

LEVELS IN ^{97}Tc , ^{99}Tc , ^{101}Tc and ^{101}Ru

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by

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SCOPE AND CONTENTS: Levels in ^{97}Tc , ^{99}Tc , ^{101}Tc and ^{101}Ru have been studied using the electron-capture decay of ^{97}Ru and the beta decays of ^{99}Mo , ^{101}Mo and ^{101}Tc , respectively. The energies and intensities of the gamma rays involved in these decays were measured with high precision by Ge(Li) detectors. Gamma-gamma coincidence experiments with Ge(Li)-Ge(Li) and Ge(Li)-NaI(Tl) spectrometers have been used to establish 9 levels in ^{97}Tc , 14 levels in ^{99}Tc , 45 levels in ^{101}Tc and 11 levels in ^{101}Ru . The low energy conversion electron spectrum from the decay of ^{101}Mo was investigated with the Chalk River $\pi\sqrt{2}$ spectrometer. From these data, energies, spins and parities for states in ^{97}Tc , ^{99}Tc , ^{101}Tc and ^{101}Ru have been established and interpreted in terms of the shell model.

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TABLE OF CONTENTS

	Page
CHAPTER I	1
NUCLEAR MODELS AND RADIOACTIVE DECAY	1
Introduction	1
1.1	1
Shell Model	1
1.2	3
Extreme Single-Particle Model	3
1.3	5
Single-Particle Model	5
1.4	7
Collective Motion	7
1.5	15
Core-Excitation Model	15
1.6	15
Beta Decay	15
1.7	20
Gamma-Ray Emission	20
1.8	22
Internal Conversion	22
CHAPTER II	25
EXPERIMENTAL TECHNIQUES	25
2.1	25
Gamma-Ray Spectrometers	25
2.2	33
The Ge(Li) Detector	33
2.3	36
Signal Processing	36
2.4	44
Source Preparation	44
2.5	50
Intensity Measurements	50
2.6	59
Energy Measurements	59
2.7	68
Peak-Fitting Procedure	68
2.8	71
Internal Conversion Measurements	71
2.9	80
Gamma-Gamma Coincidence Measurements	80
2.9.1	80
Introduction	80
2.9.2	82
Time Resolution	82
2.9.3	89
Gamma-Gamma Coincidence Measurements	89
2.9.4	105
Interpretation of the Coincidence Data	105

	Page
CHAPTER III	EXPERIMENTAL RESULTS
3.1	Levels in ^{97}Tc Populated in the Decay of ^{97}Ru
3.1.1	Introduction
3.1.2	Gamma-Ray Singles Measurements
3.1.3	Gamma-Gamma Coincidence Results
3.1.4	The Level Scheme of ^{97}Tc
3.1.5	Subsequent Investigations
3.2	Levels in ^{99}Tc Populated in the Decay of ^{99}Mo
3.2.1	Introduction
3.2.2	Gamma-Ray Singles Measurements
3.2.3	Gamma-Gamma Coincidence Results
3.2.4	Discussion of Transition Intensities
3.2.5	The Level Scheme of ^{99}Tc
3.2.6	Subsequent Investigations
3.3	Levels in ^{101}Tc Populated in the Decay of ^{101}Mo
3.3.1	Introduction
3.3.2	Gamma-Ray Singles Measurements
3.3.3	Internal Conversion Measurements
3.3.4	Gamma-Gamma Coincidence Results
3.3.5	The Level Scheme of ^{101}Tc
3.4	The Decay of 14.0 min. ^{101}Tc to Levels in ^{101}Ru
3.4.1	Introduction
3.4.2	Singles and Gamma-Gamma Coincidence Results for ^{101}Tc
3.4.3	The Level Scheme of ^{101}Ru
3.4.4	Subsequent Investigations

	Page
CHAPTER IV DISCUSSION	
4.1 Introduction	228
4.2 Discussion of the Levels in ^{101}Ru	230
4.3 Interpretation of the Technetium Isotopes ^{97}Tc , ^{99}Tc and ^{101}Tc	243
REFERENCES	267
APPENDIX A	A1
A.1 Jagspot	A1
APPENDIX B	A6
B.1 The Nucleus $^{88}_{38}\text{Sr}_{50}$.	A6
B.2 The Zirconium Isotopes	A7
B.3 The Molybdenum Isotopes	A15
B.4 The Niobium Isotopes	A25

LIST OF TABLES

Table No.		Page
2.1	List of gamma-ray energy standards	60
2.2	Example of Jagspot energy calibration	64
2.3	Explanation of Code used in Figs. 3.5 and 3.6	92
3.1	Transitions in ^{97}Tc	114
3.2	Assignment of $\log f_0 t$ and J^π values to levels in ^{97}Tc	120
3.3	Transitions in ^{99}Tc	141
3.4	Multipole Assignments in ^{99}Tc	144
3.5	Assignment of $\log f_0 t$ and J^π values to levels in ^{99}Tc	149
3.6	Low-lying transitions in ^{101}Tc	172
3.7	Transitions in ^{101}Tc from the decay of ^{101}Mo	193
3.8	The β -decay branching ratios, $\log f_0 t$ and estimated J^π values for the decay of ^{101}Mo	201
3.9	Transitions in ^{101}Ru from the decay of ^{101}Tc	218
3.10	the β -decay branching ratios, $\log f_0 t$ and J^π values for the decay of ^{101}Tc	222
4.1	Negative parity states in the technetium isotopes	251
A.1	Comparison of centers of gravity and emptiness factors for zirconium and molybdenum nuclei	A22

LIST OF FIGURES

Figure No.		Page
Fig. 1.1	The pair population v_v^2 for the single-particle levels ϵ_v (a) when no pairing force is present (i.e. $\Delta=0$) (b) with pairing included.	10
Fig. 2.1	A comparison of the resolution (FWHM)-energy relationships for several nuclear spectrometers	31
Fig. 2.2	A comparison of the practical efficiencies of various types of gamma-ray spectrometers	32
Fig. 2.3	Example of Jagspot fits to spectral peaks.	72
Fig. 2.4	The effect of Walk on time resolution.	85
Fig. 2.5	Block diagram of coincidence circuit for ^{97}Tc and ^{99}Tc experiments.	91
Fig. 2.6	Block diagram of coincidence circuit for ^{101}Tc and ^{101}Ru experiments.	97
Fig. 2.7	Coincidence projection for the 36 cm ³ detector showing the gates used in the analysis	102
Fig. 3.1	The low-energy gamma-ray spectrum of ^{97}Ru	112
Fig. 3.2	The gamma-ray spectrum of ^{97}Ru recorded with a 12 cm ³ Ge(Li) detector shielded with a graded absorber.	113
Fig. 3.3	The Ge(Li) spectrum in coincidence with a gate set on a 216 keV photopeak in the NaI(Tl) detector.	116
Fig. 3.4	The Ge(Li) spectrum in coincidence with a gate set on a 324 keV photopeak in the NaI(Tl) detector.	117
Fig. 3.5	The level scheme of ^{97}Tc	118

		Page
Fig. 3.6	The decay scheme of ^{99}Mo prior to the present investigation.	128
Fig. 3.7	The low-energy portion of the gamma-ray spectrum of ^{99}Mo .	130
Fig. 3.8	The gamma-ray spectrum of ^{99}Mo above 170 keV.	131
Fig. 3.9	A detailed presentation of the weaker peaks of Fig. 3.8.	133
Fig. 3.10	The Ge(Li) spectra in coincidence with gates set on the NaI(Tl) photopeaks of the 140 and 181 keV gamma rays.	135
Fig. 3.11	The Ge(Li) spectra in coincidence with NaI(Tl) gates set from 335-367 keV, 368-401 keV and 402-436 keV.	136
Fig. 3.12	The Ge(Li) spectra in coincidence with NaI(Tl) gates set from 437-469 keV, 470-503 keV and 619-650 keV.	138
Fig. 3.13	The level scheme of ^{99}Tc .	148
Fig. 3.14	The decay scheme of ^{101}Tc prior to the present work.	158
Fig. 3.15	The low energy portion of the gamma-ray spectrum of ^{101}Mo and ^{101}Tc taken with an 0.8 cm^3 Ge(Li) detector.	161
Fig. 3.16	The gamma-ray spectrum of ^{101}Mo and ^{101}Tc from 150 to 1360 keV recorded with a 12 cm^3 Ge(Li) detector.	162
Fig. 3.17	The gamma-ray spectrum of ^{101}Mo above 1360 keV taken with a 12 cm^3 Ge(Li) detector.	163

	Page
Fig. 3.18	165
Upper half: Expanded portions of the spectrum of Fig. 3.16.	
Lower half: Portions of the spectrum as recorded with the high resolution 0.8 cm ³ detector.	
Fig. 3.19	167
Conversion lines of the 6.281 and 9.317 keV transitions in ¹⁰¹ Tc.	
Fig. 3.20	168
Conversion lines of the 9.317, 15.606, 80.926 and 192.00 keV transitions in ¹⁰¹ Tc.	
Fig. 3.21	174
The spectrum in coincidence with the 591 keV transition.	
Fig. 3.22	175
The spectrum in coincidence with the 196 keV transition.	
Fig. 3.23	177
The spectrum in coincidence with the 500+506 keV gate. The gate also contains portions of the 510+513 keV photopeaks.	
Fig. 3.24	179
Spectra in coincidence with the 1251 and 1532 keV gates.	
Fig. 3.25	180
Spectra in coincidence with the 871+877 and 1161 keV gates.	
Fig. 3.26	181
Spectra in coincidence with the 934 keV and the 1010+1012+1019 keV gates.	
Fig. 3.27	183
The spectrum in coincidence with the 81 keV transition.	
Fig. 3.28	185
Spectra in coincidence with the 409 and 1346 keV gamma rays.	

		Page
Fig. 3.29	Coincidence spectra associated with the 212, 1184 and 1186 keV transitions.	186
Fig. 3.30	Spectra in coincidence with the 222, 1431+1433, and 1653+1662 keV gates.	188
Fig. 3.31	Spectra in coincidence with the 449 and 1064+1066 keV gates.	190
Fig. 3.32	Spectra in coincidence with the 1514+1518, 1599 and 1304 keV gates.	191
Fig. 3.33	Level scheme of ^{101}Tc showing the properties of the levels below 1775 keV.	198
Fig. 3.34	Level scheme of ^{101}Tc showing the properties of the levels between 1806 and 2418 keV.	199
Fig. 3.35	The gamma-ray spectrum of ^{101}Tc taken with the 12 cm ³ detector.	214
Fig. 3.36	Spectra in coincidence with the 127 and 180+184 keV gamma rays from the decay of ^{101}Tc .	215
Fig. 3.37	Spectra in coincidence with gates containing the 234, 238, 295, 307+311, and 531 keV transitions from the decay of ^{101}Tc .	216
Fig. 3.38	The level scheme of ^{101}Ru from the decay of ^{101}Tc .	221
Fig. 4.1	The low-lying states in the odd-mass technetium isotopes.	245

CHAPTER I

NUCLEAR MODELS AND RADIOACTIVE DECAY

Introduction

The fundamental problem of nuclear physics is the development of a general theory which will provide an explanation for observed nuclear properties. Ideally, the total nuclear wavefunction could be derived from a direct determination of the components of the nucleon-nucleon interaction. Although considerable experimental effort has been expended in trying to obtain the exact form of the nucleon-nucleon interaction, the results have been inconclusive. Instead, a number of models have been developed which attempt to approximate the real interaction. A successful model should be able to predict the properties that identify a nuclear state: energy, spin, parity, electric and magnetic moments. In addition, it should predict the transition rates between nuclear states. At the present time, there exists no single model which is able to account for all the observed nuclear properties. However, some models have provided limited descriptions of the observed phenomena for selected nuclei. The first half of the chapter discusses some of the features of these models. The latter half considers some of the aspects of radioactive decay.

1.1 Shell Model

The shell model is based on the assumption that each nucleon in the nucleus moves as an independent particle in

an orbit determined by the static central potential $V(r)$ created by the averaged interaction between the other nucleons. This potential is expected to have a shape that is similar to the charge density, but with a rather larger range. The orbital angular momentum ℓ of each nucleon is a constant of motion for a central potential. For each value of ℓ , there is a series of energy levels which can be defined by the quantum number n , the number of nodes of the radial wavefunction.

Experimentally, it has been found that there are large gaps in the spacing of energy levels with corresponding discontinuities in other nuclear properties, such as binding energy, neutron capture cross section, and angular momentum. In terms of the shell model, these gaps have been correlated with shell closures which occur when the number of protons or neutrons in the nucleus corresponds to one of the so-called "magic numbers": 2, 8, 20, 50, 82 and 126. The positions of these magic numbers cannot be predicted by using only the central potential. By using this potential and assuming a strong coupling between the spin of a nucleon and its orbital angular momentum, Mayer (1949, 1950) and Haxel et al. (1949, 1950) independently arrived at a natural explanation of the observed magic numbers. The complete potential then takes the form

$$V = V(r) - F(r) \vec{\ell} \cdot \vec{s}$$

The spin-orbit term has the effect of splitting a single $n\ell$ state into two states with $j=\ell-1/2$ and $j=\ell+1/2$; these

correspond to the two orientations of the individual particle spin with respect to ℓ . The $\ell-1/2$ states are raised from their original position, the $\ell+1/2$ states are lowered by the same amount. It should be noted that the strength of the force is sufficient to remove the degeneracy, but not large enough to destroy the good angular momentum character of the orbital angular momentum ℓ and the total angular momentum j . Since the magnitude of the splitting increases with the value of ℓ , levels with high angular momentum are lowered close to the next shell. By properly selecting the value of the interaction strength, major shell closures can be made to occur for nucleon numbers which correspond to the magic numbers. The level ordering of the first 126 nucleons in the simple shell model is given by Preston (1962). Each level is labelled by $n\ell j$ and the parity is given by $(-1)^\ell$. The Pauli exclusion principle restricts the number of neutrons or protons in each $n\ell j$ level to $2j+1$.

1.2 Extreme Single-Particle Model

The extreme single-particle model is the simplest form of the shell model. It is assumed that the neutron and proton states fill independently and that the nucleons in the ground state have dynamically paired motions (each pair couples to zero angular momentum) so that the nuclear properties are attributed to the last unpaired nucleon, whenever it is present. This model predicts the experimentally observed fact that all even-even nuclides have 0^+ ground states. The spin and parity

of the ground state of odd-A nuclei will be given by that of the last unpaired particle. The angular momentum of odd-odd nuclei cannot be precisely predicted by the model.

In the extreme single-particle model, low-lying excited states are created by exciting one or more nucleons into levels that are unoccupied. For an even-even nucleus, the simplest type of intrinsic excitation is the so-called particle-hole excitation formed by the promotion of a single nucleon into a single shell. In addition to one-particle one-hole excitations (1p-1h), two-particle two-hole excitations (2p-2h) can be created by raising a pair of particles into a higher-lying state. The lowest-lying excited states of odd-A nuclei are formed by excitation of the unpaired particle into unoccupied orbitals. The spin and parity is that of the unpaired particle. In addition, particles lying below the Fermi surface may be promoted to the surface to pair-off with the previously unpaired nucleon. A vacancy or hole is created below the Fermi surface and the J^π value of the resulting configuration is that of the unpaired hole.

The pairing energy is another phenomenon which must be considered in determining the filling of the shell-model levels. Two nucleons with the same (n, ℓ, j) quantum numbers and with their intrinsic spins aligned in opposite directions tend to interact strongly to form a "zero-coupled pair". Formation of this rather stable pair results in an increase in

nuclear binding energy over and above the "individual" binding energies. The gain in energy depends on the angular momentum, being larger for higher j values. It is found experimentally that when two shell-model states of different angular momentum are close together in energy, the higher angular momentum orbital will tend to fill in pairs before the lower angular momentum orbital has been completely filled.

1.3 Single-Particle Model

Since the extreme single-particle model ignores all nucleon correlation, it is clear that it is an oversimplified description of the nucleus. In a more realistic single-particle model, only those nucleons forming closed shells (in particular, "magic numbers") are treated as part of the inert core. The nuclear properties are then derived from a consideration of the internucleon forces between all the remaining particles outside the core. All the different states formed by k particles with the same (n, l, j) have the same energy in the extreme single-particle model. The interaction (residual interaction) in the extended model is sufficiently strong to remove the above degeneracies, but not so strong, compared with the spin-orbit force, that j ceases to be a good quantum number for each nucleon. A further assumption is that the effective interparticle forces between two particles in the same orbit are much stronger than those between nucleons in different orbits l and l' . Therefore, the angular momentum coupling within a shell, whether filled or not, is not modified by

internucleon interactions with other shells. For example, in an odd- Z , even- N nucleus, the neutrons still couple to $j=0$ and the protons to $j=j_p$. The resultant angular momentum of the whole nucleus is $J=j_p$. This explains why the simple shell model is able to predict correct angular momentum values.

In practice, it is rather difficult to determine the exact form of the residual interaction since even the interaction between two nucleons, as determined from scattering experiments and the deuteron, is not very well understood. The two-nucleon force is clearly very strong at small distances and can be expected to modify the nuclear wavefunction only where two nucleons are very close together. Shell-model wavefunctions modified by two-particle short-range correlations can be used to approximate real nuclear wavefunctions. Instead of using these modified wavefunctions with the real interaction, it is possible to use unmodified wavefunctions with interactions that are modified by the two-particle correlations. Thus, to use the model, it is necessary to consider a modified Hamiltonian which contains effective two-body interactions between nucleons. Unfortunately, this is also difficult to obtain. The effective interactions could be replaced by a phenomenological simple interaction; however, the latter must contain many adjustable parameters. An alternative procedure is to try to determine the effective interaction from the experimental data. If a definite scheme of wavefunctions, e.g. j - j coupling, is used, the energies

in the n-particle configurations can be calculated in terms of two-nucleon configurations. The experimental data are taken from a group of nuclei in which the same configurations occur. All energies can be expressed in terms of a small number of unknown parameters if it is assumed that the matrix elements of the effective nuclear interaction are the same in all nuclei of the group. A best fit to the experimental data is then used to obtain the matrix elements. These can be used in latter calculations to predict properties of other nuclei. An example is given below to illustrate the method.

The Hamiltonian for a system of nucleons outside a core nucleus (e.g. $^{90}_{40}\text{Zr}_{50}$) can be written (Bhatt and Ball 1964) as

$$H = H_{\text{core}} + \left(\sum_p H_{p\text{-core}} + \sum_n H_{n\text{-core}} \right) + \left(\sum_{\text{pairs}} H_{pp} + \sum_{\text{pairs}} H_{nn} \right) + \sum_{\text{pairs}} H_{pn}$$

The first term represents the total interaction between nucleons inside the core; the second part represents the interaction of the extra-core protons and neutrons with the core. The third and fourth terms represent the residual interactions between the extra-core proton, neutron and proton-neutron pairs. The nn, pp and pn effective interactions are determined from experimental data.

1.4 Collective Motion

The concept of pairing was introduced in a previous section. Pairing is the preferred interaction between two nucleons in similar (n, l, j) states, but with oppositely

aligned spin vectors and, as such, represents a type of collective behaviour. Since it is interactions between pairs of particles in the same state that are important, the residual force responsible for pair-correlations would appear to be of very short range. If the model is extended to include interactions between zero-coupled pairs in the same and different orbitals, a description for the nucleus is obtained in which the motion of these pairs is also correlated. This latter approach is the so-called BCS theory (Bardeen et al. 1957) which was initially introduced as a description of the phenomenon of superconductivity. Bohr et al. (1958) applied this theory to the nuclear problem. The pairing interaction results in the pairs of particles being distributed among the single (n, ℓ, j) levels of the simple shell model in a correlated fashion. These levels are no longer described as completely full or completely empty. Thus, for a given state of the nucleus, a particular (n, ℓ, j) shell model level has a probability $v_{n\ell j}^2$ of being occupied by a pair and $u_{n\ell j}^2$ of being unoccupied. In the same state, another (n, ℓ, j) level will be characterized by a different set of U^2 and V^2 values. The resulting pair state will consist of a particular distribution of pairs among the various shell-model levels.

The BCS wavefunction for the ground state of an even-even nucleus can be written as

$$|0\rangle_{\text{BCS}} = \prod_v (U_v + V_v a_v^+ a_{-v}^+) |0\rangle$$

where $|o\rangle$ represents the vacuum state, the state with no particles present

v represents a (j,m) particle

$-v$ represents a $(j,-m)$ hole

a_v^+ is a particle creation operator

a_{-v}^+ is a hole creation operator

The coefficients U and V are subject to the constraints

$$U_v^2 + V_v^2 = 1$$

and

$$\sum_v 2\Omega_v V_v^2 = N$$

where N is the total number of nucleons in the state and

$\Omega_v = j+1/2$ is the number of degenerate pairs of type j . The probabilities U_v^2 and V_v^2 are given by

$$U_v^2 = \frac{1}{2} \left| 1 + \frac{\epsilon_v - \lambda}{\sqrt{(\epsilon_v - \lambda)^2 + \Delta^2}} \right|$$

$$V_v^2 = \frac{1}{2} \left| 1 - \frac{\epsilon_v - \lambda}{\sqrt{(\epsilon_v - \lambda)^2 + \Delta^2}} \right|$$

where ϵ_v is the single-particle energy ignoring the pairing forces and λ is the average Fermi level. The gap parameter $\Delta = G \sum_v \Omega_v U_v V_v$ provides a measure of the strength of the pairing (G being a strength parameter).

In the absence of pairing, all states below the Fermi level λ are completely occupied while those states above the Fermi level are completely vacant. This situation is shown in Fig. 1.1(a). The addition of the pairing force causes a

"smearing" of the Fermi level as shown in Fig.1.1(b).

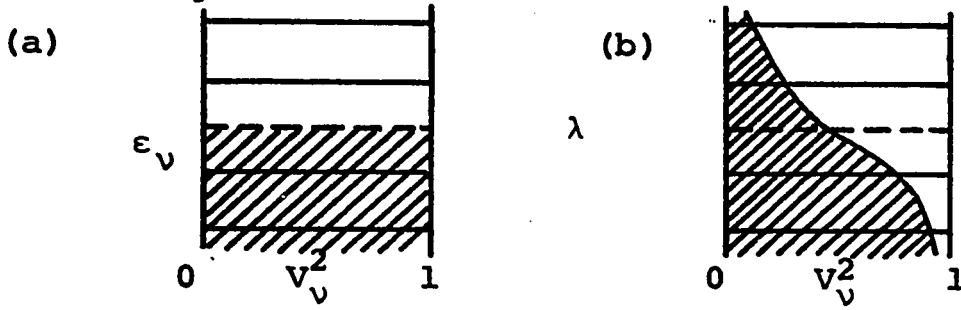


Fig. 1.1 The pair population v_v^2 for the single-particle levels ϵ_v (a) when no pairing force is present (i.e. $\Delta=0$) (b) with pairing included.

Although the high and low energy states are still almost completely empty and filled respectively, the states near the Fermi surface are only partially occupied. The inflection point of the level population curve is determined by the value of λ while the diffuseness of the Fermi surface is measured by the parameter Δ .

The description of excited states in the pairing formalism is usually done in terms of quasi-particles which have properties of both particles and holes. The quasi-particle creation and annihilation operators can be written in terms of the particle creation and annihilation operators (Bogoliubov (1958), Valatin (1958)).

$$\alpha_v^+ = U_v a_v - V_v a_{-v}^+$$

$$\alpha_{-v} = U_v a_{-v} + V_v a_v^+$$

The even-even system is transformed into its neighbouring odd nucleus by the application of one of the above operators.

The energy of the single quasi-particle state is given by

$$E_v = \sqrt{(\epsilon_v - \lambda)^2 + \Delta^2}$$

The energy difference between one-quasi-particle excitations is simply $E_v - E_v$. This leads to a compression of the low-energy spectrum relative to the single-particle picture. For an even-even nucleus, the lowest excited states are the two-quasi-particle states $\alpha_v^+ \alpha_v^+ |\tilde{0}\rangle$ ($|\tilde{0}\rangle$ being the quasi-particle vacuum). Since $|\epsilon_v - \lambda|$ is usually $\ll \Delta$ for the lowest odd state, these two-quasi-particle states correspond to an energy $(E_v + E_v) \geq 2\Delta$. This provides the energy gap observed in all even-even nuclei.

Other types of collective motion, which involve all the nucleons as a whole, may exist. For example, the interaction of a quasi-particle, which has undergone a slight perturbation from its orbit, with other quasi-particles may produce a coherent oscillation. The corresponding vibrational states may be considered as due to combinations of several quasi-particle excitations.

In the vibrational model, the nuclear core is treated as an incompressible fluid. The vibrational excitation modes are characterized by angular momentum λ and z-component μ . The surface of the vibrating nucleus can be described in terms of the coordinates of vibration $\alpha_{\lambda\mu}$ (Bohr 1952).

$$R = R_0 \left[1 + \sum_{\lambda} \sum_{\mu=-\lambda}^{\lambda} \alpha_{\lambda\mu} Y_{\lambda}^{\mu}(\theta, \phi) \right]$$

where R_0 is the equilibrium radius, Y_{λ}^{μ} is the usual spherical harmonic function and (θ, ϕ) are polar angles with respect to an arbitrary axis. Any collective motions are expressed by letting $\alpha_{\lambda\mu}$ vary in time.

The Hamiltonian for nuclear oscillations which are harmonic and of small amplitude is given by

$$H_{\text{vib}} = \sum_{\lambda} \sum_{\mu=-\lambda}^{\lambda} [1/2 \beta_{\lambda} |\dot{\alpha}_{\lambda\mu}|^2 + 1/2 C_{\lambda} |\alpha_{\lambda\mu}|^2]$$

The terms represent the kinetic and potential energies, respectively. The vibrational motion is that of a surface wave travelling around the nucleus and β_{λ} is associated with the mass transported by this wave. In order to calculate β_{λ} and C_{λ} , it is necessary to make definite assumptions about the nature of the nucleon motions. Each mode of vibration λ has associated with it a frequency

$$\omega_{\lambda} = \sqrt{\frac{C_{\lambda}}{\beta_{\lambda}}}$$

The vibrational spectrum will therefore be that of a harmonic oscillator and the excitation energies will be multiples of $\hbar \omega_{\lambda}$.

$$E \sim \sum_{\lambda} n_{\lambda} \hbar \omega_{\lambda}$$

where n_{λ} refers to the number of phonons of order λ in the excited state. These phonons carry an angular momentum λ and parity $(-1)^{\lambda}$. Terms with $\lambda=0$ describe density oscillations in a spherical nucleus; if these occur, they will be at much higher energies than the incompressible vibrations. Terms with $\lambda=1$ describe vibrations of the mass center of the nucleus and are not internal degrees of freedom. The vibrations found experimentally are the quadrupole ($\lambda=2$) and octupole ($\lambda=3$) terms. Coupling of these phonons with the ground-state

of an even-even nucleus will produce the states given below.

λ	n_λ	J^π	Energy of excited state
2	1	0^+	$\hbar\omega_2$
	2	$0^+, 2^+, 4^+$	$2\hbar\omega_2, 2\hbar\omega_2, 2\hbar\omega_2$
3	1	3^-	$\hbar\omega_3 \approx 2\hbar\omega_2$

The multiplets are degenerate for an ideal harmonic oscillator; however, the degeneracy is removed in any real nucleus and the ordering of the $0^+, 2^+, 4^+$ triplet can vary.

If the single-particle excitation energies of the odd-particle in an odd-A nucleus are fairly large, states may exist which consist of vibrational excitations of the even-even core coupled to the single-particle level. The vector coupling of the particle state of angular momentum j and the vibrational state of angular momentum R_λ gives rise to an approximately degenerate multiplet of $(2j+1)$ or $(2R_\lambda+1)$ states (whichever is smaller) having angular momentum values $J=(R_\lambda+j), (R_\lambda+j-1), \dots, |R_\lambda-j|$. The energy of the multiplet relative to the intrinsic state will be approximately equal to the phonon energy for neighbouring even-even nuclei.

Kisslinger and Sorensen (1963) have attempted to calculate the low-lying states of spherical nuclei by assuming that these states can be treated in terms of the two basic excitations discussed above, quasi-particles and phonons. For even-even nuclei, the lowest excitations are the phonons and only those were considered. For the odd-mass nuclei, both of these modes of excitation, as well as their interactions, were considered.

The particles were assumed to interact with a two-body force represented by two components, the pairing force (Bohr et al. 1958) and a long-range part represented by a quadrupole force (Elliot 1958). The pairing force was assumed to exist only for protons and neutrons separately, since the neutrons and protons are filling different shell-model states. The quadrupole force was effective for protons, neutrons, and for proton-neutron pairs. This model had moderate success in predicting nuclear properties over a wide range of mass numbers.

Other nuclear models have been proposed to interpret nuclear structure. These include the rotational model in which the nucleus is assumed to have a large stable deformation in order to account for the large ground-state quadrupole moments. This model predicts that the energy spectra of even-even nuclei can be represented by the expression

$$E_J \sim \frac{\hbar^2}{2\mathcal{I}} (J(J+1)) \quad J=0,2,4$$

where $\mathcal{I}$ is the moment of inertia. At higher excitations the vibrations described previously become energetically possible.

In odd-A nuclei, rotations and vibrations are coupled to the various single-particle excitations. These single-particles are generally described (Nilsson 1955) as moving in a non-spherical deformed potential. The Nilsson model has been very successful in predicting the properties of rotational nuclei. Since this model is not particularly

applicable to the mass region of interest in this thesis no further discussion is included here.

1.5 Core-Excitation Model

The core-excitation model was proposed by De-Shalit (1961) to account for the low-energy properties of some odd-mass nuclei. An odd-mass nucleus can be considered as an odd nucleon added to an even-even core. Normally, the lowest-lying states of the odd-mass nucleus are formed by the excitation of the unpaired particle into an unoccupied orbital. In some cases, the energy of the next single-particle state is high compared to the lowest energy of the even-even core. It is reasonable to assume that the lowest excitations of the odd-mass nucleus will be described by the odd nucleon staying in its lowest orbit (spin j) and the core excited to its first excited state (spin J_c). That is, there is a collective excitation of the core coupled to the odd nucleon. In the absence of any interaction between the core and the odd particle, all the states (the "core multiplet") with the same J_c and j , but with different values of J (and m) are degenerate. A particle-core interaction will split the degeneracy within the core-multiplet. The "center of gravity" of each multiplet coincides with the "unperturbed position". $M1$ transitions from a multiplet constructed on a 2^+ state to the ground state are forbidden; however, $M1$ transitions within the multiplet are not.

1.6 Beta Decay

The weak interaction is one of the basic forces of

nature. Its properties are revealed experimentally by the process of beta decay which may take several forms depending on whether the decaying nucleus is neutron or proton rich. Neutron-rich nuclei decay by changing a neutron into a proton in the nucleus and emitting an electron and an anti-neutrino. It is found experimentally that the electrons emitted in beta decay have a continuous energy distribution from zero to some maximum value (called the beta end-point). Pauli (1931) hypothesized the three-body character (nucleus, electron, neutrino) of the beta-decay process in order to account for this continuous energy distribution, and to satisfy angular momentum (spin) and statistics requirements. The third particle, the neutrino, was assumed to have zero charge, zero mass and intrinsic spin $1/2\hbar$. Its existence was verified through the inverse neutrino-induced reaction by Reines et al. (1960). A continuous distribution of electron energy results because of the statistical sharing of the available decay energy between the electron and the anti-neutrino. Proton-rich nuclei may decay by two competing processes: positron decay and orbital electron capture. The first of these is similar to negatron decay and involves changing a proton into a neutron with the emission of a positron and a neutrino. Emission of the lepton pair results in a continuous distribution of positron energies up to the end-point energy. The second process involves the capture of an atomic electron (K,L,M, etc.) and

the subsequent transformation of a proton into a neutron with the accompanying emission of a monoenergetic neutrino. In the case of negatron decay, a measurement of the end-point energy, the energy for which all of the decay energy is imparted to the electron, determines the total energy difference between the initial and final nuclear configuration. For positron decay, the end-point energy must be increased by $2m_0c^2$ to yield the same quantity.

The first adequate mathematical description of the beta-decay process was provided by Fermi(1934). From perturbation theory, the transition probability for beta decay is given by

$$\frac{2\pi}{\hbar} |\langle \psi_f \phi_f | H_\beta | \psi_i \phi_i \rangle|^2 \rho_f(E) \quad 1.1$$

where ψ_i and ψ_f represent initial and final nuclear states, ϕ_i and ϕ_f the initial and final lepton states, H_β the weak interaction Hamiltonian and $\rho_f(E)$ is the density of final states. Under the conditions of special relativity and momentum and angular-momentum conservation, the theory of beta decay must be invariant under Lorentz transformations. Therefore, the form of H_β is restricted to a combination of five possible forms: the scalar(S), pseudoscalar(P), vector(V), axial vector (A) and tensor(T) interactions. Experiments involving the lepton helicities (Wu et al. 1957) have shown that only the V(singlet) and A(triplet) interactions are non-zero. These interactions result in the Fermi and Gamow-

Teller selection rules, respectively. The spin change ΔJ and parity change $\Delta\pi$ for transitions between the initial state (J_i, π_i) and the final state (J_f, π_f) follow the selection rules given by

$$|J_i - J_f| \leq \Delta J \leq (J_i + J_f)$$

$$\Delta\pi = \pi_i \pi_f$$

The table given below lists the selection rules for spin and parity

<u>Transition</u>	<u>Fermi (Vector)</u>	<u>Gamow-Teller (Axial Vector)</u>
allowed	$\Delta J=0 \quad \Delta\pi=\text{No}$	$\Delta J=0, \pm 1 \quad \Delta\pi=\text{yes}$ No $0 \rightarrow 0$
first forbidden	$\Delta J=0, \pm 1 \quad \Delta\pi=\text{yes}$ No $0 \rightarrow 0$	$\Delta J=0, \pm 1 \quad \Delta\pi=\text{yes}$ No $\left\{ \begin{array}{l} 0 \rightarrow 0 \\ 0 \rightarrow 1 \\ 1/2 \rightarrow 1/2 \end{array} \right.$
first forbidden unique	$\Delta J=\pm 2 \quad \Delta\pi=\text{yes}$	$\Delta J=\pm 2 \quad \Delta\pi=\text{yes}$

If the lepton wavefunctions of Eq. 1.1 are approximated by plane waves, the product of these plane waves can be expanded as a series of terms of diminishing magnitude. The leading term in the expansion yields the term for allowed decay. If the initial and final nuclear states do not satisfy the spin and parity requirements for allowed decay, additional terms in the series must be included until the spin and parity

requirements are satisfied and a non-zero transition probability occurs. The number of vanishing terms specifies the degree of forbiddenness. For allowed transitions, the density of momentum states can be evaluated to give the beta momentum distribution

$$N(p) dp = C |M_{if}|^2 F(Z, E) p^2 (E_0 - E)^2 S_n dp \quad 1.2$$

where p is the electron momentum, $N(p)dp$ is the probability that the β -particles are emitted between p and $p+dp$, C is a constant, M_{if} is the nuclear transition matrix element, $F(Z, E)$ is the electron screening function (Fermi function) which corrects for the Coulomb interaction between the emitted electron and the nucleus, Z is the nuclear charge, E is the electron energy and E_0 is the beta end-point. S_n is a shape factor which is unity for allowed spectra and which corrects the expression when higher degrees of forbiddenness are required (n indicates the degree of forbiddenness).

The total transition probability λ for beta decay is obtained by integrating 1.2 over all possible momenta

$$\lambda = \frac{\ln 2}{T_{1/2}} = C |M_{if}|^2 f(Z, E_0) \quad 1.3$$

where $f = \int F(Z, E) p^2 (E_0 - E)^2 S_n dp$ and $T_{1/2}$ is the half-life for the particular decay mode. The integral can be calculated from the charge of the nucleus and the beta end-point.

Equation 1.3 can be rewritten as

$$f(Z, E_0) T_{1/2} = \frac{C'}{|M_{if}|^2} \quad C' \text{ a constant}$$

The quantity " ft ", called the comparative half-life, is inversely proportional to the square of the nuclear matrix element and should be an indicator of the type of beta transition. In practice, since S_n is not uniquely defined for most forbidden decays, it is taken as unity and values of $f_0 t$ are usually calculated. For the same beta end-point, a factor of ~ 100 in $|M_{if}|^2$ is introduced by each degree of forbiddenness. Therefore, it is more usual to compare $\log_{10} f_0 t$ values; the following table shows typical values which give a rough indication of the type of transition involved.

<u>Type</u>	ΔJ	$\Delta \pi$	$\log_{10} f_0 t$ (Avg. Value for Mass 100)
Super Allowed	0	no	3-4
Allowed	0, ± 1	no	4-7 (~ 5.6)
First Forbidden Non-Unique	0, ± 1	yes	6.5-9.5 (~ 7.9)
First Forbidden Unique	± 2	yes	8.5-11
Second Forbidden	$\pm 2, \pm 3$	no	11-13

It should be pointed out that these are very crude divisions and that $\log f_0 t$ values for a given type increase gradually with A .

1.7 Gamma-Ray Emission

Since beta decay does not always proceed from the ground state of the initial nucleus to the ground state of the final nucleus, the final nucleus may be left in an excited configuration. For excited states that are stable against nucleon emission, the dominant decay mode is the

emission of electromagnetic radiation. The decay of an excited state involves the changing of one nuclear configuration to another, this results in a rearrangement of one or more nucleons. Since these nucleons have charges and magnetic moments, this rearrangement is accompanied by a perturbation in the electromagnetic field surrounding the nucleus with the subsequent emission of a gamma ray. The energy associated with the gamma ray is just the energy difference between the two nuclear configurations.

The probability of gamma-ray emission can be derived from first-order time-dependent perturbation theory and is given by

$$T = \frac{2\pi}{\hbar} |\langle \psi_f^* | H_{int} | \psi_{in} \rangle|^2 \rho_f(E)$$

where ψ_i and ψ_f are the initial and final state wavefunctions, H_{int} is the interaction between the nucleus and the surrounding electromagnetic field and $\rho_f(E)$ is the density of final states. Usually H_{int} is expanded in a series of multipole operators $\phi_{\lambda\mu}^\sigma$. The transition probability can then be written as

$$T(\sigma\lambda\mu) = \frac{2\pi}{\hbar} |\langle \psi_f^* | \phi_{\lambda\mu}^\sigma | \psi_{in} \rangle|^2$$

where $T(\sigma\lambda\mu)$ represents the transition associated with the emission of a photon with angular momentum λ and projection μ . The operator may be classified as either electric ($\phi=E$) or magnetic ($\phi=M$). The following restrictions are imposed on the spins and parities of the initial and final states

by the form of these operators.

$$|J_i - J_f| \leq \lambda \leq (J_i + J_f)$$

$$\pi_i \pi_f \begin{cases} = (-1)^\lambda & \text{for electric transitions} \\ = (-1)^{\lambda+1} & \text{for magnetic transitions} \end{cases}$$

Transitions with multipole order zero cannot occur since the photon has one unit of angular momentum associated with it.

Exact evaluation of the transition probabilities is impossible unless the correct nuclear wavefunctions are known. However, order of magnitude estimates for the multipole matrix elements indicate that electric transitions are faster than magnetic ones of the same multipole order and that the transition probabilities decrease rapidly as the multipole order increases.

Weisskopf (1951) has calculated transition rates by assuming an extreme single-particle model. Experimental rates are often compared with the single-particle rates and are classified as hindered or enhanced. Normally E1 and M1 transitions are found to be hindered and E2 transitions enhanced. Therefore, the M1 and E2 strengths may compete favourably and the corresponding radiation will have a mixed multipole character.

1.8 Internal Conversion

A mode of decay which always competes with gamma-ray emission is internal conversion. In this process, the excitation

energy of the nucleus is transferred by the electromagnetic field to one of the atomic electrons. The electron is then ejected from the atom with an energy $E_\gamma - B_j$ (where E_γ is the transition energy and B_j is the atomic binding energy of the j^{th} shell). For any given gamma-ray transition, several competing conversion lines will appear, corresponding to internal conversion of electrons from the K, L_1 , L_2 , L_3 etc. shells. Since the electrons closest to the nucleus have the greatest probability of interacting with it, the K lines are stronger than the L lines, the L lines are stronger than the M lines and so on.

The probability for internal conversion is given by the transition rate

$$T = \frac{2\pi}{\hbar} |\langle \psi_f^* | H_{\text{int}} | \psi_i \rangle|^2 \rho(E)$$

where ψ_f and ψ_i involve products of the electronic and nuclear wavefunctions. As in the case of gamma-ray emission, H_{int} is the electromagnetic interaction which is written in terms of multipole operators. The electronic and nuclear matrix elements may be separated if it is assumed that the electronic wavefunctions do not penetrate the nucleus. The nuclear matrix elements for gamma-ray emission and internal conversion will therefore be equal. The competition between internal conversion and gamma-ray emission is measured by the internal conversion coefficient

$$\alpha = \frac{N_{e_i}}{N_\gamma}$$

where N_{e_i} is the probability of internal conversion in the i^{th} shell and N_λ is the probability of gamma-ray emission. Due to the equality of the nuclear matrix elements, the conversion coefficients depend only on the electronic wavefunctions and the electromagnetic multipole interaction. Their values can be theoretically predicted (e.g. Hager and Seltzer 1968). There is some penetration of the nucleus by the electrons due to the finite size of the nucleus. Thus, complete cancellation of the gamma-ray and internal conversion matrix elements does not in fact occur. The correction for these penetration effects is usually small, and has been allowed for in recent tabulations of conversion coefficients.

Since the conversion coefficients are strongly dependent on the multipolarity of the transition, a comparison of the experimental values with the theoretical predictions allows a determination of the multipolarity of the radiation.

CHAPTER II

EXPERIMENTAL TECHNIQUES

A primary task of the nuclear spectroscopist is accurate determination of the energies and intensities of the radiations given off by a decaying nucleus. Once this has been done, other measurements can be used to establish the relationship of the various transitions in order to discover the unique level scheme of the nucleus under study. This level scheme can then be used as a starting point for further investigations which seek to derive other nuclear properties such as: spins, parities, lifetimes, magnetic and quadrupole moments and perhaps even the wave-functions of the nuclear states. This chapter discusses the various experimental techniques used in the investigations of the radioactive nuclei considered in this thesis.

2.1 Gamma-Ray Spectrometers

By the early 1950's, a number of methods and instruments had been developed for detecting the radiations released in the decay of radioactive nuclei. The most versatile of these devices was the magnetic spectrometer. At low transition energies where the internal conversion coefficient is large, magnetic spectrometers were used to analyze the K,L,M, etc. electrons emitted in the internal conversion process. In this way, the energies, the intensities, and the multi-

polarities of many low-energy transitions in a large number of nuclei were determined. Magnetic spectrometers were also used for the detection of secondary electrons from processes such as the Compton effect and the photo-electric effect (external conversion). The secondary electrons were produced by the interaction of primary gamma rays with a material external to the radioactive source. These latter techniques were employed to measure the energies and intensities of many high energy transitions, which were unsuitable for study with the internal conversion method because of their small conversion coefficients. The efficiencies of magnetic spectrometers are relatively low. The solid angle of a typical double-focussing instrument is $\sim 10^{-3}$ of a sphere and this number reduces to $\sim 10^{-7}$ (at 1 MeV) when the spectrometer is used in the external conversion mode. Typical resolutions of magnetic spectrometers suitable for routine spectrum investigation range from 0.2 to 2% in momentum. Resolutions of 1 part in 10^4 have been achieved with high resolution spectrometers such as the iron-free, double-focussing $\pi\sqrt{2}$ spectrometer at Chalk River Nuclear Laboratories. Unfortunately, the achievement of very high resolution results in considerable loss in efficiency. For example, the efficiency of the Chalk River spectrometer at a resolution of 0.01% in momentum is $\sim 10^{-5}$ (Graham et al. 1960). Moreover, these spectrometers require extremely thin sources having reasonably long half-lives (typically >1 hour). Magnetic spectrometers are

single-channel devices and therefore it is only possible to scan very small regions of the total spectrum at one time.

The curved crystal diffraction spectrometer developed by Dumond (1947) was another valuable instrument whose resolving power was quite high. The line-width (full-width half maximum) of the transmission-type diffraction spectrometer may be expressed as (Alburger 1960).

$$\Delta E = kE^2/n$$

(where $k \sim 3 \times 10^{-5}$ for the usual quartz bent-crystal spectrometer and n is the order of the reflection). For first order reflection, line widths of ~ 0.3 keV (0.3%) at 100 keV and ~ 30 keV (3%) at 1 MeV were attained. About three times better resolving power has been achieved by the use of germanium crystals instead of quartz (Seppi et al. 1962). The dependence on the square of the energy means that the resolution of the diffraction spectrometer deteriorates rapidly with energy so that the crystal spectrometer cannot be considered a high-resolution instrument for gamma rays with energies above a few hundred keV. Because of its very low efficiency ($\sim 10^{-10}$ at 1 MeV), extremely strong sources (~ 1 -10 curies) are required. In addition, the crystal spectrometer is a single-channel device and therefore, only a small fraction of the energy interval can be studied at one time. Thus, the crystal spectrometer has proved to be an instrument of limited applicability and has been used chiefly for the

determination of primary gamma-ray energy standards and for the study of relatively long-lived activities.

Some of the limitations of the above methods were reduced with the discovery of the sodium-iodide, thallium-activated (NaI(Tl)) crystal by Hofstadter (1948). The NaI(Tl) scintillator provided for the first time a means of measuring all of the components of a gamma-ray spectrum simultaneously. The relatively high efficiency of this detector made it possible to investigate more thoroughly nuclei with short half-lives and nuclei having a large number of weak gamma rays. Unfortunately, the resolving power of the NaI(Tl) detector is one or two orders of magnitude lower than that of the crystal diffraction spectrometer. It was therefore rather difficult to work with nuclei having large numbers of gamma rays closely spaced in energy.

It was clear that a gamma-ray "spectrometer-gap" existed between the very high resolution instruments whose applications were limited by very low efficiency and the very efficient scintillators whose applications were limited by poor resolution. This "gap" in gamma-ray spectroscopy was virtually eliminated by the development of the semiconductor detector.

Following the development of the lithium-drift process in silicon (Pell 1960), several attempts were made to use lithium-drifted silicon Si(Li) detectors as gamma-ray detectors. Unfortunately, the photo-electric cross section

of silicon is small and therefore the absorption cross section was dominated by the Compton effect over most of the energy range of interest to gamma-ray spectroscopists. Because of the low peak-to-compton ratios obtained (0.02 at 660 keV), these detectors were found to be useful only below 100 keV. However, the fact that the intrinsic resolution of Si(Li) was much superior to that of NaI(Tl) indicated that further investigation of semiconductor materials was justified. Since $\sigma_{\text{photo}}/\sigma_{\text{compton}}$ goes approximately as Z^4 , it was clear that detector development should concentrate on higher atomic number semiconductors in order that the detectors would have high photopeak efficiency as well as high resolution. Germanium was therefore a logical candidate for a semiconductor detector.

The first lithium-drifted germanium (Ge(Li)) detectors (Freck and Wakefield 1962, Webb and Williams 1963) were small in volume and suffered from poor photopeak efficiencies and from poor peak-to-compton ratios. Soon afterward, Tavendale and Ewan (1963) were able to fabricate detectors having sufficient volumes to be useful as practical gamma-ray spectrometers. Since then, improvements in the fabrication techniques and in the quality of the germanium have made it possible to produce large volume (up to 100 cm³) detectors having excellent energy resolution and efficiency characteristics. Smaller volume (1 or 2 cm³), higher resolution detectors with very thin entrance windows have also been developed; these

detectors are able to detect radiations having energies as low as 2 keV. It should be noted that the improvements in detector resolution were due in part to major developments in preamplifier and amplifier technology.

Fig. 2.1 shows a comparison of the resolution (fwhm)-energy relationships for several nuclear spectrometers. Excluding special instruments like the iron-free $\pi\sqrt{2}$ spectrometers and energies below 200 keV, it is clear that the Ge(Li) spectrometers give superior resolving powers.

A comparison of the practical efficiencies of various types of gamma-ray spectrometers is presented in Fig. 2.2 (Fig. 10 of Hollander 1966). The assumed solid angles are given in parenthesis. The efficiency of the Ge(Li) spectrometer is several orders of magnitude greater than the corresponding efficiencies of the external conversion or crystal diffraction spectrometers. The largest Ge(Li) detectors have efficiencies (at 1332 keV) of approximately 15% of the efficiency of a 7.6 cm $\times$ 7.6 cm NaI(Tl) detector (the source is 25 cm from each detector for the comparison) and peak-to-compton ratios of 40:1.

The above discussion has shown that the efficiency and resolution characteristics of the Ge(Li) detector make it an excellent general purpose gamma-ray spectrometer. As with the scintillation detector, the entire spectrum can be accumulated simultaneously. Therefore, the investigation of nuclei having short half-lives or large numbers of weak

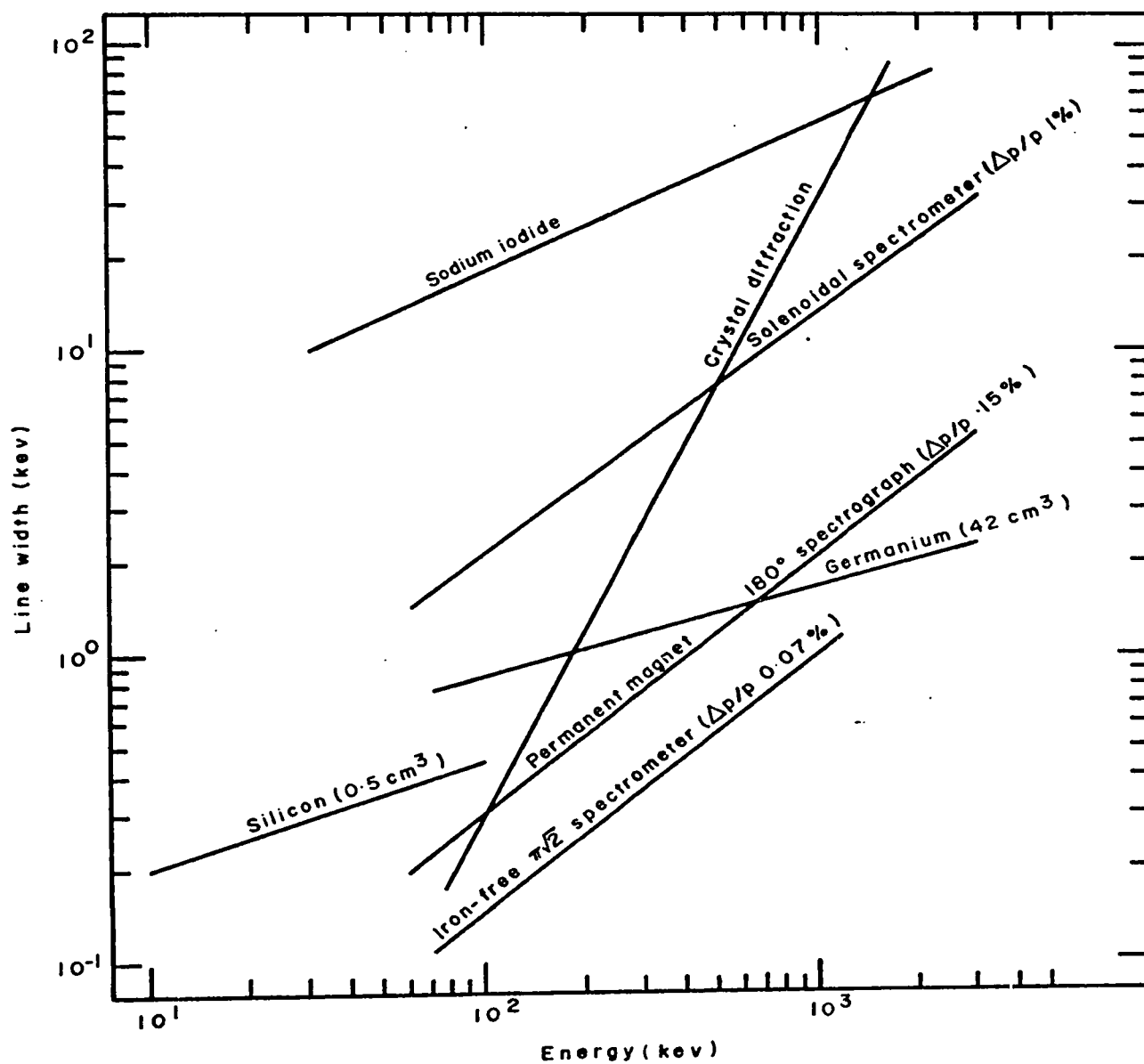


Fig. 2.1 A comparison of the resolution (FWHM)-energy relationships for several nuclear spectrometers

APPROXIMATE TOTAL FULL-ENERGY ABSORPTION EFFICIENCY
VS. ENERGY

OF SEVERAL GAMMA-RAY SPECTROMETERS

(assumed solid angles given in parentheses)

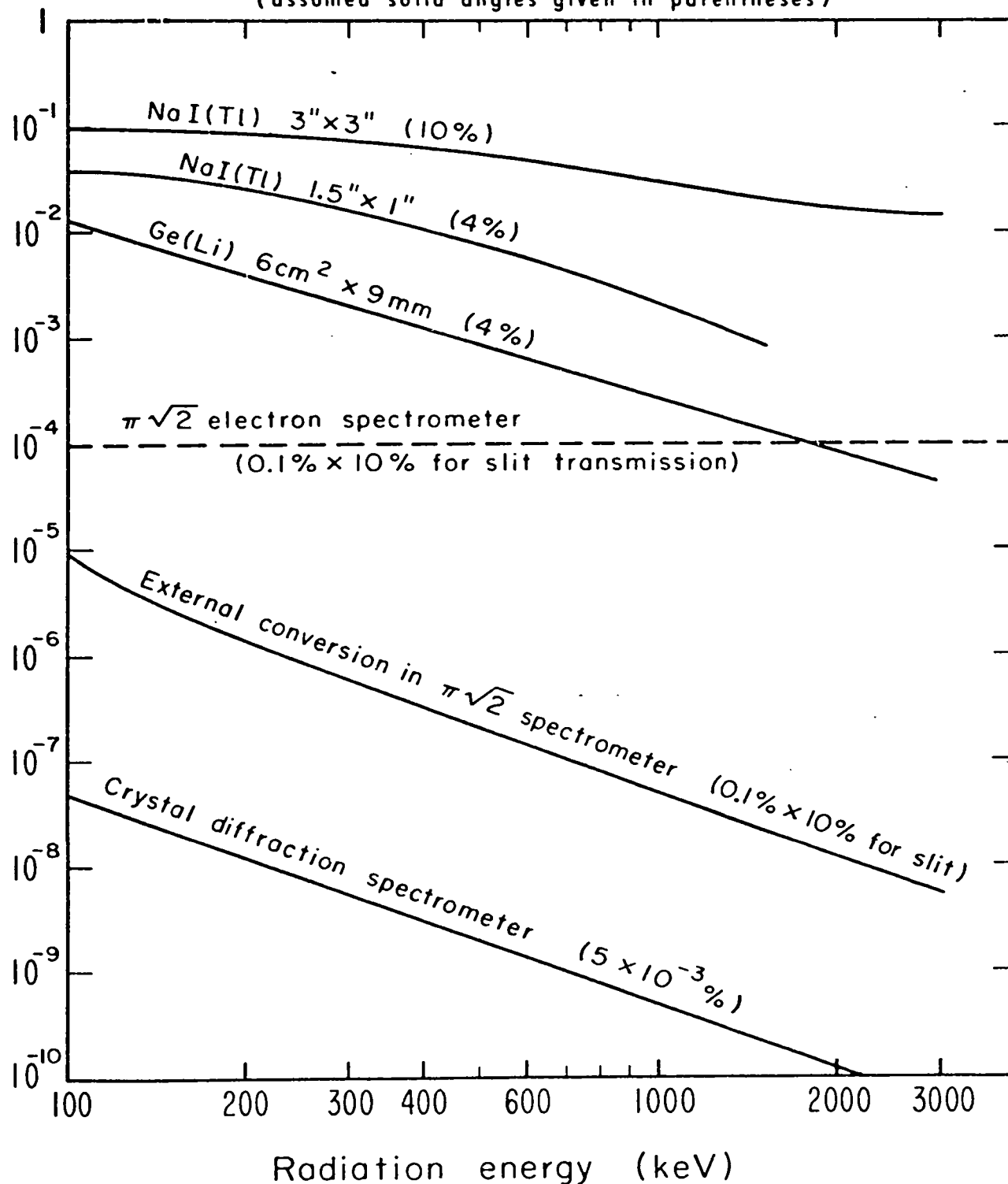


Fig. 2.2 A comparison of the practical efficiencies of various types of gamma-ray spectrometers

transitions does not constitute an insurmountable problem.

2.2 The Ge(Li) Detector

Ge(Li) detectors are normally fabricated in the following way. Donor atoms of lithium are diffused onto the front surface of a crystal of p-type germanium forming a heavily doped n^+ -type layer. Lithium ion-drift is then carried out; the lithium ions compensate the acceptor impurities in the drifted zone creating a region of high resistivity. The resulting configuration is a compensated region (sensitive region) with p-type (undrifted) material on one side and n-type (lithium) material on the other. Complete details of the fabricating procedures are readily available in the literature (see for example: Goulding 1966, Camp 1967).

A gamma ray is detected when it interacts (for example, via the photo-electric or the compton effect) with charged particles (electrons) in the compensated region of the detector and transfers energy to them. Since the probability of interaction is dependent on the volume of the compensated region, the efficiency for detection of the full energy peak will increase as the drifted depth is enlarged. The larger volume also allows more multiple interactions (such as a compton scattering event followed by a photo-electric one) to take place, with the result that there is a higher probability of absorption of the full energy of the incident gamma ray.

Fast and efficient collection of the electron-hole pairs created by the interaction of a gamma-ray with the sensitive region of the detector is influenced both by the quality of the compensated region and by the field strength applied. The field strength should be high enough so that the fraction of charge pairs collected is independent of the energy of the ionizing radiation. The quality of the sensitive region is manifested in the number of centres that either trap ion pairs or inhibit their collection. Too many trapping or recombination centres will mean that not all of the charge is collected before the pulse is clipped by the amplifier. In other words, the collection time will become too long. In practical terms, this means that the peak corresponding to full energy absorption acquires a low energy tail associated with events in which the charge was incompletely collected. Since the majority of the centres are caused by impurities or crystal imperfections, very pure samples of germanium are required for detector fabrication.

The basic limitations (aside from preamplifier noise) of the energy resolution of a Ge(Li) detector include noise sources such as surface leakage currents, thermal generation-recombination effects, which are aided by impurity centres, and the statistical fluctuations in the number of ion pairs created for a given energy E . The first two problems are alleviated by cooling the detector to liquid nitrogen temperatures (77°K) and by keeping the detector surfaces clean

by operating under high vacuum conditions (also stops condensation on the cold surfaces). The latter limit involves the Fano Factor (Fano 1947). When the energy E of a gamma-ray is completely absorbed within the compensated region, part of the energy goes into heating the crystal lattice (creating phonons) and part into the formation of ion pairs. The ratio of the variance to the yield is defined as the Fano factor. The yield is the number of electron-hole pairs produced by a given quantity of energy deposited E , while the variance (σ^2) is the mean square variation in the yield. If ϵ designates the average number of electron volts that results in the production of any ion pair in germanium, then the yield is $Y = E/\epsilon$ and the Fano factor is

$$F = \frac{\sigma^2}{(E/\epsilon)}$$

The division of energy loss between heating and ionization is essentially statistical. All of the incident energy would result in ion pair production without the competing process of lattice heating. Therefore, there would be no statistical fluctuation and thus $F = 0$. If the probability of ion-pair production is small resulting in a very large ϵ , Poisson statistics are expected to hold and $\sigma^2 \rightarrow Y, F=1$.

The average energy ϵ required to produce a hole-electron pair should be as small as possible; therefore ,

materials such as silicon (1.08 eV) and germanium (0.66 eV) with small band gaps will have good resolution characteristics. In germanium, ϵ has a value of 2.98 eV/ion pair (Camp 1967); in NaI(Tl), the value is ~ 300 eV/ion pair. This difference is responsible for the superior resolution of Ge(Li) detectors as compared to NaI(Tl). Latest estimates suggest a value of 0.15 for the Fano factor for germanium (Camp 1967).

The energy resolution obtained for a gamma ray of energy E is given by

$$R = \epsilon \sigma = \sqrt{\epsilon E F}$$

The full width at half maximum is then given by

$$\text{FWHM} = 2.35 R$$

At 1 MeV and with a Fano factor of 0.15, the best possible resolution is expected to be 1.55 keV. A comparable figure for NaI(Tl) ($F \sim 1$) is ~ 40 keV. The best detectors used in this work had resolutions of ~ 3 keV at this energy.

2.3 Signal Processing

The ultimate resolution attainable with a particular detector is determined by the detector characteristics (material, ionization potential, leakage current, Fano factor, capacitance, etc.) and by the radiation energy. Since intrinsic high resolution is an inherent property of Ge(Li) detectors, it is very important to preserve this when charge pulses formed in the detector are processed by electronic equipment such as preamplifiers and amplifiers. The nature of the final signal after passage through the electronics will depend on the shape of the input pulse,

count rate, preamplifier and amplifier noise, overall system stability, and on proper interfacing between components within the system. At low count rate and with a matched system,

$$(\text{Total Resolution})^2 = (\text{Intrinsic Resolution})^2 + (\text{Electronic Noise})^2$$

Since some noise sources are intrinsic to the electronic components being used, there is always a noise contribution which will prevent the intrinsic resolution being attained. Fortunately the effect of noise on energy resolution is a well-behaved one; noise broadens spectral peaks without changing the original peak shape. Noise is also essentially independent of count rate.

A number of techniques are employed in the design of the detector and preamplifier system in order to keep noise sources to a minimum. The first stage of the charge sensitive preamplifier consists of a very high impedance field-effect transistor (FET). These devices have excellent low noise characteristics; without them, it would be impossible to achieve the high resolutions currently available in Ge(Li) detectors. In some systems, the FET is cooled to liquid nitrogen temperatures to further reduce its noise figures. Capacitance is kept to a minimum by attaching the preamplifier directly to the cryostat and leads are kept as short as possible. Considerable improvement in resolution can be achieved by designing a preamplifier specifically for

a certain type of detector. The following example is used to illustrate this point. The resolution (W) of a preamplifier is normally stated as

$$W = W_0 + m C_p$$

where W_0 is the resolution at zero input capacitance

C_p is the capacitance in pf

m is the change in resolution per pf of input capacitance in keV/pf.

In order to obtain the highest resolution when using small volume, low capacitance detectors, it is essential to have a preamplifier that has a minimum value of W_0 . However, for larger volume detectors having considerable capacitance ($>10\text{pf}$), it is more important to design a preamplifier with a small value of m ; the value of W_0 is not as critical as above. Since the requirements of the two types of detectors are not the same, it is the usual practice to have different preamplifiers for use with them.

In addition to noise, the electronic circuitry can contribute pulse pile-up and baseline shift. Both of these factors will tend to degrade the intrinsic high resolution. Pulse pile-up and baseline shift usually distort the shape of spectral peaks; the peaks become "lopsided" or develop "skirts". These effects are closely related to count rate; in general, the higher the count rate, the poorer the spectral resolution. Pulse shape is also important with respect to

pulse pile-up. If the duration of individual pulses is too long, the probability of pile-up is much greater.

A typical preamplifier output pulse has a very fast rise time (<100 nsecs) and an exponential fall-off with a relatively long ($50 \mu\text{sec}$ - 1 msec) time constant. The rise time is determined by the detector characteristics and the exponential tail by the resistance and capacitance appearing at the detector terminals. For the best possible signal-to-noise ratio, the capacitance must be kept as low as possible and the resistance as high as possible; so high, in fact, that the tail is undesirably long from the standpoint of pile-up. Therefore, pulses pile-up on the tail of the preceeding pulses even at relatively low count rates. Normally this occurrence is not a problem, since the useful information is found only in the approximately step-function rise of each pulse. The pulse shaping circuits in the subsequent main amplifier (particularly the first differentiator), operating with time constants (~ 0.5 - $2.0 \mu\text{secs}$) much shorter than the decay time of the preamplifier signal, effectively remove the slow component of the preamplifier signal and produce individual peaks. If however, the count rate becomes very high, the amplifier will be unable to distinguish individual pulses in the output of the pre-amplifier. Portions of the first pulse will be added to the second in the amplifier. Therefore, the voltage of some of the output pulses of the amplifier will be increased. This

leads to a bulge on the high energy side of the peaks in the gamma-ray spectrum.

The differentiation network and the subsequent integration network of an amplifier also serve the purpose of enhancing the signal-to-noise ratio by removing the low and high frequency components respectively, of the noise spectrum. Unfortunately, these filtering networks are the cause of some unwanted effects on the pulse being processed. Baseline shift is caused by some RC interstage coupling network remaining which is not pole-zero cancelled (see below). This is usually a differentiation network having a time constant much longer than those of the primary pulse-shaping networks. In amplifiers without pole-zero cancellation, the dominant additional differentiation will be that of the exponential decay of the preamplifier signal. The effect of the remaining undesired differentiation in main amplifiers is an undershoot in the output pulse. If the counting rate in the amplifier is such that there is a high probability of a second pulse occurring before the undershoot of the prior pulse has fully returned to the baseline, then the apparent amplitude of the second pulse will be decreased by the amount of undershoot present at the time of the second pulse. Since this occurs in a random time distribution, the spectral peaks are smeared toward low energies.

Several improvements have been made in amplifier

design to reduce the effect of pile-up and baseline shift at high counting rates. Pole-zero cancellation, which reduces the undershoot caused by the preamplifier recovery time constant, is achieved by connecting a resistor across the capacitor in all CR differentiating circuits in the preamplifier as well as the main amplifier. The R value is such that the time constant formed by it and the differentiating capacitor is equal to the preamplifier time constant. Baseline restoration removes any residual undershoots remaining at the amplifier output by forcing the signal to return to the baseline immediately after each pulse.

The analog pulses from the amplification system associated with the detector are sent to a multichannel analyzer for further processing. The amplitude of the pulse is proportional to the energy deposited in the detector and its frequency of occurrence is related to the intensity of the radiation. The function of the analyzer is to classify the analog pulses according to their height and number.

An essential component of the analyzer is the analog-to-digital converter (ADC) which converts the analog pulse into a digital number. Present ADC's employ Wilkinson-type converters in which a capacitor is charged with the input voltage pulse. After the height of the pulse has been sensed, the capacitor is allowed to discharge linearly until a certain threshold is reached. The discharge time interval

is measured by scaling a gated high-frequency clock. The number of clock pulses are recorded by the advance channel scaler (ACS); this number is a digital representation of the amplitude of the input pulse. Since the ADC will not accept any further pulses for analysis while a pulse is being processed, the system is said to be "dead" during this time interval; a linear gate is used to block the input. The percentage dead time is a measure of the detector counting rate and can be read from a front-panel meter. The dead-time will depend on the analysis time of the converter and will vary with different ADC's.

The other major component of an analyzer is the memory unit. The number in the ACS is transferred to an address register of the memory. The contents of core memory corresponding to this address are then loaded into the memory register. The register is incremented and its contents are transferred back into core memory. Upon completion of the storage cycle, a reset signal allows the linear gate of the ADC to reopen.

Most of the gamma-ray measurements discussed in this thesis were taken with a Nuclear Data 161F 4096-channel ADC coupled with a 160M memory unit. A 16 MHz clock frequency is used by the ADC during the capacitor discharge. This is relatively "slow" by present-day standards since ADC's now use clock frequencies of up to 200 MHz .

The performance of the analyzer is chiefly determined by the characteristics and stability of the ADC. Of particular

importance is the differential non-linearity which expresses the deviation from linearity as a function of channel number. Non-linearity is discussed in section 2.6; the differential non-linearity of the present system was less than 3 channels over the region of interest. Two other important characteristics are stability over long periods of time and stability under temperature variations. The ND 161F ADC was relatively stable over the long term (2 or 3 days) if temperature variations were kept to a minimum. However, it did not perform exceptionally well under conditions of variable temperature (particularly if the temperature was increasing) and was subject to gain shifts. In fact, if the experimental room was not kept air-conditioned, the unit would cease to work at about 26°C.

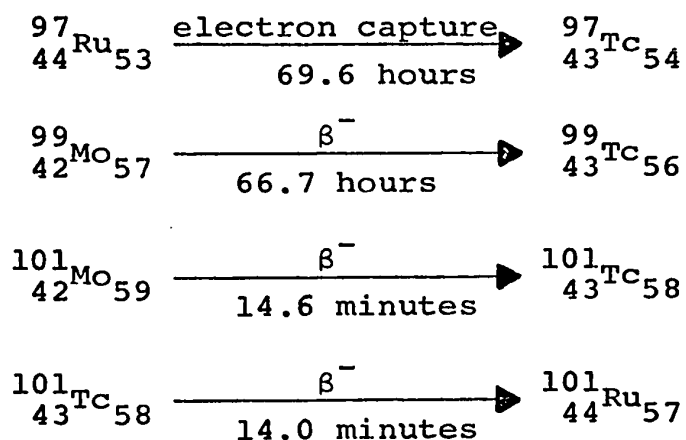
Because of these facts, a digital spectrum stabilizer (Nuclear Data Model 502) was added to the basic system to prevent gain shifts during collection of singles spectra. A precision mercury relay pulser (Canberra model 1501) was fed to the amplification system; its gain was adjusted so that the peak appeared around channel 4000 in the ADC. A window was then set around the channels containing this peak by the spectrum stabilizer. In this way, gain shifts of small amplitude were avoided.

An oscilloscope display unit was employed to monitor the memory contents. The spectrum was normally dumped onto magnetic tape after each experiment. This was accomplished through a buffer control unit which presented the memory

contents to the tape transport for transcription. If the experiment was run for an extended time period, several dumps were also made during the experiment to check for zero or gain shifts.

2.4 Source Preparation

This thesis is concerned with studying the decays of the radioactive nuclei listed below.



All of these nuclei were produced by irradiating suitable target materials in the McMaster nuclear reactor which has a flux of about 10^{13} neutrons /cm²/sec. The table given below indicates the target material, weight and container, irradiation time and the irradiation facility used for each nuclear species.

Irradiation Data

Activity	Target Material (Enrichment)	Weight	Container	Irradiation Time	Irradiation Facility
${}^{97}\text{Ru}$	${}^{96}\text{Ru}$ (94.9%)	1mg	quartz vial	10 hours	Aluminium Can
${}^{99}\text{Mo}$	${}^{98}\text{Mo}$ (98.5%)	1mg	quartz vial	4-10 hrs	Aluminium Can
${}^{101}\text{Mo}$	${}^{100}\text{Mo}$ (95.9%)	2mg	quartz vial	0.5-3 min.	Rabbit
${}^{101}\text{Tc}$	${}^{100}\text{Mo}$ (95.9%)	2mg	quartz vial	10 min.	Rabbit

Irradiations were normally carried out with the container placed in an aluminium can which was then assigned an irradiation position in the reactor core. Shorter term irradiations were carried out using the rabbit facility. This system is a series of pneumatic tubes connecting the reactor core to several laboratories which are external to the reactor building. Use of the rabbit system allowed rapid delivery of the radioactive sample to the experimental counting area. All of the irradiations were performed with 1 or 2 mgs. of target material sealed in a quartz vial. Since considerable activity was produced in the quartz by the neutron irradiation, the active sample material had to be removed from the quartz vials for counting. The main advantage of the quartz vials was that they could be broken open quickly without the sample material being contaminated by broken pieces of the container. In most cases, the source was prepared by depositing the active sample material in gelatine capsules which were then placed in a standard geometry for counting purposes. These capsules and the tape used to seal them were made of thin material with a low atomic number ($Z=6$) and therefore caused little gamma-ray absorption.

The following paragraphs discuss the various methods used to identify the transitions observed in the investigations of the four nuclei under study in this thesis.

Sources of 2.9 day ^{97}Ru were produced by thermal neutron capture in enriched ruthenium metal. In addition to the desired ^{97}Ru activity, neutron irradiation produced

several other activities: ^{103}Ru , ^{105}Ru and ^{96}Tc . The ^{103}Ru ($T_{1/2}=39.8$ days) and ^{105}Ru ($T_{1/2}=4.44$ hrs) activities were due to thermal neutron capture in ^{102}Ru and ^{104}Ru . Although these isotopes were only present in small quantities (1.22% and 0.59%, respectively) their cross sections for thermal neutron capture (1.3 and 0.5 barns, respectively) are considerably larger than that of ^{96}Ru (<0.2b). The ^{96}Tc ($T_{1/2}=4.3$ days) activity formed by the $^{96}\text{Ru}(n,p)^{96}\text{Tc}$ fast neutron reaction, was roughly 1% of the ^{97}Ru activity under study.

The following chemical procedure was carried out to remove the ^{96}Tc activity and to minimize other impurities. The quartz irradiation vials were broken under a 3 ml solution of equal parts of 1N sodium hydroxide and 1N sodium hypochlorite. After boiling for 5 minutes to dissolve all the ruthenium, a 10 ml saturated solution of sodium persulfate was added to the boiling solution. The ruthenium tetroxide formed in this way was bubbled out of the solution by means of an air stream and trapped in 3 ml of 1/2 N sodium hydroxide. A drop of ethyl alcohol was added to this active solution, which was then passed through a fine filter paper. The filter paper was subsequently dried and used as the source for counting in standard geometries. No trace of ^{96}Tc activity was found after this purification procedure was carried out. The source was allowed to decay for approximately 24 hours before the chemical procedure was begun. This delay

allowed most of the ^{105}Ru activity to decay to ^{105}Rh ($T_{1/2}=35.5\text{h}$); the chemical procedure then removed the rhodium activity. No transitions from either ^{105}Ru or ^{105}Rh were observed in sources treated in the above manner. Since the half-life of ^{103}Ru is much longer than that of ^{97}Ru , it was impossible to wait until the ^{103}Ru activity had decayed. Therefore, transitions from the decay of ^{103}Ru appeared as impurity lines in the ^{97}Ru gamma-ray spectrum. Fortunately, the ^{103}Ru spectrum is rather simple and it has been well-investigated by several workers (e.g. Manthuruthil et al. 1968). A series of spectra were taken at regular intervals over a period of 10 days. Comparison of these spectra allowed positive identification of the transitions belonging to the ^{97}Ru decay.

Sources of 67 hour ^{99}Mo were made by thermal neutron capture in targets of enriched (98.8%) molybdenum metal and molybdenum oxide (MoO_3). Counting was delayed for approximately 24 hours to allow the ^{101}Mo ($T_{1/2}=14.6$ minutes) and ^{101}Tc ($T_{1/2}=14.0$ minutes) activities produced by thermal neutron capture in ^{100}Mo (0.33% of the target material) to decay away. In addition, equilibrium was established between the decay of ^{99}Mo and the decay of the 6.1 hour isomeric state in ^{99}Tc . Thus, the intensity of the isomeric transition was able to be correctly related to the intensity of the other transitions. Since no strong impurities were observed, no chemical purification procedure was used. The decay of a number of sources was followed by collecting gamma-ray spectra

at regular intervals over a period of ten days. From these spectra it was possible to definitely assign most of the observed transitions to the decay of ^{99}Mo .

Thermal neutron capture in enriched targets ($94.9\%^{100}\text{Mo}$) of molybdenum metal was used to obtain the ^{101}Mo activity. For most of the experiments, the source was made by depositing the active material in gelatine capsules for counting in standard geometries. The collection of data was started one minute after irradiation and was continued for 30 minutes. Although a shorter counting period would have minimized interference from the ^{101}Tc daughter ($T_{1/2}=14.0$ minutes), the procedure used gave a larger number of counts per source while maintaining a reasonable ratio of parent and daughter activities. The spectra of from 20 to 30 sources were summed to improve the statistics. Several methods were employed to identify the transitions belonging to the beta decays of ^{101}Mo and ^{101}Tc and to look for possible impurities. Gamma-ray spectra were accumulated through four successive half-lives for a number of sources. Using this procedure, transitions from the decay of ^{101}Mo decreased in intensity by a factor of eight in going from spectrum 1 to spectrum 4. Similarly, transitions from the decay of ^{101}Tc decreased to 75% of their original intensity. This method determined the parentage of over 80% of the gamma-rays observed. As a further check, a simplified version of the procedure developed by Fritze (1966) to separate molybdenum from other fission

products was utilized. The activity was dissolved in 10 ml of 6N HNO_3 , and then the solution was placed in a separatory funnel. After addition of 10 ml of an acetylacetone solution, the mixture was equilibrated by shaking for 1 minute. Spectra from both the molybdenum and technetium fractions were taken. In the technetium fraction, the molybdenum activity was decreased by a factor of 20. In order to search for any long-lived contaminants, several sources were irradiated for 10 minutes, allowed to stand for several hours, and then counted. Only the 140.5 keV transition from the beta decay of ^{99}Mo was found.

In the previous section it has been shown that limitations of the electronics at very high counting rates can introduce considerable peak distortions. Therefore, it was important to determine the maximum rate (or source strength) that can be used. The source strength was monitored by means of the dead-time meter on the ADC's. Although this meter did not give a direct reading in counts/sec, the relationship between dead time (per cent losses) and source strength was determined empirically. It was found that if the dead time was kept below 20%, the peak distortion was held to a tolerable level. In the case of the shorter-lived isotopes, the source strength decreased regularly with time. It was established empirically that the probability of a gain shift was held to a minimum if the dead-time was kept above 4%. When the source strength had decreased to this amount, a new source was installed and counting was continued.

2.5 Intensity Measurements

The intensities of the gamma rays associated with the decay of a particular nuclear species can be extracted from the measured areas of the full energy peaks in the following way. The area associated with the i^{th} gamma ray is given by

$$N_i = N_o (\epsilon \omega)_i I_i$$

where N_o is the source strength, ϵ_i and ω_i are the intrinsic efficiency of the detector and the solid angle subtended by the source at the detector respectively, for the i^{th} gamma ray and I_i is the intensity of the i^{th} gamma ray. Since the source strength N_o is usually unknown in most gamma-ray intensity measurements, it is only relative gamma-ray intensities that are determined. These relative intensities can be found if the quantity $\epsilon \omega$ is known for each gamma-ray energy.

The following procedure is used to determine $\epsilon \omega$ as a function of energy (the so-called efficiency curve). A number of gamma-ray sources having well-established source strengths are counted in a "standard geometry" for the detector to be calibrated; the term "standard geometry" meaning a geometry in which all sources, standard or unknown, will subtend the same solid angle at the detector. This procedure avoids determining ω each time a source is counted.

By carefully choosing the energies of the standard sources, it is possible to generate a curve representing the

relationship of efficiency ($\epsilon\omega$) and energy (E) for the energy range of interest. Above a certain energy, which is determined by the detector and by the presence of any absorbing material between the source and the detector, it is found for most detectors that $\log(\epsilon\omega)$ decreases as a linear function of $\log E$. Below this energy, the log-log relationship is no longer satisfied and it is therefore essential to space the standards more closely to avoid any errors due to incorrect interpolation. The problem is particularly acute if thick absorbers of material having a high atomic number are employed.

The standard sources used in the present work for efficiency calibration were of three types. The primary standards were the Vienna IAEA group of calibrated sources $^{22}\text{Na}(2.602\text{y})$, $^{54}\text{Mn}(312.7\text{d})$, $^{57}\text{Co}(271.6\text{d})$, $^{60}\text{Co}(5.264\text{y})$, $^{88}\text{Y}(107.4\text{d})$, $^{137}\text{Cs}(30.6\text{y})$, $^{203}\text{Hg}(46.57\text{d})$, $^{241}\text{Am}(458\text{y})$. The half-lives have been given in brackets. The manufacturer supplied the source strength and date of calibration for each standard. The accuracy of the measurements was guaranteed to be of the order of one or two per cent. The source strength at the time of use was calculated from the calibrated value, the half-life and the time elapsed between calibration and use. A number of secondary standards were also employed: ^{24}Na , ^{27}Al , ^{37}Cl , ^{42}K , ^{51}Cr , ^{51}V , ^{198}Au , ^{207}Bi , ^{232}Th . Their source strengths were measured relative to a 7.6 cm x 7.6 cm NaI(Tl) detector. The absolute efficiency of this type of

detector as a function of energy is well known (Heath, 1964). The secondary standards had the dual purpose of supplementing the energy scale (particularly above 1500 keV) covered by the Vienna sources and of providing an independent check of the efficiency curve. Portions of the efficiency curve not covered by the standard sources were established by using sources emitting many gamma rays for which the relative intensities are well known. Among these were ^{151}Pm , ^{152}Eu , $^{180\text{m}}\text{Hf}$, and ^{192}Ir .

The intensity measurements were carried out with several detectors arranged in different geometries. This procedure avoided the systematic errors which might have arisen had only one detector and one geometry been employed. Since the qualities of highest resolution and highest efficiency are not usually possessed by a single detector, the use of a range of detectors allowed the detector (and geometry) to be matched to the particular demands of the experimental situation. A number of experimental situations are examined below.

Consider for example, what happens when the radiation from a radioactive source is incident on a detector system having some absorbing material between it and the source. As the gamma-ray energy decreases, the attenuation due to the absorber will become appreciable and the efficiency of the system (detector + absorber) will no longer follow the $(\epsilon\omega) \propto E^{-n}$ relationship applicable at higher energies. Eventually, the efficiency will decrease with decreasing energy. In this region, no simple mathematical relation exists and

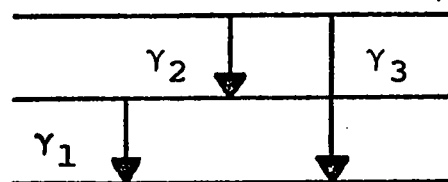
the efficiency curve is rather difficult to establish with high accuracy. The situation is made more complicated by the fact that not many low energy standard sources are readily available and therefore it is difficult to obtain enough calibration points in this energy region. If the intensities of low energy gamma rays or X-rays are to be accurately determined, a geometry with very little absorbing material between the radioactive source and the detector is a very desirable requirement. The small volume ($<1\text{cm}^3$) Ge(Li) detectors are particularly useful for low-energy investigations. The very thin entrance windows (e.g. 0.13 mm beryllium + $150\text{ \AA}$ gold) of these detectors cause little absorption of the incoming radiation. Below 25 keV the absorption is due entirely to the beryllium window; at 6 keV, the efficiency is still 90% of the maximum value.

Unfortunately, there are some limitations to the use of both the thin window detectors and the "no absorber" geometry. Normally, thin window Ge(Li) detectors have rather small volumes and small depletion depths with the result that their efficiencies decrease rapidly as the energy is increased above 60 keV. As a consequence, detectors with a large depletion depth are required for the detection of higher energy gamma rays, particularly those above 500 keV. Since more events can contribute to the full energy peak in the larger detector, the peak-to-compton ratio is 3 or 4 times greater than that for a smaller volume detector. This means

that gamma rays of weak intensity can be found more easily on the Compton background of higher energy transitions.

In order to compensate for the low intrinsic efficiency ϵ of the small Ge(Li) detectors, the solid angle ω could be increased by moving the source as close as possible to the detector and the source strength could be greatly increased. However, there are several reasons why this cannot be done. If the count rate is increased too much, then the pulse handling system (preamplifier, amplifier, analyzer) can no longer process a single pulse independent of other pulses. The result will be extensive distortion of the peaks and considerable loss of resolution. In addition, sum peaks may result when two gamma rays arrive in the detector within the processing time of the associated electronics. A photopeak will be produced at an energy equal to the sum of the incident energies. Coincidence summing or random summing can occur, as described below.

Coincidence summing is the adding together of two radiations, gamma rays or X-rays, which are in cascade; for example, γ_1 and γ_2 in the figure below.



The area in the coincidence sum peak is given by

$$N_{cs} = N_o (\epsilon\omega)_1 (\epsilon\omega)_2 C_{12}$$

where C_{12} is the probability that γ_1 and γ_2 are in coincidence.

If there is some doubt as to the existence of a true crossover transition, the presence of a coincidence sum peak can complicate the problem considerably. In any event, the intensity of crossover transitions should be corrected for possible coincidence summing effects, particularly if the cascade transitions are strong, low energy radiations.

Random summing is the adding together of two uncorrelated radiations, for example, γ_1 and γ_3 of the figure above. The area of the random sum peak is given by

$$N_{rs} = N_o^2 (\epsilon\omega)_1 (\epsilon\omega)_3 I_1 I_3^2 \tau$$

where τ is the processing time of the electronics.

It is clear that both summing effects can be reduced by increasing the source to detector distance thereby decreasing the efficiency. The amount of random summing can be lowered by using weaker source strengths. Normally, considerable deterioration of resolution occurs before random summing becomes a serious problem. However, coincidence summing is often seen if the source is counted in a close geometry.

The use of an absorber between the radioactive source and the detector can be a valuable technique in certain experimental situations. If strong low energy radiations are attenuated by the absorber, it is possible to increase the overall source strength while maintaining an effective source strength at the detector which is within acceptable limits. This means that the source strength can be increased

for the purpose of seeking very weak transitions without degrading the system resolution because of count rate effects. The absorber can also prevent summing effects since the effective efficiency of the detector for possible summing radiations will be lowered considerably. The absorber used in most of the experiments was a graded absorber of 2 mm lead, 1 mm cadmium and 0.5 mm copper. The purpose of the cadmium and copper was to absorb lead X-rays and cadmium X-rays produced by photo-electric events in the lead and cadmium, respectively.

A number of radioactive nuclei are β^- emitters; in fact, three (^{99}Mo , ^{101}Mo and ^{101}Tc) of the four nuclei discussed in this thesis are β^- emitters. Absorption of the beta rays in the detector or its adjacent surroundings can be a possible source of background radiation. If a thin window Ge(Li) detector is being used, the thin window allows direct penetration of many of the beta rays into the intrinsic region of the detector. A continuous beta-ray spectrum will result and will be superimposed on the gamma-ray spectrum. In addition, a continuous spectrum of photons, called bremsstrahlung, is produced in radiative collisions between beta rays and the atoms of the absorbing material. The maximum photon energy is equal to the maximum energy of the beta rays. According to Evans (1955), the average bremsstrahlung per beta ray is given by

$$I \approx \frac{Z}{7,000} E_0^2 \frac{\text{MeV}}{\beta \text{ disint.}}$$

where Z is the atomic number of the absorber and E_0 is the maximum beta energy. The intensity of the bremsstrahlung depends on Z ; however, the shape of the bremsstrahlung spectrum is independent of the Z of the absorber and only depends on the beta-ray emitter being considered. The bremsstrahlung intensity is 1.2% for lead ($Z=82$) and 0.086% for carbon ($Z=6$) if a beta spectrum with a 1 MeV endpoint is being considered.

It is clear that regardless of the detector system, a low Z material should be used to absorb the beta rays before they reach the detection system. Unfortunately, the lower the Z of the absorber, the greater the thickness of material needed to absorb a particular beta energy. Normally a compromise between absorber thickness and Z of the material used must be made; this will depend on the particular geometry desired, on the gamma-ray energy spectrum, and on the beta spectrum to be absorbed. For the experiments of this thesis, an absorber consisting of 12.5 mm of lucite ($Z=6$) was used to remove the beta rays. The presence of the lucite beta stopper had little effect on the transmission of gamma rays above 60 keV and essentially no effect above 100 keV. Since the beta stopper was much thicker than either the gelatine capsules used for counting or the "Scotch" tape used to seal them, any small variation in these latter quantities was not a factor in the intensities derived for low energy radiations.

The various geometries and detectors utilized in the intensity measurements are listed below

Intensity Measurement Data

Detector	Depletion Depth	Absorber	Source to Detector Distance	Resolution
0.15 cm ³	(3 mm)	beta stopper	10 cm	1.6 keV @ 662 keV
0.15 cm ³		beta stopper, standard ab- sorber	2 cm	
0.8 cm ³	(5.5 mm)	beta stopper,	10 cm	{ .7 keV @ 122 keV 1.3 keV @ 662 keV
0.8 cm ³		beta stopper, standard ab- sorber	2 cm	
12 cm ³	(11 mm)	beta stopper, standard ab- sorber	5 cm	2.9 keV @ 1332 keV

The sources, standard or unknown, were placed in a depression machined in a shelf which was then located in the appropriate slot of the source holder. With the shelf completely pushed into the holder, the source was automatically centered in front of the detector. Since the source holder was mounted securely to the detector housing, the distances were quite reproducible. By using this technique the variation in source to detector distance for all sources was kept at less than 0.1 cm.

In most experiments, the relative intensities of the stronger transitions were determined with 5-7% error. There are several reasons why the author feels that these errors could not be reduced. First of all, the shapes of the standard sources employed for efficiency calibration were

not identical. Therefore, there was some scatter in the efficiency curves, particularly at low energy. Secondly, since many sources were sometimes required to attain good statistics, some error in source positioning was unavoidable. Also because of the finely divided nature of the metal powder used as the source material, it was impossible to ensure that all of the material was at the bottom of the gelatine capsule during the counting period. Some source activity was therefore not at the standard geometry position during counting.

2.6 Energy Measurements

The energies of the transitions observed in gamma-ray spectra can be found by comparing the peak positions of these transitions with the peak positions of standard gamma rays whose energies are very accurately known (see Table 2.1). Since the number of standard gamma-rays is rather limited, their peak positions and corresponding energies are normally combined to form a calibration curve in which energy is a simple function of peak position. The calibration curve is then employed as an interpolating device to obtain the energy of unknown gamma rays. It is very important that the energy calibration curve include the full effect of any non-linearities inherent in the analyzing system; otherwise, the interpolating procedure will be incorrect.

The non-linearity of any system can be established in the following way. A spectrum containing approximately

TABLE 2.1

List of Gamma-Ray Energy Standards

Nuclide	Gamma-Ray Energy (keV)	Reference	Nuclide	Gamma-Ray Energy (keV)	Reference
^{241}Am	26.348 $\pm$ 0.010	a, b	^{203}Hg	279.191 $\pm$ 0.008	a
^{241}Am	59.543 $\pm$ 0.015	a, b	^{192}Ir	295.938 $\pm$ 0.009	d, a, b
$^{203}\text{Hg}(\text{Tl } K_{\alpha 2})$	70.833 $\pm$ 0.001	b	^{192}Ir	308.429 $\pm$ 0.010	d, a, b
$^{203}\text{Hg}(\text{Tl } K_{\alpha 1})$	72.873 $\pm$ 0.001	b	^{192}Ir	316.486 $\pm$ 0.010	d, a, b
$^{203}\text{Hg}(\text{Tl } K_{\beta 2})$	84.9 $\pm$ 0.1	b	^{51}Cr	320.080 $\pm$ 0.013	a, b
^{153}Gd	97.45 $\pm$ 0.02	c	^{198}Au	411.795 $\pm$ 0.009	d, a, b
^{182}Ta	100.104 $\pm$ 0.002	b	^{192}Ir	468.053 $\pm$ 0.014	d, a, b
^{153}Gd	103.19 $\pm$ 0.02	c	moC^2	511.006 $\pm$ 0.002	a
^{141}Ce	145.442 $\pm$ 0.006	a	^{207}Bi	569.64 $\pm$ 0.06	Avg a, b, e
^{182}Ta	152.435 $\pm$ 0.003	b	^{192}Ir	588.557 $\pm$ 0.017	d, a, b
^{182}Ta	156.387 $\pm$ 0.003	b	^{192}Ir	604.385 $\pm$ 0.017	d, a, b
^{139}Ce	165.852 $\pm$ 0.010	a	^{192}Ir	612.435 $\pm$ 0.017	d, a, b
^{182}Ta	179.393 $\pm$ 0.004	b	^{137}Cs	661.635 $\pm$ 0.076	a
^{192}Ir	205.782 $\pm$ 0.014	b, a	^{54}Mn	834.81 $\pm$ 0.03	a
^{182}Ta	222.109 $\pm$ 0.005	b	^{46}Sc	889.22 $\pm$ 0.07	Avg a, b, f
^{182}Ta	229.322 $\pm$ 0.008	b	^{88}Y	898.04 $\pm$ 0.04	a
^{182}Ta	264.072 $\pm$ 0.009	b	^{207}Bi	1063.59 $\pm$ 0.06	Avg a, b, e

Table 2.1 (continued)

Nuclide	Gamma-Ray Energy (keV)	Reference	Nuclide	Gamma-Ray Energy (keV)	Reference
^{65}Zn	1115.45 $\pm$ 0.07	Avg a, b, f	^{40}K	1460.75 $\pm$ 0.06	a
^{46}Sc	1120.44 $\pm$ 0.07	Avg a, b, f	^{207}Bi	1769.71 $\pm$ 0.13	a, b
^{60}Co	1173.226 $\pm$ 0.040	d, b	^{88}Y	1836.13 $\pm$ 0.04	a
^{22}Na	1274.55 $\pm$ 0.04	a	^{208}Tl	2614.47 $\pm$ 0.10	d, a, b
^{60}Co	1332.483 $\pm$ 0.046	d, b	^{24}Na	2753.92 $\pm$ 0.12	d, a, b
^{24}Na	1368.526 $\pm$ 0.044	d, a, b			

a Marcon (1968)

b Lederer et al. (1967)

c Alexander (1964)

d Murray et al. (1965)

e Kern (1970)

f Black and Heath (1967)

15 or 20 standard gamma rays and covering a wide energy range is collected with the system whose non-linearity is to be checked. Two of these gamma rays, well separated in energy and therefore channel position, are used to establish two channel positions which are assumed to have no deviation from linearity. Since energy is a linear function of channel position for these two data points, a linear calibration curve of energy versus channel number can be generated. The deviation from linearity for all other channel positions (relative to the two channels above) is then determined by fitting a curve through the differences between the peak positions predicted on the basis of the linear fit and the actual peak positions. The differential non-linearity for the Nuclear Data 161F ADC's used in most of the present experiments was found to have the form of a simple polynomial. It should be noted that the non-linearity characteristic shape varies with amplifier gain setting. Since the determination of a non-linearity curve is a rather time consuming process, it is somewhat impractical to perform this task for every spectrum from which gamma-ray energies are required. Instead, the experimenter would prefer to be able to account for the non-linearity effects by means of a simple analytic function. The function which approximates the shape of the non-linearity curve for the ND 161 analyzer is the cubic

$$E = a + bx + cx^2 + dx^3$$

where x is the channel number and a, b, c and d are parameters

to be determined. These parameters were calculated with a non-linear least-square fit to the standard energies using the program JAGSPOT which is discussed section 2.7. An example of a cubic fit to a group of standard energies is displayed in Table 2.2. Quadratic fits to the data were much poorer and it was felt that quartic fits were somewhat unrealistic in view of the shape of the non-linearity characteristic determined experimentally.

Energy calibrations of gamma-ray spectra can be performed either by a "mixed source" technique or by a "before and after" technique. In the latter method, an energy calibration spectrum is taken before and after the collection of the unknown spectrum. The peak positions of the two calibration runs are averaged and a resultant calibration curve is generated. This curve is then employed to calculate the energies of the unknown spectrum from the peak positions of the transitions. The above method should account reasonably well for any slowly varying system instabilities which are not too large. It is probably suitable for experiments of short duration and for experiments where counting rate variation is not a serious problem. However, this method cannot correct for system fluctuations which are rapidly varying or periodic or of a large amplitude. For many experiments, long counting times (2 or 3 days) are needed to accumulate sufficient statistics and system instabilities (e.g. those due to line voltage variations or temperature changes) may become a problem. If short half-lives are

TABLE 2.2

Example of Jagspot Energy Calibration

Standard Energy (keV)	Error (keV)	Fitted Energy (keV)	Error (keV)	Deviation (keV)
279.191	0.008	279.187	0.062	-0.004
295.938	0.009	295.917	0.071	-0.021
316.486	0.010	316.504	0.052	+0.018
411.795	0.009	411.798	0.042	+0.003
468.053	0.014	468.096	0.055	+0.043
569.640	0.060	569.696	0.064	+0.056
604.385	0.017	604.342	0.060	-0.043
612.435	0.017	612.397	0.080	-0.038
661.635	0.076	661.560	0.070	-0.075
834.810	0.030	834.850	0.053	+0.040
898.040	0.040	898.072	0.068	+0.032
1063.590	0.060	1063.615	0.173	+0.025
1115.450	0.070	1115.412	0.073	-0.038

Avg = 0.034

$$E = -(16.495 \pm 0.025) + (3.412 \pm 0.004) \times 10^{-1} C - (1.48 \pm 0.24) \times 10^{-6} C^2 + (1.50 \pm 0.39) \times 10^{-10} C^3$$

where C is the channel number

involved, the counting rate will vary widely because of the decay of the source. Multiple sources are usually needed for short-lived activities and it is normally quite difficult to avoid variation in activity from source to source. Any one of these effects can cause considerable gain shifts between the spectrum containing the unknown gamma rays and the calibration spectra. There is a modification of the "before and after" technique which improves the situation. One or two energy standards can be included with the unknown spectrum so that any instabilities arising during collection of the unknown spectrum can be monitored. This method is still unsuitable if the system fluctuations are non-linear. The technique would be useful in neutralizing any zero shifts, but any gain shifts could not be compensated for.

For the reasons discussed above, it was decided to only make use of the "mixed source" technique to obtain energies. To employ this method, it is necessary to simultaneously accumulate the energy calibration standards with the unknown gamma rays from the source that is to be calibrated. In general, this technique is to be preferred over any other, since zero shifts, gain shifts and counting rate effects such as distortion of peak shape will automatically affect both sets of gamma rays in the same way.

Because of the rather large differential non-linearity of the analyzing system used, several procedures were employed to minimize the probability that systematic errors in energy calibration would arise due to incorrect non-linearity

compensation. First of all, the energy calibrations were performed using a number of overlapping energy ranges. For example, the energy standards contained in channels 300-1500, 1000-3000 and 2000-3800 were used to generate three separate energy calibration curves, one for each region. This method reduced the problem of trying to reproduce the non-linearity of the system over an excessively wide range of channel numbers. In addition, the two sets of energies calculated from the overlapping energy ranges served as a crosscheck. In most cases excellent agreement was obtained. Many of the energy calibrations involved only the low energy portion of the gamma-ray spectrum. It was therefore possible to run these experiments at a relatively high gain so that the contribution of the error in the peak position to the error in the energy was reduced. The energy calibrations were checked for reproducibility of results by replacing some of the energy standards with different ones and repeating the measurements. In all cases, excellent agreement was obtained.

The energies of the stronger lines in the spectrum obtained from the "mixed source" technique were later employed as secondary energy standards in spectra taken for intensity purposes. These intensity spectra were collected for much longer time periods in order to increase the statistics and to look for weaker gamma rays. Of course, no energy standards were included in these experiments since they might obscur

weak transitions. In addition to the gamma rays used as secondary standards, energies were extracted for a number of weaker gamma rays from the energy calibration spectra. It was very gratifying to find that these energies agreed quite well with those calculated from the intensity measurements using the secondary standards.

A number of precautions were taken to ensure that the experimentally determined energies would be essentially independent of the peak fitting procedure (see Section 2.7). The total counting rate due to the energy standards was adjusted (by placing them at different distances from the detector) to be approximately equal to the average count rate from the unknown source. In this way, equal peak heights were attained for both the standard set of gamma rays and the stronger transitions from the unknown source. This procedure would reduce any discrepancies which might arise due to inadequacies of the peak fitting routine. In addition, the upper limit of the combined source strength was kept less than 20% on the dead time meter in order to prevent peak distortions.

Very strong sources were necessary to attain the required statistics for the high energy calibrations due to the relatively low Ge(Li) efficiencies at these energies. However, such source strengths gave large peak distortions because of the high counting rate. This problem was resolved by placing between the detector and the unknown source

(and any energy standards which had low energy gamma rays or X-rays in addition to the gamma rays used for calibration) large amounts of lead (5 to 10 mm) to attenuate the strong low energy radiations. The counting rate at the detector was then reduced to a tolerable level without having to lower the source strengths.

In order to reduce the number of sources employed in the energy calibration experiments while retaining good statistics, the procedure below was followed. A very strong source (5 to 10 minutes irradiation time) was made and then placed at a distance from the detector which gave the required counting rate. Every 10 minutes or so the source was moved closer to the detector; close watch was kept on the dead time meter so that the count rate remained within acceptable limits (20%, 15%). The lower limit was necessary so that the analyzer wasn't only accepting events from the energy standards. This method gave the same results as the many source technique.

2.7 Peak-Fitting Procedure

Since the primary function of a singles gamma-ray experiment is to provide accurate energies and intensities for the observed gamma rays, an important prerequisite is the extraction of accurate peak positions and areas from the experimental data. Therefore, the method employed to calculate these quantities should introduce no systematic errors or prejudices of its own. Ideally, one would like

to use a peak-fitting procedure which gives exact fits to the peaks in the spectrum, fits that are independent of the underlying background from higher energy gamma-rays. Since this is impossible to do in practice, it is necessary to use an analytic function, with a small number of parameters, to approximate the shape of the peaks in the spectrum together with a second function to approximate the background.

The peak-fitting routine JAGSPOT was used for determining the peak positions, energies and areas of the gamma-ray peaks found in the present investigation. This program was originally written at the Chalk River Nuclear Laboratories (Williams and McPherson 1968, Graham et al. 1972) in an ALGOL-type language called APEX. The gamma-ray spectroscopy group at McMaster changed the program language to FORTRAN IV for use on the CDC6400 computer and also extended several features of the program.

The fitting function used in the program was specifically selected to handle peaks with an excess of counts on their low energy side. This excess was present because of the properties of the overall detection system (electronics + detector). The form of the function is

$$I(x) = \alpha + \beta x + \sum_{j=1}^N \gamma_j \int_{-\infty}^{xp_j} e^{\epsilon(y-xp_j)} e^{-\delta(x-y)^2} dy$$

where x = channel number

$I(x)$ = counts in channel x

N = number of peaks in the region being considered (≤ 6)

$\alpha + \beta x$ = linear background term

γ_j = intensity of the Gaussian located at x_{p_j}

x_{p_j} = position of first Gaussian in peak j .

ϵ = constant determining the exponential fall-off
of Gaussian intensities

$\delta = 1/2\sigma^2$ of normal Gaussian in peak j

The parameters that are determined by the program's least-squares fitting routines are α , β , δ , ϵ , $(\gamma_j, x_{p_j})_{j=1}^N$.

The values of these parameters are examined after each of the 20 iterations that the routine performs. If the values indicate that the shape of the fitted function is tending towards a Gaussian form, the program automatically switches to the function

$$I(x) = \alpha + \beta x + \sum_{j=1}^N \gamma_j e^{-\delta(x-x_{p_j})^2}$$

which is derived from the more general function by setting $\epsilon=0$.

It could be argued that the Gaussian function should be sufficient for fitting all spectra. Although it is not correct, it will miss the same fraction of the peak each time, and therefore, the effect of the incorrect fitting procedure should have little influence on the final results. In fact, since the exponential distribution of Gaussians accounts only approximately for the peak shape, the same argument must be made for that fitting procedure. However, consideration of a multiple peak system shows that, unless the fitting procedure closely approximates the true peak shape, large errors can result. If a large fraction of the peak is not

accounted for by the fitting procedure, peak positions and areas of neighbouring peaks can be influenced greatly by any residuals.

A description of the use of JAGSPOT is given in Appendix A . This discussion makes it clear that the program does not solve all the problems associated with peak fitting. However, accurate and consistent fits can be achieved with reasonable care. Figure 2.3 displays a number of peaks which have been fitted using JAGSPOT; the plus signs indicate the original data points, the curve is drawn through the values generated by the fitting procedure. The plotting routine made no attempt to draw a smooth curve through the two highest values, they are just joined together. The top group of diagrams shows several strong peaks, one a doublet. All parameters were allowed to vary in these fits. The middle group displays several multiplets; δ and ϵ have been allowed to vary for the one on the left, they have been chosen from a calibration curve for the one on the right. The bottom group of diagrams shows a few weak peaks which were fit with constant (from a calibration curve) values of δ and ϵ . High and low background situations are included. The fits are seen to be excellent.

2.8 Internal Conversion Measurements

Internal conversion is a mode of decay competing with gamma-ray emission in the deexcitation of a nuclear excited state. For each transition, the internal conversion process

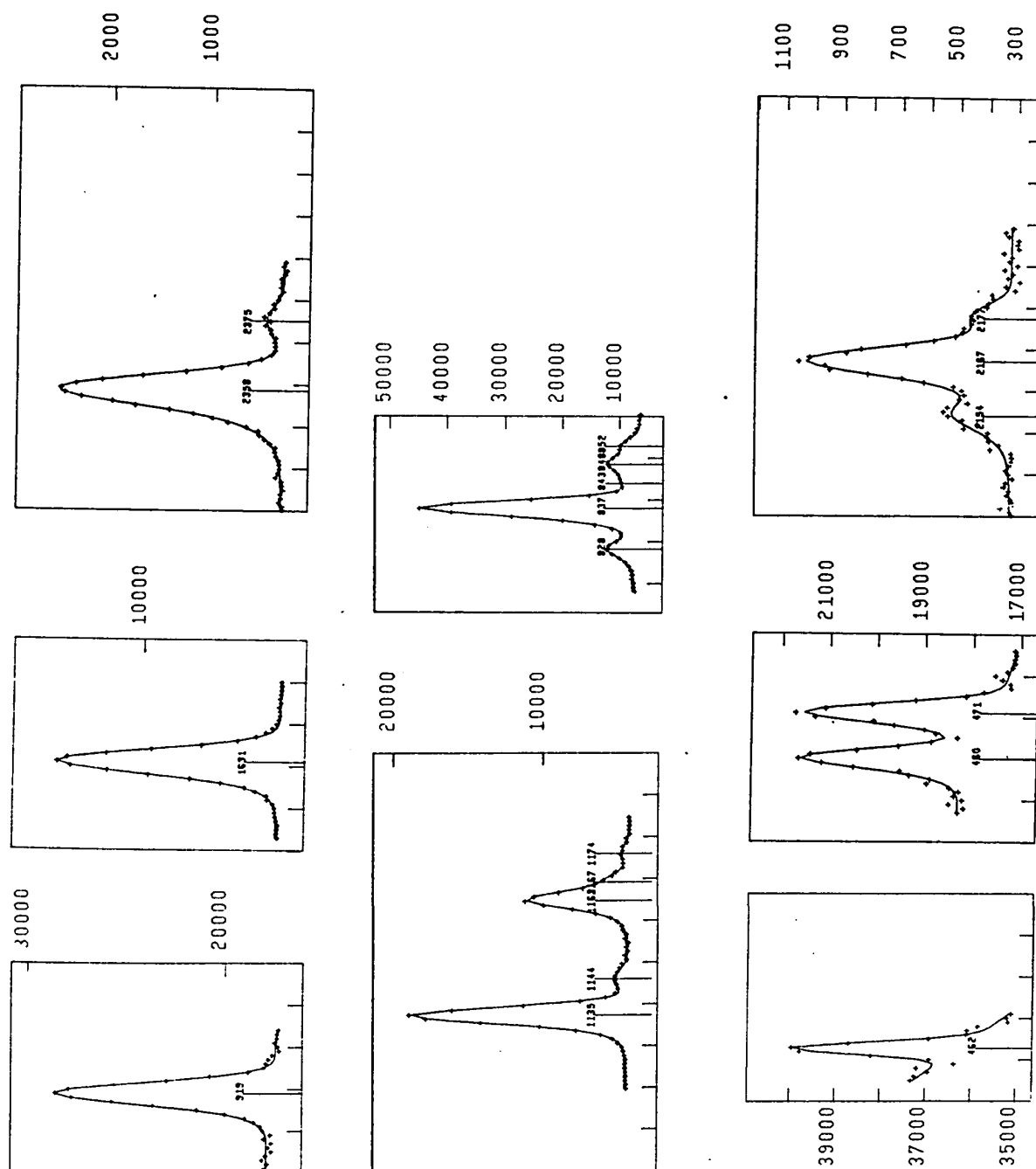


Fig. 2.3 Example of Jagspot fits to spectral peaks

gives rise to monoenergetic electron groups corresponding to conversion in each of the K, L_1 , L_2 , L_3 , etc. atomic shells. The resulting electron spectrum can therefore be quite complex demanding high energy resolution from the detection device. A unique identification of the multipolarity of a given transition can be deduced by comparing the experimental internal conversion line intensity ratios $K:L_1:L_2:L_3$ etc. with the theoretically predicted ratios (Hager and Seltzer 1968). The internal conversion coefficients are quite high for low energy transitions with the result that the internal conversion process becomes the dominant mode of deexcitation for transitions with energy less than ~ 50 keV. Therefore, for such transitions, it is particularly important to measure the internal conversion lines.

One of the most common ways of achieving high resolution for the measurement of electron spectra is through the use of a magnetic field to deflect the electrons in a suitable fashion. A magnetic field having rotational symmetry and a radial field gradient causes the electrons to describe oscillatory paths in both the radial and vertical directions. Focussing in both these directions (double-focussing) will occur after an angle of $\pi/\sqrt{2}$ radians if the field decreases radially as $1/\sqrt{r}$. The Chalk River spectrometer (Graham et al. 1960) is a one-meter radius, iron-free, double-focussing $\pi/\sqrt{2}$ spectrometer. This instrument was designed to have a resolution of from 0.01% to 1% full line width at half maximum depending on the

transmission aperture, which can be adjusted from 0.1% to 1% of 4π steradians. The magnetic field strength, and hence the momentum of the focussed electrons, is directly proportional to the excitation current in the field-producing coils, due to the iron-free construction of the spectrometer. Standard methods can then be used to measure the current to high precision. The precision with which the current through the coils can be controlled and measured will determine the precision with which momenta of focussed monoenergetic electrons lines can be measured. In addition, the accuracy of electron momentum measurements depends critically on the accuracy with which the earth's field can be compensated. For this reason, two large sets of degaussing coils are used to provide cancellation of the vertical and north-south components of the earth's magnetic field. For the Chalk River instrument, the proportionality between the momentum of the focussed electrons and the current in the coils is constant over the useful momentum range to better than 1 part in 10^5 . To achieve this stability, the IR drop in a precision resistor, connected in series with the spectrometer coils (all in series), is compared with the reference potential chosen by the experimenter. The current regulator signal is derived from the difference signal between these two potentials. Measurement of the spectrometer current, and hence the electron momentum, is achieved by measuring the IR drop across the standard resistor with a precision potentiometer. This was done before and

after each run.

In order to minimize the effects of the variation in background counting rate from the reactor-produced ^{41}Ar (Graham et al. 1968) which drifts over the spectrometer building, a rapid-scan counting technique (Gelletley et al. 1968) was employed. In this technique, a Nuclear Data 3300 analyzer was used to collect the events which occurred during repetitive rapid scans of the momentum region being considered. The upper and lower reference potentials, which corresponded to the initial and final channels of the analyzer, were selected by two separate potentiometers. The channel address was normally incremented every 0.5 seconds. At the same time, an analog voltage proportional to the channel number was generated. Since the channel address register and the current regulator are linked electrically, the β -spectrometer field strength was incremented simultaneously with the channel address. Therefore, the spectrometer field setting was caused to sweep the desired region of the spectrum. The counts recorded at any particular voltage (field) setting were deposited into the corresponding channel. For the present experiments, the region of interest (typically 5 to 10% in momentum) was covered by a scan involving 512 steps with a counting period of 0.5 seconds per step. Upon completion of a scan, the scan number and total counts accumulated during the scan were printed out and then the sequence was restarted. The elapsed time per scan, including readout time, was 270 seconds. Six scans were normally made

for each electron source. Upon completion of a run, the contents of the 512 channels were printed out on paper tape for subsequent analysis.

The internal conversion study of 14.6 min. ^{101}Mo posed some rather serious source preparation problems due to the short half-life. In order to prepare thin sources suitable for counting in the $\pi\sqrt{2}$ spectrometer, it was necessary to evaporate the ^{101}Mo activity onto a suitable source backing. Molybdenum is rather difficult to evaporate and the time elapsed during source preparation can be rather long. The resultant loss of activity prohibits any such processing of the samples after irradiation. Therefore, to minimize the time interval between the source irradiation and the start of counting, ^{100}Mo samples were prepared prior to neutron irradiation. This was done by evaporating molybdenum metal powder, enriched to 96% in ^{100}Mo , in vacuum from a hot tungsten crucible onto a 0.0001 inch thick pure iron foil.

The thickness of the ^{101}Mo deposit was measured using the X-ray fluorescence technique. A gamma-ray source of ^{131}I was used to excite fluorescent X-rays in the ^{100}Mo sample and in a standard sample of molybdenum metal of known thickness. The X-rays were detected in an 30 mm² Kevex Si(Li) detector with a resolution of approximately 250eV at 6 keV. The samples used had thicknesses in the range of 3 $\mu\text{g}/\text{cm}^2$ to 17 $\mu\text{g}/\text{cm}^2$. Small strips (~ 6 mm $\times$ 25 mm) of the foil were irradiated for ~ 30 minutes in either the NRX (flux $\sim 0.5 \times 10^{14}$ neutrons /cm²/sec) or NRU (flux $\sim 2.5 \times 10^{14}$ neutrons /cm²/sec) reactor at Chalk

River. After completion of the irradiation, the sample was delivered via pneumatic tube to a nearby building. The active foil was mounted on a source mask with a 3 mm (wide) $\times$ 25 mm aperture. Then, the source was transported the approximately 1/4 mile to the spectrometer building and mounted in the spectrometer. Approximately one half-life elapsed between removal of the source from the reactor and initiation of counting.

In addition to the desired ^{101}Mo activity, neutron irradiation of the iron foils produced the 45 day ^{59}Fe and 2.58 hour ^{56}Mn activities. The ^{101}Mo activity dominated the backing activities for the electron momentum regions of interest. In order to prevent accumulation of the background activities, new source material was used for each irradiation. Each source was counted for two half-lives and then a new one was installed. Forty sources were used in all.

The internal conversion measurements were carried out at a momentum resolution of 0.3%. Since extremely high resolution was not required, rather large source (3 mm) and counter (10 mm) apertures and a wide baffle setting were employed to increase the counting rate at the detector. The transmission was approximately 0.3%.

In the first series of experiments the L and M conversion lines of the 6.281 and 9.317 keV transitions (see section 3.3.3) were studied using a Johnson Laboratory focussed mesh electron multiplier (Johnson Laboratories Inc.,

Baltimore, Maryland, U.S.A.) as the electron detector.

The electron multiplier is a series of dynodes each of which consists of hundreds of raised electron-emissive surfaces and holes for electron passage. These dynodes are stacked so that the raised surfaces of one dynode are aligned with the holes in the dynode above it. Electrostatic focussing insures efficient transmission of secondary electrons from dynode to dynode. In this manner, electron multiplication progresses through the dynode assembly of 20 stages to give a gain $\sim 10^6$. It was assumed that the efficiency of the electron multiplier remained constant over the range of any one momentum scan (typically 0.5 keV).

In the second set of experiments the M lines of the 6.281 and the L lines of the 9.317 and 15.606 keV transitions were recorded with a flow-type proportional counter (Graham et al. 1960). It was operated with an $\sim 100 \mu\text{g}/\text{cm}^2$ laminated VYNS window and at a pressure of ~ 90 torr. The counter with its preamplifier, amplifier and discriminator was electrically insulated and could be operated at potentials of up to ± 5 kV with respect to the vacuum chamber. This made it possible to accelerate or decelerate the electrons after they had been magnetically analyzed. In order to establish the counter window absorption correction, the L-Auger complex from the decay of ^{166}Ho was studied. Scans of this region (2.8 - 10.9 keV) were taken using several values of post-accelerating voltage on the counter. The variation in the shape of the Auger spectrum under the different post-acceleration conditions

gave the window transmission curve for the energy range of 4 - 15 keV. A post-acceleration voltage of 4 kV was used for scanning the region containing the M lines of the 6.281 and the L lines of the 9.317 keV transitions. For this scan the data were corrected by 5% for the counter-window absorption. No post-acceleration was used for the 15.606 keV L lines and no absorption correction was applied. The proportional counter was operated with an ~ 1 mg/cm² Mylar window and atmospheric gas pressure for the measurements of the K-conversion lines of the 80.926 and 192.00 keV transitions. The counter window transmission of 93% at 59.9 keV was taken from the data given in Fig. 1 (curve D) of Geiger et al. (1960).

Since the conversion-electron measurements required multiple sources, normalization was provided by recording a Ge(Li) gamma-ray spectrum from each source before it was placed in the spectrometer. The relative strengths of the ¹⁰¹Mo activity were deduced from the peak areas of the 192 keV gamma ray. Since the background decayed much more slowly than the ¹⁰¹Mo activity, it was treated as a constant during the analysis. After the background was subtracted, a decay correction was applied to the data. The spectrometer was calibrated for energy using the L₁ line of the 39.85 keV transition in ²⁰⁸Tl which has a β_0 value of 534.11 ± 0.06 Gauss-cm (Graham et al. 1965).

The observed conversion line shapes were asymmetric with

a low energy tail caused partly by the focussing properties of the spectrometer and partly by energy degradation in the source. For the low energy lines, the second factor was the most serious. As a result, it was difficult to define the peak position. A useful definition was obtained by drawing a line midway between the high and low energy sides of the peak until it intercepted the curve drawn through the data points. This intercept was defined as the peak position. The asymmetrical line shape also complicated the evaluation of the areas of conversion lines which were not fully resolved. For the L-subshell lines, a first approximation to the areas was obtained by using the L_1 shape to fit the L_3 line first and then the L_2 line (this area was of course affected by the area chosen for the L_3 line). Removal of these two lines from the spectra gave a new L_1 shape. An iterative procedure then gave the final areas used to calculate the L-subshell ratios. Since the spectrometer has a constant momentum resolution $\Delta p/p$, the measured areas of the conversion lines were divided by the peak momentum (as well as corrected for the scanning increment used) to obtain a true measure of their intensity.

2.9 Gamma-Gamma Coincidence Measurements

2.9.1 Introduction

If the spectroscopist is to infer any information about nuclear structure from his data, a primary requirement is the establishment of a unique level scheme appropriate to the

particular nucleus being studied. A number of possible levels can be arrived at by considering the various energy sums that can be formed by adding together the energies of the observed transitions. Some of these sums are rejected outright since they would create levels above the energy permitted by the decay energy (for beta-decay processes). Other sums are excluded because the corresponding energy levels cannot satisfy intensity balance requirements. For any level, the total outgoing transition intensity must be equal to the total incoming transition intensity plus any intensity due to direct population of the level by beta decay. For a nucleus which emits only a few transitions, the limited number of energy sums permitted by the intensity balance often allows construction of a unique level scheme. However, if the number of observed transitions is large, there are many possible levels which are consistent with both intensity and energy sum considerations. Therefore, it becomes impossible to determine a unique level scheme in this way.

The existence of cascade transitions in nuclei can be established by the coincidence method, a powerful technique for determining the positions of transitions in a level scheme. In this method the radiations from a radioactive source are observed in two detectors oriented in a fixed position relative to the source and to each

other. By means of a suitable electronic circuit, the time correlation of events recorded in the two detectors can be measured. It is clear that the electronic circuitry must be able to distinguish between those events that are in true coincidence and those that are unrelated. Normally, portions of the pulse-height spectrum called windows or gates are chosen from one of the detectors. The spectra in coincidence with these gates are then recorded in the other detector. Transitions appearing in the latter spectra are in cascade with the gating transition.

2.9.2 Time Resolution

The signal corresponding to a detected gamma ray is examined by a discriminator circuit in such a way as to define the time of occurrence of the detected event. The leading edge of the standard logic pulse generated by the time pick-off device conveys the time information. One way that the difference in time of occurrence of the events in detectors 1 and 2 can be measured is with a time-to-amplitude converter (TAC). This measurement is accomplished by charging a capacitor from a constant current source during the time interval between the start and stop pulses. In an ideal system, the TAC output would contain only pulses of one height, those corresponding to true coincidence events. The output will be a spectrum with a peak of finite width associated with the true events and a continuous background associated with the chance events. The full width at tenth maximum

(FW 1/10 M) of this peak is the significant parameter which indicates the time resolution of the system. A number of factors determine this resolution. Some of these will be discussed below.

Since both NaI(Tl) and Ge(Li) detectors were used in the coincidence experiments, some of the timing characteristics of each will be discussed below. The time resolution attainable with scintillation detectors is determined by the properties of both the scintillator and the photomultiplier and their coupling efficiency. The thickness of the scintillator, the finite decay time of light emitting states in the phosphor, the light yield of the scintillator as a function of the energy of the detected radiation and the variations in the path length for collection of photons from the scintillator at the photomultiplier photocathode are some of the factors modifying the time resolution of the scintillator. Similarly for the photomultiplier, the factors are the efficiency of conversion of photons to emitted photoelectrons at the photocathode, the absolute number of photoelectrons released, the variation in the transit time of photoelectrons from the photocathode to the first dynode and the single electron response of the dynode gain structure. The time resolution for a photomultiplier is proportional to $\frac{1}{\sqrt{R}}$, where R is the average number of photoelectrons released per detected event (Gedcke and Williams 1968)

The time resolution obtainable with a Ge(Li) detector

depends upon properties of the detector (drift depth, detector geometry and electric field strength throughout the detector), the type of electronic system used and upon the source of radiation under investigation. Gamma rays produce ionizing events throughout the active volume of the detector. The collection time of the carriers produced will depend upon the position of each ionizing event. As a result, there must be a spread in the pulse rise times; the magnitude of this spread will be a function of the drift depth of the detector.

The time t_o at which a gamma ray is emitted from a source can be related to the time t_t at which the resulting signal from the detector causes a discriminator circuit to produce a standard timing pulse. The accuracy of timing or time resolution is determined by the variations in delay between t_t and t_o . Two characteristics strongly influence the choice of the proper time pickoff device: (a) the intrinsic time jitter of the detector signal, (b) timing walk due to amplitude or rise time variations. These effects are illustrated in Fig. 2.4. The upper left hand diagram shows two pulses of different rise times but equal amplitudes crossing a discriminator level. The slower rise time pulse reaches the discriminator level at a later time than the fast rise time pulse. The lower left hand diagram displays two pulses of different amplitudes but equal rise times crossing a discriminator level at different times. Although

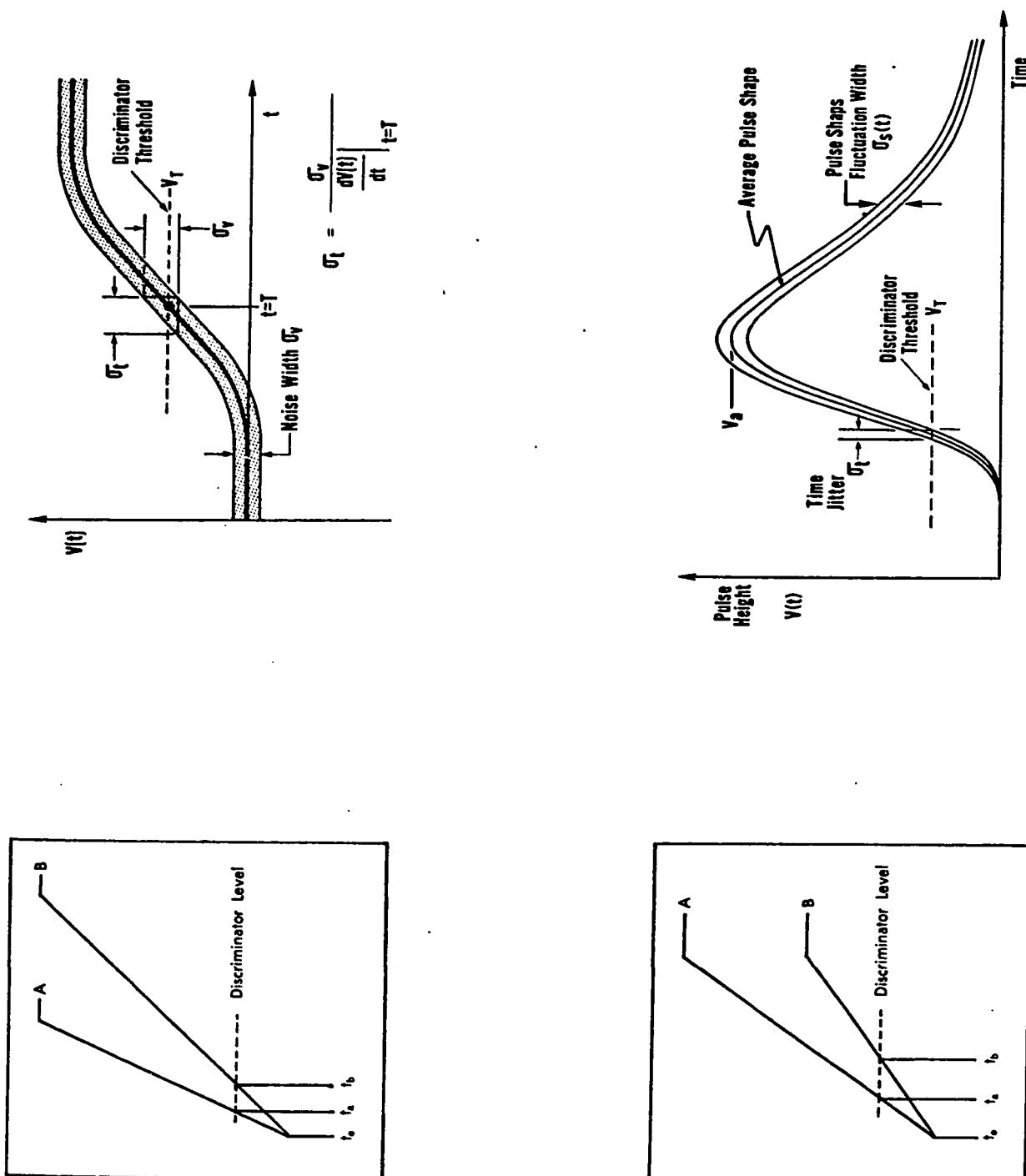


Fig. 2.4 The effect of walk on time resolution

the two signals were caused by events occurring at the same time, t_0 , the larger amplitude signal crosses the discriminator threshold before the smaller. In both of the above effects, the shift in the crossing time causes the output pulse of the discriminator to "walk" along the time axis as a function of pulse amplitude or rise time. The walk is most noticeable for pulse heights close to the discriminator threshold. The trigger circuit must accumulate the finite amount of charge sufficient to respond to the input pulse, after the discriminator level has been reached. Therefore, there will be an additional variation in delay time for each of the preceeding cases.

Independent of the walk effect is another source of timing error which is statistical in its nature. This error is due to an intrinsic time jitter of the detector signal as seen by the time pickoff device. The upper right hand portion of the figure shows the effect of system noise on the time variation. The noise can be generated in the detector itself, in the amplifying electronics, or may be the equivalent input noise of the timing discriminator. If the steady-state noise has amplitude σ_v , then the width of the timing uncertainty, σ_t , in crossing the discriminator threshold V_T is given by

$$\sigma_t = \frac{\sigma_v}{\left. \frac{dV(t)}{dt} \right|_{t=T}}$$

where $\left. \frac{dV(t)}{dt} \right|_{t=T}$ is the slope of the average signal shape as it crosses the discriminator threshold. This source of

jitter is a dominant limitation for semiconductor detectors, but is usually not important with scintillation detectors.

For scintillation-photomultiplier combinations, the statistics of emission and collection of photoelectrons dominate the timing properties. The net result is observed as fluctuations in the shape of the photomultiplier output pulse; the effect is shown in the lower right hand portion of the figure. The timing jitter, σ_t , is given by

$$\sigma_t = \left. \frac{\sigma_s(t)}{\frac{dV(t)}{dt}} \right|_{t=T}$$

where $\sigma_s(t)$ is the amplitude fluctuation width and where now both the numerator and denominator must be evaluated at the time, T , when the pulse crosses the discriminator threshold.

A number of techniques have been developed to obtain good time resolution. These include: "slow-crossover timing", "leading edge timing", "constant fraction of pulse height trigger" and "extrapolated leading edge timing". Slow-crossover timing can be used where the best time resolution is not required and where a moderately wide range of pulse heights (that is, a moderate dynamic range) must be processed. The signal from the detector is fed via the preamplifier to a shaping amplifier. A double differentiation produces a bipolar output pulse; the crossover point of this pulse represents the time to collect 50% of the total charge in the detector pulse. The output pulse is then fed to a timing

single channel analyzer (TSCA) which triggers at the zero crossing point and produces a standard logic pulse. Since the zero-crossing represents the same phase point for all pulses, almost all the walk effect is eliminated.

Leading edge timing with a fixed threshold discriminator offers the best time resolution for a narrow dynamic range. However, the fixed discriminator threshold does not have an optimum value simultaneously for all pulse heights. This together with walk effects leads to poorer time resolution for wide dynamic ranges. Due to walk effects, the FW 1/10 M increases much more rapidly than the FWHM as the dynamic range increases.

The experience gained with leading edge timing led to the design of a circuit which would provide amplitude and rise time compensation by triggering at a constant fraction of the pulse height. In this technique the input pulse is fed through two paths. Along one path, the input pulse is attenuated to a fraction of its original height. Along the other, the pulse is inverted and delayed by an interval T_D . Both signals are then added together to form a zero-crossing signal. A zero-crossing discriminator triggers at the zero-crossover point. Since the time of zero-crossing is essentially independent of pulse height, the timing walk of the system is minimized.

In extrapolated leading edge timing, the rate of rise of the input signal is measured with two discriminators, the second triggering at approximately twice the pulse height of

of the first. A time correction which is proportional to the rate of rise is then applied to the timing pulse from the upper discriminator. For pulses large enough to trigger the second discriminator, correction can be made for timing walk.

2.9.3 Gamma-Gamma Coincidence Measurements

The coincidence measurements described in this section were carried out over a three year period. During this time there were considerable improvements in both the electronics used to perform the experiments and the computer facilities used to process the experimental data. Thus the measurements described below reveal a sequence of increasing sophistication which matches the increasing complexity of the problems studied. It would have been very difficult to have interpreted the ^{101}Mo decay with the equipment used for the ^{97}Ru study.



The gamma-gamma coincidence measurements for ^{97}Ru were carried out using the 12 cm³ Ge(Li) detector and a 7.6 cm × 7.6 cm NaI(Tl) scintillation detector. The detectors were oriented at 180° to maximize the coincidence efficiency. A shaped anti-Compton shield of lead and cadmium, 3 mm thick and 8.5 cm in diameter, was inserted between the detectors to suppress inter-crystal scattering. The gamma-ray source was mounted in a 7 mm tapered hole in the middle of the absorber so that the source could "see" both detectors. In this

arrangement, gamma rays back-scattered from one detector had very little chance of reaching the other. As a consequence, the Compton and backscatter peaks, which were very prominent without the shield, were almost eliminated.

Standard crossover timing techniques (see Fig. 2.5) employing DDL's and TSCA's were used to generate timing pulses which were then fed to a TAC. Table 2.3 explains the abbreviations. The Ge(Li) detector provided the start pulse, the NaI(Tl) detector the stop pulse. A resolving time (FWHM) of 90nsecs was obtained. Two windows, each 200nsecs wide and corresponding to portions of the time spectrum on and off the prompt coincidence peak, were set on the TAC output by two SCA's. These two windows represented true + chance and chance coincidences, respectively; the former window incorporated ~95% of the prompt coincidence peak. Since source strengths were chosen to maintain a true-to-chance ratio of better than 50, it was not considered necessary to correct the coincidence data for chance, although the mechanism for doing this was provided.

Logic pulses from the SCA set on the prompt coincidence peak were fed to the coincidence inputs of the two Nuclear Data 161 ADC's. The memory of the ND 160-161 analyzer system is composed of 4096 locations, each of which is 18 bits in length. If this system is used directly as a two-dimensional analyzer, the memory size limits the experimental configuration to a 64×64 element matrix. Although such a configuration

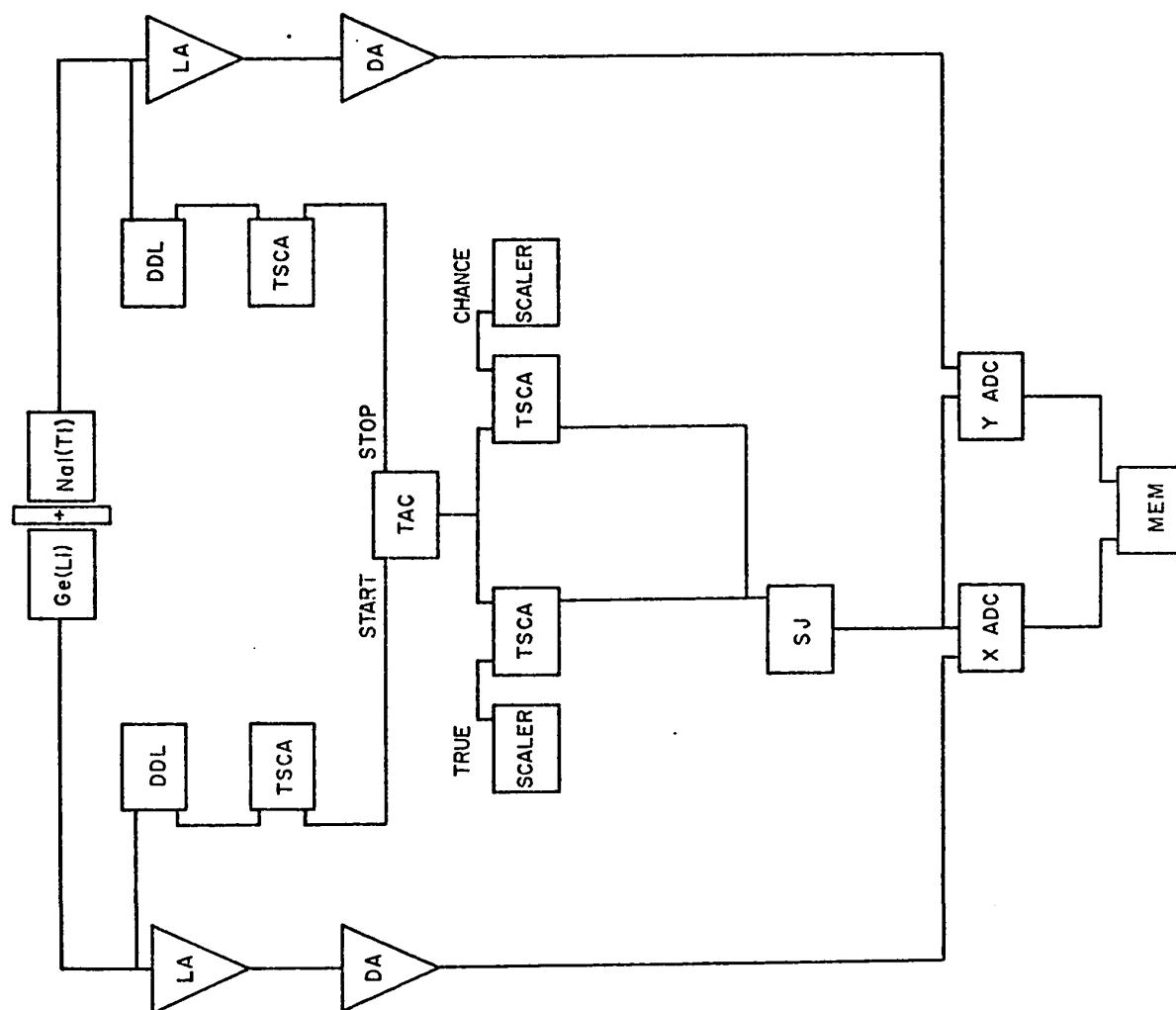


TABLE 2.3

Explanation of Code Used in Figs. 3.5 and 3.6

Abbreviation	Explanation
BR	Baseline Restorer (Tennelec Model 610)
CFD	Constant Fraction Discriminator (Ortec Model 453)
DA	Delay Amplifier (Ortec Model 427)
DDL	Double-Delay Line Amplifier (Canberra 1411)
LA	Linear Amplifier (Tennelec Model TC200 or Ortec Model 452)
LSAD	Logic Shaper and Delay (Canberra Model 1455)
MEM	Memory for the ND160-161 System or for the ND3300 system
PDP-9	Digital Equipment Corp. Computer
SJ	Summing Junction (Diodes)
TAC	Time-To-Amplitude Converter (Ortec Model 437A)
TFA	Timing-Filter Amplifier (Ortec Model 454)
TSCA	Timing Single Channel Analyzer (Canberra Model 1436)
SCALER	SCALER (Ortec Model 1473)

had proven quite acceptable for NaI(Tl) - NaI(Tl) coincidence measurements, the high resolution of a Ge(Li) detector requires more channels (e.g. 1024) along the Ge(Li) axis. Therefore, the following address-recording procedure was employed to increase the effective memory size. Each 18-bit word was divided into two parts. The lowest order 10 bits (1024 channels) were assigned to the linear pulses from the Ge(Li) detector, the highest order 8 bits (256 channels) were assigned to the linear pulses from the NaI(Tl) detector. These address pairs defined the coordinates of a two dimensional (1024×256) matrix whose axes represents the energy of events observed in each detector. The resulting 18-bit words were stored sequentially in one half of the analyzer memory. After accumulation of 2048 events, the complete block was automatically dumped onto magnetic tape and then counting was resumed. The remaining half of the memory was used to display the coincidence projections for monitoring purposes. Some 1.8×10^6 coincidence events were recorded in the course of the ^{97}Ru experiment.

The coincidence data were sorted into the desired matrix at a later time with the university IBM 7040 computer. Since the size of the 1024 × 256 matrix was too large for the computer memory to process at one time, this matrix was divided into eight 128 × 256 submatrices. A computer program written by Dr. W.H. Fleming of the Computing Center was used

to check the data tape for events with addresses corresponding to each submatrix. For each such event, a count was added to the existing counts in the appropriate element of the submatrix. Eight passes through a data tape were required to extract all the information.

The spectrum in coincidence with a single gamma ray was obtained in the following manner. The total sum of all coincidence events was projected onto both the Ge(Li) and NaI(Tl) axes. A number of adjacent channels were selected from one of the projections; such a region was called a gate. The spectra in coincidence with the channels composing each gate were obtained from the sorted matrix and then summed to form a resultant coincidence spectrum. In this way, gates of variable width at any energy were selected from the matrix of coincidence data. Background, which was due to coincidences with the Compton distribution of higher energy gamma rays underlying the peak, was removed as follows. Gates adjacent to the peak and containing an equal number of channels were subtracted from the gate containing the peak. The "difference" spectra so obtained were a good approximation to the spectrum in coincidence with a single gamma ray.

(b) $^{99}\text{Mo} \rightarrow ^{99}\text{Tc}$

The experimental setup employed in the ^{99}Mo gamma-gamma coincidence measurements was similar to that used in the ^{97}Ru measurements (see Fig. 2.5). However, a 40 cm³ coaxial Ge(Li) detector replaced the 12 cm³ Ge(Li) detector and the Nuclear

Data 3300 system of the Tandem Accelerator Laboratory was used as the analyzer instead of the Nuclear Data 160-161 system. Use of the coaxial detector improved the timing resolution and a resolving time (FWHM) of 80ns was attained. As before, true + chance and chance windows (150ns wide in this case) were determined by two SCA's; the true to chance rate was monitored by the two scalers. A second output from each SCA was fed to a summing junction consisting of two diodes with their cathodes soldered together; the resultant output was then fed to the coincidence inputs of the ADC's and served as the enable pulse to open the linear gates.

The 16-band digital band selector of the 3300 analyzing system was employed for the two-parameter coincidence experiments. This feature of the analyzer allowed up to 16 separate energy restrictions (bands) to be placed on one parameter. The corresponding coincidence spectra from the other parameter are stored in separate portions of the memory. In the present experiments, eight energy windows (~ 30 keV wide) were established on the NaI(Tl) axis. The resulting Ge(Li) spectra were stored in eight 1024-channel subsections of the 16K memory. The corresponding chance spectra were routed into the remaining eight subsections. This procedure allowed subtraction of the chance contribution in the prompt coincidence peak. No further processing of the data was necessary. Coincidences with the Compton distribution of higher energy gamma rays underlying the desired peak (peaks) in the (NaI(Tl))

window were not subtracted. Their contribution to the gate was estimated by comparison with the neighbouring spectra.

(c) $^{101}\text{Mo} \rightarrow ^{101}\text{Tc}$

Coincidence events from the decay of ^{101}Mo and ^{101}Tc were observed with two true-coaxial Ge(Li) detectors having active volumes of 36 cm^3 and 40 cm^3 . The detectors were oriented at 135° and an anti-Compton shield of lead and cadmium, 2 cm thick, was placed between them in order to minimize inter-crystal scattering. The electronics are displayed in Fig. 2.6. In order to take advantage of the good timing properties of the coaxial Ge(Li) detectors, a constant fraction timing discriminator (CFD-Ortec model 453) was used to obtain the timing signal from one of the Ge(Li) detectors. The timing filter amplifier (Ortec model 454) was used to enhance the signal-to-noise ratio by shaping the preamplifier output pulse; this procedure improved the performance of the CFD. Unfortunately, since only one CFD was available, conventional leading edge timing had to be used for the other detector. However, there was considerable improvement over the circuit shown in Fig. 2.5 and a resolving time (FWHM) of 50ns was obtained.

Since the longer propagation time of the timing signal through the DDL and TSCA modules resulted in the TSCA signal being delayed relative to the CFD signal, the latter was used as the start signal for the TAC, the former as the stop.

Two windows, each 120nsecs wide and corresponding to portions of the time spectrum on and off the prompt coincidence peak, were set on the TAC output by two TSCA's. Outputs from these TSCA's were sent via the OR circuit to a logic shaper and delay (LSAD); this latter unit provided the coincidence enable pulse for the two ADC's. The output of the TSCA set off the prompt peak was fed to a second LSAD which provided an external routing pulse for one of the special purpose inputs of the ADC's. The delays of the two LSAD's were adjusted to make the signals coincide in time. A chance event was designated by the appearance of a logic pulse at both the special purpose input and the coincidence inputs. In this way a chance spectrum was generated for each coincidence gate. The chance contribution in the prompt coincidence peak was then removed from each gate at a later time.

Due to the short half-life of ^{101}Mo , it was necessary to use approximately 70 sources to obtain the coincidence data. The source strength was not constant from source to source and each source decayed during the counting period. Therefore, there was considerable variation in the true to chance rate over the course of the experiment. The overall true to chance ratio observed in the scalers was 9:1. Each source was counted for approximately 45 minutes, a time somewhat longer than the counting period for the singles experiments. This procedure allowed collection of ^{101}Tc coincidence events

without performing a separate experiment. The resolution of the Ge(Li) detectors in the coincidence experiments was somewhat poorer than in the singles spectra due to the higher singles rates required to obtain a satisfactory coincidence rate. Some resolution was also lost through use of the delay amplifiers (DA); these amplifiers delayed the arrival of the linear signals at the linear inputs of the ADC's until processing of the timing signals was complete. The detectors used in the coincidence experiments had been subjected to sustained fast neutron bombardment during accelerator experiments. Empirical evidence from the accelerator laboratory and from a study by Kraner et al. (1969) have shown that neutron damage causes considerable loss of resolution and pronounced low energy tailing (particularly at energies > 1 MeV). These features were quite noticeable in both the 36 cm^3 and 40 cm^3 Ge(Li) detectors.

The processing of the logic and analog pulses was handled by two Nuclear Data 4096-channel ADC's interfaced to a PDP-9 computer. In addition to being used to control the ADC's and the magnetic tape drive, the computer was employed for buffer storage of the incoming coincidence data and for live display of the coincidence projections along each axis for monitoring purposes. The information conveyed to the ADC's by the electronic circuitry was preserved by a computer program which set appropriate bits of the 18-bit word pairs used for storage of the coincidence results. The twelve lowest order

bits (bits 0-11) defined the address (channel number) of the detected gamma ray. Chance events were designated by setting a value of 1 for the chance bit (bit 12). The ADC was identified with bit 14; the X ADC was 0 and the Y ADC was 1. Bits 13, 15 and 16 were not used; the 15th and 16th bits could be used to define other parameters or ADC's if more than two were required. The highest order bit (bit 17) was the time sequence bit which was set to the same value for each component of the coincidence event. Related pairs of events were alternately designated by 0,0:1,1:0,0 etc.

The data were stored in the computer as pairs of 18-bit words; 500 of these pairs constituted a full buffer. When full, the buffer was emptied automatically onto magnetic tape. The data tapes were processed at a later time with the McMaster CDC 6400 computer. In all, 6×10^6 pairs of coincidence events were recorded.

Although the two-parameter coincidence experiments were performed in a 4096×4096 channel configuration, the coincidence projections were displayed using a 1024-channel conversion gain. If the rest of the analysis were to be carried out at this gain, the coincidence projections were obtained directly from the PDP-9. If some other configuration (eg. 1024×2048) were desired, a computer program was used to extract the projections from the data tapes with the CDC 6400 computer. The size of the coincidence matrix (1024×1024) prohibited construction of the coincidence matrix in the computer memory, which contained 124,000 octal locations. If the ^{101}Mo

coincidence data were analyzed in the same manner as the ^{97}Ru data, 16 or more passes through each data tape would have been required to obtain the full 1024×1024 channel matrix.

Since considerable computer time would have been consumed by the above method, an alternative means was devised. Instead of obtaining the entire coincidence matrix, only those portions of the matrix that contained peaks and their associated background were sorted into coincidence spectra. This meant that the coincidence gates had to be chosen prior to the computer analysis. Considerable care had to be taken in selecting the gates from the coincidence projection, since they could not be changed at a later date. Normally, background windows of equal width were selected from regions of the projection above and below the peaks. However, some difficulties arose when multiplets were considered. In many cases it was possible to set equal width windows over each of the components of the multiplet. A common background was then subtracted from each component. Any contribution from the adjacent peaks was estimated by comparing the background subtracted spectra. Unfortunately, it was often necessary to have windows of different widths when unresolved multiplets were considered; therefore, it was not possible to establish a common background and a separate background had to be chosen for each component of the multiplet. A selection of peak coincidence gates for ^{101}Mo is shown in Fig. 2.7.

The memory available in the CDC 6400 computer allowed

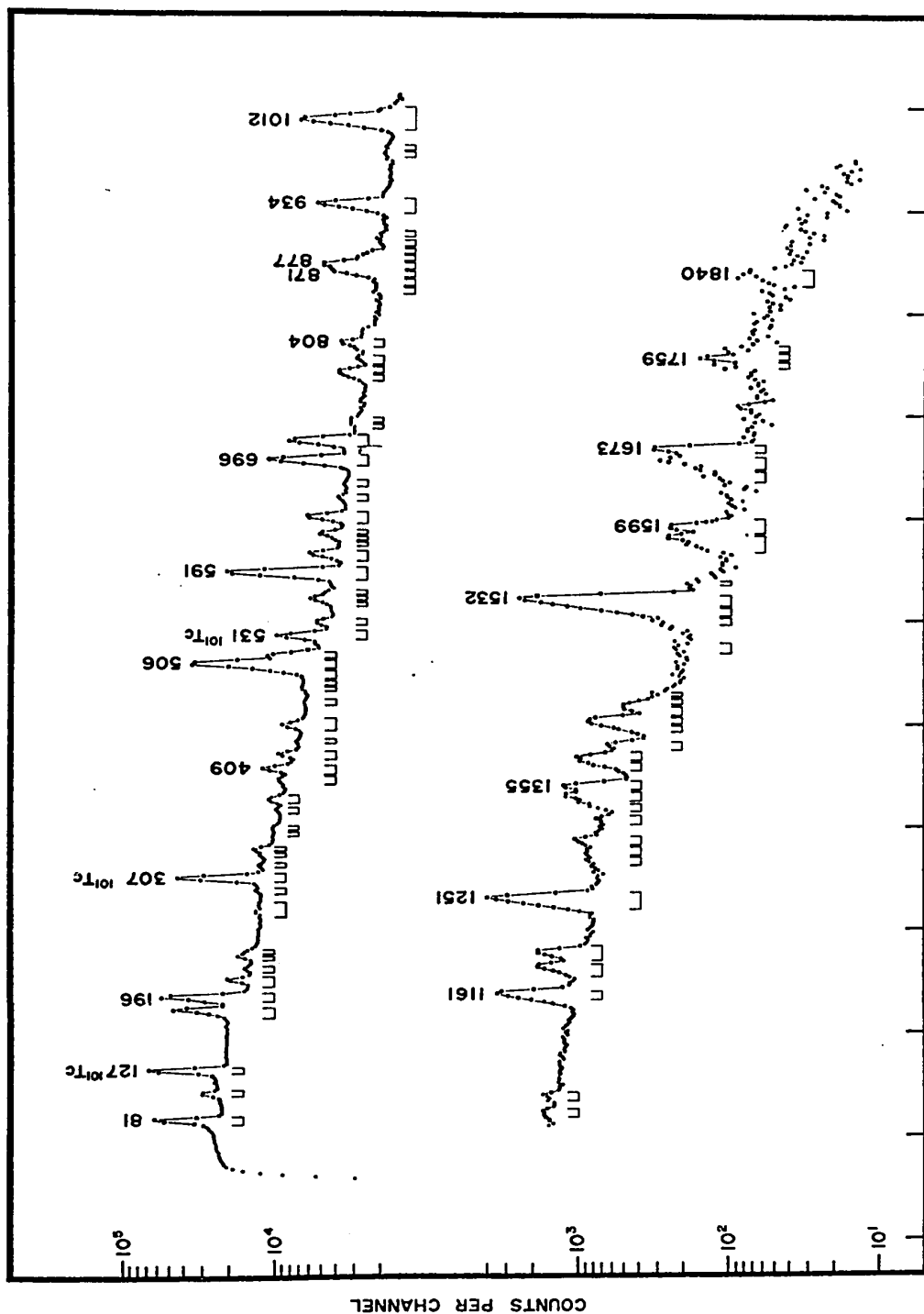


Fig. 2.7 Coincidence projection for the 36 cm³ detector showing the gates used in the analysis

storage of approximately 120 1024-channel coincidence spectra. The exact number depended on the memory space required by the computer program. Half of the storage portion of the memory was allocated to true + chance spectra, the other half to chance spectra. Thus, the maximum number of gates (windows) that could be processed during a single pass through the data was 60. Since the present experiments required more than 200 windows, at least four passes were needed. In order to save computer time, all of the coincidence data were read from tape onto disc. Since the access time of the disc was much faster than that of the tape drive, extra passes through the data required much less time.

The address-recorded coincidence data were sorted by a program written in FORTRAN and COMPASS (assembly language) by the author. The data were read from tape in buffer-sized increments (1000 18-bit PDP-9 words); each word was then subjected to several preliminary tests before being written onto disc. Bit 14 was checked to see that each event from ADC 1 (called the X ADC) was accompanied by a corresponding event from ADC 2 (called the Y ADC) and that the events were written in the proper order (X followed by Y). A second test on the time sequence bit (bit 17) confirmed that the X and Y address pair belonged to the same coincidence event. After all the data had been processed in this way, the data were read from disc under control of the second part of the program. Bit 13 of each address pair was inspected to determine whether the

coincidence event was from the true + chance or chance portion of the time spectrum. Then, the X(or Y) member of each address pair was examined in order to distribute the events among the preselected windows. This was done by comparing the addresses of the starting and stopping channels of each window with the X (or Y) address being considered. The data were stored in the computer memory as follows. If the address of the event corresponded to one of the preselected windows, the X address was set equal to the window number (i.e. a number between 1 and 60). This address designated the 1024-channel spectrum in which the coincidence event belonged. If the event was also a chance event, the X address was incremented by an amount equal to the total number of windows. This procedure allowed separate storage of the true + chance and chance events. If more than 60 windows were required, a new set of windows was read in and the whole procedure was repeated for each coincidence event.

The sorted data were printed out and were written onto tape after completion of the analysis. Another program was used to subtract the chance coincidence spectra from the true + chance coincidence spectra. The chance corrected spectra were printed out and then were written on another tape. This tape was the input for a third program which subtracted the background coincidence spectra from the peak coincidence spectra and then printed and plotted the resulting spectra. Due to statistical fluctuations in both the peak and background

coincidence spectra, "peaks" were sometimes generated in the background corrected spectra by the subtraction process. Such "peaks" were not present in the uncorrected spectra. Therefore, peaks observed in the background corrected spectra were designated as real only if they also appeared in the uncorrected spectra.

2.9.4 Interpretation of the Coincidence Data

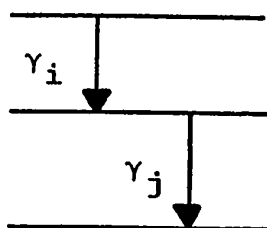
(a) Coincidence Probabilities

The areas of the peaks observed in the coincidence spectra can be used to obtain the coincidence probabilities between the gamma rays observed in the gated spectrum and the gating transition. The number of coincidence events between the full energy peaks of γ_i and γ_j is given by

$$N_{ij} = N_0 (\epsilon\omega)_i (\epsilon\omega)_j C_{ij}$$

where $(\epsilon\omega)_i$ and $(\epsilon\omega)_j$ are the corresponding efficiencies and C_{ij} is the coincidence probability of interest. The source strength N_0 was determined by normalizing to a strong cascade for which the coincidence probability was assumed to be correct. Normally, the efficiencies were determined by collecting a singles spectra for each detector in the coincidence geometry and using the measured singles intensities. In the present experiments, the efficiency curves were determined from the chance projections.

The experimental coincidence probabilities were compared with those expected from the decay scheme to test its



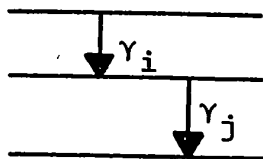
validity. The accompanying diagram shows a direct cascade involving γ_i and γ_j . The coincidence probability is defined by $C_{ij} = C_{ji} = I_i b_j$ where I_i is the intensity of the populating transition, γ_i , and b_j is the fraction of this intensity proceeding through the deexciting transition, γ_j . This fraction is equal to the intensity of γ_i divided by the total intensity (including internal conversion) deexciting the level.

In addition to the peaks corresponding to gamma rays in coincidence with the gating transition, a number of other peaks may appear in a coincidence spectrum. These peaks may be due to chance coincidences, coincidence summing or to Compton-backscatter coincidences. Under the conditions of the present experiments, these effects were not troublesome. The peaks due to coincidences with the Compton background from higher energy gamma rays underlying the gating transition have already been discussed in the previous section.

(b) Treatment of Chance

The equipment used to process the coincidence pulses has a finite resolving time. A prompt coincidence peak of finite width (50-90ns in the present experiments) superimposed on a continuous background is the time spectrum observed. The true coincidences are those in the peak, the chance coincidences are those in the continuous background. The peak-to-background ratio will depend on both the resolving time of the electronics and the source strength.

Consider the simple cascade shown below If γ_i is detected



in one detector and γ_j in the other, the counting rates will be given by

$$N_1 = N_0 (\epsilon\omega)_i I_i \quad N_2 = N_0 (\epsilon\omega)_j I_j$$

where N_0 is the source strength, $(\epsilon\omega)_i$ and $(\epsilon\omega)_j$ are the detector efficiencies, and I_i and I_j are the intensities of γ_i and γ_j . The true gamma-gamma coincidence rate is given by

$$N_{ij} = N_0 (\epsilon\omega)_i (\epsilon\omega)_j \epsilon_c C_{ij}$$

where ϵ_c is the coincidence efficiency and C_{ij} is probability of γ_i being in coincidence with γ_j . The chance rate for a resolving time of τ seconds is

$$N_{CH} = 2\tau N_1 N_2$$

The ratio of true to chance coincidences will therefore be

$$N_{ij}/N_{CH} = \epsilon_c / (2\tau N_0)$$

While this expression is only strictly valid for the simple case discussed, the true to chance ratio is always proportional to the product of τ and N_0 . Thus, for a given τ , there is an upper limit to the value of N_0 which may be used. If the window on the prompt peak is made too narrow, not all of the detected coincidence events will be recorded and a loss of coincidence efficiency will result. If N_0 is made too small, the low number of coincidence events will make the experiment impractical to do.

The coincidence experiments described in section 2.9.3 selected two "equal" portions of the time spectrum, one corresponding to true + chance coincidences, the other to

chance coincidences. In this way, it was possible to correct for the chance contribution in the prompt coincidence peak. However, the accuracy of such compensation is limited by the precision with which two equal windows can be selected and by the stability of the electronics. This sets the practical upper limit on the maximum source strength. However, long before this limit was reached in the present experiments, the detectors began losing resolution because of counting rate effects.

(c) Coincidence Summing

Another effect occurs whenever strong triple cascades are present. Since all three gamma rays are in coincidence, there is a probability that two of them may be simultaneously detected in one of the detectors at the same time the third is detected in the other detector. The presence of sum peaks can cause substantial difficulties if there is uncertainty about the existence of a true crossover transition. If the sum contribution is not removed, the measured intensity of the crossover transition will be too high; moreover, failure to consider coincidence summing may lead to an incorrect crossover assignment. The contribution of the sum peak is given by

$$N_{ijk} = N_o (\epsilon\omega)_i (\epsilon\omega)_j (\epsilon\omega)_k C_{ijk}$$

where C_{ijk} is the coincidence probability for the triple cascade.

(d) Compton-Backscatter Coincidences

Another source of spurious peaks in the coincidence

spectra is Compton-backscatter coincidences. Many gamma rays incident upon a detector will interact by means of the Compton effect; that is to say, the incident gamma ray will inelastically scatter from an electron and deposit only part of its energy in the detector. The remaining energy will be associated with the scattered gamma ray. Some of these gamma rays will be scattered at angles greater than 90° (called backscattering). When two detectors are brought close together in a coincidence geometry, there is a high probability that the second detector will absorb gamma rays backscattered from the first detector. This is particularly true when the source is between two detectors oriented at 180° to each other. A peak at the Compton energy will occur in one detector and a peak at the backscatter energy will occur in the other. Since the time involved in the scattering process is $\ll$ the resolving time of the coincidence electronics, there appears to be a direct coincidence between two gamma rays. This effect can be virtually eliminated by placing an anti-Compton shield between the two detectors. Backscattered gamma rays from one detector are prevented from entering the second detector. Of course, the thickness of the shield (normally made of lead or cadmium) will depend on the energy of the gamma rays involved.

CHAPTER III

EXPERIMENTAL RESULTS

3.1 Levels in ^{97}Tc Populated in the Decay of ^{97}Ru

3.1.1 Introduction

The electron-capture decay of 2.9- day ^{97}Ru was first studied by Sullivan et al. (1946) and the first attempt to construct a level scheme for ^{97}Tc was reported by Cork et al. (1955). The existence of a long-lived isomeric state in ^{97}Tc ($T_{1/2}=90$ days) was established by Unik and Rasmussen (1959). On the basis of conversion-electron ratios, the 96.5 ± 0.1 keV isomeric transition was assigned an M4 multipolarity. NaI(Tl) studies of the ^{97}Ru decay by Cretzu et al. (1966) led to the discovery of levels in ^{97}Tc at 216, 325, 786 and 995 keV. In a preliminary communication, Jha and Bond (1968) have reported results on the asymmetry of the 569-216, 460-324 and 645-324 keV gamma-gamma cascades. During the course of the present investigation, a further study by Graeffe (1969) was published. This work reported the population of levels at 215, 324, 784 and 969 keV. The measured values of the K-conversion coefficients of the 216 and 324 keV transitions were used in assigning J^π values of $7/2^+$ and $5/2^+$ respectively, for the 216 and 324 keV states. Nothing was known about the spins and parities of the other known levels at 784 and 969 keV.

The present investigation was undertaken as part of a program to study the level schemes of odd-mass technetium isotopes to obtain more accurate energies and intensities and to search for new levels in ^{97}Tc .

3.1.2 Gamma-Ray Singles Measurements

The low-energy region of the ^{97}Ru gamma-ray spectrum is shown in Fig. 3.1. The spectrum was obtained using a 0.15 cm^3 planar detector having a thin beryllium window; a 12.5 mm lucite disc, which normally serves as a beta stopper, was the only absorber. In addition to the K X-rays of technetium from electron capture and internal conversion, the 109, 216 and 324 keV transitions are seen. The weak peaks on the low energy side of the 216 and 324 keV lines are germanium X-ray escape peaks (Ungrin and Johns, 1969).

Fig. 3.2 shows the gamma-ray spectrum above 120 keV taken with the 12 cm^3 detector. In addition to the lucite beta stopper, a graded absorber of lead, cadmium and copper, 6 mm thick, was used to attenuate the strong 216 keV radiation. The 531, 560, 639 and 670 keV transitions have been observed for the first time. All had a half-life of the order of three days. The prominent 855 keV transition had tentatively been assigned to the ^{97}Ru decay by Graeffe (1969) on the basis of its half-life. The present work establishes its role in the decay without ambiguity. The major impurity in the spectrum is due to 40-day ^{103}Ru (see section 2.4). Other impurities arose mainly from room background (^{137}Cs , ^{228}Th) and from trace impurities in the source. They are labelled (c) in the figure.

The energies and relative intensities derived from the present work are tabulated in Table 3.1. The results of Graeffe (1969), Jha and Bond (1968), and those of Hofstetter and

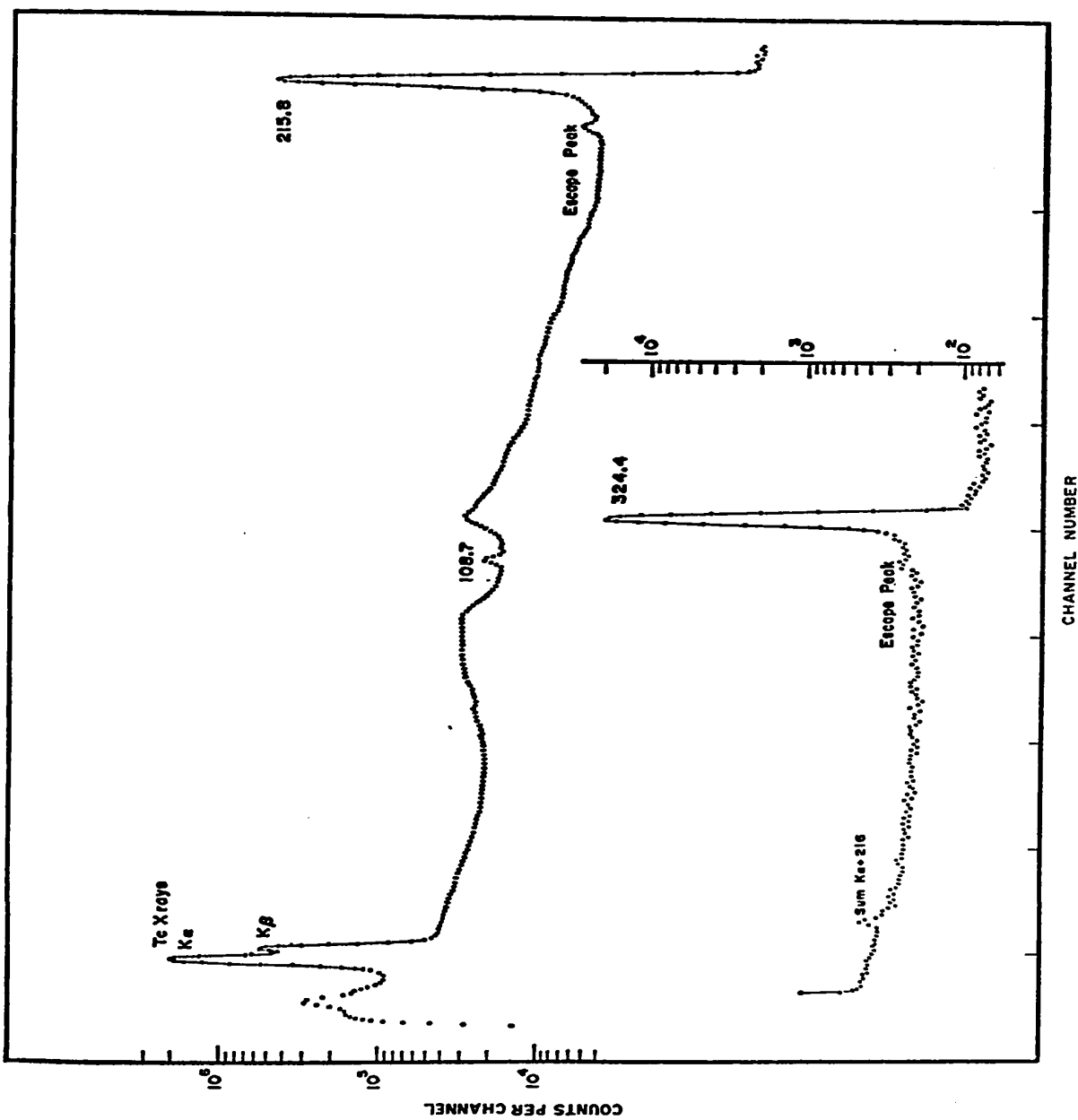


Fig. 3.1 The low-energy gamma-ray spectrum of ^{97}Ru

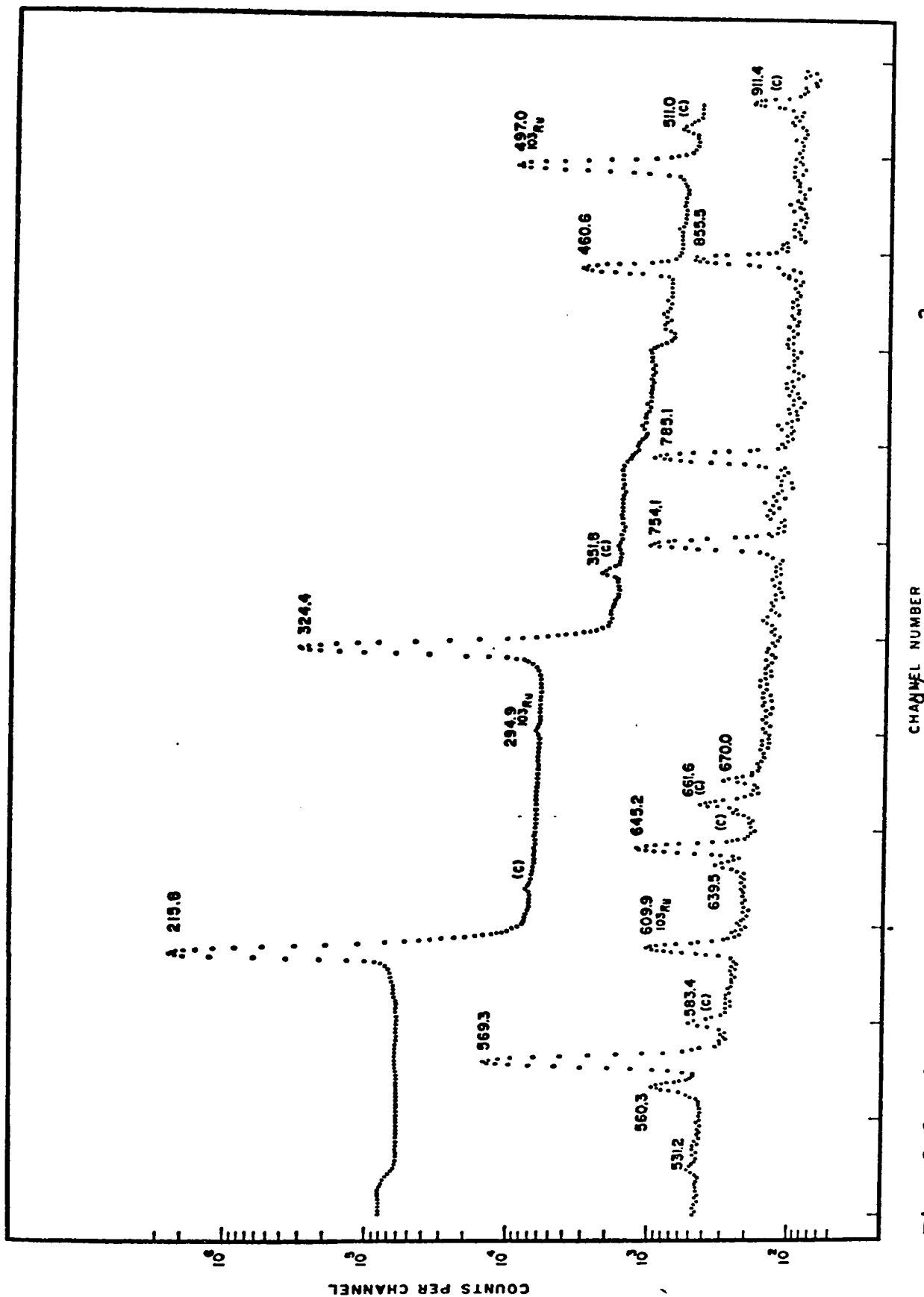


Fig. 3.2 The gamma-ray spectrum of ¹⁰³Ru recorded with a 12 cm³ Ge(Li) detector shielded with a graded absorber

TABLE 3.1
Transitions in ^{97}Tc

Graeffe (1969)		Jha and Bond (1968)		Present Work Cook et al. (1970)		Classification	Hofstetter and Sugihara (1970)		Phelps and Sarantites (1971)	
Energy (keV)	Intensity	Intensity	Energy (keV)	Energy (keV)	Intensity		Energy (keV)	Intensity $\leq 10\%$	Energy (keV)	Intensity
Tc X-rays 70 ± 7										
108.8 $\pm$ 0.5	0.085 $\pm$ 0.017		108.7 $\pm$ 0.15	0.15 $\pm$ 0.015	324.45 $\rightarrow$ 215.76	85.4	108.9 $\pm$ 0.3	0.14	108.80 $\pm$ 0.08	0.11 $\pm$ 0.009
115.2 $\pm$ 0.4	85.4*	85.4*	215.76 $\pm$ 0.08	85.4	(969.73-855.41)		215.2 $\pm$ 0.3	85.4*	114.5 $\pm$ 0.2	0.008 $\pm$ 0.004
324.2 $\pm$ 0.4	10.1 $\pm$ 0.7	9.7	324.45 $\pm$ 0.08	11.1 $\pm$ 0.6	215.76 $\rightarrow$		324.4 $\pm$ 0.3	10.1	215.71 $\pm$ 0.05	85.4*
160.2 $\pm$ 0.5	0.10 $\pm$ 0.02	0.12	460.6 $\pm$ 0.15	0.13 $\pm$ 0.015	324.45 $\rightarrow$		461.3 $\pm$ 0.3	0.11	324.48 $\pm$ 0.05	10.8 $\pm$ 0.2
					785.07 $\rightarrow$ 324.45				460.59 $\pm$ 0.09	0.13 $\pm$ 0.009
					(574.2 $\rightarrow$ 96.5)				477.7 $\pm$ 0.6	0.0023 $\pm$ 0.0011
			531.2 $\pm$ 0.5	0.0043 $\pm$ 0.0009	855.41 $\rightarrow$ 324.45				531.1 $\pm$ 0.3	0.0032 $\pm$ 0.0007
168.9 $\pm$ 0.5	0.85 $\pm$ 0.09	0.85	560.3 $\pm$ 0.2	0.031 $\pm$ 0.004	656.7 $\rightarrow$ 96.5		560.6 $\pm$ 0.3	0.035	560.32 $\pm$ 0.12	0.041 $\pm$ 0.004
			569.3 $\pm$ 0.1	0.91 $\pm$ 0.06	785.07 $\rightarrow$ 215.76		570.0 $\pm$ 0.3	0.75	569.33 $\pm$ 0.07	0.87 $\pm$ 0.02
144.4 $\pm$ 0.5	0.060 $\pm$ 0.009	0.06	639.5 $\pm$ 0.2	0.0081 $\pm$ 0.0009	855.41 $\rightarrow$ 215.76		641.5 $\pm$ 1.0	0.0068	639.73 $\pm$ 0.12	0.0085 $\pm$ 0.0017
			645.2 $\pm$ 0.15	0.073 $\pm$ 0.0075	969.73 $\rightarrow$ 324.45		646.9 $\pm$ 0.3	0.047	645.35 $\pm$ 0.08	0.058 $\pm$ 0.003
			670.0 $\pm$ 0.3	0.010 $\pm$ 0.002	994.4 $\rightarrow$ 324.45				670.30 $\pm$ 0.10	0.0085 $\pm$ 0.009
153.8 $\pm$ 0.7	0.068 $\pm$ 0.009	0.10	754.05 $\pm$ 0.15	0.079 $\pm$ 0.0075	969.73 $\rightarrow$ 215.76		754.2 $\pm$ 0.3	0.060	754.03 $\pm$ 0.06	0.074 $\pm$ 0.002
184.7 $\pm$ 0.7	0.085 $\pm$ 0.017	0.10	785.1 $\pm$ 0.15	0.080 $\pm$ 0.0075	785.07 $\rightarrow$		785.5 $\pm$ 0.3	0.079	785.04 $\pm$ 0.06	0.070 $\pm$ 0.002
					(946.5 $\rightarrow$ 96.5)				850.1 $\pm$ 0.4	0.0014 $\pm$ 0.0011
155.1 $\pm$ 0.8	0.043 $\pm$ 0.017	0.05	855.5 $\pm$ 0.2	0.047 $\pm$ 0.005	855.41 $\rightarrow$		856.3 $\pm$ 0.3	0.032	855.45 $\pm$ 0.08	0.042 $\pm$ 0.009
					(994.4 $\rightarrow$ 96.5)				898.0 $\pm$ 0.4	0.0009 $\pm$ 0.0009
			970	<0.005	(969.73 $\rightarrow$ 0		971 $\pm$ 1	0.004	970	<0.0005

* Normalized to this value for comparison purposes

Sugihara (1970) and Phelps and Sarantites (1971), which have been published since the present work was completed, are also included in Table 3.1.

3.1.3 Gamma-Gamma Coincidence Results

Figs. 3.3 and 3.4 display the Ge(Li) spectra, taken with the 12 cm³ detector, in coincidence with NaI(Tl) gates containing the 216 and 324 keV transitions, respectively. The background due to the Compton distributions of other gamma rays has been subtracted in both cases. The chance contribution, which was not subtracted, is responsible for the 216 and 324 keV peaks in these spectra. The relevant portions of the ⁹⁷Tc level scheme are shown as inserts in both figures. The presence of the 109, 569 and 754 keV peaks in Fig. 3.3 verify known levels at 324, 785 and 969 keV. The existence of the 639 keV transition, which was also seen in singles, establishes a new level at 855 keV. The 460 and 645 keV lines in Fig. 3.4 confirm the levels at 785 and 969 keV, while the 531 keV line gives further support to the 855 keV level mentioned above. The 855 keV level also accomodates the 855 keV gamma ray as a ground-state transition. The coincidence relation between the 324 and 670 keV gamma rays defines a new level at 994 keV.

3.1.4 The Level Scheme of ⁹⁷Tc

The singles and coincidence measurements of the preceding sections lead to the decay scheme of Fig. 3.5. It accomodates all 13 of the gamma-ray transitions observed. Table 3.2 gives a summary of the proposed level energies, the percent electron

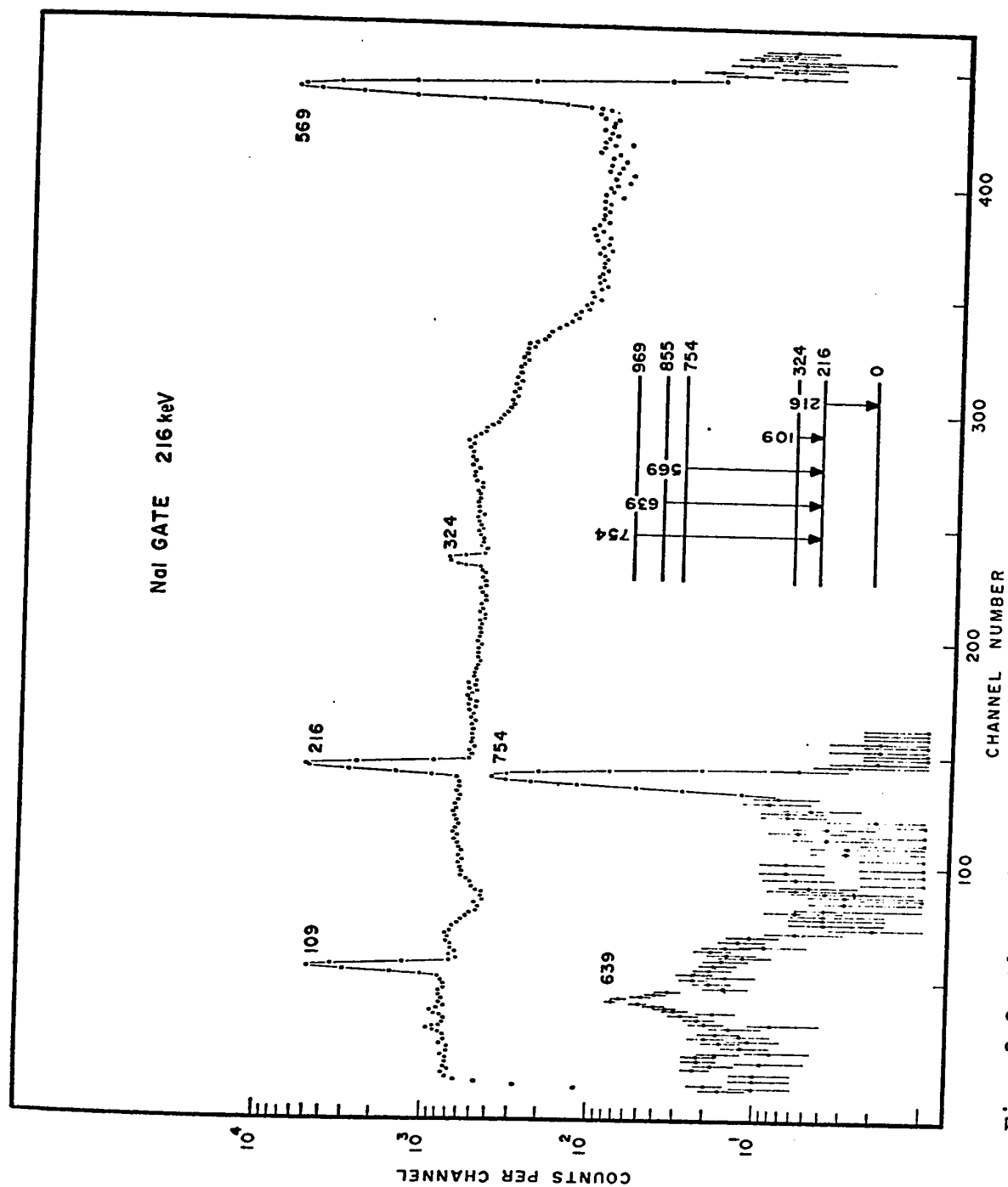


Fig. 3.3 The Ge(Li) spectrum in coincidence with a gate set on a 216 keV photopeak in the NaI(Tl) detector

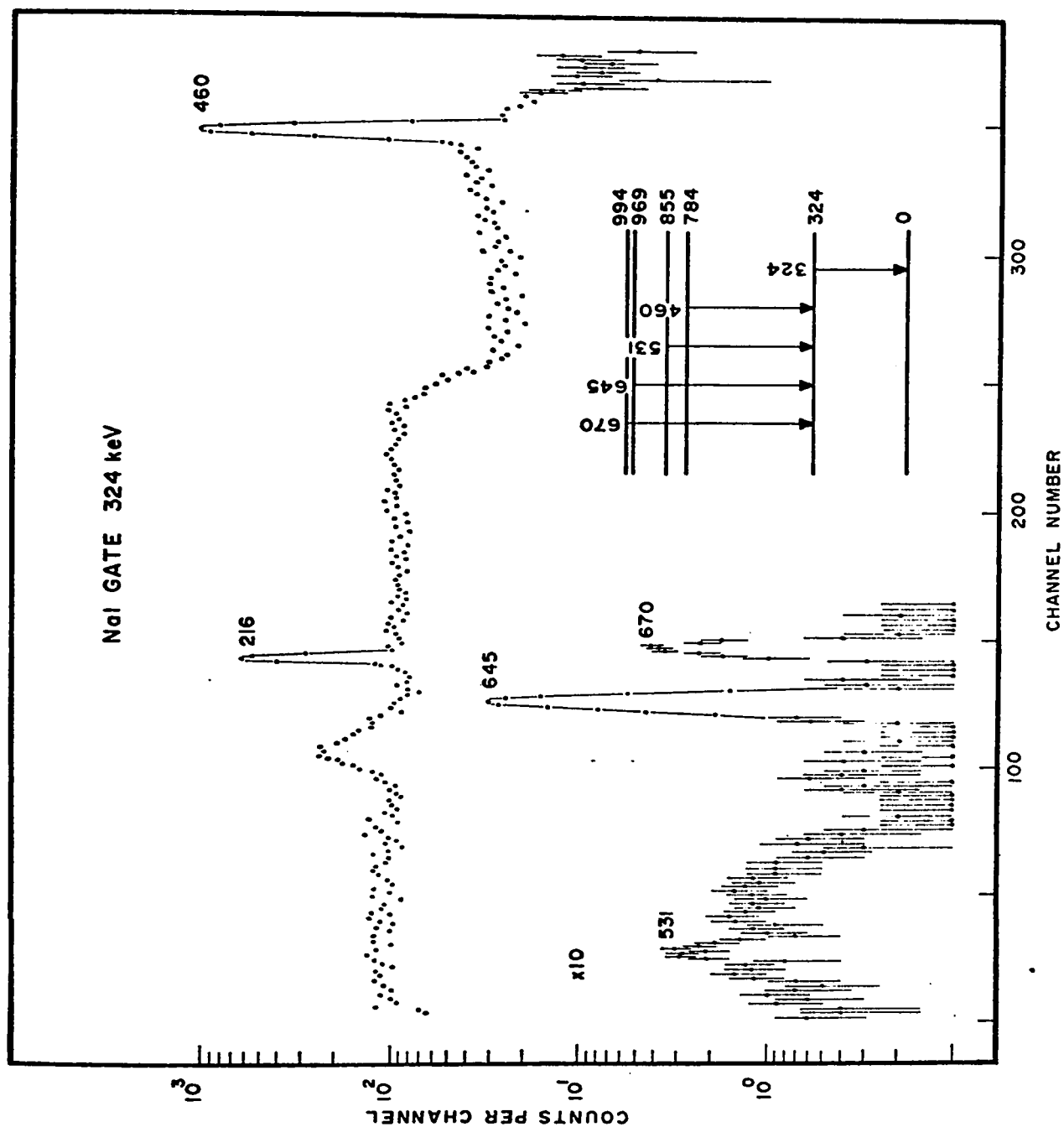
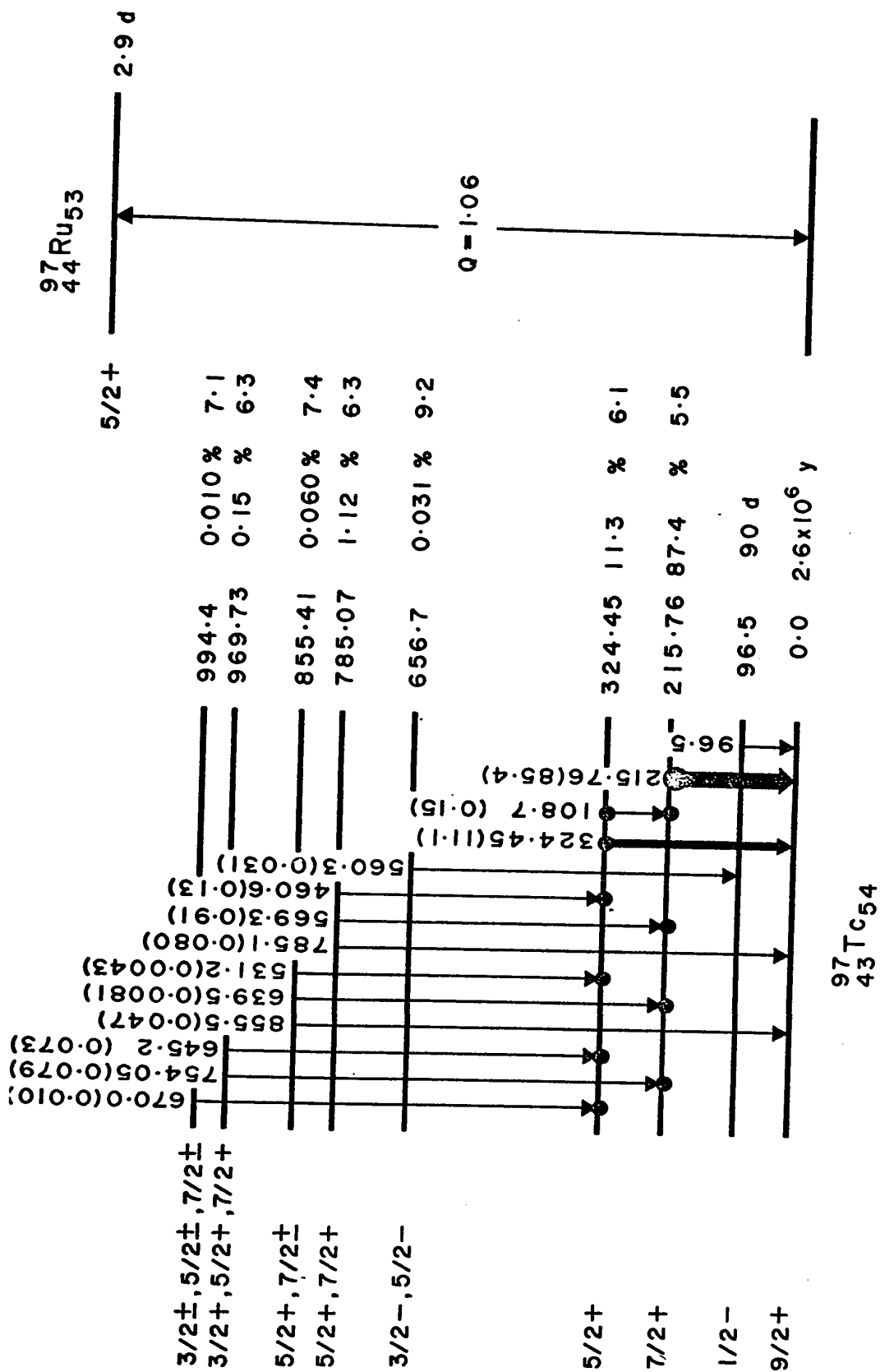


Fig. 3.4 The Ge(Li) spectrum in coincidence with a gate set on a 324 keV photopeak in the NaI(Tl) detector

Fig. 3.5 The level scheme of ^{97}Tc

capture populations, the deduced $\log f_0 t$ values and the J^π values assigned for each proposed level. The table also includes the level energies obtained from the reaction work of Kim et al. (1970) which is discussed below. The $\log f_0 t$ values listed in Cook and Johns (1970) were calculated using an estimate of 1200 keV for the decay energy (Mattauch et al. 1965), and on the assumption that there is no electron capture to the ground state. A value of 1060 keV is obtained from a simple average of the calculated values of Garvey et al. (1969), Mattauch et al. (1965) and Everling (1962). This latter value has been used in the thesis for the calculation of the $\log f_0 t$ values. The nomograms of Lederer et al. (1967) were used to make the calculations. It should be pointed out that the $\log f_0 t$ values for the levels above 800 keV are particularly sensitive to possible errors in the estimated decay energy. An indication of the $\log f_0 t$ variation is shown in the table.

The ground state spin and parity of the parent nucleus ^{97}Ru can be assumed to be $5/2^+$ in analogy with the $5/2^+$ assignment for the ^{95}Ru , ^{99}Ru and ^{101}Ru ground states. The spin and parity of ^{95}Ru was established with the $^{96}\text{Ru}(p,d)^{95}\text{Ru}$ experiments of Ball (1971). The spins of ^{99}Ru and ^{101}Ru have been determined directly by electron-spin resonance measurements (Lindgren 1964).

The ground state and the 96.5 keV isomeric state. The nucleus ^{97}Tc has 43 protons. The simple shell model predicts that the $2p_{1/2}$ subshell should be filled ($Z=40$). The last

TABLE 3.2
Assignment of $\log f_{\text{O}t}$ and J^{π} Values to Levels of ^{97}Tc

Level Energy (keV)	% Electron Capture	Q=1200	$\log f_{\text{O}t}$ Q=1100	Q=1060	J^{π}	Level Energy(keV) Kim et al. (1970)
0					$9/2^{+}$	0
96.5					$1/2^{-}$	101 ± 10
215.76 ± 0.08	87.4 ± 2.0	5.6	5.5	5.45	$7/2^{+}$	213 ± 10
324.45 ± 0.08	11.3 ± 0.6	6.45	6.2	6.1	$5/2^{+}$	328 ± 10
656.7 ± 0.25	0.031 ± 0.003	9.5	9.3	9.2	$3/2^{-}, 5/2^{-}$	575 ± 10
785.07 ± 0.09	1.12 ± 0.06	6.7	6.4	6.3	$5/2^{+}, 7/2^{+}$	662 ± 9
855.41 ± 0.15	0.060 ± 0.005	7.8	7.55	7.4	$5/2^{+}, 7/2^{+}$	785 ± 9
969.73 ± 0.12	0.15 ± 0.01	7.05	6.6	6.3	$3/2^{+}, 5/2^{+}, 7/2^{+}$	857 ± 9
994.4 ± 0.3	0.010 ± 0.0015	8.1	7.5	7.1	$3/2^{+}, 5/2^{+}, 7/2^{+}$	941 ± 9
						962 ± 9
						987 ± 9

three protons will then be found in the $1g_{9/2}$ subshell; the ground state spin and parity of ^{97}Tc is therefore expected to be $9/2^+$. Excitation of a proton from the $2p_{1/2}$ subshell to the $1g_{9/2}$ subshell should then produce a $1/2^-$ state at a relatively low excitation energy. The M4 character of the 96.5 keV transition (Unik and Rasmussen 1959) suggests that the ground state and the 96.5 keV isomeric state can be identified with the $9/2^+$ and $1/2^-$ states predicted by the shell model. Unfortunately, the internal conversion measurement cannot determine which of the two states ($9/2^+$ or $1/2^-$) forms the ground state.

Since there exists no direct measurement of the ground-state spin of ^{97}Tc , the value must be inferred from beta decay data and nuclear systematics. The ground state of ^{97}Tc decays by electron capture to the $5/2^+$ ground state of ^{97}Mo with a half-life of $(2.6 \pm 0.4) \times 10^6$ years. The associated $\log f_0 t$ value of ~ 12 is based on an estimate of ~ 300 keV for the decay energy (Lederer et al. 1967). This high $\log f_0 t$ value suggests a $9/2^+ \rightarrow 5/2^+$ rather than a $1/2^- \rightarrow 5/2^+$ electron capture transition. Supporting evidence for this choice is obtained from the neighbouring nuclei ^{95}Tc and ^{99}Tc . The decay patterns of the ground and isomeric states of ^{95}Tc (Chilosi et al. 1969) indicate J^π values of $9/2^+$ and $1/2^-$, respectively. A direct measurement of the ground state spin of ^{99}Tc (Lindgren 1964) yields a $9/2^+$ value. The $1/2^-$ state is found at 142.68 keV (see section 3.2). The existing experimental evidence is therefore consistent with an assignment of $9/2^+$ and $1/2^-$ respectively, for the ground state and the 96.5 keV isomeric state.

The 215.76 and 324.45 keV levels. The allowed character of the electron-capture feeds ($\log f_{\text{oc}}$ values of 5.5 and 6.1, respectively) to these levels suggest J^{π} values of $3/2^{+}$, $5/2^{+}$ or $7/2^{+}$. The presence of a ground state transition from each level rules out the $3/2^{+}$ choice. The Si(Li) electron measurements of Graeffe (1969) indicate that the multipolarities of the 215.76 and 324.45 keV transitions are $M1+E2$ and $E2$, respectively. Therefore, the 215.76 and 324.45 keV levels have J^{π} values of $7/2^{+}$ and $5/2^{+}$, respectively. It should be noted that 98.7% of the total electron-capture intensity is associated with these two states.

The 656.7 keV level. Since the 560.3 keV gamma ray has been seen only in singles, it can feed either the ground state or the isomeric level at 96.5 keV. The latter assignment seems more probable since transitions to the 215.76 and 324.45 keV states have not been observed. In addition, two negative parity levels, which decay only to the $1/2^{-}$ isomeric state, exist at a similar excitation energy in the neighbouring nucleus ^{99}Tc (Cook et al. 1970). The high $\log f_{\text{oc}}$ value of 9.2 is characteristic of a first forbidden beta transition implying J^{π} values of $3/2^{-}$, $5/2^{-}$ or $7/2^{-}$ for the 656.7 keV level. If the 560.3 keV transition does populate the $1/2^{-}$ isomeric state, then the spin and parity must be $3/2^{-}$ or $5/2^{-}$.

The 785.07 keV level. The existence of the 569.3–215.76 and 460.6–324.45 keV gamma-ray cascades, together with a 785.1 keV ground-state transition, confirms the previously established

785 keV level. The $\log f_0 t$ value of 6.3 indicates allowed capture to the state. This fact plus the existence of the ground-state transition restricts the spin and parity of the 785.07 keV state to $5/2^+$ or $7/2^+$.

The 855.41 keV level. The new level at 855.41 keV is established by the presence of the 639.5-215.76 and 531.2-324.45 keV gamma-ray cascades and by the existence of the 855.5 keV transition to the ground state. The $\log f_0 t$ value of 7.4 is characteristic of either allowed hindered or first-forbidden electron capture. The $\log f_0 t$ value quoted by Cook et al. (1970) was 7.8 (based on a Q-value of 1200 keV) and therefore a first-forbidden assignment was favoured, restricting the J^π values for the 855.41 keV state to $3/2^-$, $5/2^-$ or $7/2^-$. The existence of a ground-state transition further restricted the J^π to $7/2^-$. The revised $\log f_0 t$ value also allows J^π of $5/2^+$ or $7/2^+$. Therefore, the spin and parity of the 855.41 keV state may be $5/2^+$ or $7/2^+$.

The 969.73 keV level. The existence of this level, which was previously suggested by earlier workers, is confirmed by the presence of the 754.05-215.76 and 645.2-324.45 keV cascades. The revised $\log f_0 t$ value of 6.3 for the electron-capture decay to this level is characteristic of an allowed transition, thus indicating J^π values of $3/2^+$, $5/2^+$ or $7/2^+$ for the level. If the ground-state transition exists, its intensity must be $<0.006\%$. In the absence of any knowledge of the mixing ratios of the transitions deexciting the 969.73 keV level, it is not possible to use the relative reduced transition rates to make further

The 994.4 keV level. This new level is established through the 670.0-324.45 keV gamma-ray cascade. The $\log f_0 t$ value of 7.1 for the transition to this state is characteristic of either allowed or first-forbidden electron capture. The spin of the 994.4 keV level can be 3/2, 5/2 or 7/2 and the parity can be positive or negative.

3.1.5 Subsequent Investigations

At about the same time that the present work appeared in print, an investigation of the decay of ^{97}Ru by Hofstetter and Sugihara (1970) was also published. Their energies and intensities have been included in Table 3.1. Although the agreement is normally quite good, there are some serious discrepancies between their results and the results of both the present work and those of Phelps and Sarantites (1971). In addition, they used the 560 keV transition to define a level at 560 keV rather than at 657 keV, as in the present work. The reaction work described below supports the latter choice. Their observation of a ground-state transition from the 970 keV level is not confirmed by other workers with greater sensitivity.

Kim et al. (1970) have investigated the level structure of ^{97}Tc by means of the $^{97}\text{Mo}(p,n)^{97}\text{Tc}$ reaction. For the spin 5/2 target, it is expected that final states with spins ranging from 1/2 to 9/2 should be populated with s-, p- and d- wave neutrons. Below 1 MeV, levels were found at 0, 101, 213, 328, 575, 662, 775, 857, 941, 962 and possibly 987 keV; the errors

in the energies were typically 10 keV. These levels have been included in Table 3.2 and are identified with corresponding levels from the present work. Since levels were observed at 575 and 662 keV, the 560 keV transition could deexcite either level and therefore a choice between the suggestions of the present work and that of Hofstetter and Sugihara (1970) cannot be made. The energy agreement would tend to support the assignment of the present work. Unfortunately, no gamma-ray spectra were recorded by Kim et al. (1970). Also, no spin information was obtained since no angular distributions were determined.

In a subsequent publication, Kim et al. (1971) have used the fact that isospin-forbidden neutron decay to low-lying states of proton induced analogue resonances in the mass-90 region preferentially populates those residual states which are fed by p-wave neutrons. The results of these measurements indicate that their states at 575 and 662 keV have J^π values of $3/2^-$ and $5/2^-$, respectively.

Phelps and Sarantites (1971) have reinvestigated the ^{97}Ru decay using a Ge(Li) detector in an anti-Compton arrangement. A Si(Li) detector was also used in order to measure the internal conversion spectrum. Their gamma-ray results have been included in Table 3.1; the agreement with the results of the present work is excellent. In addition to the levels observed in the present work, they identified two more levels at 574.2 and 946.5 keV. These probably correspond to the levels at 575

and 941 keV found by Kim et al. (1970). The 574.2 keV state was observed to be deexcited by a 477.7 keV transition to the 96.5 keV state; no transition to the ground state was found. The placement of the 560.3 keV gamma ray as a deexciting transition from the 656.7 keV level to the 96.5 keV level was confirmed. The K-conversion coefficients of the 460.6, 569.3, 645.2, 754.05, 785.1 and 855.5 keV transitions indicated M1 and/or E2 multipolarities. Therefore, the 785.07, 855.41 and 969.73 keV states must have positive parity. If the $7/2^-$ choice for the J^π value of the 855.41 keV level is eliminated, the J^π assignments of the present work are consistent with the data of Phelps and Sarantites (1971). The identification of a weak 898.0 keV transition from the 994.4 keV level to the $1/2^-$ state at 96.5 keV further restricts the J^π value of the former level to one of $3/2^\pm$ or $5/2^-$.

3.2 Levels in ^{99}Tc Populated in the Decay of ^{99}Mo

3.2.1 Introduction

The beta decay of 67-hour ^{99}Mo was first identified by Seaborg and Segrè (1939) who also showed that this activity was in equilibrium with a 6-hour isomeric state at approximately 140 keV in ^{99}Tc . Medicus et al. (1949, 1950) demonstrated that the ~ 140 keV level was a close doublet ($\Delta E \sim 1.8$ keV) with the higher energy member being the isomeric state, which decayed to the lower energy member and to the ground state. Bunker and Canada (1950) were among the first to attempt construction

of a level scheme for ^{99}Tc and established new levels at 180 and 920 keV in addition to the ~ 140 keV doublet. The doublet nature of the ~ 140 keV state was confirmed by Mihelich et al. (1951) with internal conversion measurements of $^{99\text{m}}\text{Tc}$. The multipolarities of the 142 and 140 keV transitions were found to be M4 and M1 respectively, from K/L electron intensity ratios; in addition, the branching ratio for the decay of the 142 keV level to the 140 keV level and to the ground state was estimated to be ~ 70 . Since the ground state of ^{99}Tc was $9/2^+$ (Kessler and Meggers 1950), J^π values of $1/2^-$ and $7/2^+$ respectively, were proposed for the 142 and 140 keV states. In addition to these early investigations, a number (Nuclear Data Sheets 1961) of other studies using scintillation and/or magnetic spectrometers has been published, the most recent being those of Cretzu et al. (1965), Crowther et al. (1965) and Dorikens-Vanpraet et al. (1967). These results and those discussed above are summarized in the decay scheme shown in Fig. 3.6.

While the present investigation was in progress, Van Eijk et al. (1969) published results obtained with Ge(Li) detectors, magnetic and crystal spectrometers, and a Ge(Li)-NaI(Tl) coincidence spectrometer. A number of previously unreported levels were established at 672, 761, 1004, 1105, 1169 and 1197 keV.

The low-energy portion of the internal conversion spectrum has been investigated very carefully in recent years (Geiger 1968; Van Eijk et al. 1969; Ranakumar et al. 1969; and Bashandy 1969). The previous results for the multipolarities

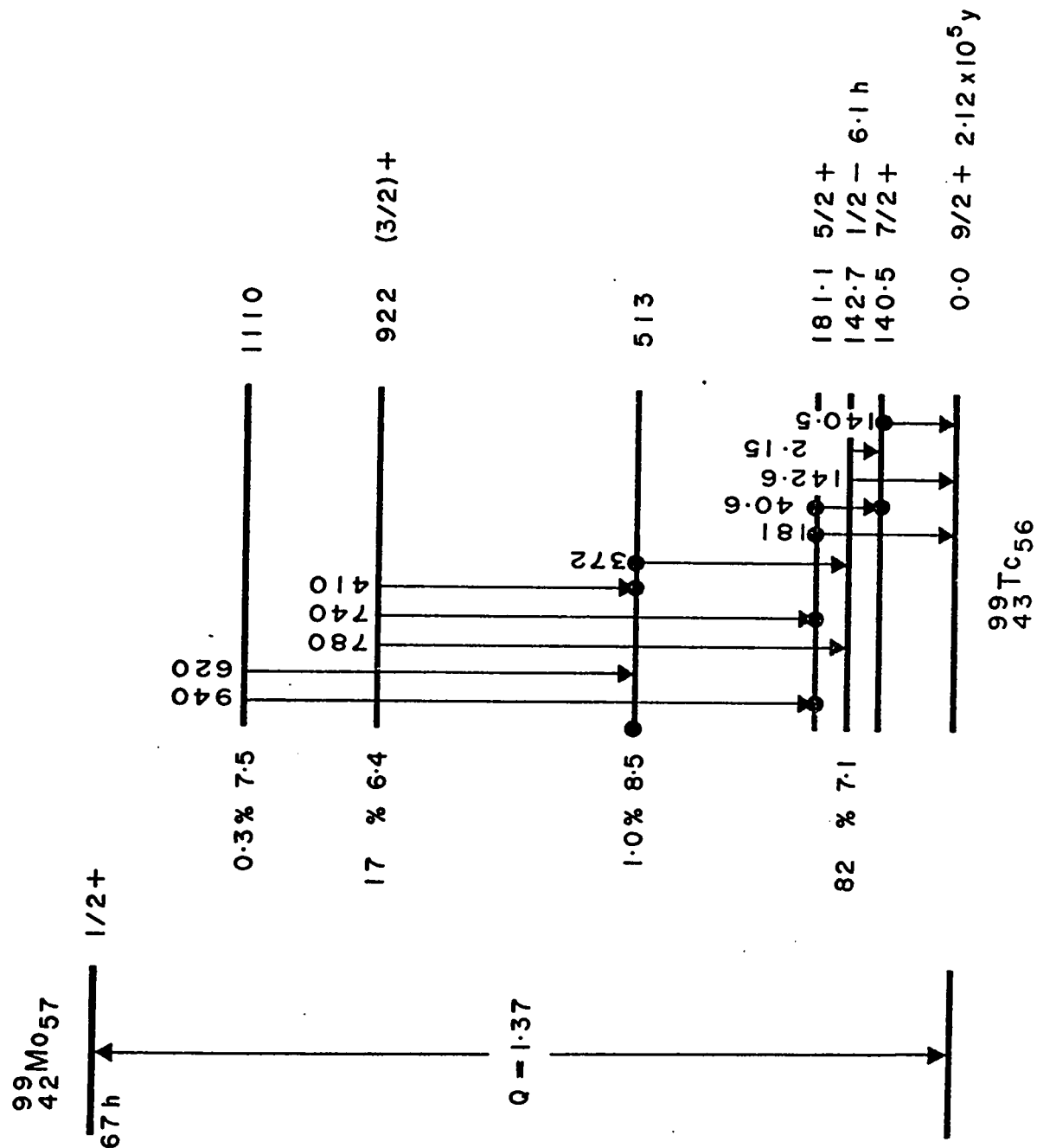


Fig. 3.6 The decay scheme of ^{99}Mo prior to the present investigation

of the 140 and 142 keV transitions were verified and M1 and E2 multipolarities respectively, were determined for the 40 and 181 keV transitions. The ratio of the intensities of the 142.63 and the 140.51 keV gamma rays and K-conversion lines have been found to be $(30 \pm 6) \times 10^{-6}$ (Van Eijk et al. 1969) and 0.0746 ± 0.0011 (Geiger 1968), respectively. The properties of the low-lying states up to 181 keV are therefore very well established.

The present investigation was undertaken as part of a program to study the level structure of odd-mass technetium nuclei. The principal aim was to establish the role of the weak transitions in the decay.

3.2.2 Gamma-Ray Singles Measurements

The ^{99}Mo gamma-ray spectrum for the energy range below 200 keV is presented in Fig. 3.7. The 0.15 cm^3 thin window Ge(Li) detector with a 12.5 mm lucite beta stopper, was used for the measurement. The intense 140 and 181 keV peaks dominate the spectrum. The other peaks correspond to the K_α and K_β X-rays of technetium from internal conversion and the 40.5 keV gamma ray. The weak peak below the 140 keV transition is the germanium X-ray escape peak (Ungrin and Johns 1969).

Fig. 3.8 displays the higher energy portion of the gamma-ray spectrum taken with the 12 cm^3 detector. A graded absorber of lead, cadmium and copper, 6 mm thick, was used to attenuate the strong 140 and 181 keV transitions. The above spectra are representative of a number of different runs in various geometries

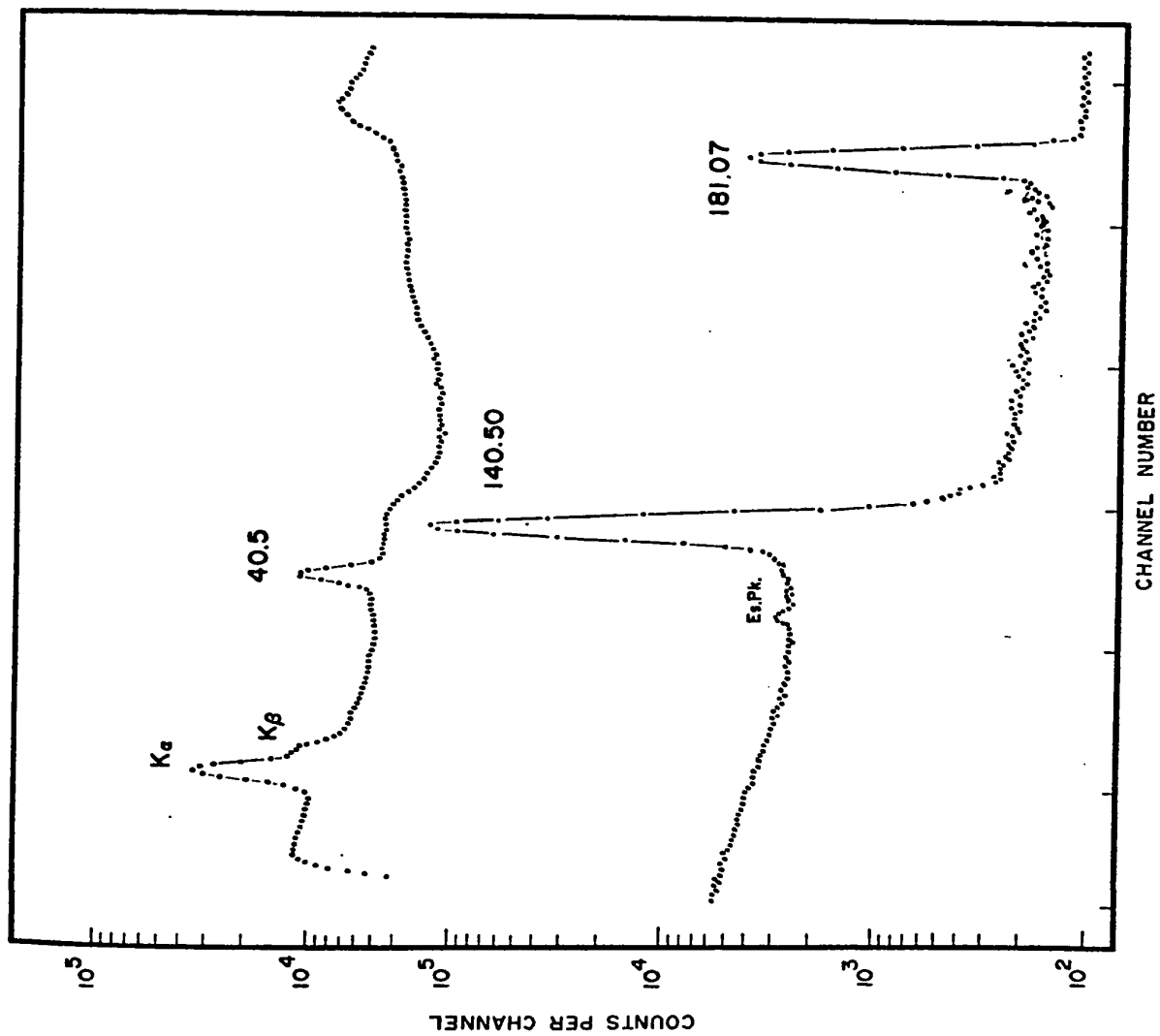


Fig. 3.7 The low-energy portion of the gamma-ray spectrum of ^{99}Mo

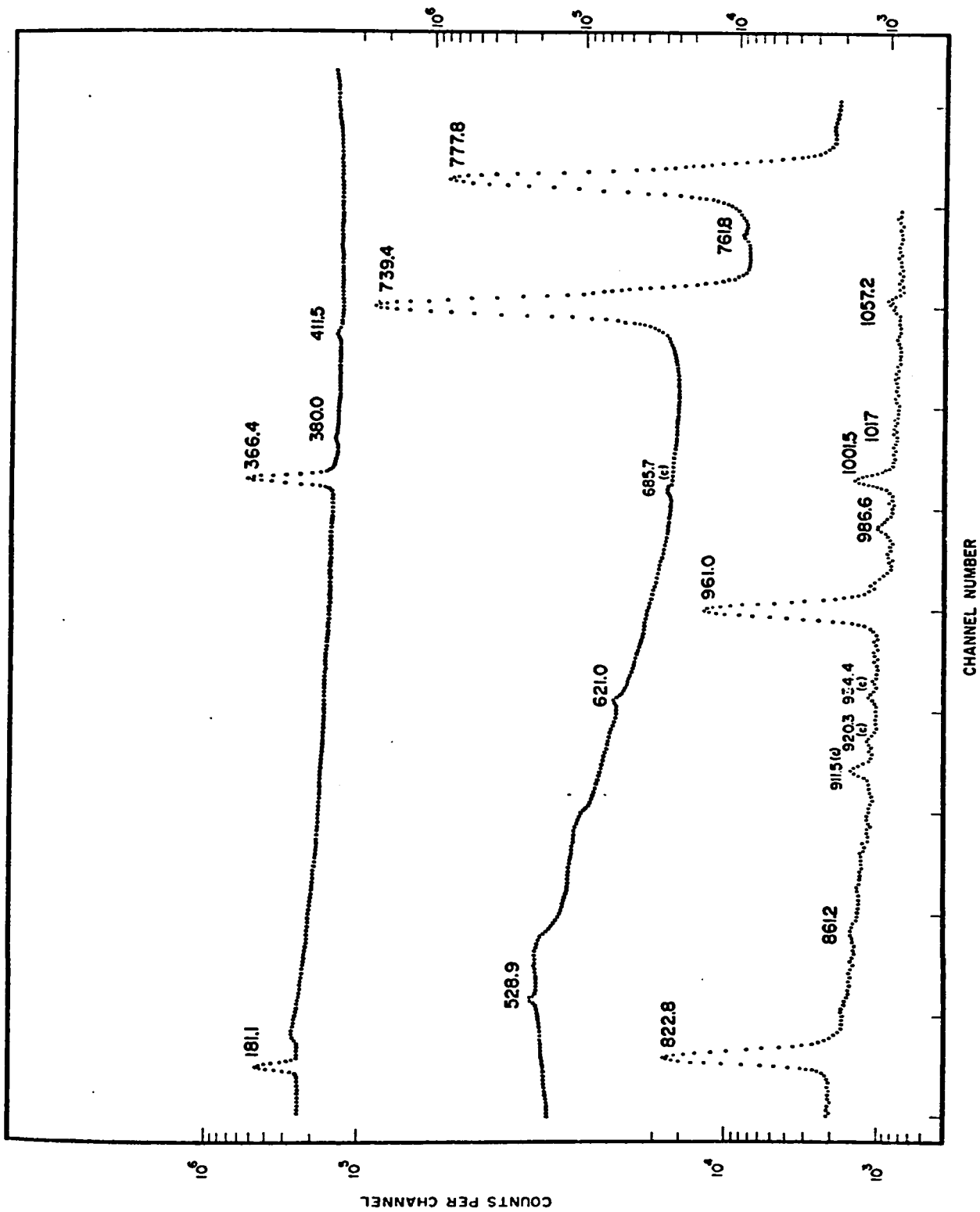


Fig. 3.8 The gamma-ray spectrum of ^{99}Mo above 170 keV

using a range of absorbers. For a number of sources, spectra were recorded as a function of time in order to check the half-life of the gamma rays observed. In most cases, it was possible to definitely assign the transition to the decay of ^{99}Mo . Contaminant lines present in the source are indicated by the letter (c) in the figure. Some of the weaker peaks have been replotted on a linear scale in Fig. 3.9. The peaks indicated have been seen consistently in several runs.

Since rather strong sources were used to obtain the spectrum of Fig. 3.8, the strong peaks show indications of counting rate effects (see section 2.3). These effects distort the high energy side of the peaks near the base (<5% of the height). Other runs, taken with weaker sources and therefore little peak distortion, indicate that no features on the high energy side of the strong peaks have been hidden by the counting rate distortion. The weak peak on the high energy side of the 961 keV peak is a contaminant and should be labelled (c).

3.2.3 Gamma-Gamma Coincidence Results

The details of the Ge(Li)-NaI(Tl) coincidence spectrometer used for these measurements are discussed in section 2.9. A number of Ge(Li) spectra in coincidence with selected regions (approximately 30 keV wide) of the NaI(Tl) gamma-ray spectrum are shown in Figs. 3.10, 3.11 and 3.12. The corresponding chance spectrum has been subtracted in each case. The background contribution, which is due to Compton distributions from higher energy peaks underlying the desired peak in the gate, has not been subtracted. However, its contribution was estimated

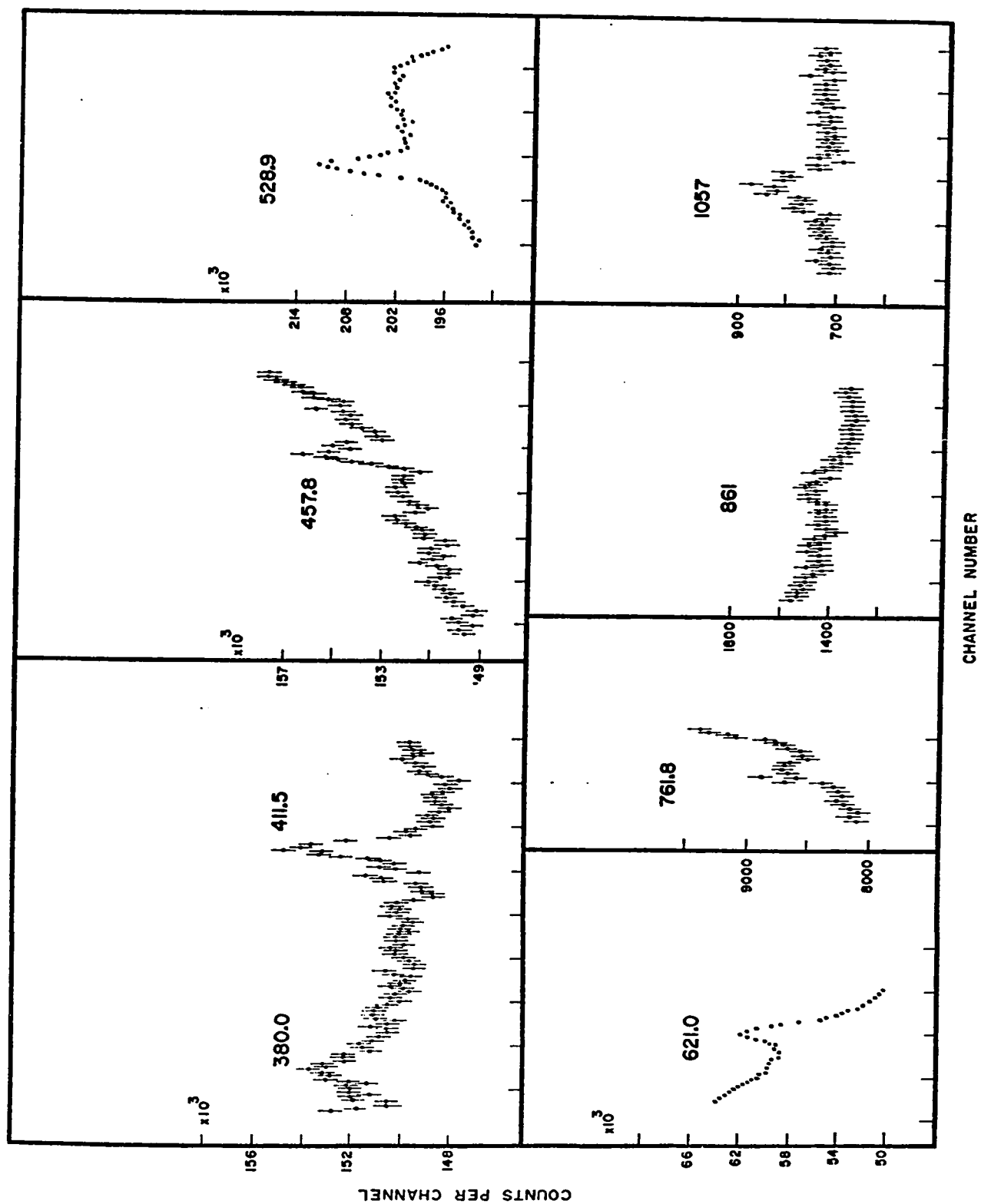


Fig. 3.9 A detailed presentation of the weaker peaks of Fig. 3.8

by comparing the adjacent spectra (gates) shown in the figures.

Fig. 3.10 presents the Ge(Li) spectra in coincidence with NaI(Tl) gates set on the photopeaks of the 140 and 181 keV radiations, respectively. Since these transitions were not completely resolved in the NaI(Tl) spectrum, the 181 keV gate was set to include $\approx 35\%$ of the 181 keV photopeak and none of the 140 keV peak while the 140 keV gate included all of the 140 keV photopeak and about 65% of the 181 keV photopeak. The relative intensities of the 739, 823 and 961 keV peaks in the two spectra are consistent with the conclusions arrived at by other workers and need no further comment. A comparison of the two spectra gives clear evidence for the 621-140 and 1002-140 keV cascades established by Van Eijk et al. (1969). The strong 778 keV transition feeds the isomeric state at 142 keV, its presence in the 140 keV gated spectrum is due to incomplete chance subtraction. Van Eijk et al. (1969) reported transitions of 988 and 1017 keV in coincidence with the 181 keV gamma ray. These transitions are not found in the 181 keV gated spectrum of Fig. 3.10, although the total number of coincidence counts in the 739 and 823 keV photopeaks in the present work was comparable to that of their work.

The spectra associated with portions of the NaI(Tl) spectrum containing the 366, 380 and 411 keV photopeaks are displayed in Fig. 3.11. The 411 keV peak appears in the first two spectra with an intensity consistent with that expected

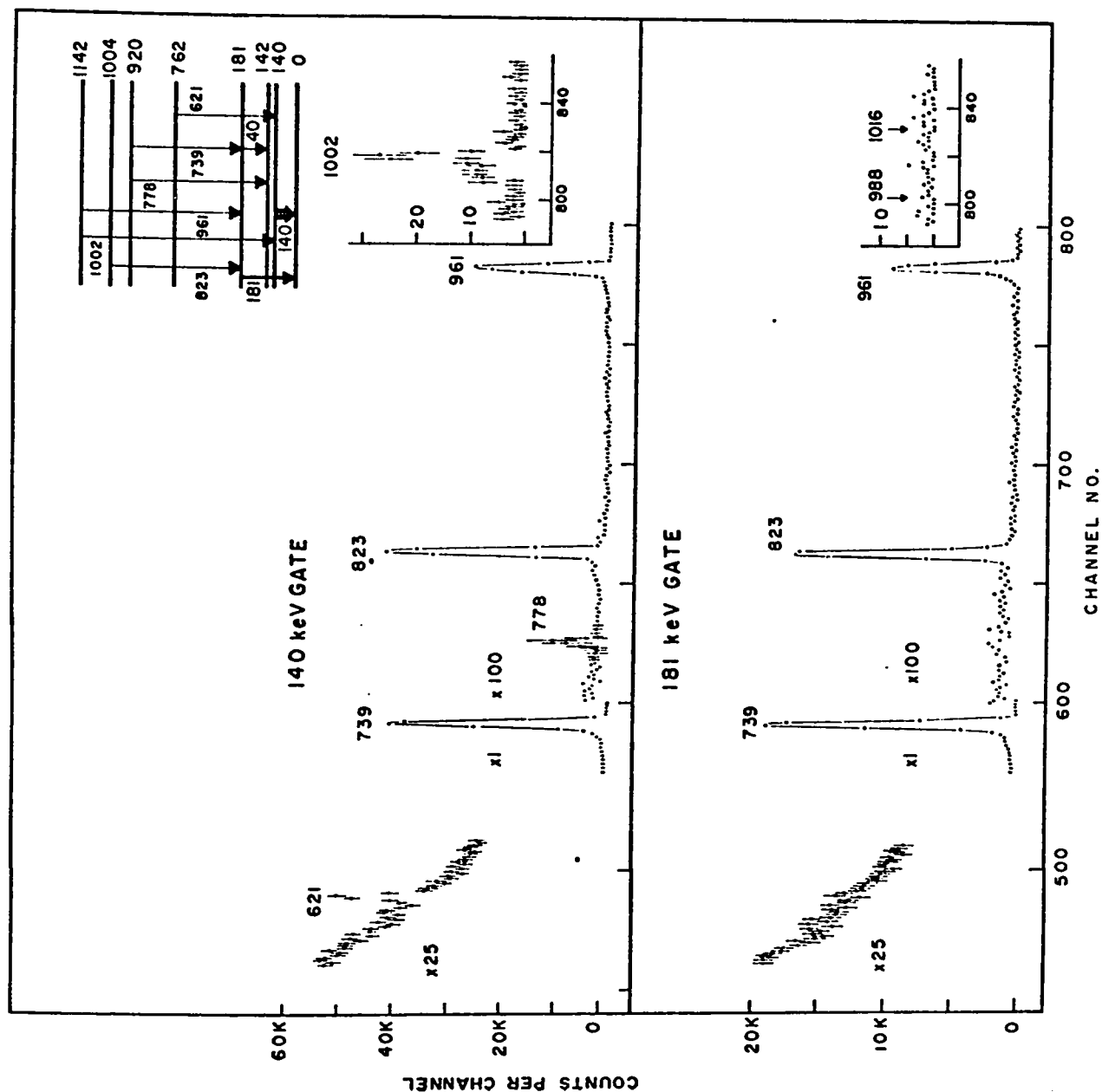


Fig. 3.10 The Ge(Li) spectra in coincidence with gates set on the NaI(Tl) photopeaks of the 140 and 181 keV gamma rays. The 40 keV transition should feed the 140 keV state

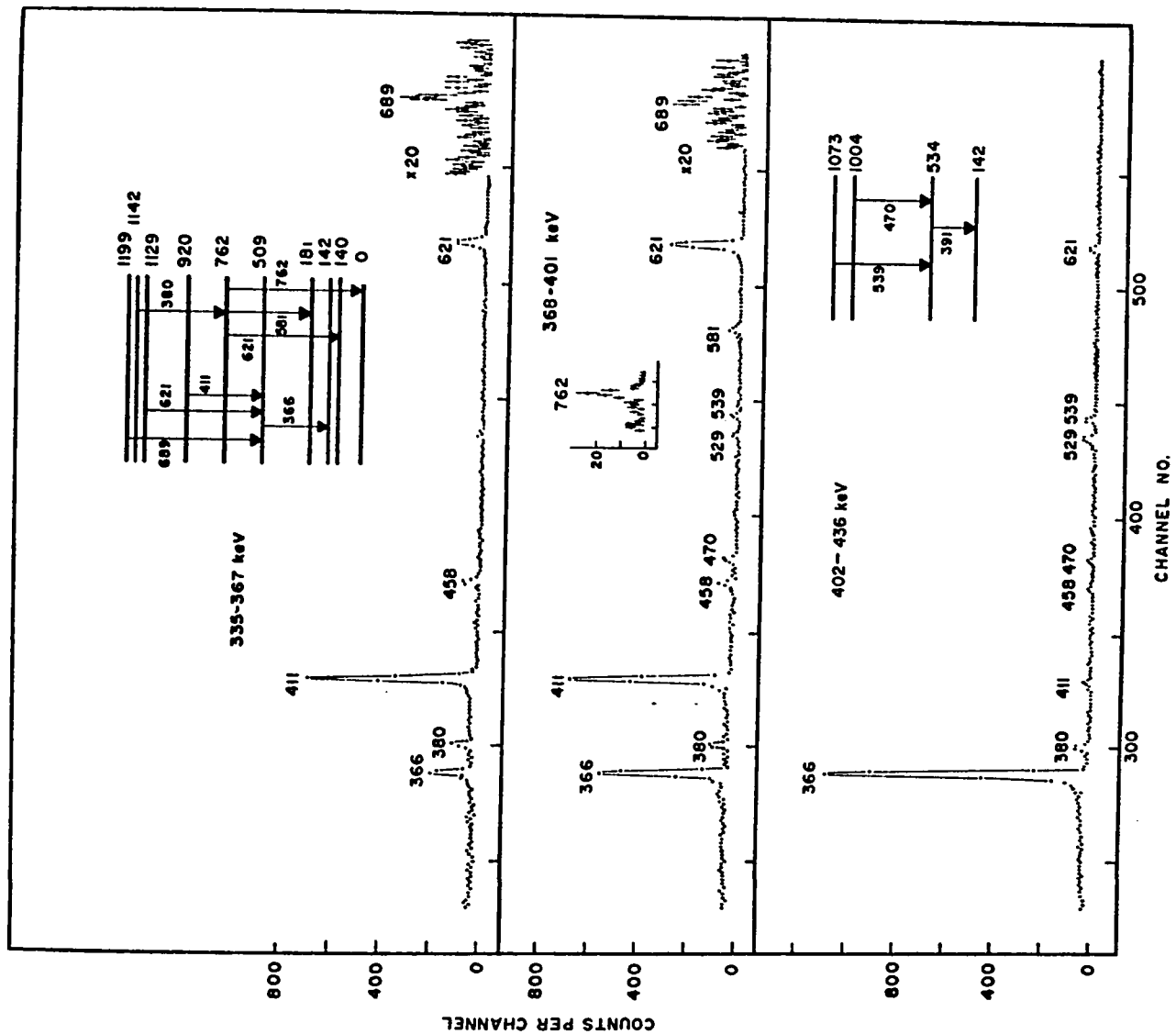


Fig. 3.11 The Ge(Li) spectra in coincidence with NaI(Tl) gates set from 335-367 keV, 368-401 keV and 402-436 keV

for the well-established 411-366 keV cascade. The presence of a 689 keV peak in both spectra implies a 689-366 keV coincidence cascade and establishes a level at 1199 keV. The 621 keV peak is due to two independent cascades and will be discussed below. The 366, 380 and 458 keV peaks in the 335-367 keV gate are all due to coincidences with Compton distributions from higher energy gamma rays.

The 581, 621 and 762 keV peaks in the 380 keV gated spectrum deexcite the 762 keV level which is fed by a 380 keV gamma ray from the 1142 keV state. These data confirm the findings of Van Eijk et al. (1969). The 581 keV line was masked in the spectrum in coincidence with the 181 keV line by the Compton background from the strong 739 keV gamma ray. The spectrum in coincidence with the 411 keV gamma ray shows strong 411-366 keV coincidences and nothing else. The 366 and 411, 689 keV peaks in the 368-401 keV gate are in coincidence with portions of the 411 and 366 keV photopeaks respectively, contained in the gate. The origin of the 470 and 539 keV peaks in the last two spectra will be discussed below. The weak peaks at 380, 458 and 529 keV are due to coincidences with the Compton background.

Fig. 3.12 presents three spectra in coincidence with gates set on the 400 to 650 keV region of the NaI(Tl) spectrum. The main components in the gate shown in the upper spectrum are the 458 and 470 (about 1/2 the photopeak) keV gamma rays which account for the 529 and 391 keV peaks, respectively. The

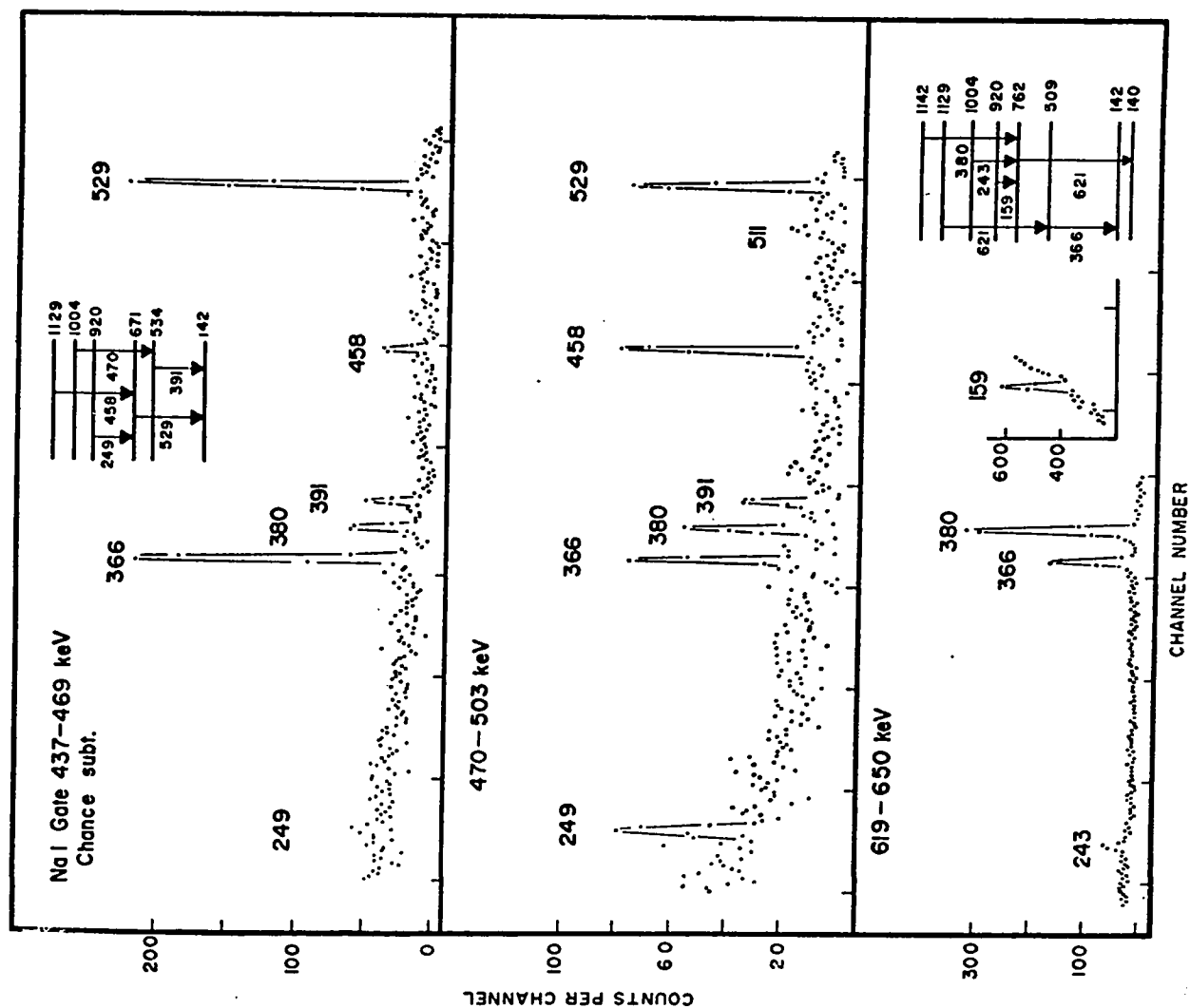


Fig. 3.12 The Ge(Li) spectra in coincidence with NaI(Tl) gates set from 437-469 keV, 470-503 keV and 619-650 keV.

366 keV peak is due to the portion of the 411 keV photopeak contained in the 437-469 keV gate. The remaining peaks are in coincidence with the Compton background.

The gate for the middle spectrum contains almost equal amounts of the 458 and 529 keV photopeaks and the remaining half of the 470 keV photopeak. The 470-391 keV cascade will be discussed below. The 249-529 and 458-529 keV cascades verify the 671 and 1129 keV levels.

The gate for the lower spectrum contains only the 621 keV doublet. The three peaks at 159, 243 and 380 keV are due to gamma rays feeding the 762 keV state from the well-established 920, 1004 and 1129 keV levels, respectively. The 159 and 243 keV transitions have not previously been reported. The 366 keV peak arises from the 621-366 keV cascade. Thus, there is clear evidence for two 621 keV transitions having intensities of 0.24 and 0.035, respectively. These intensities were deduced from the ratio of the 380 and 366 keV peaks observed in the coincidence spectrum. The weak lines reported by Van Eijk et al. (1969) at 344 and 409 keV in their 621 keV gated spectrum are not evident in the 621 keV gated spectrum of Fig. 3.12. The present measurements set upper limits on their intensities which are a factor of ten lower than the intensities quoted by Van Eijk et al. (1969).

Evidence for a 470-391 keV cascade appears in Figs. 3.11 and 3.12. Their energy sum of 861.8 keV corresponds to the energy difference between the well-established levels at 1004

and 143 keV. There is also evidence for a 539-391 keV cascade in Fig. 3.11. This cascade determines the order of the 470-391 keV cascade and defines levels at 534 and 1073 keV. The intensities of these three weak transitions have been established from the coincidence data using the coincidence probabilities for stronger cascades in the same gates.

3.2.4 Discussion of Transition Intensities

The energies and intensities of the transitions associated with the decay of ^{99}Mo are given in Table 3.3. The table includes the gamma-ray data of both Van Eijk et al. (1969) and Cook et al. (1969). In the latter paper, the gamma-ray intensities were normalized to the intensity of the strong 739.4 keV transition. Recent measurements of conversion coefficients and mixing ratios in other laboratories have made it possible to revise the published intensities and express them as intensities/1000 decays. The intensities quoted for the 2.17, 40.55, 140.51, 142.68 and 181.07 keV transitions are total transition intensities (the gamma-ray intensities are given in brackets).

Since transient equilibrium was maintained between ^{99}Mo ($T_{1/2} = 66.7$ hours) and its daughter activity $^{99\text{m}}\text{Tc}$ ($T_{1/2} = 6.1$ hours), the decay rate of the transitions de-exciting the 142.68 keV level exceeded that of the parent by a factor

$$\frac{\lambda_2}{\lambda_2 - \lambda_1} = 1.10$$

where λ_2 and λ_1 are the decay constants of the daughter and

TABLE 3.3 99
Transitions in ⁹⁹Tc

Van Eijk et al. (1969)			Present Work	
Energy	Intensity a)	Energy	Intensity b)	Classification
40.584±0.002		2.17	T ₈₆₁ ±60 ^{c)}	142.68→140.51
		40.55±0.10	T 57.0 ± 6.0	181.07→140.51
140.511±0.006	T ₈₅₀	140.51±0.05	(8.85 ± 1.0)	
	(760)		T ₉₁₈ ±60 ^{d)}	140.51→ 0.0
142.63 ±0.03	10.1 ±2.0	142.68	(823 ±60)	
	(0.25 ±0.05)		T 8.20 ± 0.70 ^{e)}	142.68→ 0.0
181.06 ±0.04	71.9 ±1.8	158.9 ±0.3*	(0.20 ± 0.02)	
	(62.3 ±1.6)	181.07±0.05	T 0.14 ± 0.04*	920.5 →761.6
			T 73.6 ± 5.1	181.07→ 0.0
			(63.9 ± 4.4)	
		242.7 ±0.4*	0.009 ± 0.003	1003.9 →761.6
249 ±1*	0.05 ±0.025*	249.1 ±0.4*	0.065 ± 0.025	920.5 →671.6
344 ±1*	0.010 ±0.004*		<0.0025	
366.4 ±0.1	13.6 ±0.9	366.45±0.1	13.7 ± 0.7	509.13→142.68
380.7 ±0.2*	0.090 ±0.025*	380.0 ±0.4	0.090 ± 0.025	1142.0 →761.6
		391.7 ±0.4*	0.020 ± 0.005	534.4 →142.68
				141

Table 3.3 (continued)

Energy	Intensity a)	Energy	Intensity b)	Classification
409 $\pm 2^*$	0.013 $\pm 0.005^*$		<0.0015	
411.5 ± 0.5	0.23 ± 0.025	411.5 ± 0.4	0.165 ± 0.025	920.5 $\rightarrow$ 509.13
458 $\pm 1^*$	0.05 $\pm 0.025^*$	457.8 ± 0.4	0.10 ± 0.025	1129.3 $\rightarrow$ 671.6
		450.1 $\pm 0.5^*$	0.0075 $\pm 0.002^*$	1003.9 $\rightarrow$ 534.4
528.9 ± 0.2	0.50 ± 0.065	528.9 ± 0.2	0.63 ± 0.06	671.6 $\rightarrow$ 142.68
		538.6 $\pm 0.5^*$	0.013 $\pm 0.003^*$	1073.0 $\rightarrow$ 534.4
		580.9 $\pm 0.4^*$	0.033 $\pm 0.009^*$	761.6 $\rightarrow$ 181.07
620.7*	0.040 $\pm 0.015^*$	620*	0.035 $\pm 0.011^*$	1129.3 $\rightarrow$ 509.13
620.7 ± 0.2	0.23 ± 0.04	621.0 ± 0.2	0.24 ± 0.03	761.6 $\rightarrow$ 140.51
		689.6 $\pm 0.9^*$	0.0045 $\pm 0.002^*$	1199.4 $\rightarrow$ 509.13
739.7 ± 0.1	128 ± 6	739.4 ± 0.1	128 ± 6	920.5 $\rightarrow$ 181.07
		761.8 ± 0.5	0.024 ± 0.006	761.6 $\rightarrow$ 0.0
778.2 ± 0.1	44.9 ± 2.5	777.8 ± 0.1	44.7 ± 2.5	920.5 $\rightarrow$ 142.68
823.1 ± 0.1	1.33 ± 0.09	822.8 ± 0.1	1.42 ± 0.10	1003.9 $\rightarrow$ 181.07
		861.2 ± 0.9	0.0075 ± 0.0025	1003.9 $\rightarrow$ 142.68
961.0 ± 0.2	1.01 ± 0.08	961.0 ± 0.1	1.01 ± 0.08	1142.0 $\rightarrow$ 181.07
988.2 ± 0.5	0.017 ± 0.006	986.6 ± 0.6	0.018 ± 0.005	1129.3 $\rightarrow$ 142.68

Table 3.3 (continued)

Energy	Intensity ^{a)}	Energy	Intensity ^{b)}	Classification
1001.7±0.2	0.055 ±0.015	1001.5±0.3	0.046 ±0.008	1142.0→140.51
1016 ±2*	0.0075±0.004		<0.004	
		1057.2±0.8	0.010 ±0.0025	1199.4→142.68

T Total transition intensity (Gamma-ray intensity in brackets)

* Transition seen only in coincidence

a) Renormalized to intensities/1000 decays

b) Renormalized from Cook et al. (1969) to intensities/1000 decays

c) Transient equilibrium transition intensity 947±65

d) Transient equilibrium transition intensity 1005±65; Photon intensity 901±60

e) Transient equilibrium transition intensity 9.0±0.75; Photon intensity 0.22±0.02

TABLE 3.4

Multipole Assignments in ^{99}Tc

Transition	Multipolarity	α_T	Method	Reference
2.17	E3	∞	M subshell ratios	Amtey et al. (1966)
40.55	M1 (<1% E2)	<4.0	K X-ray/ γ ratio	Lacasse and Hamilton (1971)
	98.6%M1+1.4%E2	4.17	L subshell ratios	Rankumar et al. (1969)
	94.1%M1+5.9%E2	5.45	γ intensity, intensity balance at 181.07 keV level	McDonald et al. (1971)
140.51	(97.2 $\pm$ 0.3)%M1+(2.8 $\pm$ 0.3)%E2	0.115	L subshell ratios	Present work
	96%M1+(4 $\pm$ 2)%E2	0.120	L subshell ratios	Geiger et al. (1968)
142.68	M4	40.2	K/L ratio	McDonald et al. (1971)
				Mihelich et al. (1951)
				Ravien et al. (1961)
181.07	E2	0.151	α_K	Van Eijk et al. (1969)
366	M1		α_K	Bashandy (1969)
411	E1		α_K	
380	M1+E2		α_K	
458	M1+E2		α_K	

Table 3.4 (continued)

Transition	Multipolarity	α_T	Method	Reference
529	E2		α_K	
621	M1		α_K	
739	M1+E2		α_K	
778	E1		α_K	
823	E1		α_K	
961	M2		α_K	
1002	M2		α_K	

parent activities, respectively. Therefore, it was necessary to correct the observed intensities (which were too large) of the 2.17 (and therefore that part of the 140.51 keV transition intensity due to the 2.17 keV feeding) and 142.68 keV transitions by this transient equilibrium factor. This correction has been included in the results of Table 3.3.

The transition intensities in the table have been deduced on the reasonable assumption that there is no beta feed to the 0.0, 140.51 and 181.07 keV states. Any such feeding would require second, or higher order forbidden beta transitions, if the spin and parity assignments of Fig. 3.13 are correct. Table 3.4 tabulates the conversion electron data which have been published recently.

The intensity balance at the 181.07 keV level indicates that the total intensity of the 40.55 keV transition is 57.0. When this value is combined with the gamma-ray intensity of 8.85, an $\alpha_T = 5.45$ is obtained for the 40.55 keV transition; this α_T corresponds to a multipole mixture of 94.1% M1 + 5.9% E2. This value is in reasonable agreement with the 1.4% E2 admixture deduced from L-subshell ratios by McDonald et al. (1971). In view of the fact that the L-subshell lines were incompletely resolved in their spectra, the present measurement is of comparable precision.

The intensity balance at the 140.51 keV level leads directly to the intensity of the unobserved 2.17 keV transition. The total intensity feeding the level is $I_{2.17} + 57.2$.

By using the α_T value for the 140.51 keV transition obtained by Geiger (1968) and the measured gamma-ray intensity of 901, one deduces that the total transition intensity of the 140.51 keV line when in transient equilibrium is 1005. This leads to a transient equilibrium intensity of 947 for the 2.17 keV transition. The associated intensities, which have been corrected for transient equilibrium, are given in Table 3.3.

A knowledge of the 142.68 keV transition intensity is required in order to obtain the beta feed to the 142.68 keV state from the intensity balance at the level. Using a source of ^{99m}Tc , Geiger (1968) measured the ratio of the K-conversion line intensities of the 142.68 and 140.51 keV transitions to be $(7.46 \pm 0.11) \times 10^{-2}$. These data, when combined with the K-conversion coefficients of the two transitions, lead to a 142.68/140.51 gamma-ray intensity ratio of $(2.57 \pm 0.04) \times 10^{-4}$. The gamma-ray and total intensities of the 142.68 keV transition are given in Table 3.3. The results obtained from the measured ratio of the 140.51 and 142.68 keV gamma rays by Van Eijk et al. (1969) are somewhat higher; however, the errors in this measurement are greater than in Geigers' (1968) due to the high background involved. The intensity balance at the 142.68 keV level leads to a beta feed of 81.0% to the level.

3.2.5 The Level Scheme of ^{99}Tc

The results of the measurements described in the previous sections are incorporated into the level scheme of Fig. 3.13. Table 3.5 gives a summary of the proposed level energies, the

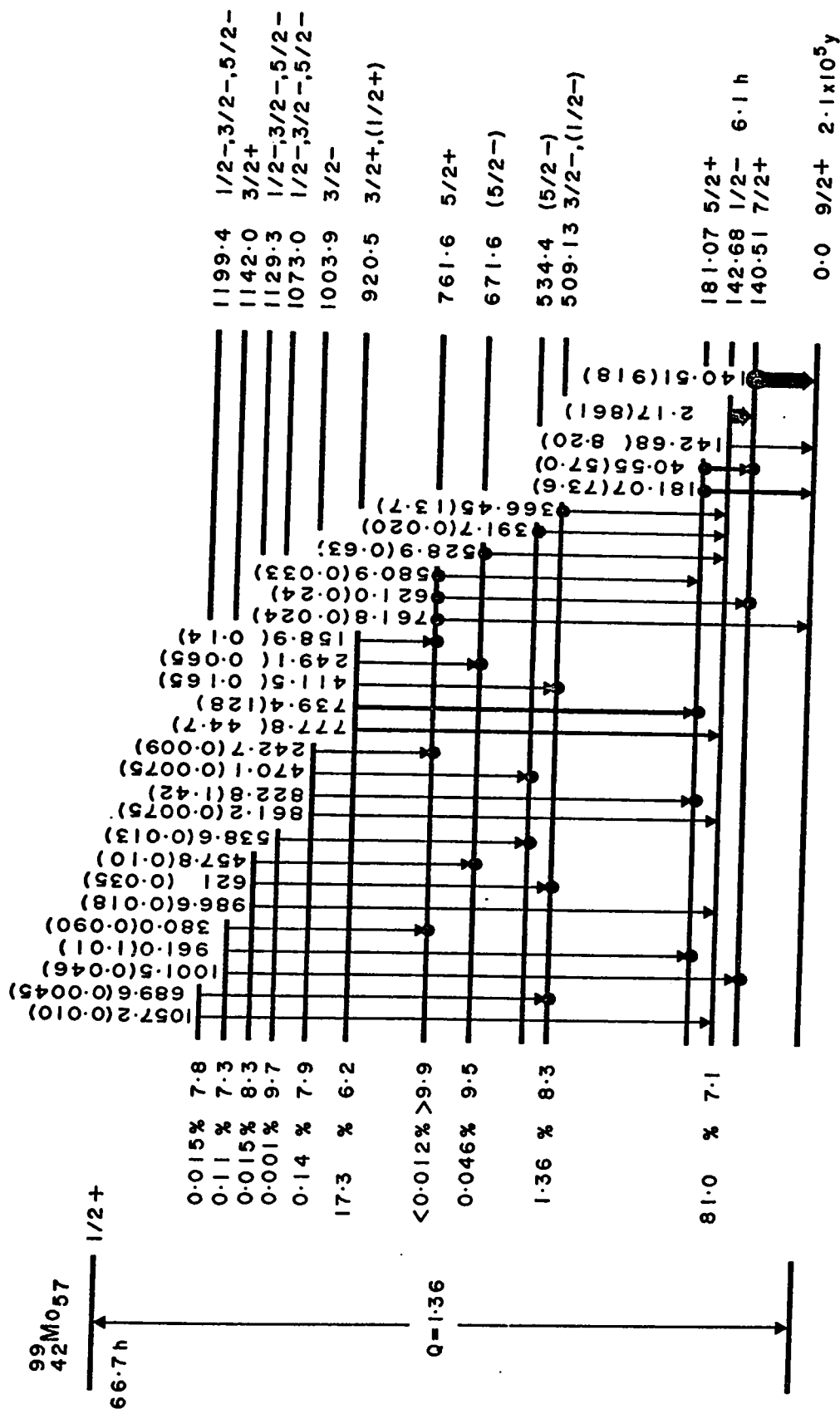
Fig. 3.13 The level scheme of ⁹⁹Tc

TABLE 3.5

Assignment of $\log f_{\text{O}t}$ and J^π values to levels of ^{99}Tc

Level Energy keV	% Beta Feed	$\log f_{\text{O}t}$	J^π
0			$9/2^+$
140.51 $\pm$ 0.05			$7/2^+$
142.68 $\pm$ 0.05	81.1 $\pm$ 6.0%	7.1	$1/2^-$
181.07 $\pm$ 0.05			$5/2^+$
509.13 $\pm$ 0.11	1.36 $\pm$ 0.07%	8.3	$3/2^-, (1/2^-)$
534.4 $\pm$ 0.4			$(5/2^-)$
671.6 $\pm$ 0.2	0.046 $\pm$ 0.007%	9.5	$(5/2^-)$
761.6 $\pm$ 0.2	<0.012%	>9.9	$5/2^+$
920.5 $\pm$ 0.1	17.3 $\pm$ 0.7%	6.2	$3/2^+, (1/2^+)$
1003.9 $\pm$ 0.1	0.14 $\pm$ 0.01%	7.9	$3/2^-$
1073.0 $\pm$ 0.6	0.001 $\pm$ 0.0003%	9.7	$1/2^-, 3/2^-, 5/2^-$
1129.3 $\pm$ 0.35	0.015 $\pm$ 0.003%	8.3	$1/2^-, 3/2^-, 5/2^-$
1142.0 $\pm$ 0.1	0.11 $\pm$ 0.01%	7.3	$3/2^+$
1199.4 $\pm$ 0.6	0.015 $\pm$ 0.003%	7.8	$1/2^-, 3/2^-, 5/2^-$

percent beta feeds, the deduced $\log f_0 t$ values and the J^π values assigned for each of the proposed levels. The nomograms of Lederer et al. (1967) were used to calculate the $\log f_0 t$ values.

The $^{98}\text{Mo}(d,p)^{99}\text{Mo}$ reaction data of Hjorth and Cohen (1964) and Moorhead and Moyer (1969) have established that the ground state spin and parity of ^{99}Mo is $1/2^+$. Allowed beta transitions will therefore populate states in ^{99}Tc with J^π values of $1/2^+$ or $3/2^+$. Similarly, states with J^π values of $1/2^-$, $3/2^-$ or $5/2^-$ will be populated with first forbidden beta transitions.

The ground state, the 140.51 and the 142.68 keV levels.

The properties of these levels have been thoroughly established by many workers (see Table 3.4 and introduction for some of the references). The ground spin and parity is known to be $9/2^+$. The $7/2^+$ and $1/2^-$ assignments for the 140.51 and 142.68 keV levels respectively, follow directly from the properties of the deexciting transitions (see Table 3.4). These results are all consistent with the measurements of the present work.

The 181.07 keV level. The intensity balance, the existence of transitions to this state from higher-lying, low-spin levels and the multipolarities of the deexciting transitions are consistent with a $5/2^+$ assignment for the 181.07 keV level. The present measurements again confirm the conclusions of earlier workers.

The 509.13 keV level. The strong 366.45 keV gamma ray is the only transition to depopulate the 509.13 keV level

which is established by the 411.5-366.45 keV cascade. The $\log f_0 t$ value of 8.3 is characteristic of a first forbidden beta transition and implies a J^π value of $1/2^-$ or $3/2^-$ for the level. On theoretical grounds (see section 4.3), the $3/2^-$ choice is to be preferred.

The 534.4 keV level. This level has been proposed on the basis on the weak 470.1-391.7 and 538.6-391.7 keV cascades. The transitions involved were too weak to be observed in singles. The rather meagre evidence at hand suggests that the 534.4 keV level possesses a spin of 5/2 and has negative parity.

The 671.6 keV level. The 457.8-528.9 and 249.1-528.9 keV cascades define the 671.6 keV level, which decays only to the $1/2^-$ state at 142.68 keV. The beta transition to the 671.6 keV level has a $\log f_0 t$ value of 9.5 characteristic of a first forbidden beta transition. The parity of the 671.5 keV state is therefore negative, the spin may be 1/2, 3/2 or 5/2. The high value of the $\log f_0 t$ favours a $5/2^-$ assignment.

The 761.6 keV level. Van Eijk et al. (1969) established this level mainly on the basis of the 621-140.51 keV gamma-ray cascade. Five gamma rays (158.9, 242.7, 580.9, 621.0 and 761.8 keV) are associated with the 761.6 keV state in the present investigation. The existence of the 761.8 keV ground-state transition rules out J^π values of $1/2^\pm$, $3/2^\pm$ or $5/2^-$. There appears to be no beta feeding to the 761.6 keV level and the $\log f_0 t$ must be >9.9 ; again J^π values of $1/2^\pm$ or $3/2^\pm$

are ruled out. The low energy transitions populating the 761.6 keV state from higher energy levels with low spin ($1/2$ or $3/2$) suggest spin values of $5/2$ or less. The data are consistent with a spin and parity of $5/2^+$ for the 761.6 keV state.

The 920.5 keV level. The 777.8 and 739.4 keV transitions to the 142.68 and 181.07 keV levels respectively, define the 920.5 keV level. The $\log f_{\text{t}}$ value of 6.2 is characteristic of an allowed beta transition; the spin and parity of the 920.5 keV state must be $1/2^+$ or $3/2^+$. The strong 739.4 keV transition to the $5/2^+$ state at 181.07 keV and the existence of low energy transitions to the 671.5 and 761.6 keV states favour a $3/2^+$ assignment. A number of inconclusive attempts (Cappeller et al. 1954; Raboy and Krohn 1958; Bodenstedt et al. 1959; Estulin et al. 1959; Andrade et al. 1965) have been made to measure the angular correlation of the 739.4-181.07 keV cascade in order to determine the spin of the 920.5 keV level. The errors involved were too large to yield definite assignments.

The 1003.9 keV level. The 822.8-181.07 and 242.7-621.0 keV gamma-ray cascades establish this level. The weak 861.2 keV gamma ray has been placed between the 1003.9 and 142.68 keV levels on the basis of energy fit. The first forbidden character of the feeding beta transition ($\log f_{\text{t}} = 7.9$) indicates J^{π} values of $1/2^-$ or $3/2^-$ for the 1003.9 keV level. In view of the low $\log f_{\text{t}}$ value, a $5/2^-$ assignment has not been considered. The presence of a strong transition to the 181.07 keV state rules out the $1/2^-$ choice. The spin and parity of the 1003.9 keV

state must therefore be $3/2^-$.

The 1073.0 keV level. The weak 538.6-391.7 keV coincidence cascade is the only evidence for the 1073.0 keV level. The high $\log f_0 t$ of 9.9 for the beta feed to the 1073.0 keV level indicates negative parity and a spin of $1/2$, $3/2$ or $5/2$.

The 1129.3 keV level. This level is established with the 621-366.45 and 457.8-528.9 keV gamma-ray cascades. In addition, the 986.6 keV line can be placed between the 1129.3 and 142.68 keV levels on the basis of energy fit. The $\log f_0 t$ value of 8.3 for the beta feed to this level is consistent with J^π values of $1/2^-$, $3/2^-$, or $5/2^-$.

The 1142.0 keV level. The 1001.5-140.51, 961.0-181.07 keV cascades and the 380.0 keV transition to the 761.6 keV state define the 1142.0 keV level. The beta transition ($\log f_0 t = 7.3$) is either allowed or first forbidden in character. The spin and parity of the 1142.0 keV level can therefore be $1/2^\pm$, $3/2^\pm$ or $5/2^-$. The existence of the 1001.5 keV transition to the $7/2^+$ state at 140.51 keV rules out the $1/2^+$ and $1/2^-$ choices. A $3/2^-$ assignment would require a M2 multipolarity for the 1001.5 keV line; the intensity of such a transition would be a factor $\sim 10^7$ less than the intensity of the much faster 961.0 keV transition (if an E1 multipolarity is assumed) rather than the factor of 0.046 observed experimentally. A J^π value of $5/2^-$ for the 1142.0 keV level would require a first forbidden unique beta transition; however, the $\log f_0 t$ value is rather low for such an assignment. A spin and parity of $3/2^+$ is therefore suggested for the 1142.0

keV state.

The 1199.4 keV level. Van Eijk et al. (1969) introduced this level on the basis of the weak 1016-181 keV coincidence cascade. There is no clear evidence for this cascade in the present coincidence data and no 1016 keV gamma ray was observed in the singles data. The present experiments reveal a 689.6-366.45 keV cascade which would deexcite a level at 1199.4 keV. In addition, Table 3.3 includes a 1057.2 keV gamma ray whose energy is appropriate to a transition from the 1199.4 keV state to the 142.68 keV isomeric state. The $\log f_{\text{ot}}$ value of 7.8 is characteristic of a first forbidden beta transition indicating J^{π} values of $1/2^{-}$, $3/2^{-}$ or $5/2^{-}$ for the 1199.4 keV level. If the 1017 keV transition exists, the $1/2^{-}$ choice would be ruled out.

Van Eijk et al. (1969) used a 344-620.7 keV cascade to define a level at 1105 keV. There is no evidence in the present work to confirm this level. They also used the 988 keV line and a weak 409 keV gamma ray to define a level at 1169 keV. There is no evidence for the 409 keV line; its intensity is at least a factor of ten lower than the value quoted by Van Eijk et al. (1969). Moreover, there is no evidence for the 988-181 keV cascade, in a coincidence experiment which had comparable statistics to their experiment. The data of the present work indicate that the 1169 keV state does not exist.

3.2.6 Subsequent Investigations

Several workers (Bashandy (1969), Lacasse and Hamilton (1971)) have reinvestigated portions of the internal conversion

spectrum from the ^{99}Mo decay; their results have been included in Table 3.4. The assignments made by Bashandy (1969) are not very convincingly documented. The K-conversion peaks are very weak and are found on rather high backgrounds due to the strong beta feeds to the 509.13 and 142.68 keV states. It is therefore expected that the α_K 's have rather large errors with the result that the data should only be used to distinguish E1 from M1 and/or E2 transitions.

The angular correlation of the 739.4-181.07 keV gamma-ray cascade has been remeasured by Gfirtner et al. (1970) using a liquid molybdenum source and Ge(Li) detectors. Their data indicate a spin of $3/2^+$ for the 920.5 keV level and a multipolarity of $(2.8 \pm 0.3)\% \text{M1} + (97.2 \pm 0.3)\% \text{E2}$ for the 739.4 keV transition.

Nagarajan et al. (1971) have used the fact that both the shape and the end-point energy can be accurately determined from a study of the shape factor of a beta transition as a function of the end-point energy. Such a study of the first forbidden beta transition to the 142.68 keV level gives a value of 1214 ± 1 keV for the end-point energy. This leads to a decay energy of 1357 keV.

Coulomb excitation of ^{99}Tc with ^{35}Cl ions (Bond et al. 1971) has resulted in the excitation of levels at 140.5(0.108), 181.1(0.027), 726.3(0.065) and 761.8(0.120) keV. The $B(\text{E2})$'s ($e^2 \text{b}^2$) are given in brackets. Measurements of the gamma-ray angular distributions are consistent with the known $7/2^+$ and $5/2^+$ assignments for the 140.5 and 181.1 keV levels and restrict

the possible spin and parity assignments of both the 726.3 and 761.8 keV levels to $7/2^+$ or $9/2^+$. Due to the high spins, these latter two levels should not be populated in the decay of ^{99}Mo , which has a ground state spin of $1/2^+$. However, a level at 761.6 keV has been observed in the present work and has been assigned a spin and parity of $5/2^+$. Both the Coulomb excited state and the state populated in beta decay are deexcited by 621.6(621.0) and 761.8(761.6) keV transitions to the 140.5 keV level and the ground state, respectively. The ratio of the intensities of the 621.0 and 761.8 keV transitions found in beta decay (10.0) is considerably different than the ratio (0.16) found for the 621.6 and 761.8 keV lines from Coulomb excitation. The data make it quite clear that the beta decay and Coulomb excitation experiments have not populated the same level. Therefore, two levels at 761.6($5/2^+$) and 761.8 ($7/2^+$ or $9/2^+$) keV exist in ^{99}Tc . Due to the small energy difference, it is impossible to say whether or not the Coulomb excitation experiments also populated the 761.6 keV level seen in beta decay. However, an upper limit of $0.018 e^2 b^2$ can be set on the $B(E2)$ to this state if the 621.6 keV line is assumed to be completely due to deexcitation of the 761.6 keV state (beta decay) rather than the 761.8 keV state (Coulomb excitation). Easterday et al. (1972) have confirmed the existence of two levels at 761.56($5/2^+$) and 762.3($7/2^+$) keV. In addition, they have identified a level at 726.7 keV ($9/2^+$ or $7/2^+$).

3.3 Levels in ^{101}Tc Populated in the decay of ^{101}Mo

3.3.1 Introduction

The beta decay of 14.6 min ^{101}Mo was identified by

Sagane et al. (1940) and Maurer and Ramm (1941). Although preliminary investigations had been carried out by Rutledge et al. (1952) and Wiles (1954), the first detailed study was performed by Martin et al. (1956). Singles, gamma-gamma and beta-gamma coincidence experiments with scintillation spectrometers and magnetic spectrographs established the existence of 22 gamma rays, ranging in energy from 80 to 2080 keV, and at least 6 beta components. The maximum energy beta-ray group of 2.2 MeV was observed to be in coincidence with a 590 keV gamma ray. A strong 191 keV transition was found to decay from an isomeric state having a half-life of $\sim 9 \times 10^{-4}$ secs. From this half-life and the value of the K-conversion coefficient, an M2 multipolarity was assigned to the 191 keV transition. The results of their measurements were incorporated into a level scheme involving eleven excited states of ^{101}Tc . A further study using scintillation methods by O'Kelly et al. (1957) determined an accurate value for the half-life ($T_{1/2} = 14.61 \pm 0.07$ min.) and measured a value of 2.23 ± 0.04 MeV for the maximum energy beta-ray group. The decay scheme resulting from the above data is presented in Fig. 3.14 (taken from Nuclear Data Sheets, 1961). The half-life of the isomeric state has been remeasured by several groups: Brandi et al. (1964) - $710 \pm 120 \mu\text{sec}$, Gabsdil et al. (1968) - $850 \pm 120 \mu\text{sec}$, and Uyttenhove et al. (1970) - $760 \pm 50 \mu\text{sec}$. The weighted mean of these results is $765 \pm 45 \mu\text{sec}$.

The present investigation of the beta decay of ^{101}Mo to levels in ^{101}Tc is a continuation of a study of the odd-mass technetium isotopes begun several years ago at McMaster. Since

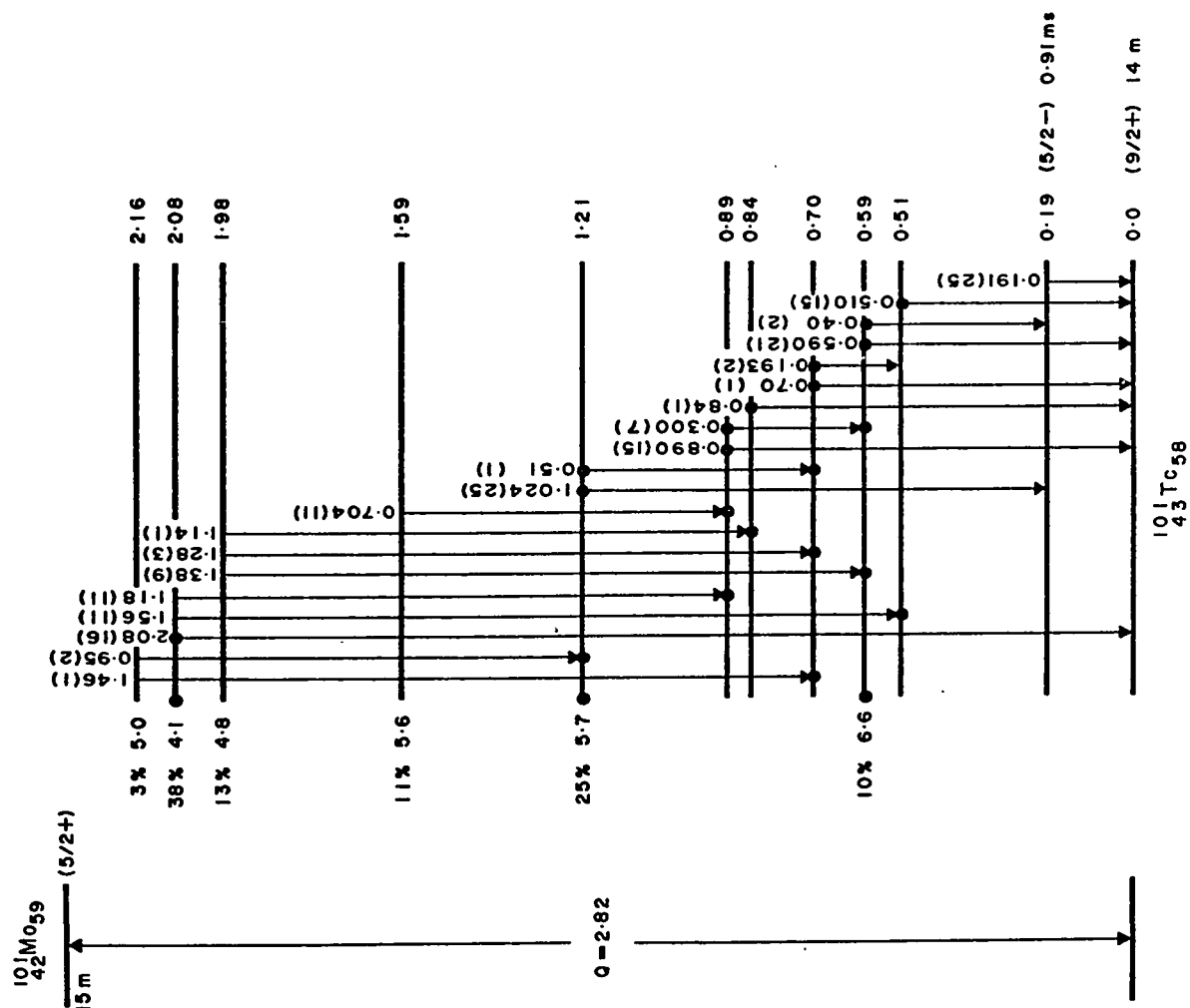


Fig. 3.14 The decay scheme of ^{101}Tc prior to the present work

there has been no previously published study of the decay of ^{101}Mo using Ge(Li) detectors, it was expected that this new investigation would find many previously unobserved transitions and would establish a number of new levels.

A systematic comparison of the levels in the neighbouring isotopes ^{97}Tc and ^{99}Tc , which are populated with the electron-capture decay of ^{97}Ru (Section 3.1 (see also Cook et al. 1970); Phelps and Sarantites 1971) and with the beta decay of ^{99}Mo (section 3.2 (see also Cook et al. 1969)) respectively, is hampered by the fact that the spins of the parent nuclei, ^{97}Ru ($5/2^+$) and ^{99}Mo ($1/2^+$) are different. In fact, it is found that $5/2$ and $7/2$ states are populated in the decay of the former, and $1/2$ and $3/2$ states in the decay of the latter. The ground state spins and parities of the odd-mass technetium isotopes are known to be $9/2^+$ (Lederer et al. 1967). The only other states where spins and parities have been positively identified in both ^{97}Tc and ^{99}Tc occur at relatively low excitation and have J^π values of $5/2^+$, $7/2^+$ and $1/2^-$.

The $^{100}\text{Mo}(d,p)^{101}\text{Mo}$ reaction data of Hjorth and Cohen (1964) have shown that the spin and parity of the ground state of ^{101}Mo is $1/2^+$. Therefore, the strong beta transitions in the decay of ^{101}Mo must proceed to states of low spin ($1/2$ or $3/2$) in ^{101}Tc . However, the large decay energy of 2.8 MeV should allow the population of higher spin states through gamma-ray cascades. Therefore, it is expected that the levels in ^{101}Tc , populated in the beta decay of ^{101}Mo , will include

many of the high and the low spin states observed in the investigations of ^{97}Tc and ^{99}Tc , respectively.

Unfortunately, the ground state of ^{101}Tc is not stable, but beta decays with a half life of 14.0 min. Therefore, transitions from this decay appear as impurities in the gamma-ray spectra of the ^{101}Mo decay.

3.3.2 Gamma-Ray Singles Measurements

The low energy region of a gamma-ray spectrum containing both ^{101}Mo and ^{101}Tc is displayed in Fig. 3.15. The spectrum was taken with the 0.8 cm^3 thin window detector; 12.5 mm of lucite served as a beta stopper, but no other absorber was used. The X-rays of technetium and ruthenium are due to internal conversion; those of molybdenum are caused by the absorption of beta radiation in the source. Absorption of low energy gamma rays in the gold layer at the front of the detector gives rise to the gold X-rays. The lines at 127, 175, 180 and 184 keV belong to the decay of ^{101}Tc ; the remainder of the gamma rays are from the decay of ^{101}Mo .

Fig. 3.16 shows the gamma-ray spectrum between 150 and 1300 keV taken with the 12 cm^3 detector. The high energy gamma rays above 1300 keV are presented in Fig. 3.17. In addition to the lucite beta stopper, a graded absorber of lead, cadmium and copper, 2 mm thick, was used to attenuate the low energy radiations from ^{101}Mo and ^{101}Tc . Gamma rays from the decay of ^{101}Tc are so labelled; gamma rays from the room background (^{137}Cs , ^{232}Th) are designated by the letter (c).

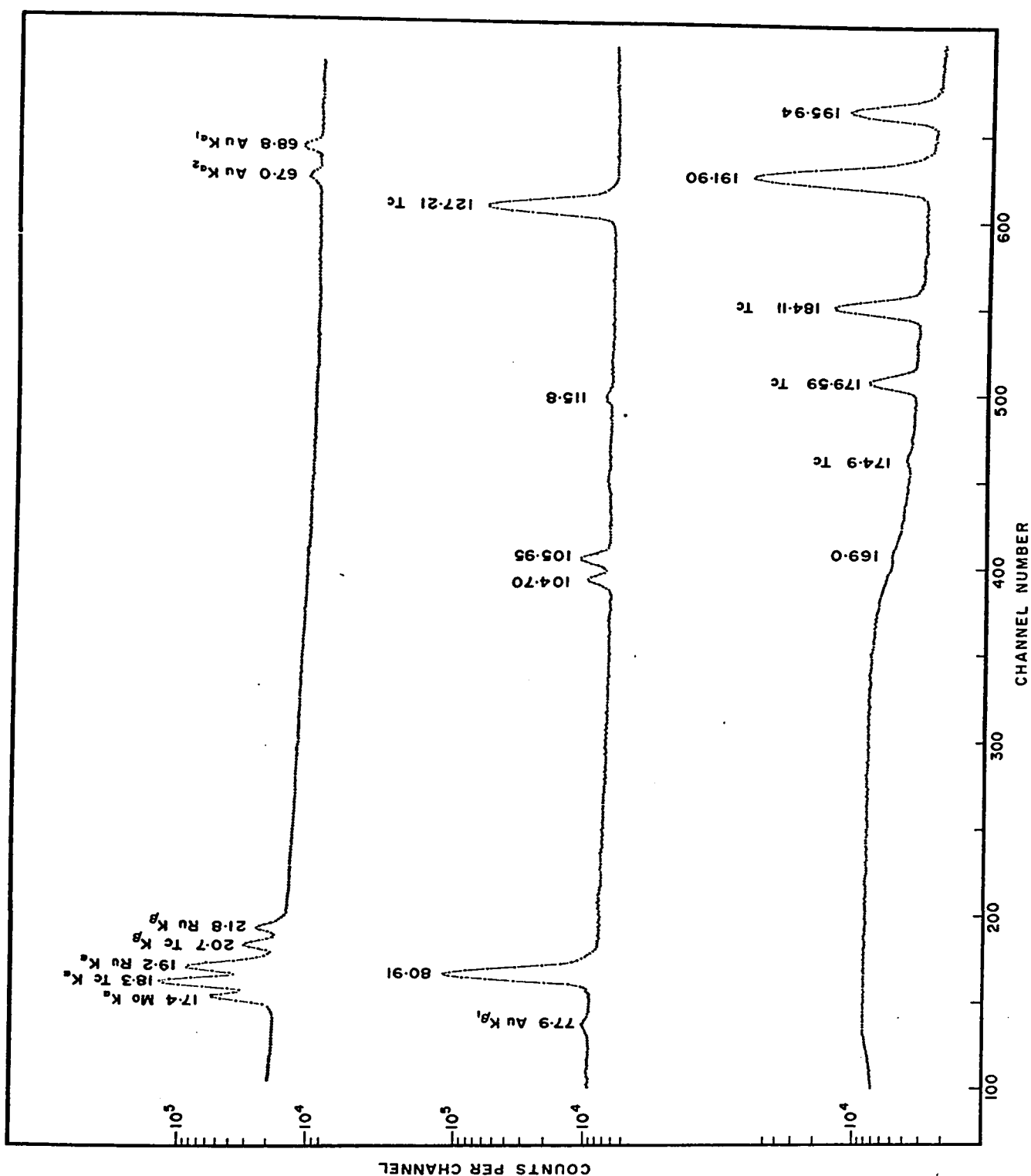


Fig. 3.15 The low energy portion of the gamma-ray spectrum

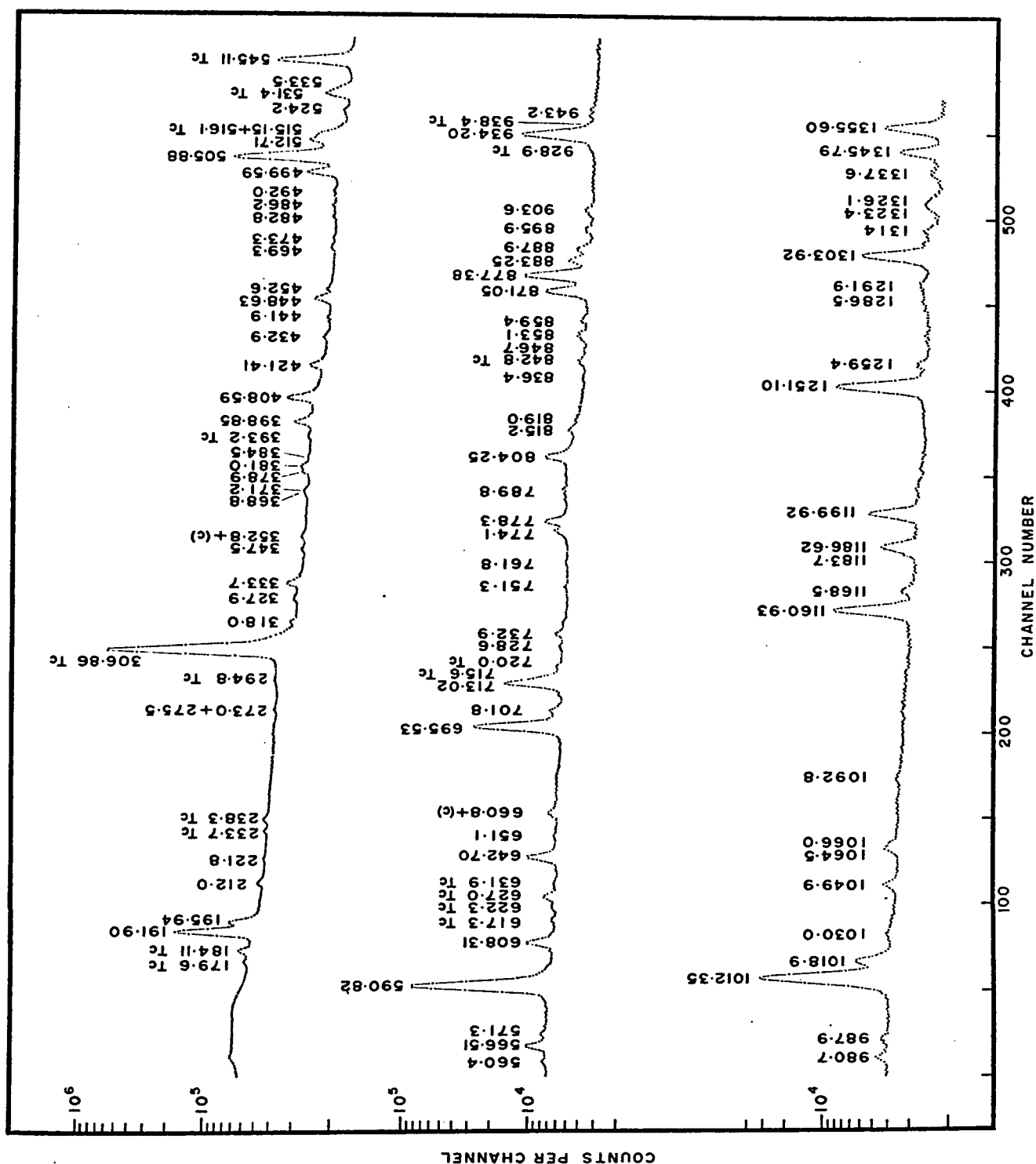


Fig. 3.16 The gamma-ray spectrum of ^{101}Mo and ^{101}Tc from 150 to 1360 keV recorded with a $12\text{ cm}^3\text{ Ge(Li)}$ detector. Contaminants from room background are labelled "c"

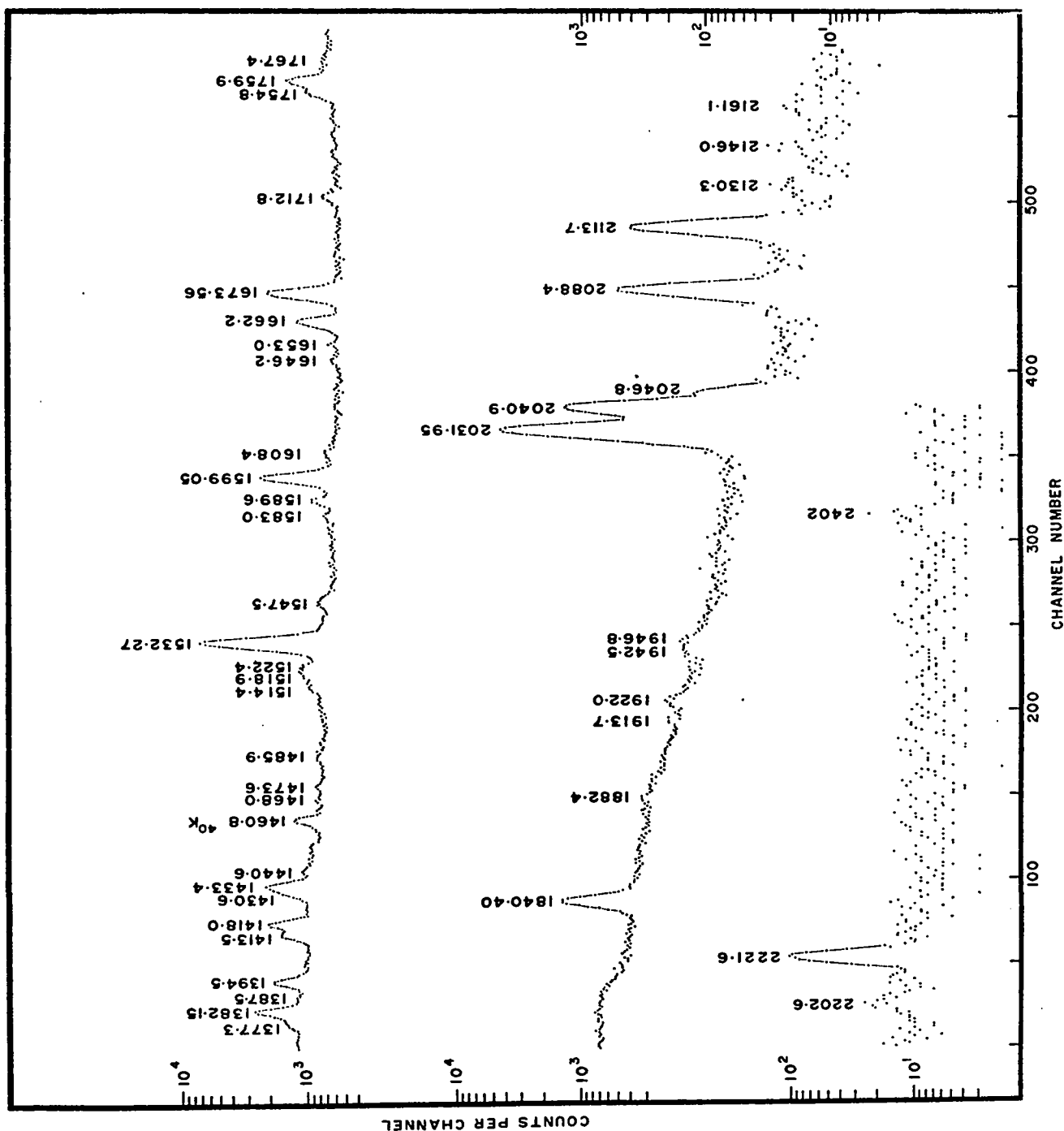


Fig. 3.17 The gamma-ray spectrum of ^{101}Mo above 1360 keV taken with a 12 cm^3 Ge(Li) detector

Several regions of Fig. 3.16 have been replotted on a linear scale in Fig. 3.18 to indicate more clearly some of the weaker lines. Portions of a spectrum taken with the 0.8 cm³ detector are included in the lower half of the figure. The doublet character of the 422 and 1012 keV gamma rays is very clear. The gamma-ray complex near 500 keV shows a 510 keV line in addition to the 500, 506, 513 and 515 keV transitions observed with the 12 cm³ detector. No 511 keV gamma ray was seen.

3.3.3 Internal Conversion Measurements

Although the internal conversion measurements were carried out after both the gamma-ray singles and coincidence experiments, the discussion of the electron data has been included here in order to simplify the interpretation of the coincidence data which follows in the next section. The motivation for performing these electron measurements was provided by several special features of the gamma-ray data. First of all, the singles spectra of Figs. 3.16 and 3.17 contain five pairs of gamma rays differing in energy by ~ 6.3 keV. In addition, the decay scheme that evolved required an excited state rather close to the ground state if both the coincidence data and the constraints imposed by intensity balance arguments were to be satisfied. From this indirect evidence, the existence of an excited state at 6.3 ± 0.1 keV was postulated. Electron measurements were therefore carried out with the iron-free, $\pi\sqrt{2}$ beta-ray spectrometer at the Chalk River Nuclear

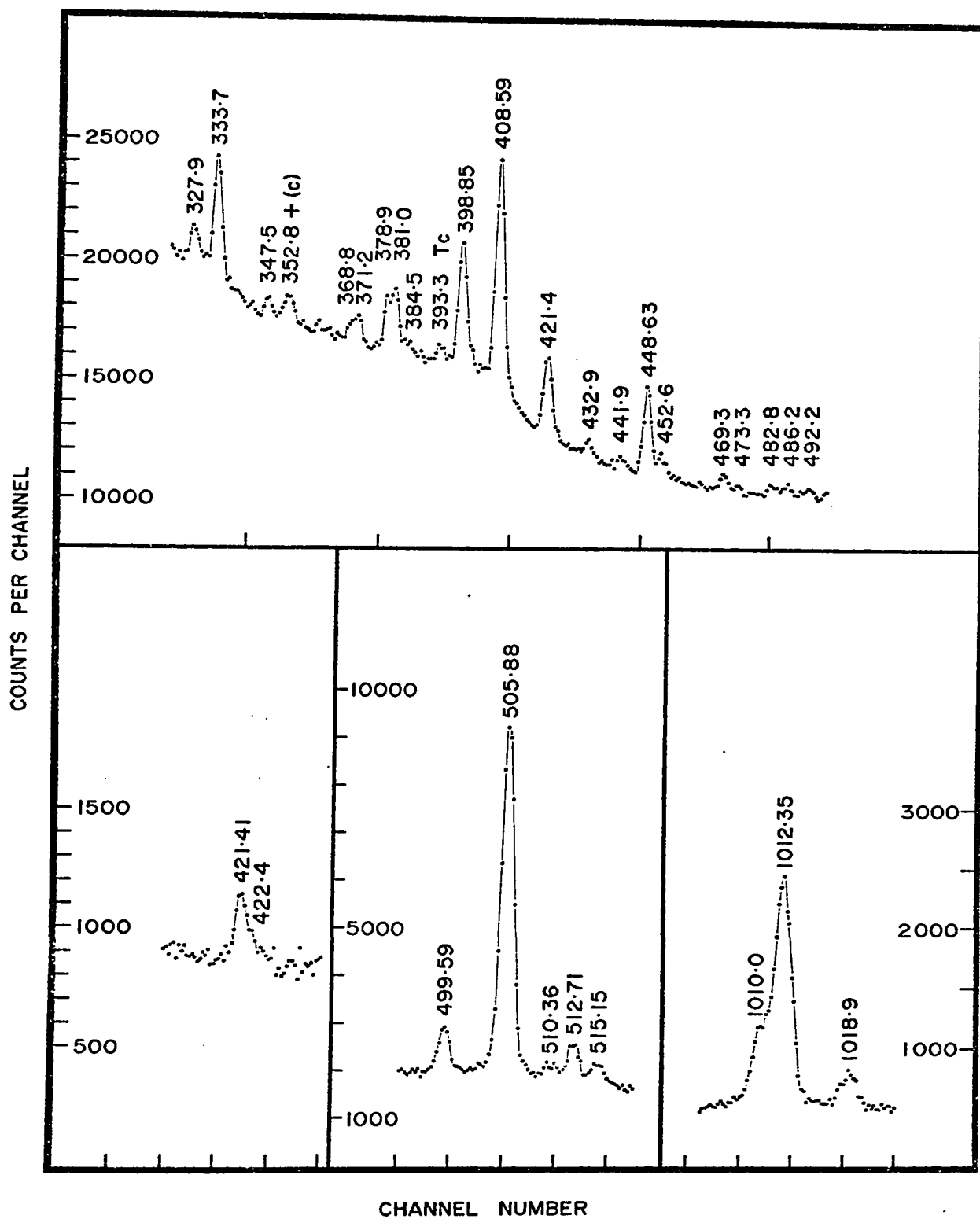


Fig. 3.18 Upper Half: Expanded portions of the spectrum of Fig. 3.17

Lower half: Portions of the spectrum as recorded with the high resolution 0.8 cm³ detector

Laboratories in order to obtain a direct measurement of the energy and multipolarity of the 6.3 keV transition. In searching for the M lines of this transition, the L lines of an even stronger transition with an energy of 9.317 keV were observed. In a subsequent experiment, the L lines of a 15.606 keV transition were also found.

A complete description of the experimental procedures is found in section 2.8 and therefore only a few details have been included here. The internal conversion lines observed in the electron experiments are displayed in Figs. 3.19 and 3.20. For this presentation and for the purpose of analysis, the raw data were summed in 5 channel groups. In addition to the ^{101}Mo activity, the neutron irradiation produced considerable activity in the iron source-backing. Most of the background in the electron spectrum was due to beta activity from this backing. The background was found to decay much more slowly than the conversion lines and was treated as a constant in making the decay corrections. The decay of the L_1 peak of the 9.317 keV transition was followed for about six half-lives. A half-life of 14.77 ± 0.16 minutes was obtained in good agreement with the accepted value of 14.61 ± 0.07 minutes (O'Kelley et al. 1957) for the ^{101}Mo decay. The half-life of the L_1 peak of the 6.281 keV transition was found to be 15.0 ± 1.0 min., the larger error reflecting the poorer statistics in this determination. No attempt was made to measure the half-life of the 15.606 keV transition. The energy separations of

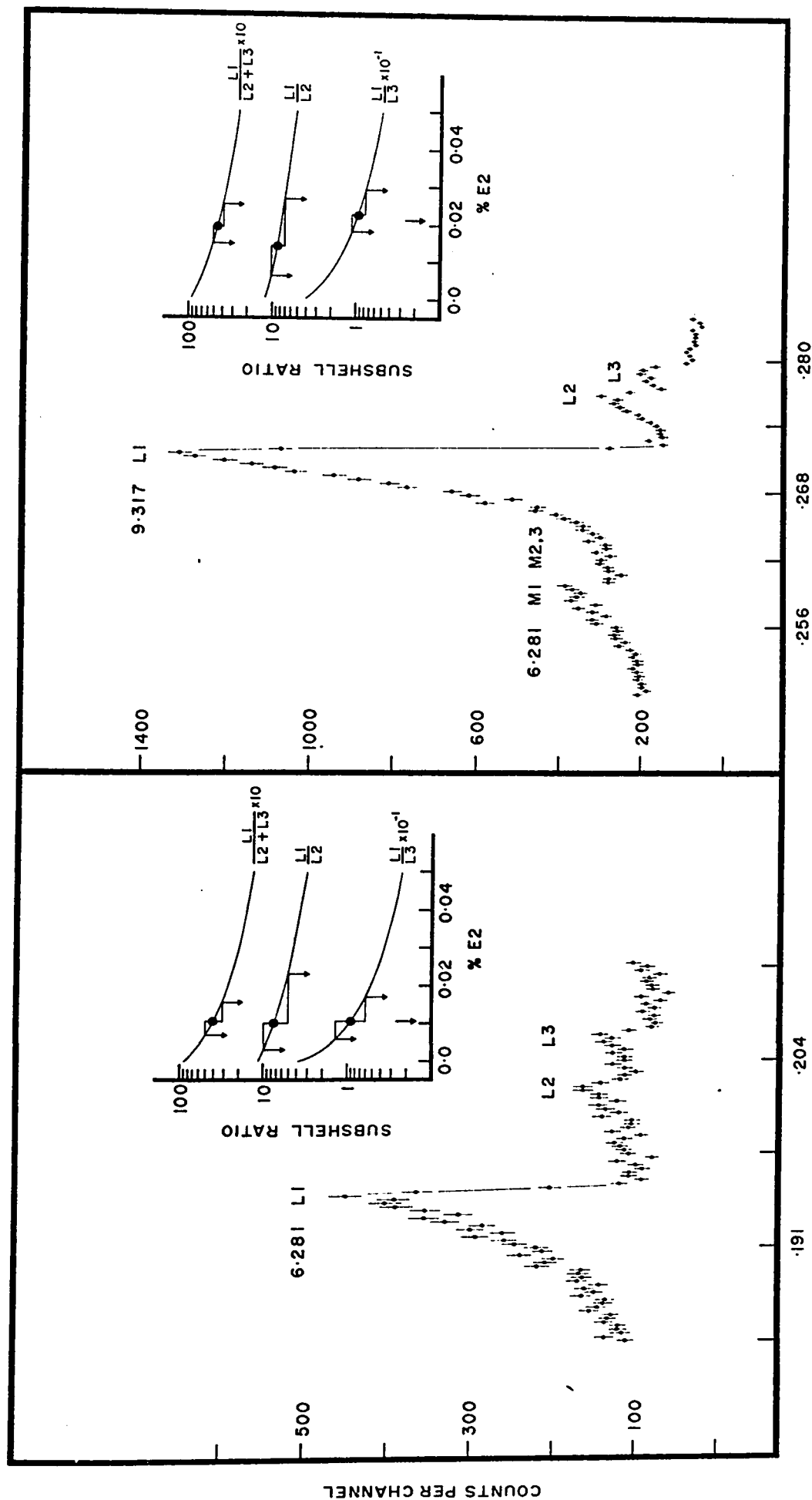


Fig. 3.19 Conversion lines of the 6.281 and 9.317 keV transitions in ^{101}Tc . Inserts show the theoretical L subshell ratios as a function of E2 strength and the multipolarities deduced from the measured ratios for each transition. The heavy arrows indicate the adopted E2 admixtures.

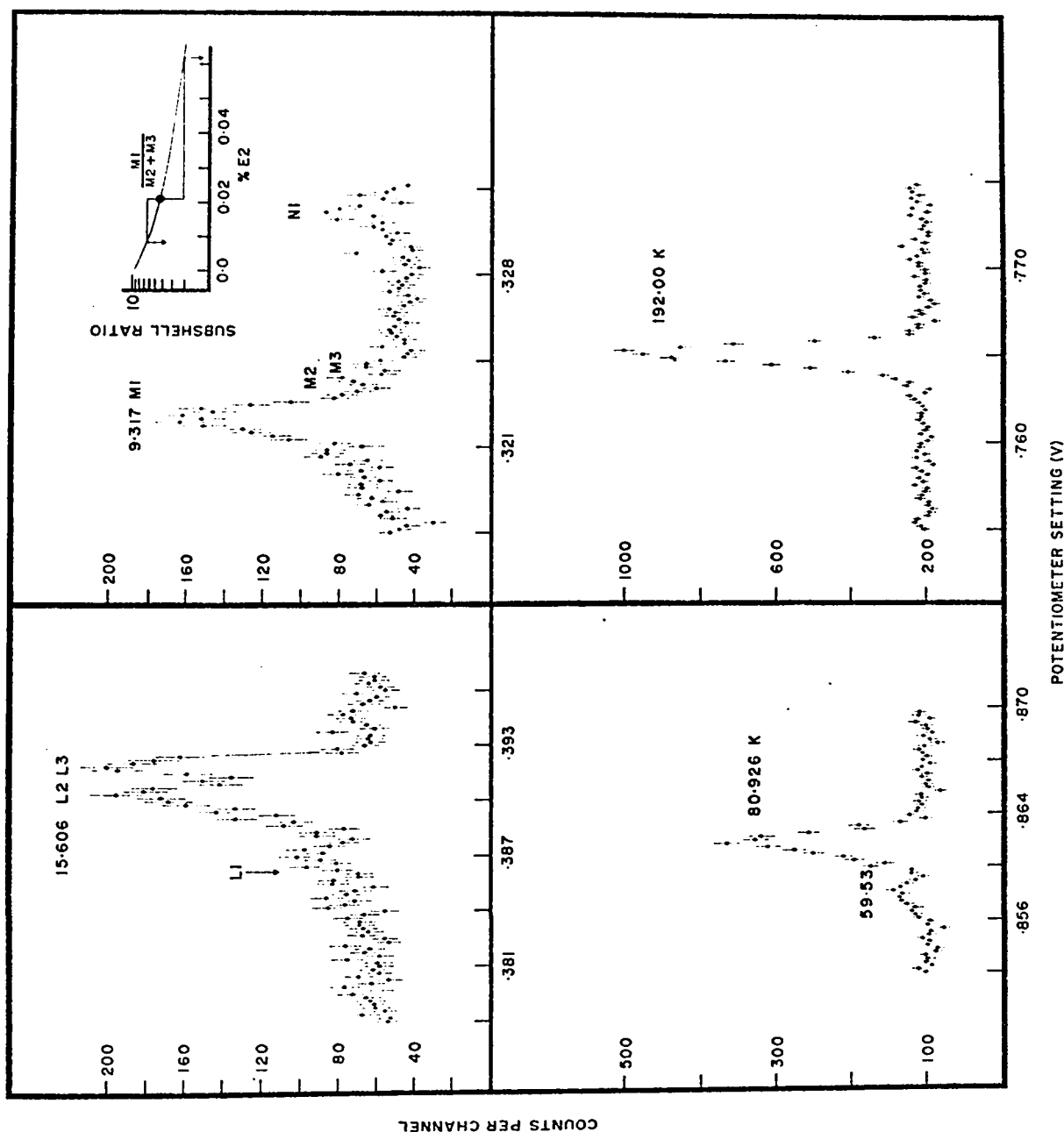


Fig. 3.20 Conversion lines of the 9.317, 15.606, 80.926 and 192.00 keV transitions in ^{101}Tc . The absence of the L_1 component for the 15.506 keV transition establishes its multipolarity as E2. The weak peak at an electron energy of 59.53 keV is unidentified and is probably due to an impurity

the L_1-M_1 and L_1-N_1 conversion lines of the 6.281 and 9.317 keV transitions are consistent only with the transitions being converted in technetium ($Z=43$).

An analysis of the spectra of Figs. 3.19 and 3.20 leads to the transition energies and conversion line intensities of Table 3.6. The 6.281 and 9.317 keV transitions are almost pure M1 in character; the 15.606 keV crossover transition is predominantly E2. Inserts in Figs. 3.19 and 3.20 show the theoretical L subshell ratios as a function of the E2 admixture in the 6.281 and 9.317 keV transitions. From these inserts, it is seen that the measured L subshell ratios correspond to an E2 component of $0.010 \pm 0.005\%$ for the 6.281 keV transition and an E2 component of $0.021 \pm 0.006\%$ for the 9.317 keV transition.

The right hand spectrum of Fig. 3.19 displays the M conversion lines of the 6.28 keV transition together with the L lines of the 9.317 keV transition. From the areas of these lines and the appropriate theoretical conversion coefficients (Hager and Seltzer 1968), it was established that the intensity of the 6.281 keV transition is $75.4 \pm 9\%$ of that of the 9.317 keV transition. A similar analysis of the areas of the L conversion lines of the 9.317 (Fig. 3.19) and the 15.606 (Fig. 3.20) keV transitions was used to determine their relative intensities. The relative strength of the ^{101}Mo activity on each source was deduced from the peak areas of the 192.00 keV gamma ray observed in the Ge(Li) gamma-ray spectrum taken for each source.

The bottom half of Fig. 3.20 displays the K-conversion lines of the 80.926 and 192.00 keV transitions. Both these transitions have been observed in the gamma-ray spectrum of Fig. 3.15. Previous work (see introduction) has shown that the 192.00 keV transition has an M2 multipolarity. Assuming the theoretical M2 conversion coefficient for the 192.00 keV gamma-ray ($\alpha_K = 0.25$) and using the measured $I_{K80.926}/I_{K192.00}$ and $I_{\gamma80.926}/I_{\gamma192}$ ratios, the K-conversion coefficient of the 80.926 keV line was found to be 0.52 ± 0.08 . This corresponds to the 80.926 keV transition being M1+E2 with an E2 component of less than 9%. When the electron energies were averaged with the gamma-ray energies, transition energies of 80.92 and 191.93 keV were obtained.

The energies and the relative intensities of the three low energy transitions observed in the internal conversion data establish the existence of excited states in ^{101}Tc at 9.317 and 15.601 keV. This conclusion is strongly supported by the coincidence data of section 3.3.4 which define a level at 515.2 keV that populates the ground state and both of the excited states. Many other levels are observed which decay to the 9.317 and 15.601 keV excited states. The intensity balance from the gamma-ray measurements requires that 82.0% of the gamma-ray intensity feeds the 15.601 keV state, 16.9% feeds the 9.317 keV state and 1.1% proceeds directly to the ground state. Since 98.9% of all the decays proceed through one or both of these excited states and the 15.606 keV crossover is quite weak, the absolute intensities of the 6.281 and 9.317 keV transitions can be deduced

almost independently from either the gamma intensity balance or the conversion electron intensities. The results of these deductions are shown in Table 3.6. The agreement between the two sets of intensities is excellent.

The low energies of these levels in ^{101}Tc accounts for the fact that these states have eluded detection for so long. The failure to identify the ground state of ^{101}Tc means that previously quoted values for the $^{101}\text{Mo} \rightarrow ^{101}\text{Tc}$ decay energy were in error by 15.6 keV.

3.3.4 Gamma-Gamma Coincidence Results

A sample of the gamma-gamma coincidence results for ^{101}Mo is presented in Figs. 3.21-3.32. These figures show the spectra recorded in one detector in coincidence with a selection of full energy peaks in the other. Since a complete discussion of the experimental details and analysis procedures is found in section 2.9, only a brief outline is included here. Contributions to the gated spectra from chance coincidences and from coincidences with the Compton distribution of higher energy gamma rays have been subtracted. Since the subtraction process sometimes created a peak, peaks have been considered as "real" only if they also appeared in the uncorrected data. Because of the necessity of processing pulses associated with gamma rays with energies from 80 to 2400 keV, it was not possible to completely avoid loss of coincidence efficiency due to "timing walk" in the electronics. The coincidence efficiency of the 36 cm^3 gating detector, which detected gamma rays in the

Table 3.6 Low Lying Transitions in ^{101}Tc

All intensities are quoted in percent

Energy (keV)	6.281±0.007	9.317±0.010	15.606±0.015	80.926±0.036	192.00±0.08
Intensity a) From C.E. (1)					
b) From γ (2)					
Photon Intensity (2)					
C.E. Intensities K	-	-	-	-	-
L_1	50.2±8.4	64.7±1.8	<0.16		
L_2	6.7±2.4	7.2±1.3	0.6±0.2		
L_3	5.6±2.4	6.7±1.1	1.0±0.2		
M_1	9.1±1.6	13.4±2.2	} 0.3 (5)		
M_2+M_3	2.5±0.8	3.1±1.5			
N	1.4 (4)	1.9±0.7			
				3.48±0.23	18.2±1.3
				1.8±0.3	4.6±0.3 (3)

- (1) deduced from conversion electron intensities, assuming 98.9% of transitions decay through the 15.601 and 9.317 keV states.
- (2) deduced from the γ -ray measurements of Cook and Johns (1972) assuming an intensity of 1.9% for the 15.606 keV line.
- (3) deduced from γ intensity, assuming M2 multipolarity from previous work ($\alpha_K = 0.25$ theoretical)
- (4) Assumed using M/N ratio observed for 9.317 keV transition.
- (5) Assumed using theoretical L/M ratios.

Deduced Multipolarities

Transition	
6.281 keV	M1+(0.010±0.005)% E2
9.317 keV	M1+(0.021±0.006)% E2
15.606 keV	E2, (<50% M1)
80.926 keV	M1, (< 9% E2)
192.00 keV	M2 (assumed)

60-1700 keV range, was 100% over the entire range of interest. However, the coincidence efficiency of the 40 cm³ detector, which covered an energy range from 60 to 2400 keV varied from 100% for the region above 170 keV to 50% at 81 keV. This variation must be considered in interpreting the spectra of Figs. 3.21-3.32. Due to the large number of coincidence relationships involved, no attempt will be made to present all the coincidence data used in constructing the decay schemes of Figs. 3.33 and 3.34. However, it is hoped that the samples presented will serve to document the conclusions made in this study. In the discussion that follows, the partial level schemes, which accompany each coincidence spectrum, are used as an aid in interpreting the coincidence cascades.

The spectrum in coincidence with the 591 keV transition is displayed in Fig. 3.21. Since the gamma rays in coincidence with the 591 keV line supply only 40% of its intensity, the 591 keV transition must depopulate a level which is strongly beta fed and which therefore has low spin ($1/2$ or $3/2$). The 591 keV line could feed either the ground state or one of the excited states at 9.3 or 15.6 keV. The evidence given below shows that the 591 keV transition populates the 15.6 keV state and therefore defines a level at 606 keV.

A comparison of the spectra in coincidence with the 196 keV (Fig. 3.22) and 696 keV (not shown) gamma rays reveals that these two spectra are identical except for the presence of the strong 500, 506 and 515 keV transitions in the former

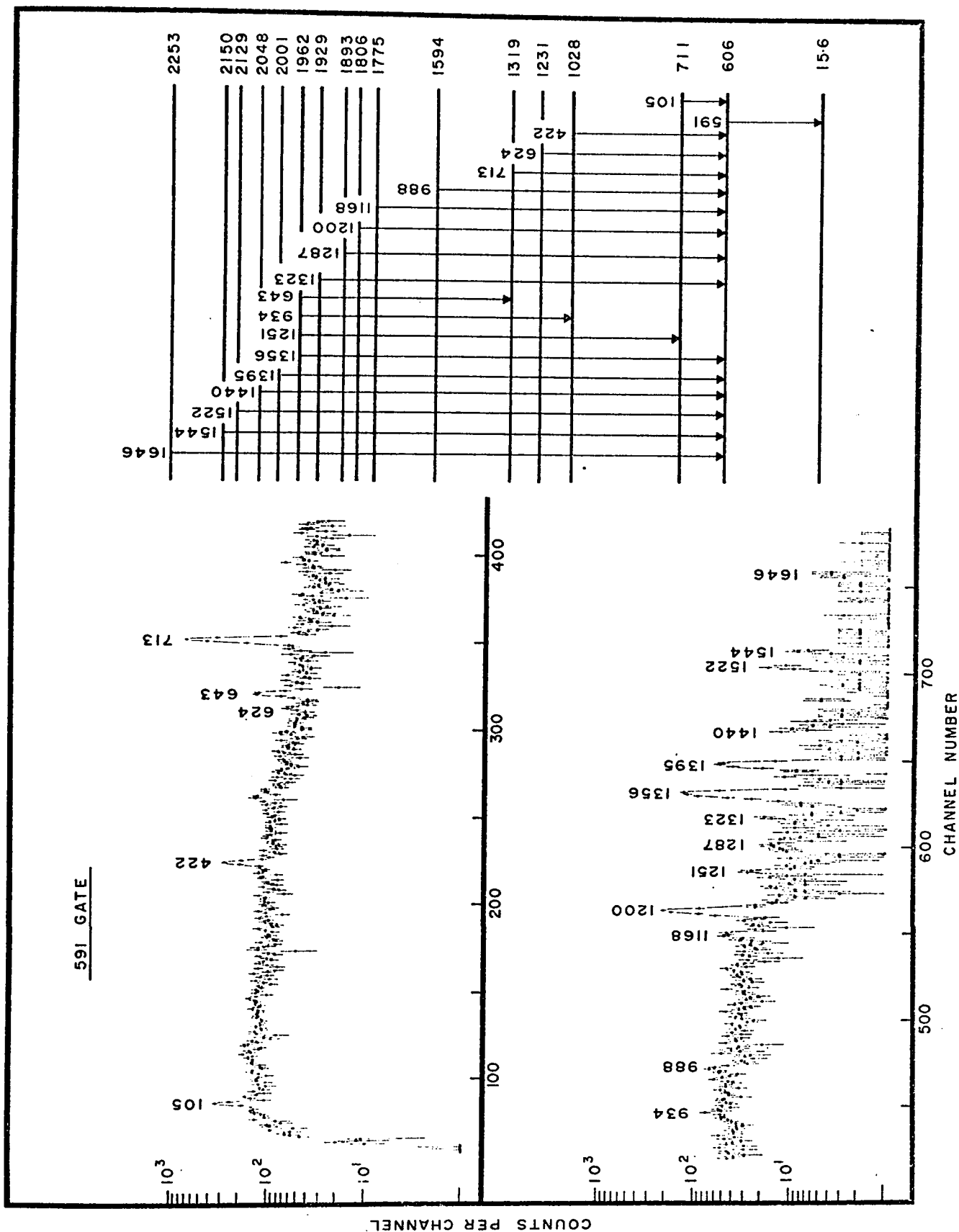


Fig. 3.21 The spectrum in coincidence with the 591 keV transition

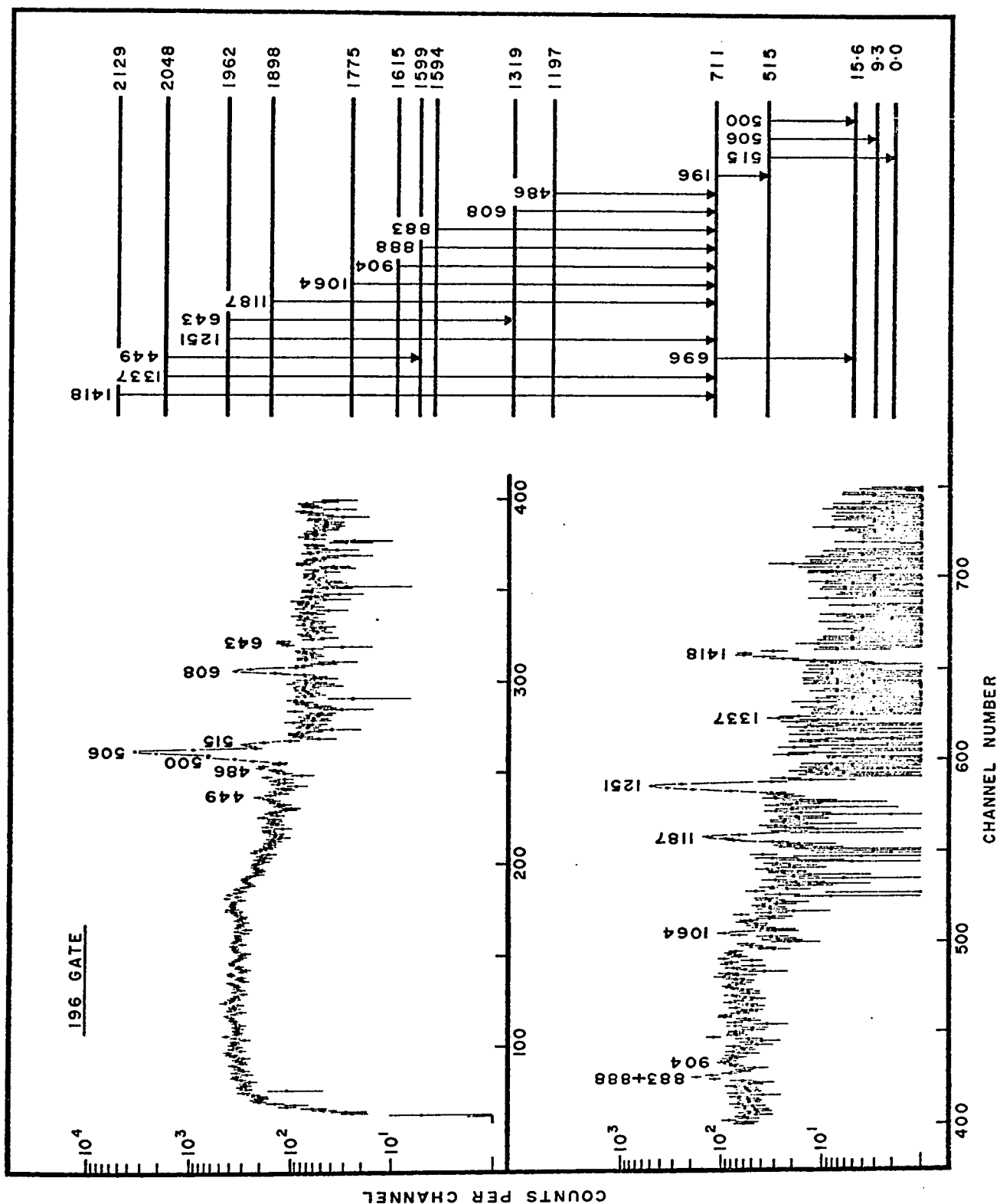


Fig. 3.22 The spectrum in coincidence with the 196 keV transition

gate. This evidence demonstrates that the 196 and 696 keV gamma rays depopulate the same level. The 696 keV transition and the 196-500, 196-506 and 196-515 keV cascades can all be accounted for by establishing levels at 515 and 711 keV. The 500, 506 and 515 keV gamma rays are transitions from the 515 keV level to the 15.6 and 9.3 keV states and to the ground state, respectively. A number of other levels are defined by the coincidence data of Fig. 3.22. Many of these levels can be identified with the levels established in Fig. 3.21 if it is assumed that the 591 keV transition populates the 15.6 keV state. Strong support for this assumption is provided by the 106-591 keV cascade, one mode of decay of the 711 keV state.

Since gates set on the 500 and 506 keV lines yielded identical coincidence spectra, the two spectra have been combined to improve the statistics. The resultant spectrum, which is presented in Fig. 3.23, also includes the 510 keV gamma ray together with a portion of the 513 keV transition. These latter peaks are responsible for the 506 keV peak in the coincidence spectrum. Unfortunately, the 510 keV peak is masked in this spectrum by the stronger 513 keV peak; however, the 510 keV transition is believed to populate the 506 keV level from a 1025 keV level which is established by energy fit. In addition to the 196, 510 and 513 keV transitions, a number of other gamma rays populate the 515 keV level directly from levels established with the 591 and 196 keV coincidence

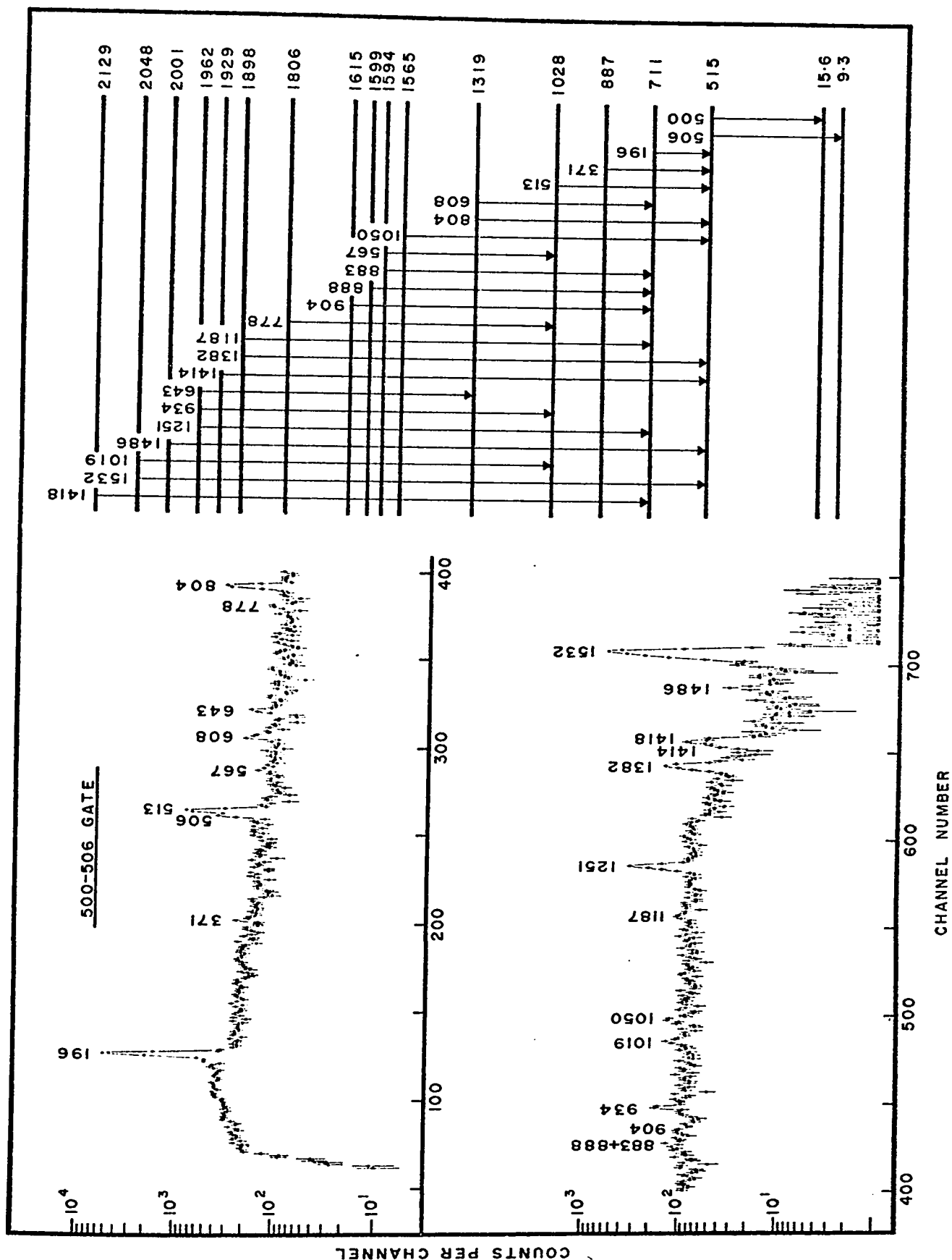


Fig. 3.23 The spectrum in coincidence with the 500+506 keV gate. The gate also contains portions of the 510+513 keV photopeaks

spectra.

The spectrum in coincidence with the 1532 keV gate, which is displayed in the bottom half of Fig. 3.24, gives further support to the decay pattern of the 515 keV level. The weak 877 peak arises from a second 1532 keV gamma ray which feeds the 877 keV state. The spectrum in coincidence with the 1251 keV gate, shown in the upper half of the figure, clearly demonstrates the relationship between the 711, 606, 515, 15.6 and 9.3 keV levels.

Fig. 3.25 presents the spectra in coincidence with the 871+877 and the 1161 keV gates associated with the 887 keV state. Since gates set separately on the 871 and 877 keV peaks yielded identical spectra, the two spectra were combined for the purpose of analysis. The weak 196 keV peak in the upper spectrum is due to a portion of the 883 keV peak contained in the 877 keV gate. Since the 888 and 1010 keV peaks in the 1161 keV gated spectrum do not account for the full intensity of the gamma rays seen in singles, both of these transitions are doublets. The weaker 888 keV component feeds the 711 keV state (see Fig. 3.22).

The spectrum in coincidence with the 934 keV gate, displayed in Fig. 3.26, shows the transitions which depopulate the 1028 keV level. The lower half of the figure displays the spectrum in coincidence with a gate composed of the 1010, 1012 and 1019 keV gamma rays. The 1012-1019 keV coincidence relationship accommodates only 31% of the 1019 keV singles intensity. The other 1019 keV line probably feeds the 9.3 keV

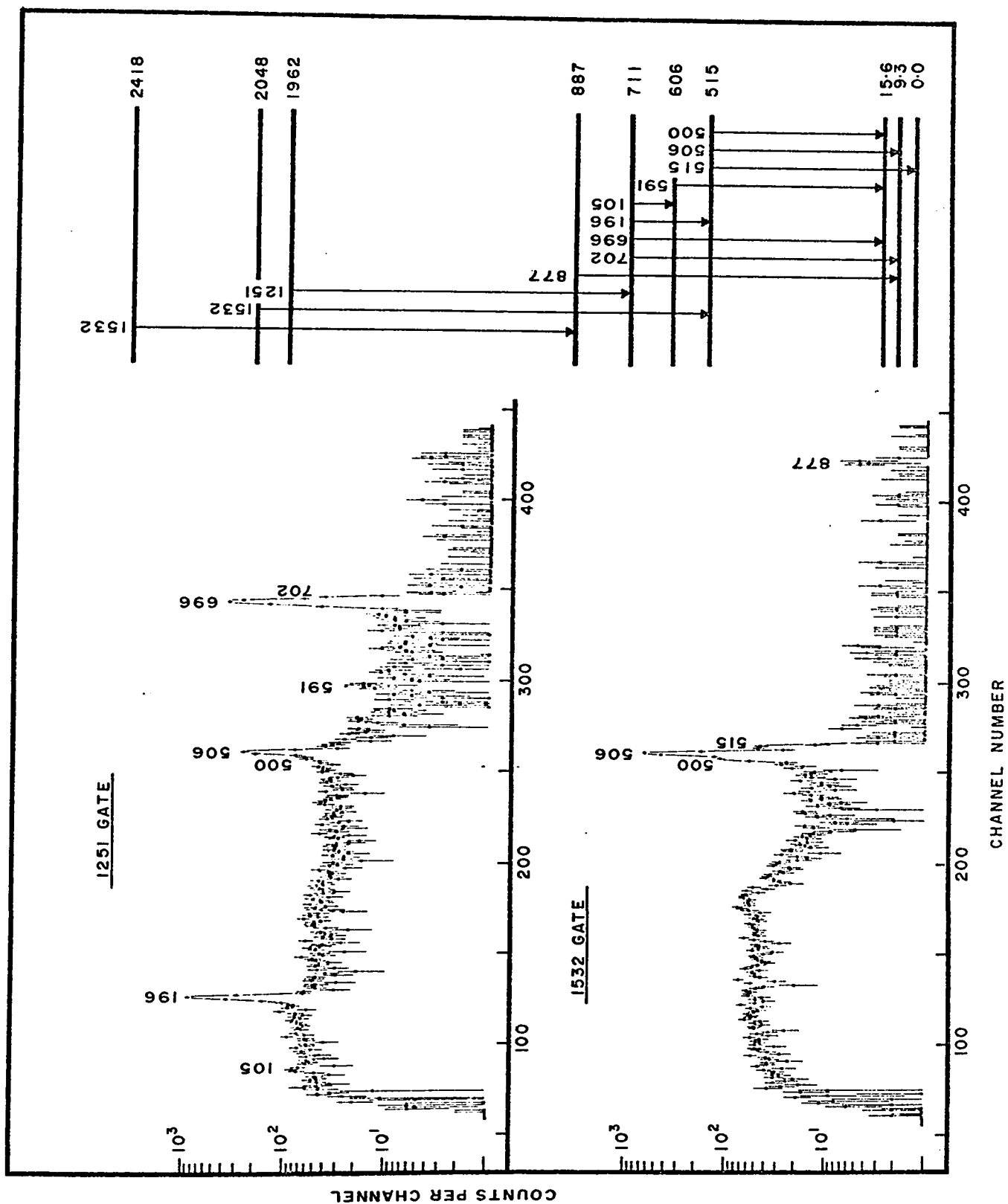


Fig. 3.24 Spectra in coincidence with the 1251 and 1532 keV gates

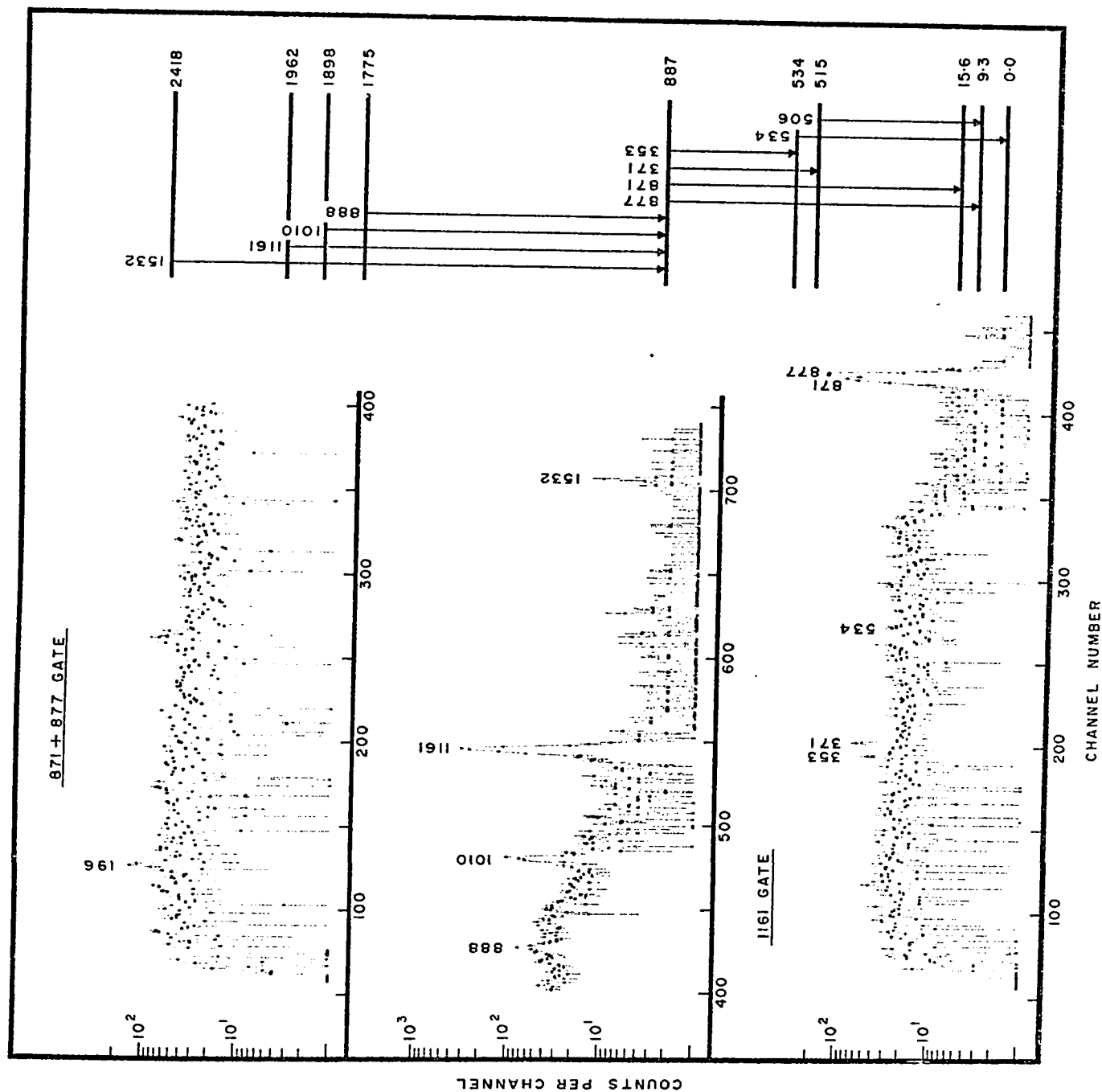


Fig. 3.25 Spectra in coincidence with the 871+877 and 1161 keV gates

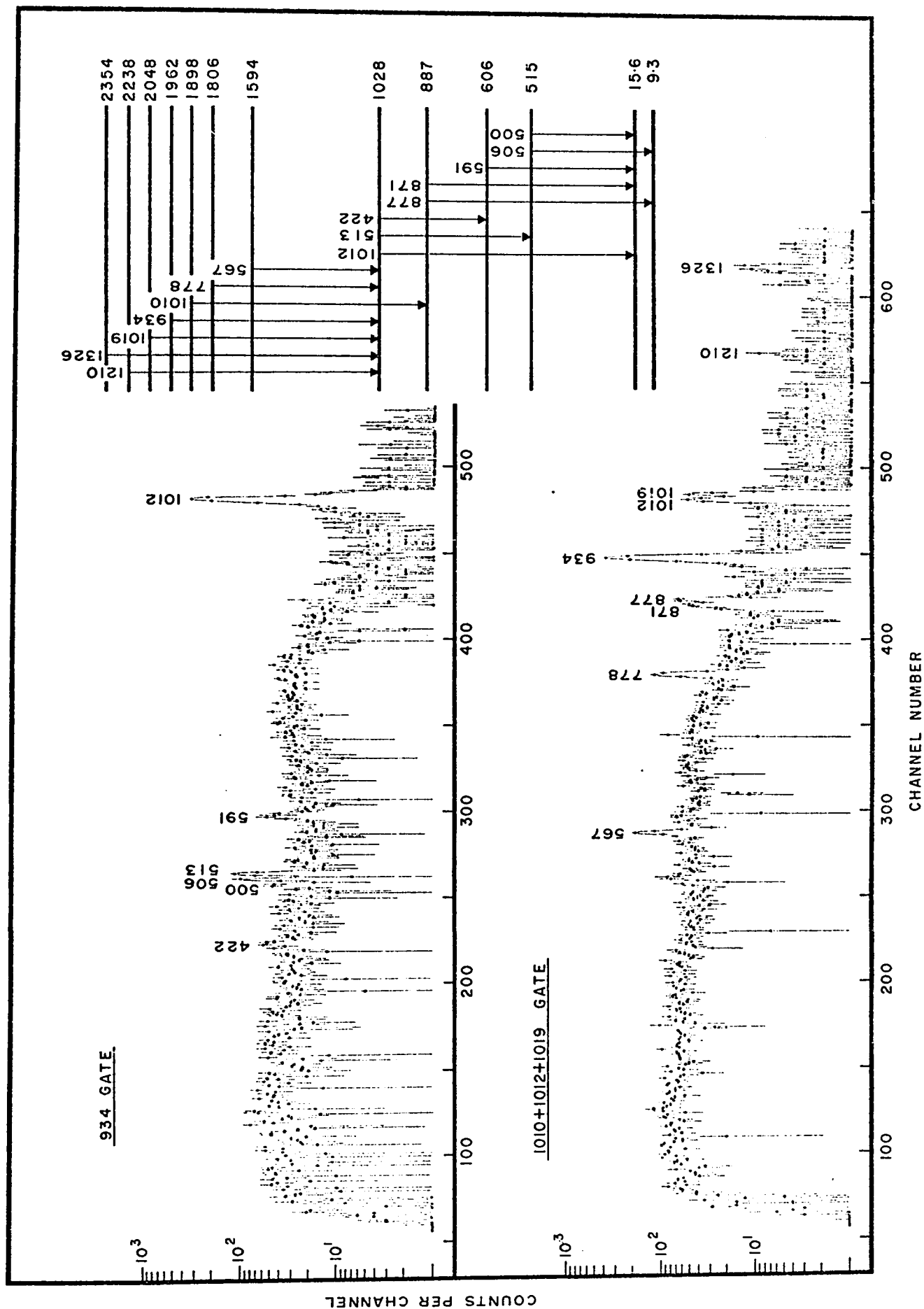


Fig. 3.26 Spectra in coincidence with the 934 keV and the 1010+1012+1019 keV gates

level; there is some evidence for a weak component on the high energy side of the 1012 keV peak in the 934 keV gate. Only 22% of the intensity of the 1010 keV gamma ray observed in singles feeds the 887 keV state. The stronger component does not appear to exhibit any coincidences, and therefore must feed one of the levels at 9.3, 15.6 or 207 keV. It has been used together with the weak 510 keV transition discussed above to define the 1025 keV level.

Fig. 3.27 presents the spectrum in coincidence with the 81 keV gate. This spectrum reveals 17 transitions in direct coincidence with the 81 keV line plus 8 others which are in indirect coincidence. The partial level scheme derived from this spectrum is consistent with evidence from other gates only if the 81 keV transition feeds the 207 keV isomeric state. This level is in turn depopulated by the 192 keV M2 transition. The assignment of the isomeric level at 207 keV is confirmed by the large number of gamma rays which can be placed by energy fit between the 207 keV level and levels established by coincidence information; most of these gamma rays are not seen in coincidence. The 1759 and 1840 keV lines observed in the 81 keV gate account for only 36% and 12% respectively, of the singles intensity. The other components of these gamma rays make no coincidences and have been placed on the basis of energy fit between well-established levels. The 1507 keV line has been identified only in the coincidence spectra and forms part of the 1507-334-81 keV triple cascade.

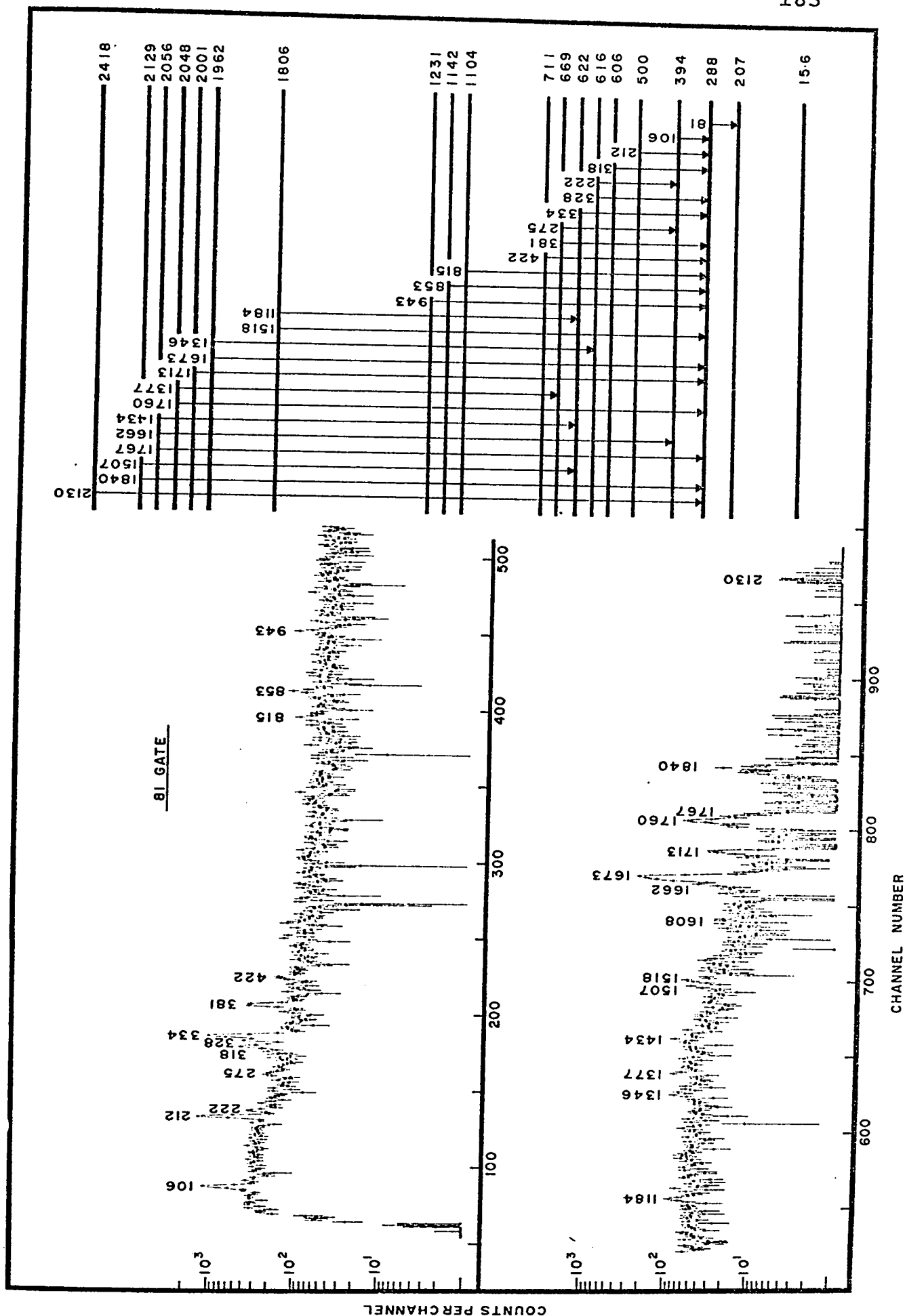


Fig. 3.27 The spectrum in coincidence with the 81 keV transition

The spectra in coincidence with the 409 and 1346 keV gamma rays (Fig. 3.28) show that the 409 keV transition connects the 616 keV state with the isomeric level. The 591 keV peak, which appears in the 1346 keV gate, is due to a portion of the 1355 keV line included in the gate; the 1356-591 keV lines form a strong cascade.

Fig. 3.29 displays the spectrum in coincidence with the 212 keV gamma ray. The 169, 602, 641, 819, 1429 and 1629 keV gamma rays are seen only in coincidence, but all of them are transitions between the 500 keV state and other well-established levels. The weak 567 and 778 keV peaks do not fit into the decay scheme, but could conceivably define levels at 1067 and 1278 keV. The weak 651 and 762 keV gamma rays, which are unclassified in Table 3.7, could be transitions from the 1278 keV state to the well-established 616 and 515 keV levels, respectively. Since all of these transitions are very weak, the 1067 and 1278 keV levels have not been included in the decay scheme. The lower half of Fig. 3.29 shows the spectrum associated with the 1184 + 1186 keV gate. The 81 and 334 keV peaks are in coincidence with the 1184 keV transition and the other three with the 1186 keV gamma ray.

In addition to revealing the transitions associated with the decay of the 394 keV level, Fig. 3.30 presents evidence that both the 379 and 1662 keV transitions are unresolved doublets. The spectrum in coincidence with the 222 keV gate shows four low energy transitions which deexcite the 394 keV

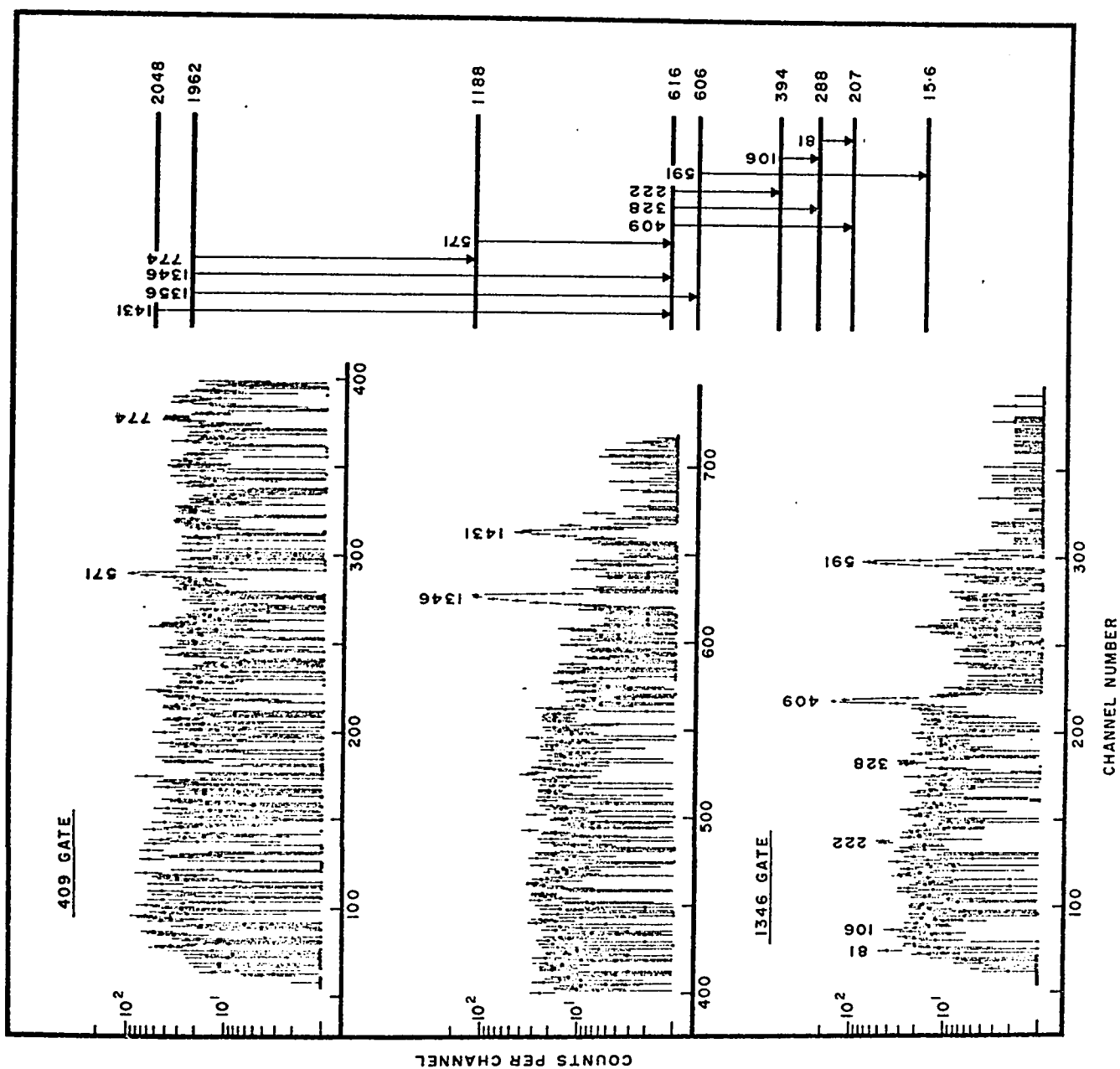


Fig. 3.28 Spectra in coincidence with the 409 and 1346 keV gamma rays

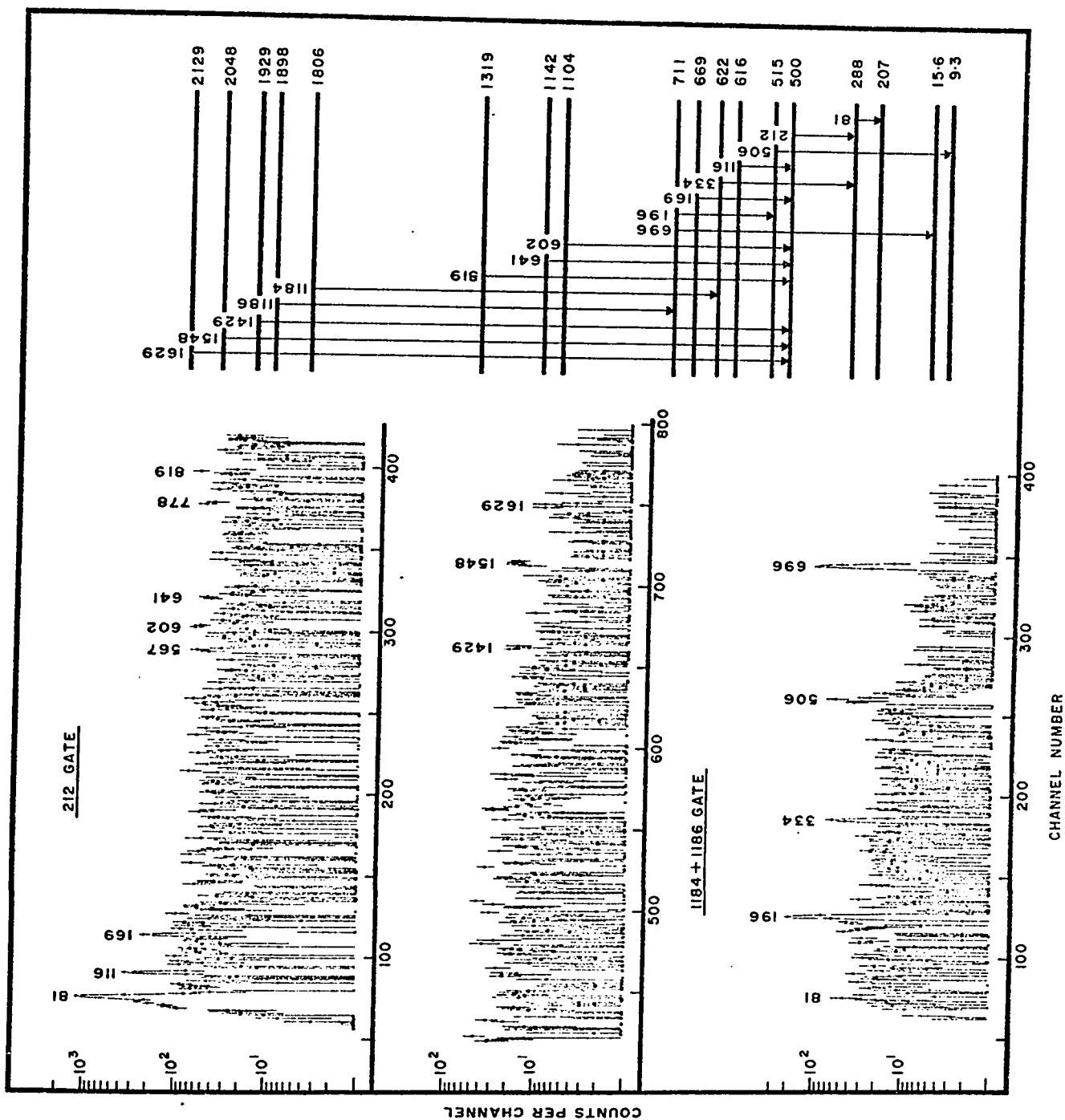


Fig. 3.29 Coincidence spectra associated with the 212, 1184 and 1186 keV transitions

state. The 385 keV gamma ray seen in coincidence was masked in the singles data by a line from the decay of ^{101}Tc (see Table 3.9) and therefore a reliable singles intensity could not be assigned. The coincidence intensity of the 379 keV transition appears to be too weak to account for the full intensity seen in singles. Unfortunately, the high error in the 222 keV singles intensity, the effect of internal conversion, and the loss of coincidence efficiency for the 81 and 106 keV transitions make it rather difficult to define an intensity for the 379 keV gamma ray observed in coincidence with the 222 keV transition. The 1653+1662 keV gated spectrum shows three peaks (81, 106 and 379 keV) from coincidences with the 1653 and 1662 keV gamma rays which feed the 394 keV level, and three peaks (370, 379 and 560 keV) which arise from transitions in coincidence with a second 1662 keV gamma ray. The latter three peaks define a level at 1678 keV. The ratio of the intensities of the 1653 and 1662 keV peaks in the spectra in coincidence with the 106 and 379 keV lines (not shown) indicate that less than 25% of the 1662 keV intensity is associated with the 394 keV level. Analysis of all these data shows that 70% of the 379 keV singles intensity is associated with the transition deexciting the 394 keV level and 30% with the transition populating the 1678 keV level. The 81 keV coincidence intensity in the 1653+1662 keV gate is too strong due to a portion of the 1673 keV transition being included in the gate; the 1673 and 81 keV gamma rays form a strong cascade (Fig. 3.27).

The coincidence measurements show that the unresolved

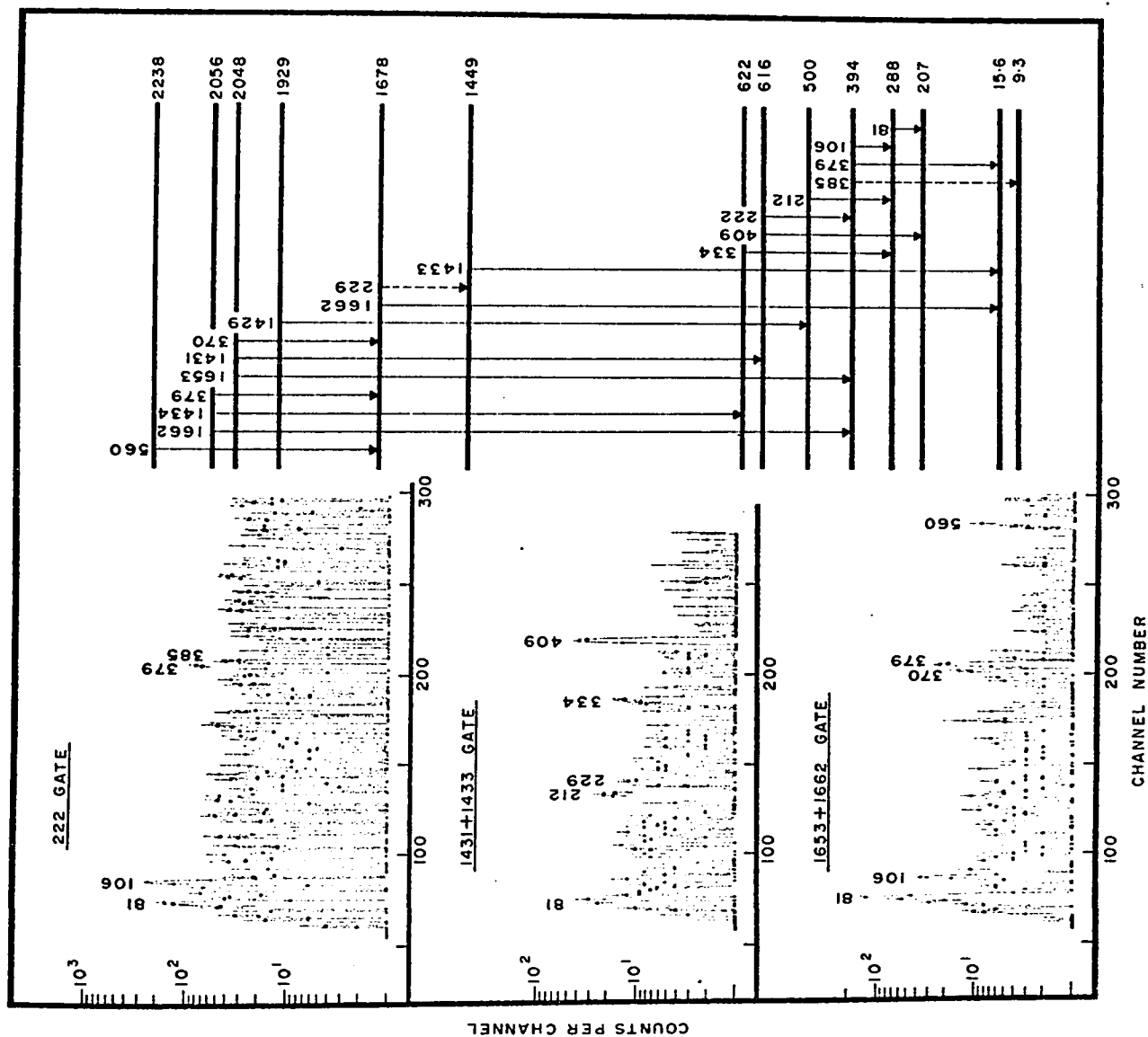


Fig. 3.30 Spectra in coincidence with the 222, 1431+1433, and 1653+1662 keV gates

structure at ~ 1432 keV in the singles data of Fig. 3.17 is composed of four components. It has already been shown in Fig. 3.29 that a weak 1429 keV transition feeds the 500 keV state. Confirmation of this coincidence relationship is also seen in the 1431+1433 keV gated spectrum of Fig. 3.30. The presence of a strong 409 keV peak in this spectrum arises from the 1431 keV gamma ray connecting the 2048 and 617 keV states. About 10% of the 1433 keV intensity can be accounted for by the 2056 keV to 622 keV transition which brings up the 334 keV peak in the coincidence spectrum. The remaining 90% of the 1433 keV intensity must arise from a transition to the 9.3, 15.6 or 207 keV states. The presence of a weak 229 keV peak in the coincidence spectrum and energy sums suggest that a 229-1433 keV cascade deexcites the 1678 keV level and defines a state at 1449 keV.

Fig. 3.31 displays the spectra in coincidence with the 449 and 1064+1066 keV gates. The former clearly shows that the 449 keV transition feeds the 1599 keV level and defines several cascades from that level. The ordering of the 1066-534 keV cascade is determined from the 1161 keV gated spectrum of Fig. 3.25 which shows that the 353 keV gamma ray originating in the 887 keV level also feeds the 534 keV state. The 888 keV peak in Fig. 3.31 is due to the weaker member of the 888 keV doublet of Table 3.7. The 1064+1066 gated spectrum confirms the conclusions reached above.

The upper spectrum of Fig. 3.32 is associated with a

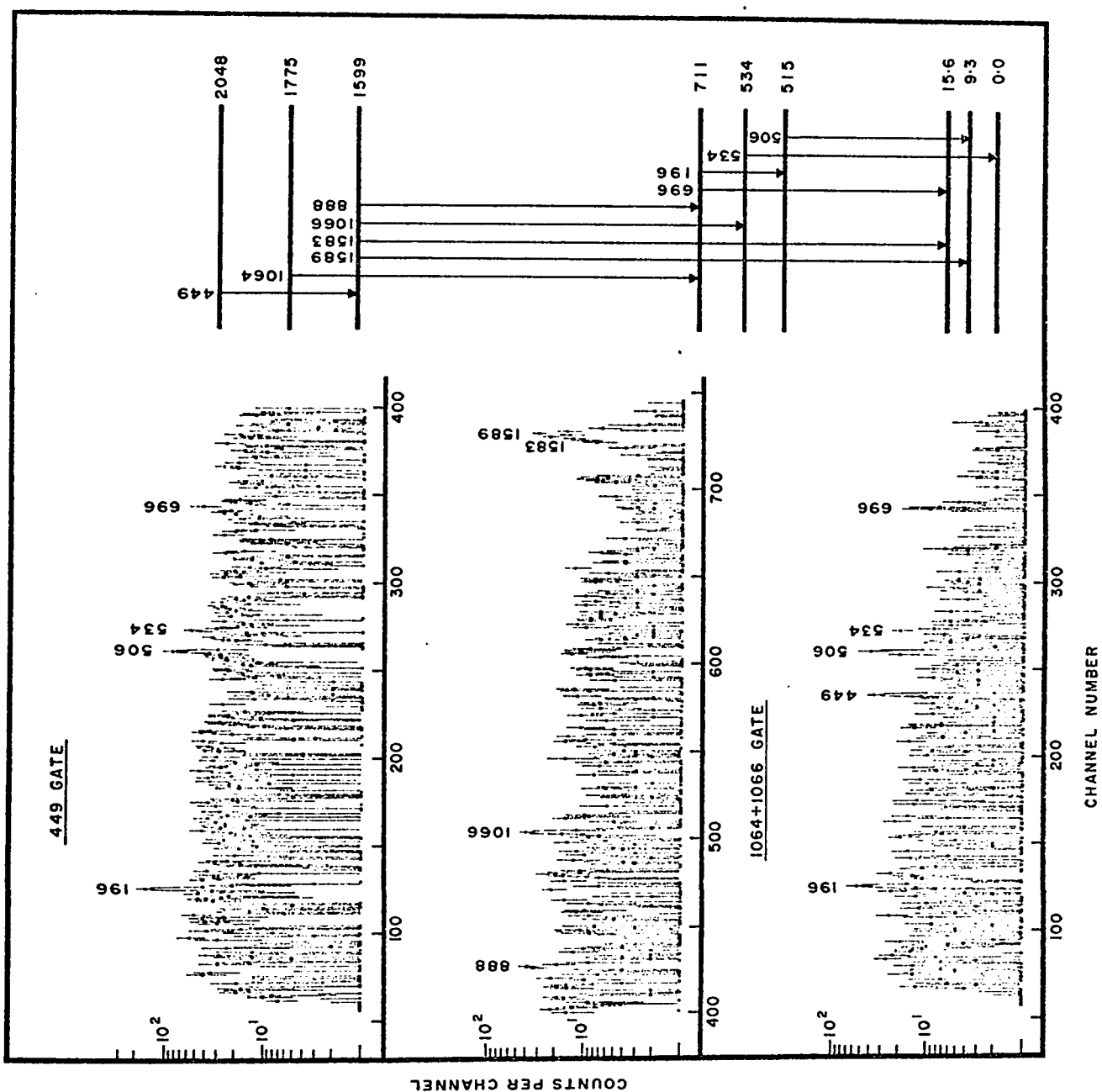


Fig. 3.31 Spectra in coincidence with the 449 and 1064+1066 keV gates

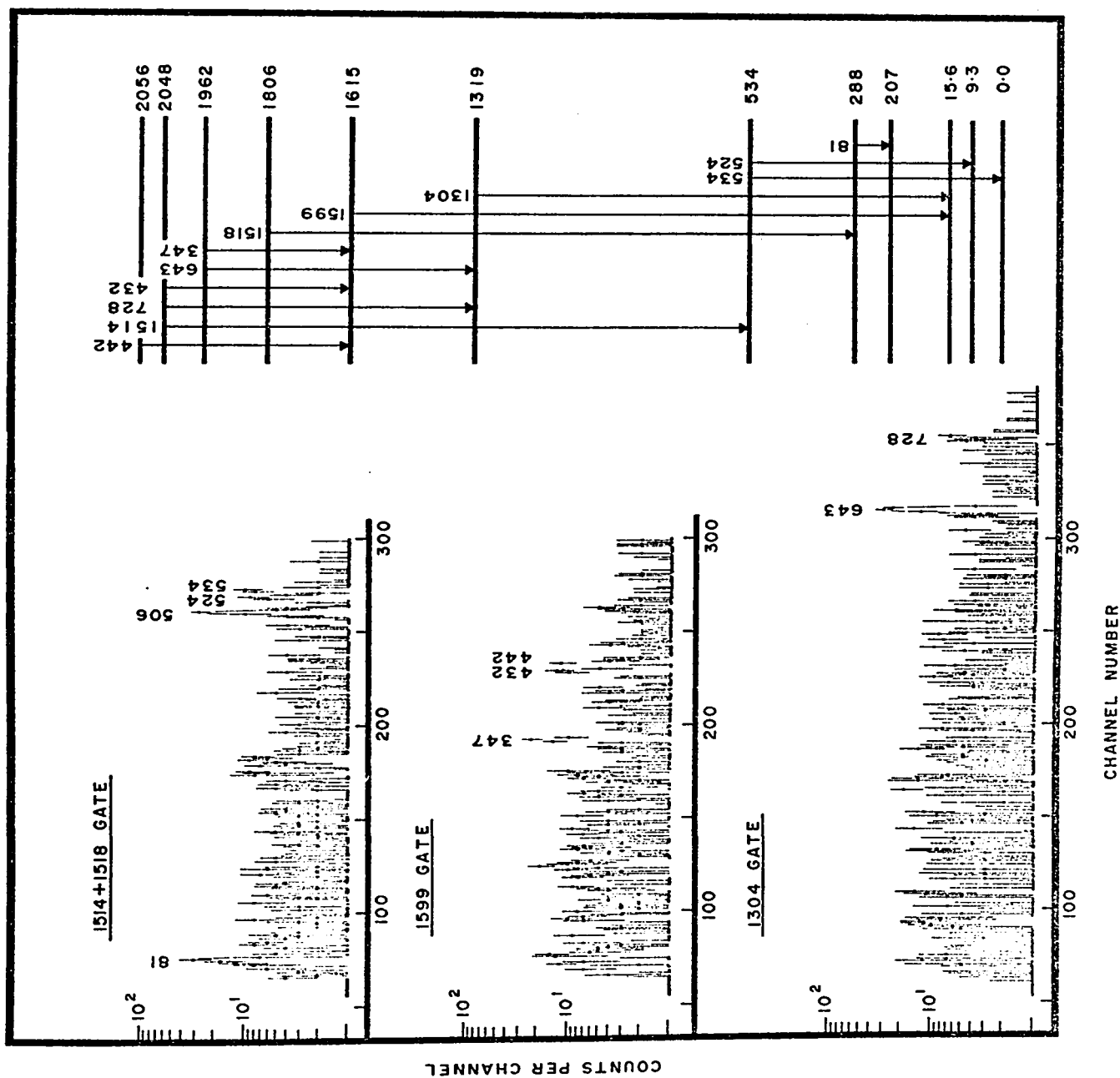


Fig. 3.32 Spectra in coincidence with the 1514+1518, 1599 and 1304 keV gates

gate containing the 1514 and 1518 keV peaks and a trace of the strong 1532 keV peak; the latter is responsible for the 506 keV peak in the spectrum. The 1514 keV transition feeds the 534 keV level which is deexcited with the 524 and 534 keV lines. The 1518 keV gamma ray is in coincidence with the 81 keV transition. The spectra in coincidence with the 1304 and 1599 keV gamma rays require no comment.

3.3.5 The Level Scheme of ^{101}Tc

The energies and intensities of the transitions in ^{101}Tc which have been derived from the singles and coincidence measurements of the previous section are presented in Table 3.7. The intensities quoted are per 1000 decays; the normalization necessary was obtained from the gamma-ray intensity balance, assuming that there is no beta feeding to the ground state or the first two excited states at 9.3 and 15.6 keV. The table also includes the classification of each transition on the basis of the decay scheme given in Figs. 3.33 and 3.34. A number of transitions, which appeared to be single gamma rays in the spectra of Figs. 3.16 and 3.17., were shown to be doublets in the coincidence data. The "singles" intensity of such doublets has been divided among the components on the basis of the coincidence measurements; a double asterisk is used to designate these transitions. In addition, a number of weak transitions were identified only in the coincidence spectra and are designated by a single asterisk. The superscript letter T is used to identify transitions for which the

TABLE 3.7

Transitions in ^{101}Tc from the Decay of ^{101}Mo

Energy (keV)	Intensity (per 1000 decays)	Classification	Energy (keV)	Intensity (per 1000 decays)	Classification
6.281 $\pm$ 0.007	800 $\pm$ 27 ^T	15.601 $\rightarrow$ 9.317	347.5 $\pm$ 0.4	0.9 $\pm$ 0.2	1962.12 $\rightarrow$ 1614.66
9.317 $\pm$ 0.010	970 $\pm$ 10 ^T	9.317 $\rightarrow$ 0	352.8 $\pm$ 0.5	1.07 $\pm$ 0.29	886.65 $\rightarrow$ 533.55
15.606 $\pm$ 0.015	19 $\pm$ 3 ^T	15.601 $\rightarrow$ 0	368.8 $\pm$ 0.7	0.95 $\pm$ 0.4	2047.58 $\rightarrow$ 1677.8
80.92 $\pm$ 0.03	54.9 $\pm$ 5.3 ^T	288.45 $\rightarrow$ 207.53	371.2 $\pm$ 0.35	1.91 $\pm$ 0.27	886.65 $\rightarrow$ 515.19
104.70 $\pm$ 0.08	1.54 $\pm$ 0.14	711.12 $\rightarrow$ 606.41	378.85 $\pm$ 0.20	[0.95 $\pm$ 0.3* 2.25 $\pm$ 0.70*	2056.42 $\rightarrow$ 1677.8
105.95 $\pm$ 0.05	2.28 $\pm$ 0.19	394.40 $\rightarrow$ 288.45			
115.8 $\pm$ 0.15	0.29 $\pm$ 0.06	616.20 $\rightarrow$ 500.45	381.0 $\pm$ 0.15	3.50 $\pm$ 0.45	394.40 $\rightarrow$ 15.601
169.0 $\pm$ 0.3	0.23 $\pm$ 0.08**	669.45 $\rightarrow$ 500.45	384.5 $\pm$ 0.7	0.6 $\pm$ 0.2**	669.45 $\rightarrow$ 288.45
191.93 $\pm$ 0.04	223 $\pm$ 16 ^T	207.53 $\rightarrow$ 15.601	398.85 $\pm$ 0.10	7.86 $\pm$ 0.68	394.40 $\rightarrow$ 9.317
195.94 $\pm$ 0.05	27.6 $\pm$ 2.3	711.12 $\rightarrow$ 515.19	408.59 $\pm$ 0.10	14.0 $\pm$ 1.2	606.41 $\rightarrow$ 207.53
212.00 $\pm$ 0.08	4.30 $\pm$ 0.45	500.45 $\rightarrow$ 288.45	421.41 $\pm$ 0.14	5.23 $\pm$ 0.53	616.20 $\rightarrow$ 207.53
221.8 $\pm$ 0.4	0.95 $\pm$ 0.3	616.20 $\rightarrow$ 394.40	422.4 $\pm$ 0.5	0.8 $\pm$ 0.3**	1027.90 $\rightarrow$ 606.41
229 $\pm$ 1	0.19 $\pm$ 0.08*	1677.8 $\rightarrow$ 1449.0	432.9 $\pm$ 0.4	1.07 $\pm$ 0.20	711.12 $\rightarrow$ 288.45
273.0 $\pm$ 1.0	0.4 $\pm$ 0.2	288.45 $\rightarrow$ 15.601 [†]	441.9 $\pm$ 0.6	0.65 $\pm$ 0.15	2047.58 $\rightarrow$ 1614.66
275.5 $\pm$ 0.8	0.6 $\pm$ 0.2	669.45 $\rightarrow$ 394.40	448.63 $\pm$ 0.12	6.22 $\pm$ 0.56	2056.42 $\rightarrow$ 1614.66
318.0 $\pm$ 0.2	2.37 $\pm$ 0.37	606.41 $\rightarrow$ 288.45	452.6 $\pm$ 0.45	1.40 $\pm$ 0.19	2047.58 $\rightarrow$ 1598.94
327.85 $\pm$ 0.2	1.89 $\pm$ 0.25	616.20 $\rightarrow$ 288.45	469.3 $\pm$ 0.45	1.09 $\pm$ 0.19	2047.58 $\rightarrow$ 1594.39
333.70 $\pm$ 0.10	6.57 $\pm$ 0.64	622.15 $\rightarrow$ 288.45			

Energy (keV)	Intensity (per 1000 decays)	Classification	Energy (keV)	Intensity (per 1000 decays)	Classification
473.3 $\pm$ 0.8	0.35 $\pm$ 0.15		641 $\pm$ 1	0.4 $\pm$ 0.1*	1141.5 $\rightarrow$ 500.45
482.8 $\pm$ 0.6	0.75 $\pm$ 0.15	2047.58 $\rightarrow$ 1565.09	642.70 $\pm$ 0.11	11.0 $\pm$ 1.0	1962.12 $\rightarrow$ 1319.46
486.2 $\pm$ 0.6	0.85 $\pm$ 0.15	1197.3 $\rightarrow$ 711.12 [†]	651.1 $\pm$ 1.0	0.6 $\pm$ 0.2	1278 $\rightarrow$ 616.20 [†]
492.0 $\pm$ 0.6	0.70 $\pm$ 0.15	500.45 $\rightarrow$ 9.317 [†]	660.8 $\pm$ 0.3	2.57 $\pm$ 0.54**	669.45 $\rightarrow$ 9.317
499.59 $\pm$ 0.10	13.4 $\pm$ 1.1	515.19 $\rightarrow$ 15.601	695.53 $\pm$ 0.07	66.3 $\pm$ 5.4	711.12 $\rightarrow$ 15.601
505.88 $\pm$ 0.05	114 $\pm$ 8	515.19 $\rightarrow$ 9.317	701.8 $\pm$ 0.3	4.03 $\pm$ 0.49	711.12 $\rightarrow$ 9.317
510.36 $\pm$ 0.20	3.23 $\pm$ 0.68	1025.59 $\rightarrow$ 515.19	713.02 $\pm$ 0.10	31.1 $\pm$ 2.7	1319.46 $\rightarrow$ 606.41
512.71 $\pm$ 0.13	12.0 $\pm$ 1.2	1027.90 $\rightarrow$ 515.19	728.6 $\pm$ 0.6	0.90 $\pm$ 0.25	2047.58 $\rightarrow$ 1319.46
515.15 $\pm$ 0.20	7.29 $\pm$ 0.78**	515.19 $\rightarrow$ 0	732.9 $\pm$ 0.15	2.29 $\pm$ 0.29	742.22 $\rightarrow$ 9.317
524.2 $\pm$ 0.4	1.17 $\pm$ 0.21	533.55 $\rightarrow$ 9.317	751.3 $\pm$ 1.0	0.45 $\pm$ 0.15	
533.55 $\pm$ 0.15	3.81 $\pm$ 0.54	533.55 $\rightarrow$ 0	761.8 $\pm$ 0.8	0.45 $\pm$ 0.15	1278 $\rightarrow$ 515.19 [†]
560.4 $\pm$ 0.45	0.8 $\pm$ 0.2	2237.8 $\rightarrow$ 1677.8	774.1 $\pm$ 0.2	3.52 $\pm$ 0.39	1962.12 $\rightarrow$ 1188.04
566.51 $\pm$ 0.10	[7.80 $\pm$ 0.68 0.25 $\pm$ 0.1*]	1594.39 $\rightarrow$ 1027.90 1067 $\rightarrow$ 500.45 [†]	778.30 $\pm$ 0.10	[8.46 $\pm$ 0.78 0.25 $\pm$ 0.1*]	1806.28 $\rightarrow$ 1027.90
571.3 $\pm$ 0.35	1.59 $\pm$ 0.23	1188.04 $\rightarrow$ 616.20	789.8 $\pm$ 0.6	1.23 $\pm$ 0.23	1278 $\rightarrow$ 500.45 [†]
590.82 $\pm$ 0.05	194 $\pm$ 10	606.41 $\rightarrow$ 15.601	804.25 $\pm$ 0.13	8.62 $\pm$ 0.88	1319.46 $\rightarrow$ 515.19
603 $\pm$ 1	0.55 $\pm$ 0.2*	1103.5 $\rightarrow$ 500.45	815.2 $\pm$ 0.4	1.40 $\pm$ 0.25	1103.5 $\rightarrow$ 288.45
608.31 $\pm$ 0.10	9.80 $\pm$ 0.88	1319.46 $\rightarrow$ 711.12	819.0 $\pm$ 0.8	0.6 $\pm$ 0.2**	1319.46 $\rightarrow$ 500.45
624.6 $\pm$ 0.7	0.6 $\pm$ 0.2*	1231.3 $\rightarrow$ 606.41	836.4 $\pm$ 0.8	0.45 $\pm$ 0.15	

Energy (keV)	Intensity (per 1000 decays)	Classification	Energy (keV)	Intensity (per 1000 decays)	Classification
846.7 $\pm$ 1.0	0.2 $\pm$ 0.1		1030.0 $\pm$ 0.8	0.65 $\pm$ 0.15	
853.1 $\pm$ 0.25	2.43 $\pm$ 0.29	1141.5 $\rightarrow$ 288.45	1049.9 $\pm$ 0.2	3.21 $\pm$ 0.33	1565.09 $\rightarrow$ 515.19
859.35 $\pm$ 0.45	1.38 $\pm$ 0.23	2047.58 $\rightarrow$ 1188.04	1064.5 $\pm$ 0.6	2.75 $\pm$ 0.60**	1775.3 $\rightarrow$ 711.12
871.05 $\pm$ 0.11	16.6 $\pm$ 1.4	886.65 $\rightarrow$ 15.601	1066.0 $\pm$ 1.0	1.45 $\pm$ 0.40**	1598.94 $\rightarrow$ 533.55
877.38 $\pm$ 0.10	31.3 $\pm$ 2.7	886.65 $\rightarrow$ 9.317	1092.8 $\pm$ 0.5	0.95 $\pm$ 0.15	
883.25 $\pm$ 0.16	6.44 $\pm$ 0.62	1594.39 $\rightarrow$ 711.12	1160.93 $\pm$ 0.12	36.0 $\pm$ 3.1	2047.58 $\rightarrow$ 886.65
887.9 $\pm$ 0.25	[1.45 $\pm$ 0.40* 2.15 $\pm$ 0.60*]	1598.94 $\rightarrow$ 711.12 1775.3 $\rightarrow$ 886.65	1168.5 $\pm$ 0.5	1.75 $\pm$ 0.35	1775.3 $\rightarrow$ 606.41
895.9 $\pm$ 0.5	1.46 $\pm$ 0.27	1103.5 $\rightarrow$ 207.53	1183.7 $\pm$ 1.0	1.85 $\pm$ 0.50	1806.28 $\rightarrow$ 622.15
903.6 $\pm$ 0.4	1.87 $\pm$ 0.27	1614.66 $\rightarrow$ 711.12	1186.62 $\pm$ 0.15	9.68 $\pm$ 0.93	1897.54 $\rightarrow$ 711.12
934.20 $\pm$ 0.09	37.5 $\pm$ 3.3	1962.12 $\rightarrow$ 1027.90	1199.92 $\pm$ 0.13	16.2 $\pm$ 1.4	1806.28 $\rightarrow$ 606.41
943.2 $\pm$ 0.7	1.46 $\pm$ 0.39	1231.3 $\rightarrow$ 288.45	1210.2 $\pm$ 0.6	1.32 $\pm$ 0.29	2237.8 $\rightarrow$ 1027.90
980.7 $\pm$ 0.25	2.64 $\pm$ 0.31	1188.04 $\rightarrow$ 207.53	1251.10 $\pm$ 0.09	42.2 $\pm$ 3.5	1962.12 $\rightarrow$ 711.12
987.9 $\pm$ 0.35	1.79 $\pm$ 0.27	1594.38 $\rightarrow$ 606.41	1259.4 $\pm$ 0.7	0.9 $\pm$ 0.3	1775.3 $\rightarrow$ 515.19
	[29.8 $\pm$ 3.7** 8.6 $\pm$ 1.7*]	1025.59 $\rightarrow$ 15.601 1897.54 $\rightarrow$ 886.65	1286.5 $\pm$ 0.8	0.9 $\pm$ 0.2	1892.9 $\rightarrow$ 606.41 [†]
1010.0 $\pm$ 0.12	114 $\pm$ 10	1027.90 $\rightarrow$ 15.601	1291.9 $\pm$ 0.7	1.48 $\pm$ 0.23	1962.12 $\rightarrow$ 669.45
1012.35 $\pm$ 0.10	[10.2 $\pm$ 1.8** 4.6 $\pm$ 0.9*]	1027.90 $\rightarrow$ 9.317 2047.58 $\rightarrow$ 1027.90	1303.92 $\pm$ 0.11	23.9 $\pm$ 2.1	1319.46 $\rightarrow$ 15.601
1018.9 $\pm$ 0.15			1314 $\pm$ 1	1.26 $\pm$ 0.40*	2056.42 $\rightarrow$ 742.22
			1323.4 $\pm$ 0.8	1.35 $\pm$ 0.40**	1928.88 $\rightarrow$ 606.41
			1326.1 $\pm$ 0.6	2.70 $\pm$ 0.70**	2353.9 $\rightarrow$ 1027.90

Energy (keV)	Intensity (per 1000 decays)	Classification	Energy (keV)	Intensity (per 1000 decays)	Classification
1337.1 $\pm$ 0.8	1.94 $\pm$ 0.29	2047.58 $\rightarrow$ 711.12	1522.4 $\pm$ 0.5	2.65 $\pm$ 0.50	2129.18 $\rightarrow$ 606.41
1345.79 $\pm$ 0.15	9.49 $\pm$ 0.88	1962.12 $\rightarrow$ 616.20	1532.27 $\pm$ 0.17	[0.8 $\pm$ 0.25* 54.4 $\pm$ 4.7	2418.3 $\rightarrow$ 886.65
1355.60 $\pm$ 0.13	16.2 $\pm$ 1.4	1962.12 $\rightarrow$ 606.41			2047.58 $\rightarrow$ 515.19
1377.3 $\pm$ 0.45	2.14 $\pm$ 0.29	2047.58 $\rightarrow$ 669.45	1544.2 $\pm$ 0.7	0.7 $\pm$ 0.2	2150.3 $\rightarrow$ 606.41
1382.15 $\pm$ 0.15	11.1 $\pm$ 1.0	1897.54 $\rightarrow$ 515.19	1547.5 $\pm$ 0.5	1.46 $\pm$ 0.23	2047.58 $\rightarrow$ 500.45
1387.4 $\pm$ 0.8	0.9 $\pm$ 0.3**	2056.42 $\rightarrow$ 669.45	1583.0 $\pm$ 1.0	0.7 $\pm$ 0.2	1598.94 $\rightarrow$ 15.601
1394.5 $\pm$ 0.2	5.70 $\pm$ 0.53	2001.01 $\rightarrow$ 606.41	1589.6 $\pm$ 0.3	2.39 $\pm$ 0.27	1598.94 $\rightarrow$ 9.317
1413.5 $\pm$ 0.25	4.20 $\pm$ 0.39	1928.88 $\rightarrow$ 515.19	1599.05 $\pm$ 0.14	15.7 $\pm$ 1.4	1614.66 $\rightarrow$ 15.601
1418.0 $\pm$ 0.20	7.95 $\pm$ 0.68	2129.18 $\rightarrow$ 711.12	1608.4 $\pm$ 0.8	0.7 $\pm$ 0.2	1897.54 $\rightarrow$ 288.45
1428.5 $\pm$ 0.7	0.55 $\pm$ 0.2*	1928.88 $\rightarrow$ 500.45	1629 $\pm$ 1	0.25 $\pm$ 0.1*	2129.18 $\rightarrow$ 500.45
1430.6 $\pm$ 0.4	2.76 $\pm$ 0.43	2047.58 $\rightarrow$ 616.20	1646.2 $\pm$ 0.65	0.95 $\pm$ 0.2	2252.6 $\rightarrow$ 606.41
1433.4 $\pm$ 0.25	[7.50 $\pm$ 0.95** 0.9 $\pm$ 0.2*	1449.0 $\rightarrow$ 15.601	1653.0 $\pm$ 0.75	0.7 $\pm$ 0.15	2047.58 $\rightarrow$ 394.40
		2056.42 $\rightarrow$ 622.15	1662.2 $\pm$ 0.2	[1.35 $\pm$ 0.5* 5.05 $\pm$ 0.8**	2056.42 $\rightarrow$ 394.40
1440.6 $\pm$ 0.7	1.00 $\pm$ 0.25	2047.58 $\rightarrow$ 606.41			1677.8 $\rightarrow$ 15.601
1468.0 $\pm$ 0.6	0.95 $\pm$ 0.3		1673.56 $\pm$ 0.13	15.3 $\pm$ 1.3	1962.12 $\rightarrow$ 288.45
1473.6 $\pm$ 0.7	0.85 $\pm$ 0.2		1712.8 $\pm$ 0.3	1.94 $\pm$ 0.23	2001.01 $\rightarrow$ 288.45
1485.9 $\pm$ 0.6	0.90 $\pm$ 0.15	2001.01 $\rightarrow$ 515.19	1754.8 $\pm$ 0.3	4.24 $\pm$ 0.47	1962.12 $\rightarrow$ 207.53
1507 $\pm$ 1	0.8 $\pm$ 0.35	2129.18 $\rightarrow$ 622.15	1759.9 $\pm$ 0.2	3.7 $\pm$ 0.9*	2047.58 $\rightarrow$ 288.45
1514.4 $\pm$ 0.7	1.65 $\pm$ 0.40	2047.58 $\rightarrow$ 533.55		6.4 $\pm$ 1.3**	1775.3 $\rightarrow$ 15.601
1518.9 $\pm$ 0.7	2.50 $\pm$ 0.50	1806.28 $\rightarrow$ 288.45	1767.4 $\pm$ 0.7	1.17 $\pm$ 0.25	2056.42 $\rightarrow$ 288.45

Energy (keV)	Intensity (per 1000 decays)	Classification	T	Total Transition Intensity
1840.40 $\pm$ 0.20	11.1 $\pm$ 1.1**	2047.58 $\rightarrow$ 207.53		
	1.5 $\pm$ 0.5*	2129.18 $\rightarrow$ 288.45	*	Seen Only in Coincidence
1882.4 $\pm$ 0.6	0.7 $\pm$ 0.15	1897.54 $\rightarrow$ 15.601	**	Intensity Derived from Consideration of Singles and Coincidence Data
1913.7 $\pm$ 1.0	0.3 $\pm$ 0.1	1928.88 $\rightarrow$ 15.601		
1922.0 $\pm$ 0.6	0.7 $\pm$ 0.1	2129.18 $\rightarrow$ 207.53	†	Tentative Assignment
1942.5 $\pm$ 0.7	0.6 $\pm$ 0.15	2150.3 $\rightarrow$ 207.53		
1946.8 $\pm$ 0.6	0.65 $\pm$ 0.15	1962.12 $\rightarrow$ 15.601		
2031.95 $\pm$ 0.20	60.7 $\pm$ 5.1	2047.48 $\rightarrow$ 15.601		
2040.9 $\pm$ 0.20	20.4 $\pm$ 1.7	2056.42 $\rightarrow$ 15.601		
2046.8 $\pm$ 0.5	1.15 $\pm$ 0.20	2056.42 $\rightarrow$ 9.317		
2088.4 $\pm$ 0.25	6.75 $\pm$ 0.58			
2113.7 $\pm$ 0.25	5.52 $\pm$ 0.49	2129.18 $\rightarrow$ 15.601		
2130.3 $\pm$ 0.7	0.29 $\pm$ 0.06	2418.3 $\rightarrow$ 288.45		
2146.0 $\pm$ 1.0	0.18 $\pm$ 0.04	2353.9 $\rightarrow$ 207.53		
2161.1 $\pm$ 0.9	0.18 $\pm$ 0.04			
2202.6 $\pm$ 0.7	0.25 $\pm$ 0.04			
2221.6 $\pm$ 0.5	1.52 $\pm$ 0.16	2237.8 $\rightarrow$ 15.601		
2402 $\pm$ 1	0.19 $\pm$ 0.04	2418.3 $\rightarrow$ 15.601		

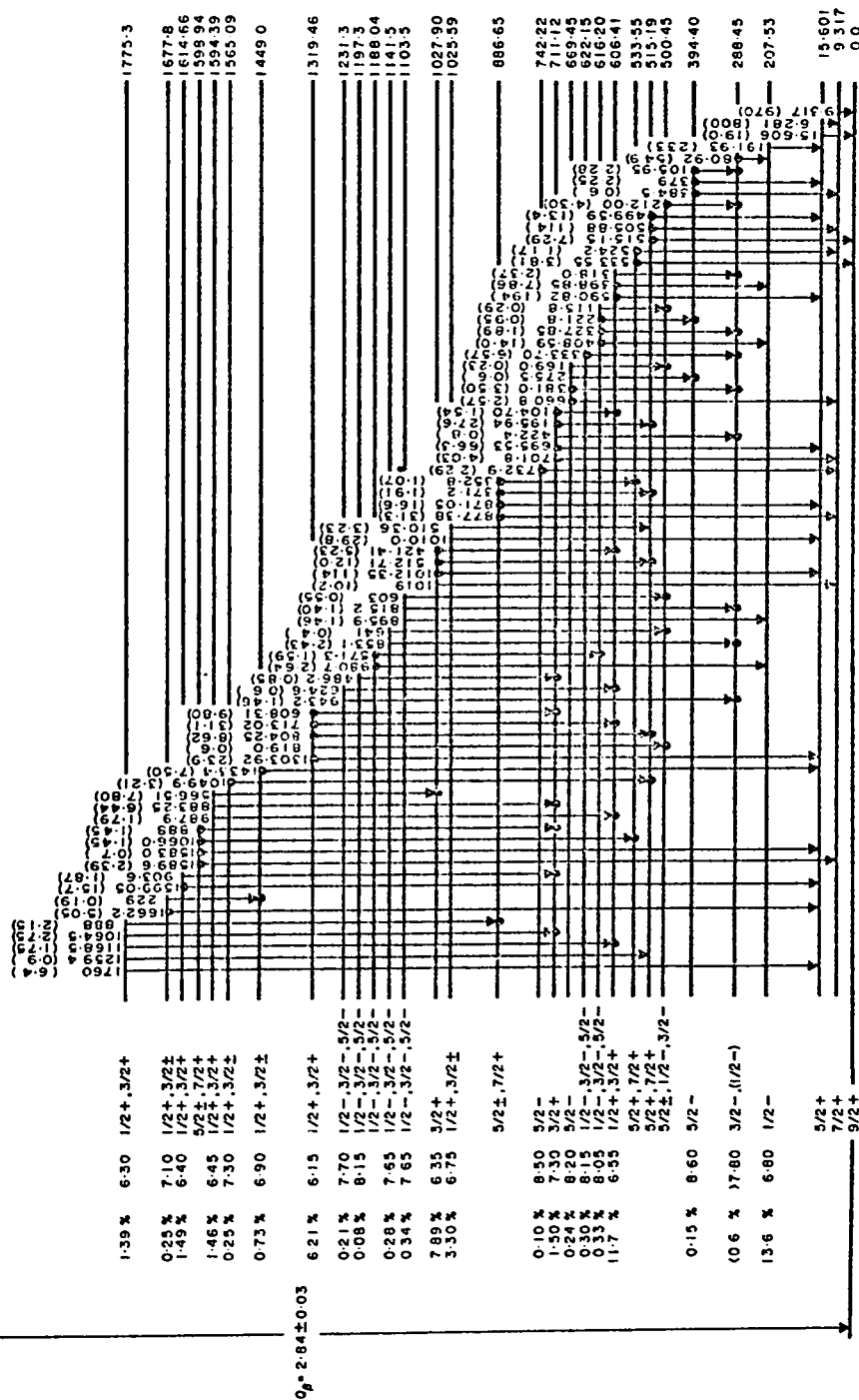


Fig. 3.33 Level scheme of ^{101}Tc showing the properties of the levels below 1775 keV. A dot at the end of a transition feeding a level indicates that it is in coincidence with one or more of the transitions deexciting the level. The spin and parity assignments shown are our best estimates from the available data

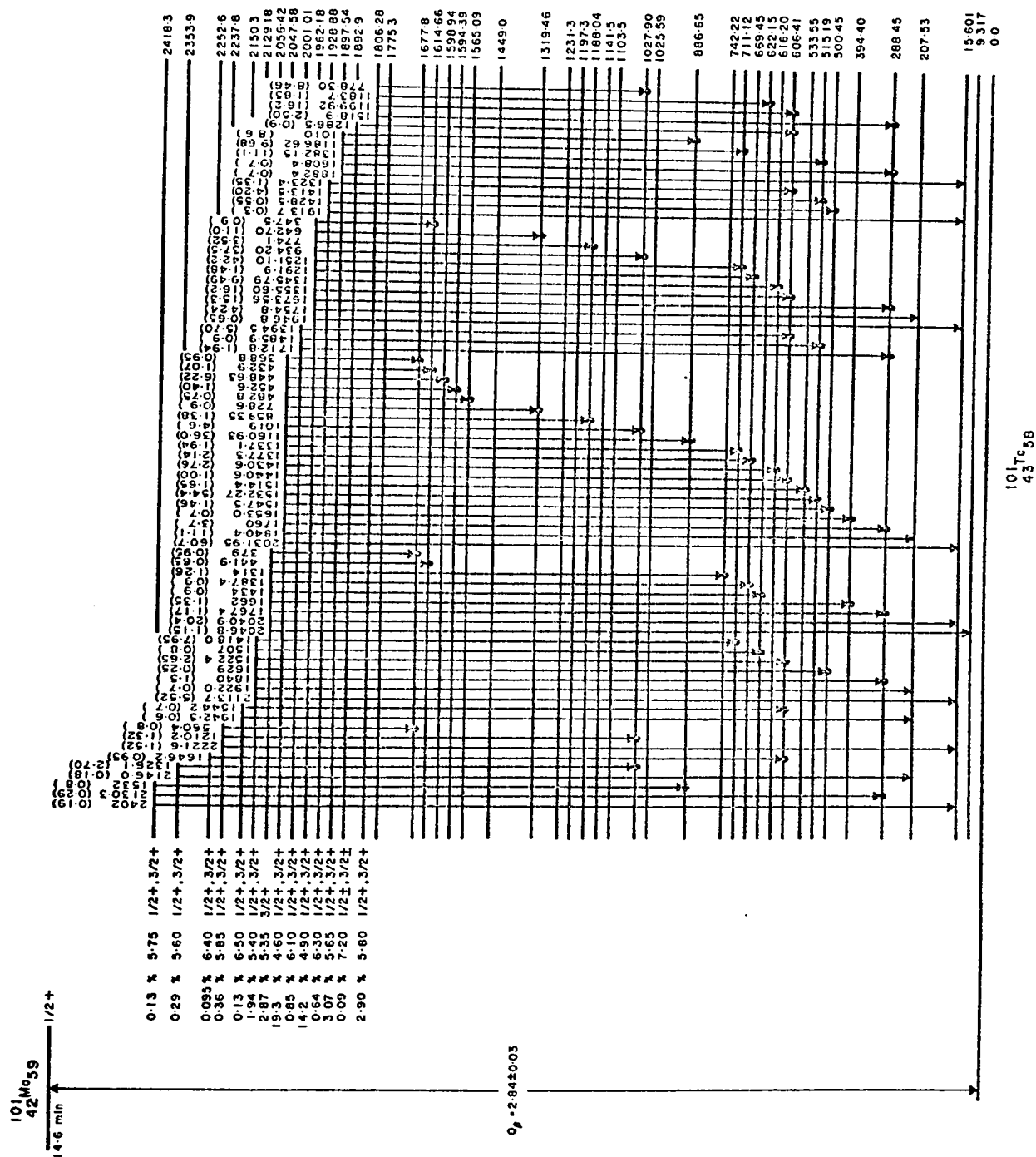


Fig. 3.34 Level scheme of ^{101}Tc showing the properties of the levels between 1806 and 2418 keV

total intensity (gamma-ray + electron) has been given.

The data of Table 3.7 have also been incorporated into the level scheme displayed in Figs. 3.33 and 3.34, where, in accordance with common usage, transitions located on the basis of coincidence evidence are indicated by heavy dots on the end of the transition arrow. The proposed level energies, the intensity (%) of the beta feeds, the deduced $\log f_0 t$ values and the J^π values assigned to each level are summarized in Table 3.8. Since the derivation of the $\log f_0 t$ values often involved evaluating small differences between rather large total intensities, some of the $\log f_0 t$ values are subject to large uncertainties. The nomograms of Lederer et al. (1967) have been used in the calculation.

Since the ground state spin of ^{101}Mo is $1/2^+$ (Hjorth and Cohen 1964), the allowed beta feeds in the decay of ^{101}Mo will populate $1/2^+$ or $3/2^+$ states in ^{101}Tc . These states can then gamma decay directly, via E1, M1 or E2 transitions, to states having J^π values of $1/2^\pm$, $3/2^\pm$, $5/2^\pm$, $7/2^\pm$. Thus, the existence of a direct gamma-ray transition from one of the $1/2^+$ or $3/2^+$ beta-fed states to a level will restrict the latter's spin and parity to one of the values listed above. This argument is applicable to many of the lower lying levels of ^{101}Tc and will be assumed without specific comment in the discussion which follows.

The ground state, the 9.317 and 15.601 keV levels. The multipolarities of the transitions connecting these states were

TABLE 3.8

The β -decay branching ratios, $\log f_0 t$ and estimated J^π values for the decay

Level Energy	% Beta Feed of ^{101}Mo	$\log f_0 t$	J^π
0.0			$9/2^+$
9.317 ± 0.010			$7/2^+$
15.601 ± 0.009			$5/2^+$
207.53 ± 0.04	13.6 ± 1.6	6.80	$1/2^-$
288.45 ± 0.05	<0.6	>7.80	$3/2^-, (1/2^-)$
394.40 ± 0.07	0.15 ± 0.10	8.60	$5/2^-$
500.45 ± 0.09	<0.065	>8.90	$1/2^-, 3/2^-, 5/2^+$
515.19 ± 0.04	<1.6	>7.50	$5/2^+, 7/2^+$
533.55 ± 0.15	<1.4	>8.50	$5/2^+, 7/2^+$
606.41 ± 0.04	11.7 ± 1.1	6.55	$1/2^+, 3/2^+$
616.20 ± 0.08	0.33 ± 0.17	8.05	$1/2^-, 3/2^-, 5/2^-$
622.15 ± 0.11	0.30 ± 0.09	8.15	$1/2^-, 3/2^-, 5/2^-$
669.45 ± 0.2	0.24 ± 0.07	8.20	$5/2^-$
711.12 ± 0.04	1.50 ± 0.70	7.30	$3/2^+$
742.22 ± 0.15	0.10 ± 0.05	8.50	$5/2^-$
886.65 ± 0.07	<0.8	>7.45	$5/2^+, 7/2^+$
1025.59 ± 0.10	3.30 ± 0.38	6.75	$1/2^+, 3/2^+$
1027.90 ± 0.07	7.9 ± 1.1	6.35	$3/2^+$
1103.5 ± 0.3	0.34 ± 0.04	7.65	$1/2^-, 3/2^-, 5/2^-$
1141.5 ± 0.25	0.28 ± 0.03	7.65	$1/2^-, 3/2^-, 5/2^-$
1188.04 ± 0.14	<0.03	>8.60	$1/2^-, 3/2^-, 5/2^-$
1197.3 ± 0.6	0.08 ± 0.02	8.15	$1/2^-, 3/2^-, 5/2^-$
1231.3 ± 0.6	0.21 ± 0.05	7.70	$1/2^-, 3/2^-, 5/2^-$
1319.46 ± 0.06	6.21 ± 0.41	6.15	$1/2^+, 3/2^+$

Level Energy	% Beta Feed	$\log f_{\beta t}$	J^{π}
1449.0 $\pm$ 0.25	0.73 $\pm$ 0.10	6.90	$1/2^{+}, 3/2^{+}$
1565.09 $\pm$ 0.20	0.25 $\pm$ 0.04	7.30	$1/2^{+}, 3/2^{+}$
1594.39 $\pm$ 0.09	1.46 $\pm$ 0.10	6.45	$1/2^{+}, 3/2^{+}$
1598.94 $\pm$ 0.13	<0.11	>7.60	$5/2^{+}, 7/2^{+}$
1614.66 $\pm$ 0.13	1.49 $\pm$ 0.15	6.40	$1/2^{+}, 3/2^{+}$
1677.8 $\pm$ 0.2	0.25 $\pm$ 0.09	7.10	$1/2^{+}, 3/2^{+}$
1775.3 $\pm$ 0.3	1.39 $\pm$ 0.16	6.30	$1/2^{+}, 3/2^{+}$
1806.28 $\pm$ 0.09	2.90 $\pm$ 0.19	5.80	$1/2^{+}, 3/2^{+}$
1892.9 $\pm$ 0.8	0.09 $\pm$ 0.02	7.20	$1/2^{+}, 3/2^{+}$
1897.54 $\pm$ 0.11	3.07 $\pm$ 0.23	5.65	$1/2^{+}, 3/2^{+}$
1928.88 $\pm$ 0.22	0.64 $\pm$ 0.06	6.30	$1/2^{+}, 3/2^{+}$
1962.12 $\pm$ 0.05	14.2 $\pm$ 0.6	4.90	$1/2^{+}, 3/2^{+}$
2001.01 $\pm$ 0.16	0.85 $\pm$ 0.06	6.10	$1/2^{+}, 3/2^{+}$
2047.58 $\pm$ 0.07	19.3 $\pm$ 0.9	4.60	$1/2^{+}, 3/2^{+}$
2056.42 $\pm$ 0.17	2.87 $\pm$ 0.19	5.35	$3/2^{+}$
2129.18 $\pm$ 0.14	1.94 $\pm$ 0.12	5.40	$1/2^{+}, 3/2^{+}$
2150.3 $\pm$ 0.5	0.13 $\pm$ 0.025	6.50	$1/2^{+}, 3/2^{+}$
2237.8 $\pm$ 0.3	0.36 $\pm$ 0.04	5.85	$1/2^{+}, 3/2^{+}$
2252.6 $\pm$ 0.65	0.095 $\pm$ 0.02	6.40	$1/2^{+}, 3/2^{+}$
2353.9 $\pm$ 0.5	0.29 $\pm$ 0.07	5.60	$1/2^{+}, 3/2^{+}$
2418.3 $\pm$ 0.6	0.13 $\pm$ 0.025	5.75	$1/2^{+}, 3/2^{+}$

determined by the electron measurements of section 3.3.3.

The data restrict the spin sequence for the three states to either $5/2^+, 7/2^+, 9/2^+$ or $9/2^+, 7/2^+, 5/2^+$. Both the 9.317 and the 15.601 keV states are directly populated by gamma rays from states fed by allowed beta transitions and hence can not have a spin of $9/2$. The ground state spin must therefore be $9/2^+$ and the spin sequence for the three low-lying levels must be $9/2^+, 7/2^+, 5/2^+$.

The 207.53 keV level. The assignment of a level at this energy was inferred from the many energy sums involving the 191.93 keV line and other gamma rays which originate from levels well established by the coincidence data. Because this state has a lifetime of $\sim 765 \mu\text{sec}$, no direct coincidence data are possible. The level is deexcited by the 191.93 keV M2 transition which feeds the $5/2^+$ state at 15.601 keV. The spin and parity of the 207.53 keV state must therefore be $1/2^-$. The first forbidden character of the beta decay to this state ($\log f_0 t = 6.8$) supports this assignment.

The 288.45 keV level. This level is established by the coincidence evidence presented in Fig. 3.27 and by the 80.92 keV transition which feeds the 207.53 keV state. Internal conversion measurements (see section 3.3.3) have shown that the 80.92 keV line is M1 with an E2 admixture of $< 9\%$. The intensity of the gamma radiation feeding this state is 1.55 times that of the gamma radiation deexciting it. This intensity difference requires an M1 assignment for the 80.92 keV line in good agree-

ment with the conversion electron data. Therefore the 288.45 keV state has negative parity and a spin of $1/2$ or $3/2$. If the 80.92 keV transition is pure M1, the intensity balance precludes any beta feed to the 288.45 keV level; if a 9% E2 admixture is permitted, a weak beta group of $\log f_0 t \sim 7.8$ might be present. The very weak 273 keV gamma ray, which appears in Table 3.7, might be the 288.45 to 15.601 keV transition; if this were so, the $1/2^-$ choice would be ruled out.

The 394.40 keV level. The 105.96-80.92 keV cascade and the 379 keV transition to the 15.601 keV level establish the 394.40 keV level. A possible transition to the 9.317 keV level is obscured in the singles data by the 385 keV gamma ray from the decay of ^{101}Tc (see Table 3.9); however, there is evidence for such a transition in the coincidence data (Fig. 3.30). The gamma-ray deexcitation pattern excludes J^π values of $1/2^\pm$, $3/2^-$ and $7/2^+$. The presence of a weak beta feed ($\log f_0 t \sim 8.6$) rules out the $3/2^+$ and $5/2^+$ choices. The only assignment consistent with all the data is $5/2^-$.

The 500.45 keV level. The 212.00-80.92 keV cascade and the coincidence data of Fig. 3.29 are used to define the level. The gamma-ray intensity balance sets a lower limit of 8.9 for the $\log f_0 t$ value associated with the beta feed to this state and thus rules out J^π values of $1/2^+$ or $3/2^+$. The presence of a transition to the 288.45 keV level excludes a J^π value of $7/2^+$. Table 3.7 includes a weak 492 keV transition which could be placed between the 500.45 and 9.317 keV levels

on the basis of energy fit. If this assignment is correct, the 500.45 keV state must have $J^\pi = 5/2^\pm$. Because this is such a weak transition, the $1/2^-$ and $3/2^-$ choices cannot be rejected.

The 515.19 and 533.55 keV levels. The 515.19 keV level is identified by deexciting transitions to the ground state and the 9.317 and 15.601 keV states and by the coincidence data of Fig. 3.23. The presence of a ground state transition rules out J^π values of $1/2^\pm$, $3/2^\pm$ and $5/2^-$. Therefore the spin and parity of the 515.19 keV state must be $5/2^+$ or $7/2^+$. The 533.55 keV state is established by transitions to the ground and first excited states (Fig. 3.32). The spin and parity of the 533.55 keV level is therefore $5/2^+$ or $7/2^+$.

The 606.41 keV level. This level is firmly supported by transitions to the 15.601, 207.53 and 288.45 keV states and by the coincidence data of Fig. 3.21. The strong beta feed to this level ($\log f_0 t = 6.55$) indicates a spin and parity of $1/2^+$ or $3/2^+$.

The 616.20 and 622.15 keV levels. The 616.20 keV level is identified by means of a number of coincidence cascades (Fig. 3.28). The presence of a strong 408.59 keV transition from this level to the 207.53 keV isomeric state should also be noted. The 622.15 keV state is established by the strong 333.70-80.92 keV cascade and several others involving transitions from higher states (Figs. 3.29 and 3.30).

The beta feeds to these two states are very weak and the $\log f_0 t$ values of 8.0 and 8.1 respectively, are characteristic of first forbidden transitions. The data permit one to restrict the J^π values for the two levels to $1/2^-$, $3/2^-$ or $5/2^-$.

The 669.45 keV state. This level is well established by a number of coincidence cascades (Fig. 3.27). The weak beta feeding ($\log f_0 t = 8.2$) is suggestive of a first forbidden beta transition. The presence of a gamma-ray transition to the $7/2^+$ first excited state eliminates the $1/2^-$ and $3/2^-$ choices and suggests a J^π value of $5/2^-$.

The 711.12 keV level. The 711.12 keV state is defined by many cascades (Figs. 3.21, 3.22, 3.23, 3.24). Although the errors are large, there appears to be some beta feeding; the $\log f_0 t$ value (7.30) is characteristic of either an allowed or a first forbidden beta transition. The gamma-ray transition to the $7/2^+$ state at 9.317 keV rules out J^π values of $1/2^\pm$ and $3/2^-$. The spin and parity of the 711.12 keV state is probably $3/2^+$.

The 742.22 keV level. The 1314-733 keV cascade between the well-known 2056.42 and 9.317 keV states is used to establish the 742.22 keV state. The order of the cascade is determined by the difference in intensity of its members; this difference yields a $\log f_0 t$ of 8.5 characteristic of a first forbidden beta transition. The only J^π value consistent with the data is $5/2^-$.

The 886.65 keV level. A number of coincidence cascades (Fig. 3.25) are used to define this level. The gamma-ray deexcitation pattern excludes J^π values of $1/2^\pm$ and $3/2^-$,

but allows values of $3/2^+$, $5/2^{\pm}$ or $7/2^+$. There appears to be no beta feeding to this state ($\log f_0 t = 7.75$); for this reason, the $3/2^+$ choice is ruled out. The data restrict the J^{π} of the 886.65 keV level to values of $5/2^{\pm}$ or $7/2^+$.

The 1025.59 keV level. The strong 1010.0 keV transition to the 15.601 keV level and the 510.36 keV transition to the 515.19 keV state are used to establish the 1025.59 keV level on the basis of energy fit. The beta feed ($\log f_0 t = 6.75$) is indicative of either an allowed or a first forbidden beta transition; the existence of a strong transition to the $5/2^+$ state at 15.601 keV does not allow a J^{π} value of $1/2^-$. The spin and parity of the 1025.59 keV state may be $1/2^+$ or $3/2^{\pm}$.

The 1027.90 keV level. This level is identified by the coincidence cascades displayed in Fig. 3.26. The strong beta feed ($\log f_0 t = 6.35$) has an allowed character. The stronger 1019 keV component has been placed between the 1027.90 and 9.317 keV levels on the basis of energy fit. The only J^{π} value consistent with the data is $3/2^+$.

The 1067 and 1278 keV levels. These levels are very tentatively proposed on the basis of the weak coincidence data of Fig. 3.29 and therefore have not been shown in Figs. 3.33 and 3.34. The 1278 keV level is also supported by energy fit involving two very weak transitions.

The 1103.5, 1141.5, 1188.04, 1197.3 and 1231.3 keV levels. Transitions to the 207.53, 288.45 and 500.45 keV levels define the 1103.5 keV level. The 853.1-80.92 and

641-212.00 keV cascades are used to establish the 1141.5 keV level. The 1188.04 keV state is confirmed by the 980.7 keV transition to the 207.53 keV level and by the 571.3-408.59 keV cascade. The 486.2-195.94 keV cascade (Fig. 3.22) is the only evidence for the 1197.3 keV level. The 1231.3 keV state is identified by the 624.6 and 943.2 keV transitions to the 606.41 and 288.45 keV levels, respectively. Each of the levels listed above is weakly populated with a $\log f_0 t$ value characteristic of a first forbidden beta transition. The spin and parity of the levels may be $1/2^-$, $3/2^-$ or $5/2^-$.

The 1319.46 keV level. This level is strongly supported by a number of coincidence cascades (Figs. 3.21, 3.22, 3.23, 3.32). The strong beta feed with an allowed $\log f_0 t$ value (6.15) restricts the J^π values to $1/2^+$ or $3/2^+$.

The 1449.0, 1565.09 and 1677.8 keV levels. The stronger 1433 keV transition makes coincidences only with the weak 229 keV gamma ray. The 1433+229 keV energy sum corresponds to the difference between the well-established (Fig. 3.30) 1677.8 and 15.601 keV states. These data define a level at 1449.0 keV. The 1565.09 keV level is established by the 1049.9-505.88 keV cascade. The beta feeds to the three levels listed above are either allowed or first forbidden in character. Since the gamma-ray decay patterns rule out the $1/2^-$ choice, J^π values are restricted to $1/2^+$ or $3/2^+$.

The 1598.94 keV level. The transitions to the $5/2^+$ and $7/2^+$ states at 15.601 and 9.317 keV together with the 888 and 1066 keV transitions to the 711.12 and 533.55 keV levels

respectively, are used to identify the 1598.94 keV level. The presence of a transition to the $7/2^+$ first excited states excludes J^π values of $1/2^\pm$ and $3/2^-$. There appears to be no beta feeding to this level; the $\log f_0 t$ value is >7.6 and would appear to rule out a $3/2^+$ assignment. The spin and parity of the 1598.94 keV state can therefore be $5/2^\pm$ or $7/2^+$.

The 1594.39, 1614.66, 1775.3 and 1806.28 keV levels.

These four states are all well established by coincidence data and the beta transitions to them have $\log f_0 t$ values characteristic of allowed transitions. Their J^π values must be $1/2^+$ or $3/2^+$.

The 1892.9 keV level. This level has been tentatively included on the basis of the weak 1286.5-590.82 keV coincidence cascade. J^π values of $1/2^\pm$ or $3/2^\pm$ are permitted.

The 1897.54, 1928.88, 1962.12, 2001.01, 2047.58, 2056.42 and 2129.18 keV levels. All of these levels are well established by the coincidence data. The beta feeds to these states are characteristic of allowed transitions; the spin and parity of each of the levels may be $1/2^+$ or $3/2^+$. In the case of the 2056.42 keV level, there is a transition to the $7/2^+$ state at 9.317 keV which further restricts its J^π value to $3/2^+$.

The 2150.3, 2237.8, 2252.6, 2353.9 and 2418.3 keV levels.

The 1544.2-590.82 keV cascade together with the 1942.5 keV transition to the 207.53 keV state are used to establish the 2150.3 keV level. Transitions to the 1677.8, 1027.90 and 15.601 keV levels provide evidence for the 2237.8 keV level.

The 2252.6 keV state is suggested by the 1646.2-590.82 keV coincidence relationship. The 2353.9 keV level is established by transitions to the 1027.90 and 207.53 keV levels. Transitions to the 886.65, 288.45 and 15.601 keV levels give firm support to the 2418.3 keV state. The beta feeds to all of these levels are characteristic of allowed transitions thus restricting the spin and parity of each of these states to $1/2^+$ or $3/2^+$.

The only strong line which has not been classified is the 2088.4 keV transition of intensity 6.75. Since this gamma ray produced no observable coincidence, it must therefore be directly beta fed and populate one of the 9.3, 15.6 or 207.5 keV states. This transition is evidence that another low spin positive parity state must exist at 2097.7, 2104.0 or 2295.9 keV.

The thirteen unclassified transitions in Table 3.7 have a total intensity of 1.5% (of which nearly half is accounted for by the 2088.4 keV gamma ray discussed above). None of these lines were observed in coincidence. Since 98.5% of the total gamma ray intensity feeding the three low lying positive parity states has been accounted for in the decay scheme, the absolute beta and gamma-ray intensity scales used in this work cannot be in error by more than $\sim 1.5\%$.

3.4 The Decay of 14.0 min. ^{101}Tc to Levels in ^{101}Ru

3.4.1 Introduction

The decay of ^{101}Tc was first identified as a daughter

activity of ^{101}Mo by Sagane et al. (1940). The existence of an excited state of ^{101}Tc at 307 keV was established with the beta-gamma coincidence experiments of Rutledge et al. (1952) and Wiles (1954). O'Kelley et al. (1957) used scintillation spectroscopy techniques to define levels in ^{101}Tc at 130, 307, 316, 410, 545, 635, 720, 846 and 939 keV. The half-life of ^{101}Tc was measured to be 14.0 ± 0.1 minutes and an energy of 1.32 ± 0.03 MeV was determined for the strong beta group feeding the 307 keV level. The Coulomb excitation experiments of McGowan and Stelson (1957), which were reported in the Nuclear Data Sheets (1961) complemented the decay studies by populating levels at 127, 307, 324, 410, 626, 720 and 850 keV. With the advent of Ge(Li) detectors, both the beta decay of ^{101}Tc (Aras et al. 1965) and the Coulomb excitation of ^{101}Ru (Kistner and Schwarzschild 1967) were reinvestigated. These new studies determined more accurate energies and intensities, confirmed the existing level scheme, and added new levels at 616.3 and 928.9 keV. The existence of an isomeric state in ^{101}Ru at ~ 525 keV has been proposed by Krehbiel and Meyer-Berkhout (1961) and McCarthy et al. (1965); half-lives of 24.5 and 17.5 μsec respectively, were determined. This level, which does not appear to be populated in the beta decay of ^{101}Tc , decays only to the 307 keV state.

The levels in ^{101}Tc can also be populated from the electron capture decay of ^{101g}Rh and ^{101m}Rh . Evans et al. (1965) have shown that the J^π values of the ground and isomeric

states are $1/2^-$ and $9/2^+$, respectively; therefore, the decay of ^{101m}Rh should populate the same (at least those energetically possible) states found in the decay of ^{101}Tc , which also has a spin and parity of $9/2^+$ for its ground state. The experiments of Evans and Naumann (1965) and Aras et al. (1966) have established that levels at 127.2(g), 306.8(m), 311.2, 325.4(g), 422.1(g), 523 and 545.2(m) keV in ^{101}Tc are populated from the decays of ^{101g}Rh and ^{101m}Rh .

Since the level scheme of ^{101}Ru seems to be fairly well defined, it is not obvious why further experimental work should be included in this thesis. The present study was carried out in order to make sure that no gamma rays from the decay of ^{101}Tc were masking transitions from the decay of ^{101}Mo . Otherwise incorrect energies and intensities might be measured for ^{101}Mo . In view of the improvements in Ge(Li) fabrication and signal processing, it was expected that more accurate energies and intensities could be determined. Moreover, the Ge(Li)-Ge(Li) coincidence experiments for ^{101}Mo gave ^{101}Tc coincidence data as a bonus, so that confirmation of the existing level scheme could be obtained without doing any additional experiments. The interpretation of the previous coincidence data, which was taken with NaI(Tl) detectors, was hampered by the poor resolution and therefore construction of previous level schemes relied primarily on energy fit considerations.

During the course of the present work, papers by Siwamogsatham and Easterday (1971) and Aras et al. (1971)

on the decay of ^{101}Tc were published; these investigations confirmed the earlier works and established new levels at 674, 912, 938 and 1001 keV. However, neither group performed coincidence experiments with Ge(Li) detectors.

3.4.2 Singles and Gamma-Gamma Coincidence Results for ^{101}Tc

A gamma-ray spectrum of ^{101}Tc for the energy range between 100 and 1100 keV is shown in Fig. 3.35. The spectrum was recorded with the 12 cm³ detector; the 12.5 mm lucite beta stopper was the only absorber used. In this spectrum, the lines of ^{101}Tc have been enhanced by removal of much of the ^{101}Mo activity with the chemical procedure outlined in Section 2.4. However, a number of the stronger lines from the decay of the ^{101}Mo parent are still present as contaminants. Other impurities arose from the room background (^{232}Th , ^{137}Cs) and are labelled (c) in the figure. A comparison between this spectrum and that of Fig. 3.16 allows identification of the transitions belonging to the beta decay of ^{101}Tc .

A selection of coincidence spectra from the decay of ^{101}Tc is presented in Figures 3.36 and 3.37. Figure 3.36 displays the spectra in coincidence with the 127 keV transition and the 180-184 keV doublet. A comparison of the two spectra indicates clearly that the 233, 531, 617 and 627 keV transitions do not feed the 127 keV level directly, but cascade through the 184 keV gamma ray. The 180 and 184 keV lines feed the 127 keV level from the 307 and 311 keV levels, respectively.

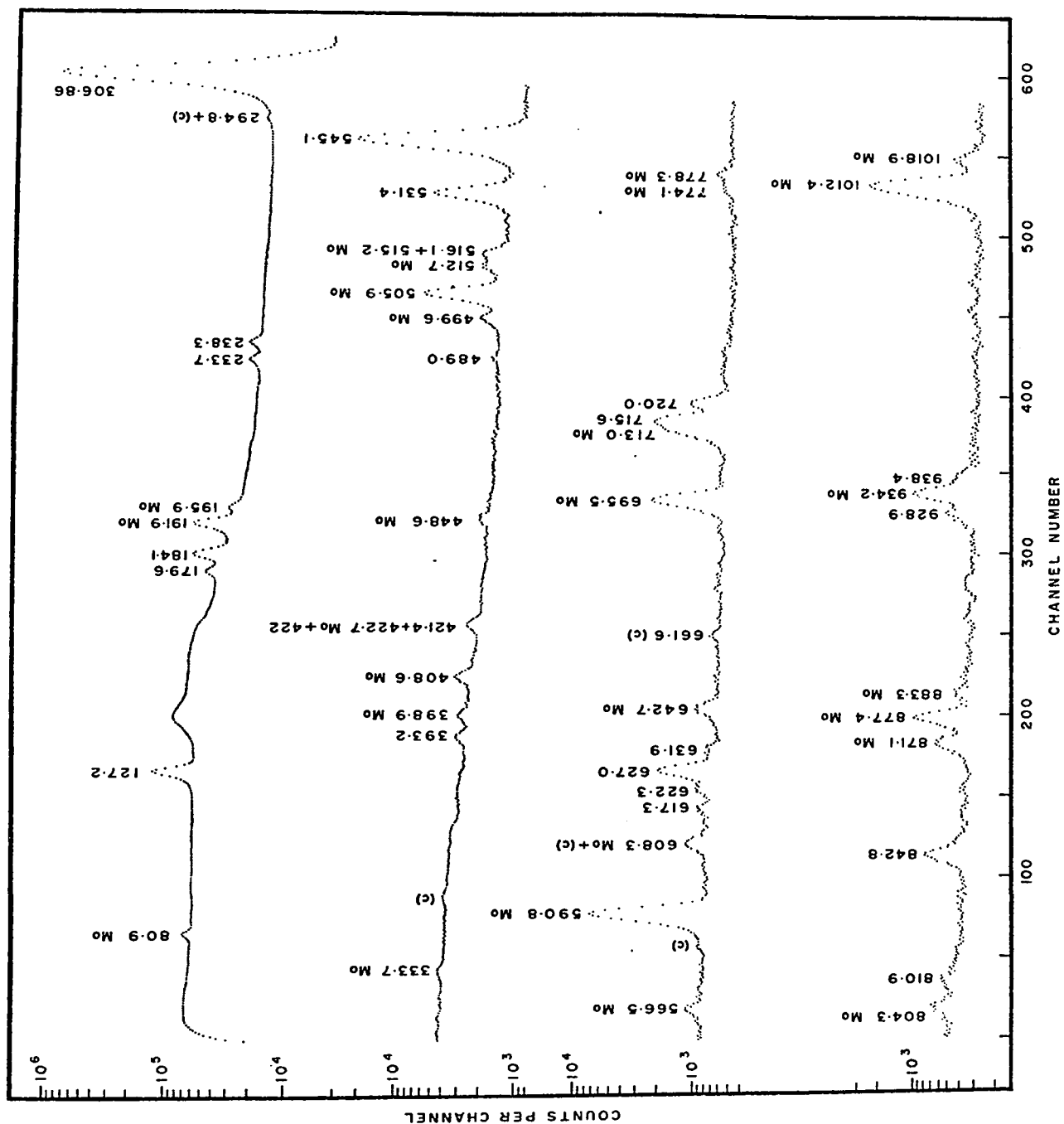


Fig. 3.35 The gamma-ray spectrum of ^{101}Tc taken with the 12 cm^3 detector

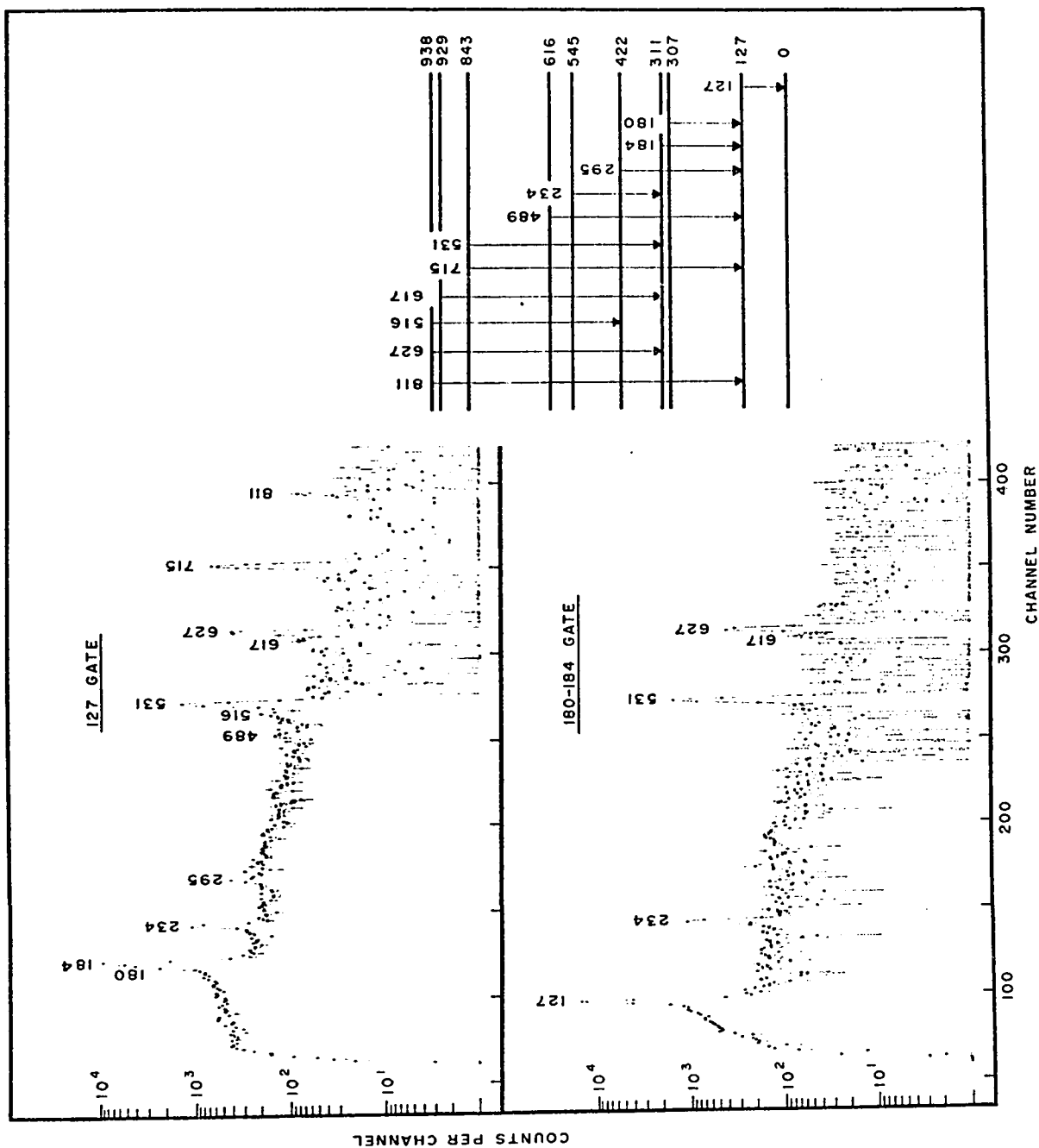


Fig. 3.36 Spectra in coincidence with the 127 and 180+184 keV gamma rays from the decay of ^{101}Tc

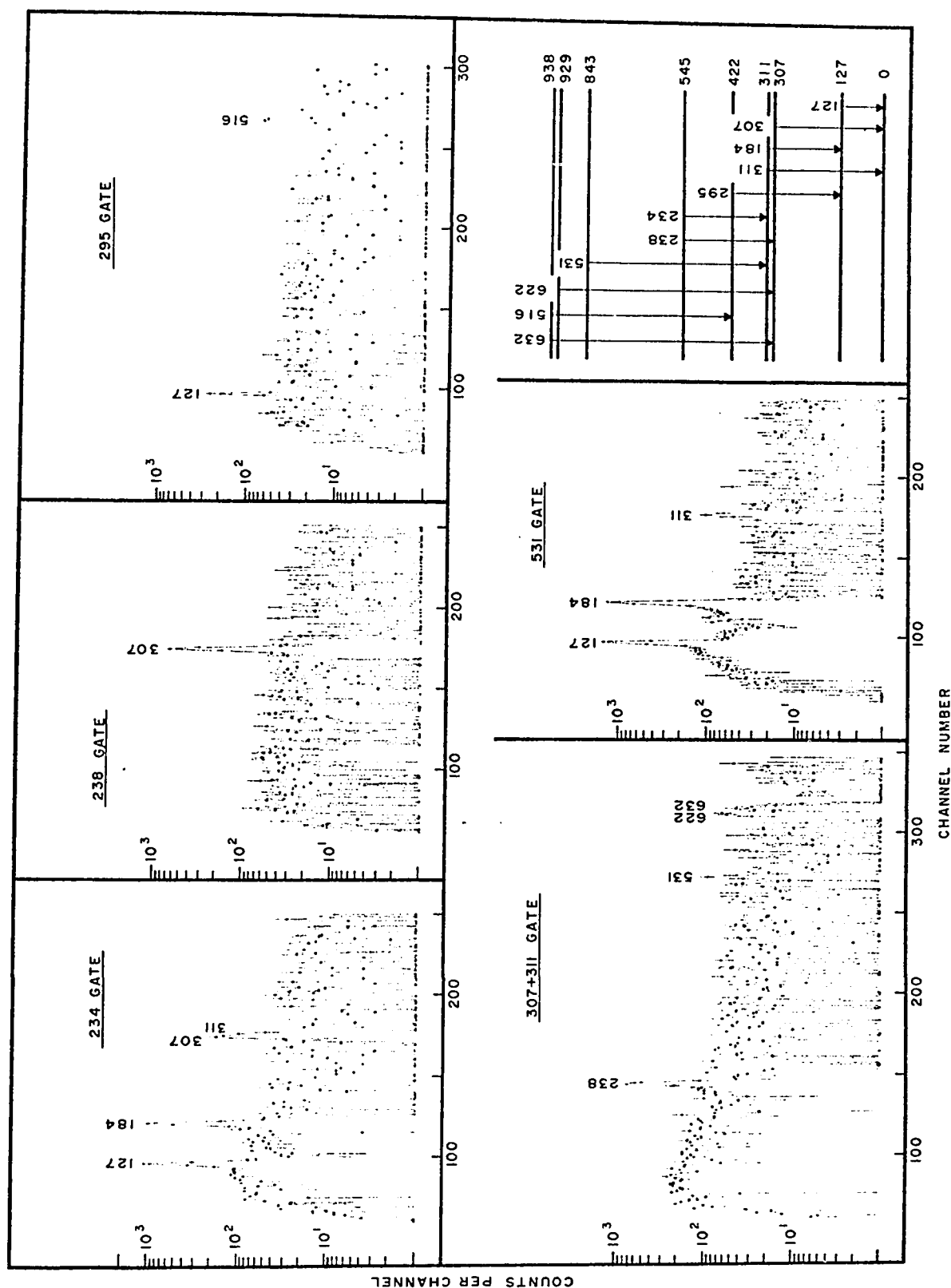


Fig. 3.37 Spectra in coincidence with gates containing the 234, 238, 295, 307+311, and 531 keV transitions from the decay of ^{101}Tc . The 234 keV gate contains a portion of the 238 keV peak

A further selection of coincidence gates is shown in Fig. 3.37. The coincidence spectra associated with gates containing the 233 and 238 keV lines confirm the presence of transitions from the 545 keV level to the 311 and 307 keV levels, respectively. The 233 keV gate included a portion of the 238 keV photopeak. The 295 keV gated spectrum shows that the 516, 295 and 127 keV lines form a triple cascade. The coincidence spectrum associated with the 516 keV gamma ray (not shown) shows a 422 keV transition in addition to the 127 and 295 keV lines and confirms the 422 keV level previously established by other workers. Similarly, the levels at 843, 929 and 938 keV are verified by the coincidence data displayed in the two figures. The presence of a 311 keV line in both the 233 and 531 keV gated spectra proves the existence of a ground-state transition from the 311 keV level. This transition has not previously been observed in decay studies and its measured intensity of 0.31% is considerably larger than the values estimated by other workers on the basis of intensity balancing. Corrections for coincidence summing effects in the spectrum associated with the 531 keV gate of Fig. 3.37 were $\sim 10\%$.

The energies and intensities of the gamma rays assigned to the decay of ^{101}Tc are tabulated in Table 3.9 together with the recent results obtained by Siwamogsatham and Easterday (1971) and Aras et al. (1971). All three sets of data agree quite well; however, some of the intensities quoted by Aras et al. (1971) are inconsistent with those of the other two workers and even with their own published spectrum. Some weaker transitions

Transitions in ^{101}Ru From the Decay of ^{101}Tc

S & E (1)

Aras et al (1971)

Present Work

E_γ (keV)	I_γ (%)	E_γ (keV)	I_γ (%)	E_γ (keV)	I_γ (%)	Classification
127.2 ± 0.1	2.31 ± 0.27	127.1 ± 0.2	1.5 ± 0.2	127.21 ± 0.05	2.73 ± 0.19	$127.21 \rightarrow 0$
179.3 ± 0.2	0.55 ± 0.05	179.4 ± 0.3	0.3 ± 0.03	174.9 ± 0.3	0.035 ± 0.008	$720.0 \rightarrow 545.09$
183.8 ± 0.2	1.51 ± 0.13	184.0 ± 0.3	1.4 ± 0.02	179.59 ± 0.05	0.57 ± 0.045	$306.85 \rightarrow 127.21$
233.6 ± 0.2	0.26 ± 0.03	233.7 ± 0.3	0.1 ± 0.02	184.11 ± 0.05	1.58 ± 0.11	$311.32 \rightarrow 127.21$
238.1 ± 0.2	0.29 ± 0.03	238.3 ± 0.3	0.3 ± 0.03	233.71 ± 0.07	0.26 ± 0.03	$545.09 \rightarrow 311.32$
295	0.06	295.2 ± 0.5	0.04 ± 0.01	238.26 ± 0.07	0.27 ± 0.03	$545.09 \rightarrow 306.85$
306.8 ± 0.1	88.0	306.7 ± 0.2	90.0 ± 9.0	294.8 ± 0.3	0.06 ± 0.02	$422.1 \rightarrow 127.21$
311.2	0.009			306.86 ± 0.04	88.7 ± 4.4	$306.85 \rightarrow 0$
321.9 ± 0.4	0.027 ± 0.009	321.7 ± 0.5	0.04 ± 0.006	311.3 ± 0.4	$0.31 \pm 0.06^*$	$311.32 \rightarrow 0$
384.0 ± 0.4	0.027 ± 0.009	383.1 ± 0.5	0.02 ± 0.003			$938.32 \rightarrow 616.2$
393.4 ± 0.3	0.10 ± 0.02	393.0 ± 0.3	0.06 ± 0.007			$928.96 \rightarrow 545.09$
413.5 ± 0.5	≤ 0.009	413.0 ± 0.5	0.02 ± 0.004	393.2 ± 0.2	0.09 ± 0.02	$938.32 \rightarrow 545.09$
421.9 ± 0.3	0.035 ± 0.009	422.5 ± 0.5	0.03 ± 0.005			$720.0 \rightarrow 306.85$
489.1 ± 0.3	0.027 ± 0.009	488.1 ± 0.5	0.03 ± 0.005	422.4 ± 0.4	0.05 ± 0.02	$422.1 \rightarrow 0$
516.1 ± 0.2	0.10 ± 0.02	515.9 ± 0.2	0.08 ± 0.009	489.0 ± 0.45	0.035 ± 0.01	$616.2 \rightarrow 127.21$
531.3 ± 0.1	1.01 ± 0.06	531.1 ± 0.2	0.8 ± 0.08	516.1 ± 0.25	0.11 ± 0.03	$938.32 \rightarrow 422.1$
544.8 ± 0.1	5.95 ± 0.53	544.6 ± 0.2	4.8 ± 0.5	531.41 ± 0.10	1.05 ± 0.09	$842.76 \rightarrow 311.32$
		602.8 ± 0.5	0.01 ± 0.005	545.11 ± 0.07	6.05 ± 0.43	$545.09 \rightarrow 0$

S & E⁽¹⁾

E_Y (keV)	I_Y (%)
617.3 ± 0.4	0.053 ± 0.009
622.0 ± 0.4	0.035 ± 0.009
626.9 ± 0.2	0.44 ± 0.07
631.5 ± 0.5	0.027 ± 0.009
674.2 ± 0.4	≤ 0.009
694.3 ± 0.2	0.044 ± 0.009
715.3 ± 0.1	0.62 ± 0.09
720.1 ± 0.2	0.21 ± 0.05
810.9 ± 0.2	0.06 ± 0.02
842.5 ± 0.1	0.19 ± 0.035
911.9 ± 0.2	0.018 ± 0.009
928.8 ± 0.1	0.080 ± 0.009
938.1 ± 0.1	0.071 ± 0.009

Aras et al (1971)

E_Y (keV)	I_Y (%)
616.9 ± 0.3	0.03 ± 0.005
621.7 ± 0.2	0.07 ± 0.008
626.7 ± 0.2	0.4 ± 0.04
631.7 ± 0.5	0.04 ± 0.003
674.0 ± 0.5	0.02 ± 0.003
694.3 ± 0.2	0.06 ± 0.007
715.6 ± 0.2	0.06 ± 0.006
720.0 ± 0.2	0.2 ± 0.02
811.1 ± 0.2	0.06 ± 0.007
826.2 ± 0.5	0.002 ± 0.001
842.7 ± 0.2	0.2 ± 0.002
911.9 ± 0.3	0.01 ± 0.002
928.5 ± 0.2	0.09 ± 0.03
938.1 ± 0.2	0.08 ± 0.008

Present Work

E_Y (keV)	I_Y (%)	Classification
617.3 ± 0.4	0.075 ± 0.02	$928.96 \rightarrow 311.32$
622.25 ± 0.2	0.09 ± 0.02	$928.96 \rightarrow 306.85$
626.98 ± 0.13	0.48 ± 0.045	$938.32 \rightarrow 311.32$
631.9 ± 0.4	0.05 ± 0.015	$938.32 \rightarrow 306.85$
715.59 ± 0.13	0.71 ± 0.07	$842.76 \rightarrow 127.21$
720.0 ± 0.2	0.23 ± 0.03	$720.0 \rightarrow 0$
810.9 ± 0.4	0.060 ± 0.015	$938.32 \rightarrow 127.21$
842.75 ± 0.13	0.22 ± 0.03	$842.76 \rightarrow 0$
928.9 ± 0.25	0.12 ± 0.02	$928.96 \rightarrow 0$
938.4 ± 0.3	0.09 ± 0.02	$938.32 \rightarrow 0$

1. Sivamogsatham and Easterday (1971)

*Seen Only in Coincidence

seen by the other workers were missed in the present work because of the higher background from the decay of ^{101}Mo .

3.4.3 The Level Scheme of ^{101}Ru

The singles and coincidence data of the preceding section lead to the level scheme displayed in Fig. 3.38. Table 3.10 summarizes the proposed level energies, the percent beta feeds, the deduced $\log f_0 t$ values, and the J^π values assigned for each proposed level. The $\log f_0 t$ values were calculated using the nomograms given in Lederer et al. (1967).

The ground state spin and parity of ^{101}Tc was shown to be $9/2^+$ in section 3.3.3. Allowed beta transitions will therefore populate states having J^π values of $7/2^+$, $9/2^+$ or $11/2^+$. Since the ground state spin of ^{101}Ru is $5/2^+$ (see below), any level in ^{101}Ru that is directly beta fed with an allowed transition cannot decay directly to the ground state unless the value of J^π is either $7/2^+$ or $9/2^+$. This argument is used several times in the following discussion without specific comment. The discussion of the J^π assignments often uses arguments based on the Coulomb excitation experiments of Kistner and Schwarzschild (1967) and on the electron capture decay of ^{101g}Rh and ^{101m}Rh (Evans and Naumann 1965; Aras et al. 1966); these references are not mentioned each time.

The ground state. A value of $5/2$ has been determined for the ground state spin of ^{101}Ru using the techniques of optical spectroscopy and electron spin resonance (Lindgren 1964). The parity is positive (Nuclear Data Sheets 1961); this

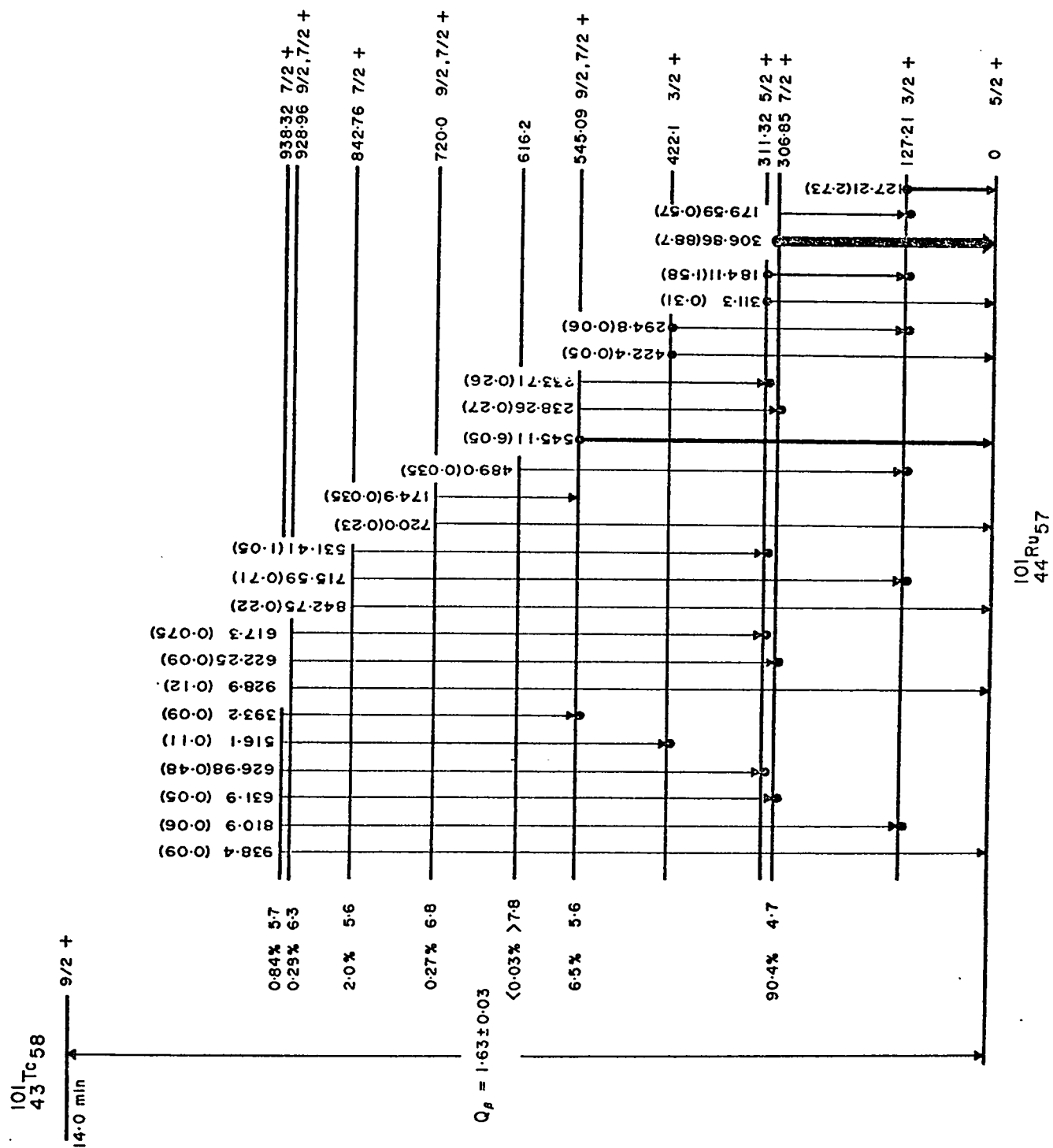


Fig. 3.38 The level scheme of ^{101}Ru from the decay of ^{101}Tc

TABLE 3.10

222

The β -decay branching ratios, $\log f_0 t$ and J^π values for the decay of ^{101}Tc

Level Energy	% Beta Feed	$\log f_0 t$	J^π
0.0			5/2
127.21 $\pm$ 0.05			3/2
305.85 $\pm$ 0.03	90.4 $\pm$ 4.4	4.7	7/2
311.32 $\pm$ 0.07			5/2
422.1 $\pm$ 0.25			3/2
545.09 $\pm$ 0.04	6.45 $\pm$ 0.43	5.6	9/2+, 7/2+
616.2 $\pm$ 0.45	< 0.03	> 7.8	
720.0 $\pm$ 0.14	0.27 $\pm$ 0.05	6.8	7/2+, 9/2+
842.76 $\pm$ 0.08	2.00 $\pm$ 0.11	5.6	7/2+
928.96 $\pm$ 0.15	0.28 $\pm$ 0.04	6.3	7/2+, 9/2
938.32 $\pm$ 0.10	0.88 $\pm$ 0.06	5.7	7/2+

result is also predicted by the shell model since the nearest single-particle states ($2d_{5/2}$, $2d_{3/2}$, $3s_{1/2}$, $1g_{7/2}$) above the closed $N=50$ shell all have positive parity.

The 127.21 keV level. This level is Coulomb excited and is directly populated in the electron capture decay of $^{101}\text{gRh}(1/2^-)$. These facts suggest that it has positive parity and a spin of $1/2$ or $3/2$. Since a number of high spin ($7/2$ or $9/2$) excited states decay to this level, the $1/2^+$ choice is ruled out. Therefore, the spin and parity of the 127.21 keV level is $3/2^+$. This assignment is confirmed by the M1 character of the 127.21 keV transition (Kistner and Schwarzschild 1967). Moreover, the intensity balance gives an α_T consistent with an M1 transition.

The 306.85 keV level. Allowed decays from $^{101}\text{Tc}(9/2^+)$ and $^{101\text{m}}\text{Rh}(9/2^+)$ populate the 306.85 keV level; it is also observed in Coulomb excitation. Therefore, J^π values of $7/2^+$ or $9/2^+$ are suggested. Since there exists a deexciting transition to the $3/2^+$ state at 127.21 keV, the only J^π value consistent with the data is $7/2^+$.

The 311.32 keV level. Transitions to the 127.21 keV state and the ground state are used to establish this level. Although it is observed in Coulomb excitation, the 311.32 keV level is not directly populated in the decays of ^{101}Tc , ^{101}Rh or ^{101}gRh . These facts suggest a spin and parity of $5/2^+$ for the level.

The 325.4 keV level. This level is only observed in the decay of ^{101}gRh and has been included here for completeness. Its spin and parity is most likely $1/2^+$.

The 422.4 keV level. Transitions to the ground and first excited states (see 531 keV gate of Fig. 3.37) define the 422.4 keV level. It is fed directly in the decay of ^{101g}Rh , but not in the decay of ^{101}Tc or ^{101m}Rh ; the level is also Coulomb excited. These facts suggest a spin and parity of $1/2^+$ or $3/2^+$. Since the 422.4 keV level is populated from high spin ($7/2$ or $9/2$) states, the only J^π value consistent with the data is $3/2^+$.

The 545.09 keV level. This level is established by transitions to the ground state and the levels at 306.85 and 311.32 keV. The 545.09 keV state is fed directly by allowed transitions in the decay of both ^{101}Tc and ^{101m}Rh . The spin and parity can be $7/2^+$ or $9/2^+$.

The 616.2 keV level. The 489.0-127.21 keV coincidence cascade defines the 616.2 keV level. The $\log f_0 t$ value (>7.8) suggests a first forbidden beta feed and a negative parity assignment; however, this level is also observed in the Coulomb excitation experiments and therefore the parity must be positive. Aras et al. (1971) and Siwamogsatham and Easterday (1971) have proposed that approximately 50% of the 617.3 keV intensity be placed as a transition deexciting the 616.2 keV level. The coincidence data of section 3.4.2 indicates that nearly all of the 617.3 keV intensity deexcites the 928.96 keV state. These workers have also placed a 322 keV gamma ray between the 938 and 616 keV levels; the intensity of this line is sufficient to account for much of the 489.0 keV intensity and would make the $\log f_0 t$ value even higher. It would appear

that J^π values of $3/2^+$ or $5/2^+$ are suggested with the $5/2^+$ assignment being the most likely choice; the level is not populated in the decay of ^{101}gRh .

The 720.0 keV level. A level of this energy was populated in the Coulomb excitation experiments. Therefore the 720.0 and 174.9 keV transitions of the present work have been assigned to the 720.0 keV level. The spin and parity can be $7/2^+$ or $9/2^+$.

The 842.76 keV level. Transitions to the ground state and the levels at 127.21 and 311.32 keV are used to define the 842.76 keV level. The $\log f_0 t$ value (5.6) is characteristic of an allowed beta transition. Since the presence of a transition to the $3/2^+$ first excited state rules out the $9/2^+$ choice, the spin and parity of the 842.76 keV state must be $7/2^+$.

The 928.96 keV level. This level is established by transitions to the 311.32 and 306.85 keV states and the ground state and is observed in Coulomb excitation. The beta feed has an allowed character ($\log f_0 t = 6.3$) and therefore J^π values of $7/2^+$ and $9/2^+$ are suggested for the 928.96 keV level.

The 938.32 keV level. Six deexciting transitions are used to define the 938.32 keV level. The $\log f_0 t$ value of 5.7 indicates an allowed beta transition. Due to the presence of a decay to the 127.21 keV level, the spin and parity of the 938.32 keV state must be $7/2^+$.

The sensitivity of the present work was insufficient to draw any conclusions regarding the 674, 912 and 644 or 1001 keV levels tentatively suggested by other workers.

3.4.4 Subsequent Investigations

Sieniawski et al. (1971) have reinvestigated the electron capture decays of ^{101m}Rh and ^{101g}Rh with Ge(Li) detectors. The results of earlier workers were confirmed and a number of new transitions were observed. In addition, the internal conversion spectrum of ^{101m}Rh was measured with a 50-cm radius double-focussing beta-ray spectrometer. The multipolarities of the 127 (M1+ $\leq$ 12% E2), 180 (E2), 184 (M1+E2), 234 (M1+E2), 238 (M1+E2), 307 (M1+ $\leq$ 30%E2) and 545 (M1+E2) keV transitions were determined. Their experiments have confirmed the existing level scheme of ^{101}Ru and have established new levels at 344.1 ($1/2^+-5/2^+$), 462.3 ($1/2^+$ or $3/2^+$), and 643.8 ($7/2^+$) keV. The 616.3 and 623.5 (not observed in the present work) keV states were both found to decay to the ground state and the $3/2^+$ first excited state and a J^π values of $7/2^+$ was assigned to each. The intensity of the 616 keV transition from the 616.3 keV level was $\sim 1/3$ the intensity of the 489 keV transition. Since the intensity of the 489 keV transition was 0.035 in the present work, this branching ratio would imply an intensity of ~ 0.012 for the 616 keV transition. This is $\sim 1/6$ of the 617.3 keV intensity and not 50% as quoted by other workers. The results of their internal conversion data established a value of $7/2^+$ for the spin and parity of the 545.09; keV state; previously, J^π values of $7/2^+$ or $9/2^+$ were allowed.

The level scheme of ^{101}Ru has also been investigated with the $^{100}\text{Mo}(\alpha, 3n\gamma)^{101}\text{Ru}$ reaction (Lederer et al. 1971). The

proposed $11/2^-$ isomeric state at 528 keV was confirmed. A spin of $5/2$ was suggested for the 311 keV state by their angular distribution data in agreement with the $5/2^+$ assignment made in the present work. In addition, J^π values of $7/2$ and $9/2$ are favoured for the 545 and 720 keV states, respectively. Evidence for a level at 1001 keV was also presented; this confirms the level assignment of Siwamogsatham and Easterday (1971).

CHAPTER IV

DISCUSSION

4.1 Introduction

The nuclei ($^{97}_{43}\text{Tc}_{54}$, $^{99}_{43}\text{Tc}_{56}$, $^{101}_{43}\text{Tc}_{58}$ and $^{101}_{44}\text{Ru}_{57}$) which are subject to investigation in this thesis have proton and neutron numbers characteristic of a mass region relatively far away from a major shell closure (magic number). The protons are filling the $1g_{9/2}$ subshell. The large energy gap due to the closed proton shell at $Z = 50$ is expected to isolate the $1g_{9/2}$ subshell from the unfilled positive parity subshells such as the $2d_{5/2}$, $2d_{3/2}$, $3s_{1/2}$ and $1g_{7/2}$. The neighbouring more stable subshells ($2p_{1/2}$, $2p_{3/2}$, $1f_{5/2}$) all have negative parity and it is generally assumed that the filling of the $2p_{3/2}$ subshell results in a minor shell closure at $Z = 38$. Thus, $^{88}_{38}\text{Sr}_{50}$ is often used as the core nucleus in shell model calculations for nuclei with $Z \leq 38 < 50$.

The filling of the $1g_{9/2}$ neutron subshell results in the major shell closure at $N = 50$. Because of the large energy gap, the negative parity $2p_{1/2}$, $2p_{3/2}$ and $1f_{5/2}$ neutron subshells below the gap are not expected to influence the behaviour of low energy states in this mass region. The positive parity $2d_{5/2}$, $2d_{3/2}$, $3s_{1/2}$ and $1g_{7/2}$ neutron subshells are all expected to contribute to the wavefunctions of nuclei with $50 \leq N < 82$. The core nucleus for such nuclei is usually chosen to be $^{90}_{40}\text{Zr}_{50}$.

Unfortunately, there has been little theoretical interpretation of the odd-mass technetium and ruthenium isotopes. The relatively large number of particles outside the closed proton and neutron shells prohibits realistic shell-model calculations. Some work in this mass region has been done by Auerbach and Talmi (1965), Cohen et al. (1964) and Vervier (1966). Their results give reasonable agreement with experiment near ${}^{90}_{40}\text{Zr}_{50}$; however, due to the small configuration space used in the calculations, the agreement is much poorer farther away from the closed shells (i.e. for technetium). Moreover, there are few published data on single-particle transfer reactions for the nuclei of interest. This means that the distribution of single-particle strength in these nuclei is unknown.

Before attempting to use the shell model for interpreting the energy levels of ${}^{97}_{43}\text{Tc}$, ${}^{99}_{43}\text{Tc}$, ${}^{101}_{43}\text{Tc}$ and ${}^{101}_{44}\text{Ru}$, its validity for describing nuclei in the proximity of the closed shells at $Z = 38$ and $N = 50$ should be investigated. For this reason, a survey of the known properties of ${}^{88}_{38}\text{Sr}_{50}$, ${}^{90}_{40}\text{Zr}_{50}$, some of the even-even and odd-even isotopes of zirconium and molybdenum, and some of the odd-odd and odd-even isotopes of niobium was made. The results of this survey has been included in Appendix B and has been used as an aid in interpreting the level structures of the nuclei which are the subject of this thesis. In the discussion that follows, no serious attempt is made to give the amplitudes

of the components contributing to the configuration which has been suggested.

4.2 Discussion of the levels in $^{101}_{44}\text{Ru}_{57}$

The experimental data discussed in section 3.4 have provided evidence for the existence of 16 levels in ^{101}Ru . Nine of these states have high spin ($7/2^+$ or $9/2^+$) and are populated with allowed transitions in the beta decay of $^{101}_{43}\text{Tc}_{58} (9/2^+)$ and/or the electron-capture decay of $^{101\text{m}}_{45}\text{Rh}_{56} (9/2^+)$. If a $^{88}_{38}\text{Sr}_{50}$ core is assumed, ^{101}Tc has five protons distributed among the $2p_{1/2}$ and $1g_{9/2}$ subshells and ^{101}Rh has seven. Previous results (Auerbach and Talmi 1965) have indicated that the $9/2^+$ ground state of $^{93}_{43}\text{Tc}_{50}$ is a mixture of the $[(\pi p_{1/2})^2_0 (\pi g_{9/2})^3]_{9/2+}$ and the $[(\pi g_{9/2})^5]_{9/2+}$ configurations; the same proton structure is assumed for the $9/2^+$ ground state of ^{101}Tc . Similarly, the $9/2^+$ first excited state of ^{101}Rh will contain the proton configurations $[(\pi p_{1/2})^2_0 (\pi g_{9/2})^5]_{9/2+}$ and $(\pi g_{9/2})^7]_{9/2+}$. In both cases the radioactive decay must leave the protons of $^{101}_{44}\text{Ru}_{57}$ in either the $(\pi p_{1/2})^2_0 (\pi g_{9/2})^4_0$ or the $(\pi g_{9/2})^6_0$ configuration. Therefore, an allowed transition in the electron-capture decay of $^{101\text{m}}\text{Rh}$ will involve the decay of a $g_{9/2}$ proton to a $g_{7/2}$ neutron leaving the daughter nucleus ^{101}Ru in some $g_{7/2}$ neutron configuration. Similarly, an allowed transition in the beta decay of ^{101}Tc will involve the decay of a $g_{7/2}$ neutron to a $g_{9/2}$ proton; again, the daughter nucleus will be left in some $g_{7/2}$ neutron configuration. It is clear

that any theoretical interpretation of ^{101}Ru must be able to account for the $g_{7/2}$ configurations which are preferentially populated in the radioactive decays.

Many of the levels found in ^{101}Ru by radioactive decay have also been observed in the Coulomb excitation experiments of Kistner and Schwarzschild (1967). Thus, it would appear that the core-excitation strength is shared among many levels rather than confined to the five states expected from a simple weak-coupling theory. Since reference is made to this work throughout the following discussion, the authors are not referred to explicitly each time. It should be noted that similar results were also obtained for ^{99}Ru in the same paper. The levels of ^{101}Ru have not been investigated with single-particle transfer reactions such as (d,p) and (d,t) .

The neighbouring nuclei, ^{103}Ru and ^{105}Ru , cannot be studied by Coulomb excitation, but many levels in these nuclei have been identified by the (d,p) and (d,t) reaction work of Fortune et al. (1971). The level spectra of the two nuclei are found to be quite similar and there is clear evidence of single-particle strength from the $\nu d_{5/2}$, $\nu d_{3/2}$, $\nu s_{1/2}$, $\nu g_{7/2}$ and $\nu h_{11/2}$ subshells. This strength is fractionated among many states. In a number of cases, it is possible to make comparisons between the levels (in ^{103}Ru and ^{105}Ru) observed in reactions and the levels in ^{101}Ru . These authors will not be referred to explicitly each time

their results are quoted.

It is also interesting to compare the level scheme of $^{101}_{44}\text{Ru}_{57}$ with that of $^{105}_{46}\text{Pd}_{59}$. The latter nucleus has been investigated by all three methods: Coulomb excitation (Bolotin and McClure 1971), beta decay (Pierson and Rengan 1967) and single-particle transfer reactions (Cujec 1963). Although ^{105}Pd has two more neutrons and two more protons than ^{101}Ru , the low energy portion of the two level schemes is rather similar. In fact, the spin sequence for the five lowest positive parity states [$0(5/2^+)$, $280(3/2^+)$, $306(7/2^+)$, $319(5/2^+)$ and $344(1/2^+)$ keV] is identical to ^{101}Ru . However, in contrast to ^{101}Ru , most of the other states in ^{105}Pd populated in Coulomb excitation have J^π values of $3/2^+$ or $5/2^+$. It is found that the states strongly excited in Coulomb excitation are only weakly populated in the transfer reactions and vice versa. The data for ^{105}Pd are consistent with an interpretation which involves configuration mixing of both single-particle and core-excited states.

The odd neutron isotope $^{101}_{44}\text{Ru}_{57}$ has seven neutrons outside the closed $N = 50$ neutron shell. The level order for neutron numbers greater than 50 is $2d_{5/2}$, $1g_{7/2}$, $3s_{1/2}$, $2d_{3/2}$, $1h_{11/2}$, etc. (Preston 1962). Thus, according to the extreme single-particle model, the first six neutrons would completely fill the $\nu d_{5/2}$ subshell and the seventh would be found in the $\nu g_{7/2}$ subshell. However, due to the pairing energy, it is expected that the higher angular

momentum $vg_{7/2}$ subshell will begin filling in pairs before the filling of the $vd_{5/2}$ subshell is completed. Therefore, the last two neutrons will be in the $vg_{7/2}$ subshell giving rise to the configuration $(vd_{5/2})^{-1}(vg_{7/2})_0^2$. This configuration results in a ground state spin and parity of $5/2^+$, the value observed experimentally. Since it is only a difference in pairing energy that results in the lowering of the $(vd_{5/2})^{-1}(vg_{7/2})_0^2$ configuration relative to the $(vd_{5/2})_0^6(vg_{7/2})$ configuration, the latter should form a single-particle state at low excitation. Similarly, other single-particle states such as $(vd_{5/2})_0^6(vd_{3/2})$, $(vd_{5/2})_0^6(vs_{1/2})$ or $(vd_{5/2})_0^6(vh_{11/2})$ might be expected to occur at low energy in ^{101}Ru . Although, the extreme single-particle model serves as a useful starting point for further discussion, it is clear that this simple model is inadequate to describe the complex level structure of ^{101}Ru . Therefore, an alternative description must be devised.

The nucleus $^{100}_{44}\text{Ru}_{56}$ is normally chosen as the even-even core for ^{101}Ru . Since ^{100}Ru has six neutrons outside the closed $N = 50$ shell, it is interesting to inquire whether or not this nucleus exhibits good closure of the $vd_{5/2}$ neutron subshell. Evidence from the zirconium and molybdenum isotopes (particularly $^{96}_{40}\text{Zr}_{56}$ and $^{98}_{42}\text{Mo}_{56}$) would suggest that good closure will not be a characteristic property of ^{100}Ru . Thus, it is anticipated that other neutron subshells such as $vg_{7/2}$, $vs_{1/2}$ and $vd_{3/2}$ will make important contributions to

the low-energy level structure of ^{100}Ru . By analogy with the molybdenum isotopes, it is expected that the ground state of ^{100}Ru will contain significant 2p-2h admixtures. The following is suggested as the predominant ground-state configuration of ^{100}Ru .

$$(\nu d_{5/2})_{0+}^6 + (\nu d_{5/2})_{0+}^{-2} \left[(\nu g_{7/2})_{0+}^2 + (\nu d_{3/2})_{0+}^2 + (\nu s_{1/2})_{0+}^2 \right]$$

which can be written as

$$(\nu d_{5/2})_{0+}^6 + \left[(\nu d_{5/2})_{0+}^{-2} (2 \text{ particle})_{0+} \right]$$

Similarly, the first 2^+ state in ^{100}Ru can be represented by the configuration

$$(\nu d_{5/2})_{0+}^{-2} \left[(\nu g_{7/2})_2^2 + (\nu d_{3/2})_2^2 \right]_{2+} + (\nu d_{5/2})_{2+}^{-2} \left[(2 \text{ particle})_{0+} \right]$$

Thus, it follows that the simple ground-state configuration of ^{101}Ru , suggested by the extreme single-particle model, should be modified to include the $(\nu d_{5/2})^{-1}(\nu d_{3/2})_0^2$ and $(\nu d_{5/2})^{-1}(\nu s_{1/2})_0^2$ configurations. The ^{101}Ru ground state can then be described by the configuration

$$\left[(\nu d_{5/2})^{-1} (2 \text{ particle})_0 \right]_{5/2+}$$

As indicated above, a $7/2^+$ excited state with the configuration $(\nu d_{5/2})_0^6(\nu g_{7/2})$ is expected to occur at low excitation in ^{101}Ru . This proposal is supported by the observation of low-lying $7/2^+$ states at 213 and 208 keV respectively, in ^{103}Ru and ^{105}Ru ; these states were populated by strong $\ell = 4$ transitions in the (d,p) and (d,t) experiments.

It should be noted that in the neighbouring nucleus ^{105}Pd , an analogous $7/2^+$ state at 306 keV is strongly populated in the (d,p) and (d,t) experiments, but very weakly excited in Coulomb excitation.

The $7/2^+$ state at 306.85 keV in ^{101}Ru is only weakly excited in Coulomb excitation; however, it is very strongly populated in the decay of ^{101}Tc with a beta transition having a very low $\log f_0 t$ value. Both of these facts suggest that the state has a predominantly single-particle nature. There is no reason to believe that this state will be composed of only one single-particle configuration. In fact, a more general configuration arises when a $g_{7/2}$ neutron is coupled to the ^{100}Ru ground state, which contains 2p-2h components. Therefore, the suggested configuration for the 306.85 keV state is given by

$$\begin{aligned} & \left[(v d_{5/2})_0^6 (v g_{7/2}) \right]_{7/2^+} + \left[(v d_{5/2})_0^{-2} \left[(v d_{3/2})_0^2 + (v s_{1/2})_0^2 \right] (v g_{7/2}) \right]_{7/2^+} \\ & + \left[(v d_{5/2})_0^{-2} (v g_{7/2})_{7/2}^3 \right]_{7/2^+} + \left[(v d_{5/2})^{-1} (2^+) \right]_{7/2^+} \end{aligned}$$

The third term arises from the coupling of the odd $g_{7/2}$ particle with the $(g_{7/2})_0^2$ component present in the ^{100}Ru ground state. The $(g_{7/2})_J^3$ configuration forms a number of states with spins ranging from $3/2^+$ to $15/2^+$. The last term, which represents a coupling of the $d_{5/2}$ hole with the 2^+ one-phonon state of the core, has been included to account for the weak Coulomb excitation probability and is expected

to have a rather small amplitude. The remaining strength will be discussed below.

The transition from the 306.85 keV level to the ground state is predominantly M1 (Sieniawski et al. (1971); there is no evidence of this level having an appreciable lifetime. If both the ground state ($5/2^+, \ell=2$) and the 306.85 keV level ($7/2^+, \ell=4$) were pure single-particle states, then the M1 transition between them should be ℓ -forbidden (since $\Delta\ell = 2$) resulting in a lifetime of $\sim 20 - 100$ nsecs for the 306.85 keV state. It would appear that there is some admixture in either the ground state or the 306.85 keV level which leads to a decrease in the hindrance of the deexciting M1 transition. If the ground state contained a portion of the $[(\nu d_{5/2})_0^{-2} (\nu g_{7/2})_{5/2}^3]_{5/2^+}$ configuration, then a transition $(\nu g_{7/2})_{7/2}^3 \rightarrow (\nu g_{7/2})_{5/2}^3$, which is not ℓ -forbidden, could proceed. This transition is seniority forbidden; however, this hindrance should be considerably less than the hindrance due to ℓ -forbiddenness. Alternately, an unhindered M1 transition can proceed between the core-excitation components $[(\nu d_{5/2})^{-1}(2^+)]_{3/2^+}$ and $[(\nu d_{5/2})^{-1}(2^+)]_{5/2^+}$ if some of the latter is present in the ground state.

The next level to be considered is the $1/2^+$ state at 325.2 keV. Although this level is not populated in the decay of ^{101}Tc , it is observed in Coulomb excitation and in the decay of $^{101g}_{45}\text{Rh}_{56}$. The neutron configuration of the ^{101}Rh ground state is expected to be very similar to that of ^{100}Ru , while the proton configuration is $[(\pi p_{1/2})^{-1}(\pi g_{9/2})_0^6]_{1/2^-}$.

Thus, the first-forbidden electron-capture decay of the ground state of ^{101}Rh to the 325.2 keV level can be explained by the transition $(\pi p_{1/2}) \rightarrow (\pi s_{1/2})$. The following is therefore, suggested as the configuration of the 325.2 keV state.

$$\begin{aligned} & \left[(v d_{5/2})_0^{-2} \left[(v g_{7/2})_0^2 + (v d_{3/2})_0^2 \right] (v s_{1/2}) \right]_{1/2^+} + \left[(v d_{5/2})_0^6 (v s_{1/2}) \right]_{1/2^+} \\ & \quad \left[(v d_{5/2})^{-1} (2^+) \right]_{1/2^+} \end{aligned}$$

The first two terms are formed by the addition of a single $s_{1/2}$ neutron to the ^{100}Ru ground state. It is through these terms that the electron-capture decay proceeds. The third component is necessary to explain the Coulomb excitation results. Since strong $l = 0$ transitions have been observed to low-lying states at 170 and 140 keV in ^{103}Ru and ^{105}Ru respectively, it is expected that the single-particle components in the 325.2 keV level are quite large. The lowest $1/2^+$ state in ^{105}Pd is also strongly populated in reaction studies.

The foregoing discussion has presented evidence for levels having significant admixtures from the $g_{7/2}$ and $s_{1/2}$ single-particle states. Therefore, one might also expect to find a level in ^{101}Ru with a large $d_{3/2}$ single-particle component. The $3/2^+$ states at 127.21 and 422.1 keV are possible candidates. The stripping data for ^{103}Ru and ^{105}Ru show that the lowest $3/2^+$ state in each isotope is very weakly populated; therefore, the single-particle amplitudes are probably small. In addition, the 127.21 keV state is

strongly excited in Coulomb excitation. This would indicate that this level has a large core-excitation component, $[(\nu d_{5/2})^{-1}(2^+)]_{3/2^+}$. There is also a possibility of some amplitude from the $[(\nu d_{5/2})_0^{-2}(\nu g_{7/2})_{3/2}^3]_{3/2^+}$ configuration. The following is suggested as the configuration of the $3/2^+$ state at 127.21 keV

$$[(\nu d_{5/2})^{-1}(2^+)]_{3/2^+} + [(\nu d_{5/2})_0^6(\nu d_{3/2})]_{3/2^+} \\ + [(\nu d_{5/2})_0^{-2}(2 \text{ particle})_0(\nu d_{3/2})]_{3/2^+} + [(\nu d_{5/2})_0^{-2}(\nu g_{7/2})_{3/2}^3]_{3/2^+}$$

It is interesting to compare the 0.77 nsec transition from this level to the ground state of ^{101}Ru with the analogous transition in ^{99}Ru . In the first case, the M1 part of the transition is hindered by a factor of ~ 12 , which seems reasonable in view of the relatively small rearrangement required to proceed from the configuration above to the $[(\nu d_{5/2})^{-1}(2 \text{ particle})_0]_{5/2^+}$ ground-state configuration. The analogous transition in ^{99}Ru has an energy of 89.6 keV and a lifetime of 20.7 nsec, corresponding to an M1 hindrance factor of $\sim 10,000$. The ground state of ^{99}Ru is most likely composed of the two single-particle components

$$(\nu d_{5/2})^{-1} + [(\nu d_{5/2})_{5/2}^3(2 \text{ particle})_0]_{5/2^+}$$

while the $3/2^+$ first excited state is probably

$$[(\nu d_{5/2})_{5/2}^3(2^+)]_{3/2^+} + [(\nu d_{5/2})_{3/2}^3(2 \text{ particle})_0]_{3/2^+} \\ + [(\nu d_{5/2})_0^{-2}(\nu d_{3/2})]_{3/2^+} + [(\nu d_{5/2})_0^2(2 \text{ particle})_0(\nu d_{3/2})]_{3/2^+}$$

The most striking difference between the descriptions of the two $3/2^+$ states is the presence of the

$\left[(v d_{5/2})_{3/2}^3 (2 \text{ particle})_0 \right]_{3/2^+}$ component in the ^{99}Ru case.

This component can easily be produced by adding a single $d_{5/2}$ neutron to the ^{98}Ru core, $\left[(v d_{5/2})_0^4 + (v d_{5/2})_0^2 (2 \text{ particle})_0 \right]$.

However, for the ^{101}Ru case, this component cannot be easily formed from the ^{100}Ru core, $\left[(v d_{5/2})_0^6 + (v d_{5/2})_0^{-2} (2 \text{ particle})_0 \right]$.

The $(v d_{5/2})_{3/2}^3 \rightarrow (v d_{5/2})_{5/2}^3$ ground-state transition in ^{99}Ru
 $\quad \quad \quad v=3 \quad \quad \quad v=1$

involves a change of two in seniority and therefore the transition may be expected to show considerable retardation.

The fact that the M1 part of this transition is strongly retarded suggests that the $\left[(v d_{5/2})_{3/2}^3 (2 \text{ particle})_0 \right]_{3/2^+}$

component is present with a fairly large amplitude. The

Coulomb excitation data show that the 127.21 keV state is

less collective than its 89.6 keV counterpart in ^{99}Ru ,

thus suggesting that the core-excitation amplitude in ^{99}Ru

is even stronger than in ^{101}Ru .

Although it is expected that the structure of the $3/2^+$ state at 422.1 keV will contain configurations similar to those of the 127.21 keV level, the amplitudes involved will be quite different. For example, the Coulomb excitation probability for the 422.1 keV level is much smaller than that for the 127.21 keV level so that the core-excitation amplitude will be correspondingly smaller. The $\log f_0 t$ values of the electron-capture transitions to the 127.21 and 422.1 keV states are identical (Sieniawski et al. 1971); this would suggest that the single-particle amplitudes are similar.

The 3-quasi-particle amplitudes may be quite different.

The transition from the 325.2 keV level to the 127.21 keV level is predominantly M1 (Sieniawski et al. 1971).

Since the former state is $s_{1/2}$ in character ($\ell=0$) and the latter $d_{3/2}$ or $d_{5/2}$ ($\ell=2$), this transition should be ℓ -forbidden. However, an unhindered M1 transition can

proceed between the core-excitation components

$$\left[(v d_{5/2})^{-1} (2^+) \right]_{1/2^+} \text{ and } \left[(v d_{5/2})^{-1} (2^+) \right]_{3/2^+}.$$

The 311.32 keV level has a spin and parity of $5/2^+$.

This state is populated with moderate strength in the Coulomb excitation experiments and therefore contains a portion of the $\left[(v d_{5/2})^{-1} (2^+) \right]_{5/2^+}$ core-excited component. The remaining strength is not accounted for; however, it should be noted that much of the expected strength for the $5/2^+$ component has been found in ^{105}Pd . The following is suggested for the structure of the 311.32 keV level.

$$\begin{aligned} & \left[(v d_{5/2})_0^{-2} (v g_{7/2})_{5/2}^3 \right]_{5/2^+} + \left[(v d_{5/2})^{-1} (2^+) \right]_{5/2^+} \\ & + \left[(v d_{5/2})^{-1} (2 \text{ particle})_0 \right]_{5/2^+} \end{aligned}$$

The last term accounts for the amplitude not present in the ground state.

The two states at 545.09 and 720.0 keV are strongly populated in the Coulomb excitation experiments with $B(E2)$ values approaching the $B(E2)$ value for the 2^+ state in ^{100}Ru . They are also connected by a "fast" M1 transition. The 545.09 keV level has a spin and parity of $7/2^+$ and can be

associated with the $\left[(\nu d_{5/2})^{-1}(2^+)\right]_{7/2^+}$ member of the core-excited multiplet. Similarly, if the $9/2^+$ choice is assumed for the 720.0 keV state, this level can be identified with the $\left[(\nu d_{5/2})^{-1}(2^+)\right]_{9/2^+}$ member. It should be noted that other configurations must be admixed with the above in order to explain the $\log f_0 t$ values for the beta transitions to the 545.09 and 720.0 keV levels. Since the 545.09 keV level has a relatively low $\log f_0 t$ of 5.6, it would appear that some of the $g_{7/2}$ single-particle strength (see the 306.85 keV level) is present as an admixture in its wavefunction. Of course, the amplitude is considerably smaller than that present in the wavefunction of the 306.85 keV state. The $\log f_0 t$ value for the 720.0 keV state is somewhat higher since there are no single-particle states with a spin and parity of $9/2^+$ to mix with its unperturbed wavefunction. One configuration through which the beta decay can proceed is the $9/2^+$ component of the $(\nu g_{7/2})^3_J$ configuration, $\left[(\nu d_{5/2})^{-2}_0 (\nu g_{7/2})^3_{9/2}\right]_{9/2^+}$.

A number of other high-spin states have been observed in the decays of ^{101}Rh and $^{101\text{m}}\text{Rh}$. These include the following: 616.2 ($7/2^+$), 623.5 ($7/2^+$), 643.8 ($7/2^+$ or $9/2^+$), 842.76 ($7/2^+$), 928.96 ($7/2^+$ or $9/2^+$), 938.32 ($7/2^+$), and 1001 keV ($11/2^+$). Several of the levels (the 616.2, 623.5 and 928.96) have also been seen in the Coulomb excitation experiments indicating that they possess some core-excitation strength. The limited number of configurations discussed above are unable to explain all of the observed levels and therefore

it is necessary to consider other types of configurations which can be found at relatively low excitation. Lederer et al. (1971) have observed states which they have attributed to the non-zero coupling of the $(d_{5/2})^{-2}$ configuration. For example, the states at 1001 and 1861 keV can be identified with the $\left[(\nu d_{5/2})_J^{-2} (\nu g_{7/2})_{7/2}^3\right] 11/2^+$ or $15/2^+$ configuration ($J = 2$ or 4). Since the $(\nu d_{5/2})_0^{-2} (\nu g_{7/2})_J^3$ configuration also produces a component of spin and parity $11/2^+$, there is a possibility of some configuration mixing in the 1001 keV wavefunction. It should be pointed out that the $(\nu d_{5/2})_2^{-2} (\nu g_{7/2})_J^3$ configuration leads to at least 21 states, four of which have a spin of $7/2$ and three a spin of $9/2$. Thus, it is clear that the existence of such a configuration in ^{101}Ru would easily produce the number of states required by the experimental data. The states produced by this configuration will probably mix with the configurations previously assigned to the low-lying levels. This configuration can also account for the many $3/2^+$ and $5/2^+$ states observed in ^{105}Pd .

Until now, no mention has been made of the $h_{11/2}$ orbital. An $11/2^-$ state has been found at 528 keV in ^{101}Ru by several workers (see section 3.4). In addition, a single-particle state, populated with an $\ell=5$ transition, has been observed in the stripping experiments for ^{103}Ru and ^{105}Ru . Moreover, Lederer et al. (1971) have found levels at higher excitation which can be attributed to the coupling of an $h_{11/2}$ neutron to the $(\nu d_{5/2})_{0+,2+,4+}^2$ configuration. The

528 keV state is expected to contain an admixture of the $(\nu d_{5/2})_0^6 (\nu h_{11/2})$ and $(\nu d_{5/2})_0^{-2} (2 \text{ particle})_0 (\nu h_{11/2})$ configurations.

Other configurations such as $\left[(\nu d_{5/2})_4^{-2} (\nu g_{7/2})_{3/2^+}^3 \text{ to } 15/2^+ \right]_{J'}$, $\left[(\nu d_{5/2})_J^3 (\nu g_{7/2})_0^4 \right]_J$ ($J = 3/2^+, 5/2^+$ or $9/2^+$), and $(\nu d_{5/2})_J^{-2}$ ($J = 2$ or 4) coupled to $(\nu d_{3/2})$ or $(s_{1/2})$ might be expected to give rise to other levels at somewhat higher excitation than the configurations already considered. It is probably that there will be mixing between configurations of similar spin.

From the preceding discussion, it would appear that there is considerable fragmentation of the available single-particle and core-excitation strength among many states. In addition to these configurations, a number of other configurations of a particle-hole type must be introduced to explain the experimental results. There is considerable evidence for mixing of the various types of configurations. Until further experimental work has been carried out, the conclusions reached above must be regarded as rather tentative. However, in view of the complexity of the low-lying level schemes of the odd-neutron nuclei in this region, it is clear that if a realistic shell model calculation is to be performed, it should take into account many of the configurations that have been suggested.

4.3 Interpretation of the Technetium Isotopes

The technetium nuclei have 43 protons. If $^{88}_{38}\text{Sr}_{50}$ is considered as the core nucleus, the $1f_{5/2}$ and $2p_{3/2}$ subshells will be completely filled and the five extra-core protons will be distributed among the $2p_{1/2}$ and $1g_{9/2}$ subshells. The extreme single-particle model predicts that the forty-

third proton will be found in the $1g_{9/2}$ subshell. Because of pairing, considerable energy can be gained if the unpaired proton is in the $2p_{1/2}$ rather than the $1g_{9/2}$ subshell. Therefore, the $g_{9/2}$ and $p_{1/2}$ single-particle states should be found close together in the level schemes of the technetium isotopes.

(a) The Ground States of ^{97}Tc , ^{99}Tc and ^{101}Tc

The low-energy levels of several of the odd-mass technetium isotopes are shown in Fig. 4.1. The levels in ^{95}Tc are those populated in the decay of ^{95}Ru (Tucker and Hein 1970); the data for the other three nuclei come from the present work. The ground states are all $9/2^+$ and can be identified with the configuration $(\pi p_{1/2})_0^2 (\pi g_{9/2})_{9/2+}^3$. It has been suggested in the case of ^{93}Tc (Cohen et al. (1964) and Auerbach and Talmi (1965)) that the ground state will also contain admixtures of the configuration $(\pi g_{9/2})_{9/2+}^5$. This suggestion was confirmed by the $(^3\text{He}, d)$ and $(d, ^3\text{He})$ reaction experiments on ^{92}Mo (Ohnuma and Yntema 1968). These data are consistent with the theoretical predictions of Auerbach and Talmi (1965) that the ^{92}Mo ground state is 70% $(\pi p_{1/2})_0^2 (\pi g_{9/2})_0^2$ and 30% $(\pi g_{9/2})_0^4$ and that the ground state of ^{93}Tc is 80% $(\pi p_{1/2})_0^2 (\pi g_{9/2})_{9/2+}^3$ and 20% $(\pi g_{9/2})_{9/2+}^5$. Cohen et al. (1964) showed that the admixture of higher seniority (e.g. $v=3$) configurations is very small in the ground-state wavefunctions so that the above configurations are predominantly seniority one ($v=1$). The ground states of all the technetium isotopes can thus be represented by the configuration

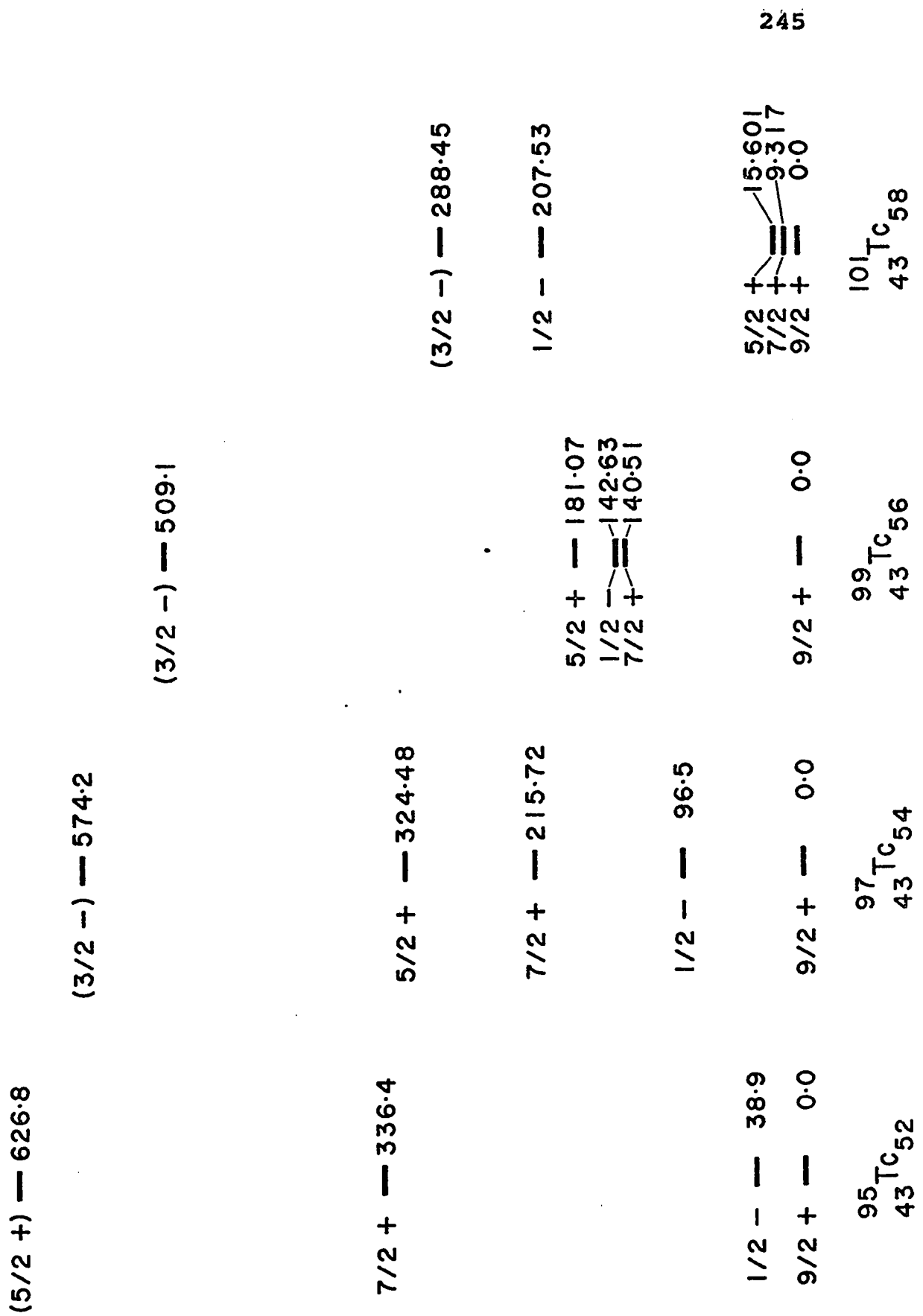


Fig. 4.1 The low-lying states in the odd-mass technetium isotopes

$$(\pi p_{1/2})_0^2 (\pi g_{9/2})_{v=1}^3 + (\pi g_{9/2})_{v=1}^5$$

Although there are (d, ^3He) data available for ^{94}Mo , ^{96}Mo and ^{98}Mo (Ohnuma and Yntema 1968), no (^3He , d) results have been published. In the (d, ^3He) data, no significant variation in the amplitudes of the two configurations was observed for this series of nuclei.

(b) Low-Lying Negative Parity States

The excitation energy of the $1/2^-$ isomeric state shows a smooth variation with neutron number as shown in Fig. 4.1. These states are identified with the particle-hole configuration $(\pi p_{1/2})^{-1}(\pi g_{9/2})_0^4$. As in the case of the niobium isotopes, the upward trend of the excitation energy would indicate that the $2p_{1/2}$ orbital is becoming increasingly full as neutrons are added.

Similar particle-hole configurations, e.g.

$(\pi p_{3/2})^{-1}(\pi p_{1/2})_0^2(\pi g_{9/2})_0^4$, with holes in the $2p_{3/2}$ and $1f_{5/2}$ subshells are expected to produce single-particle $3/2^-$ and $5/2^-$ states at higher excitation energy in the technetium isotopes. In fact, the (^3He , d) results for ^{93}Tc (Kozub and Youngblood 1971) indicate that a realistic ground-state configuration for ^{92}Mo should include sizeable

$$(\pi p_{3/2})_0^{-2}(\pi p_{1/2})_0^2(\pi g_{9/2})_0^4 \text{ and } (\pi f_{5/2})_0^{-2}(\pi p_{1/2})_0^2(\pi g_{9/2})_0^4$$

components. While these will be discussed below, it is perhaps worth noting that levels containing significant $p_{3/2}$ and $f_{5/2}$ strength have been observed in the neighbouring

niobium isotopes where the experimentally observed strengths represent $\sim 100\%$ and $\sim 50\%$ of the sum rule limits, respectively. The total $f_{5/2}$ strength is probably fragmented among more levels than the $p_{3/2}$ strength. The data indicate that there is considerable mixing between pure single-particle states and states which are not single-particle in nature.

An alternative description of the low-energy $3/2^-$ and $5/2^-$ states in technetium is provided by the particle-core weak-coupling model (De-Shalit 1961) in which the $p_{1/2}$ hole state is coupled to the 2^+ state of the even-even core. If this is a valid interpretation, then there should also be other negative parity states at higher excitation due to the coupling of the $p_{1/2}$ proton hole to the two-quadrupole phonon states of the core.

Before discussing the technetium isotopes, it is informative to consider some of the silver ($Z = 47$) and rhodium ($Z = 45$) isotopes. The experimental evidence indicates that many of the low-lying negative parity states can be explained within the framework of the particle-core weak-coupling model.

In ^{107}Ag and ^{109}Ag , the ground state is $1/2^-$ and the lowest-lying negative parity states are $3/2^-$ and $5/2^-$. Robinson et al. (1970) have shown that the excitation energies of the $3/2^-$ and $5/2^-$ states and the $B(E2)$ values for the $3/2^- \rightarrow 1/2^-$ and $5/2^- \rightarrow 1/2^-$ transitions are in reasonable agreement with those expected from the De-Shalit model. That is, the $B(E2)$ values are nearly equal and are similar to the $B(E2; 2 \rightarrow 0)$ of the even-mass cores. They have also proposed candidates

for the states attributed to the coupling of the $p_{1/2}$ hole state to the second 2^+ state of the core and showed that the model gave qualitative agreement with experiment. The average energy (the excitation energies are weighted by $2J+1$) of this two-phonon multiplet is $\sim 67\%$ of the energy of the 2^+ state in agreement with the average energy of the lowest $3/2^-$ and $5/2^-$ states relative to the 2^+ state. The $B(E2)$ values for the higher-lying states do not agree with the predicted values. The agreement is improved by considering an admixture of the one- and the two-quadrupole phonon states. However, such admixtures cannot explain the large $B(M1)$ values for the ground-state transitions from the $3/2^-$ states. Also the inelastic scattering results (Ford et al. 1970) indicate that only part of the total wavefunction of the second $3/2^-$ and $5/2^-$ states can be accounted for with the one- and two-phonon wavefunctions. It would appear that the remainder is due to admixtures of the $p_{3/2}$ and $f_{5/2}$ hole states.

As in the silver isotopes, the experimental results for the rhodium isotopes, ^{101}Rh and ^{103}Rh , show the existence of a number of core-excited levels. These nuclei have $1/2^-$ ground states and low-lying $3/2^-$ and $5/2^-$ states. The average energy of the ^{103}Rh (^{101}Rh) multiplet is 65% (61%) of the average energy of the 2^+ state in the neighbouring even-even nuclei ^{102}Ru and ^{104}Pd (^{100}Ru and ^{102}Pd). Since ^{103}Rh is stable, it has been investigated with Coulomb excitation

(Black et al. (1969) and Sayer et al. (1972)). The $B(E2; 3/2^- \rightarrow 1/2^-)$ and $B(E2; 5/2^- \rightarrow 1/2^-)$ values are approximately equal and their average compares favourably with the average $B(E2; 2 \rightarrow 0)$ value for ^{102}Ru and ^{104}Pd . Again, the $B(M1; 3/2^- \rightarrow 1/2^-)$ is quite large instead of being zero as predicted by the model. Higher energy negative parity states may arise from the coupling of the $p_{1/2}$ hole state to the 0^+ , 2^+ and 4^+ two-phonon states of the core. There is good evidence for the 2^+ two-phonon multiplet ($5/2^-$ and $3/2^-$) at 880 and 1107 keV in ^{103}Rh . Two other levels, which are populated almost exclusively by double E2 Coulomb excitation, appear to be candidates for the 4^+ two-phonon multiplet ($7/2^-$ (848 keV) and $9/2^-$ (920 keV)). A weakly excited state at 803 keV might be the $1/2^-$ level arising from the 0^+ two-phonon state. Sayer et al. (1972) also believe that they have identified a $p_{3/2}$ single-particle state at 1277 keV. It should be noted that the $3/2^-$ and $5/2^-$ core-excited states in ^{101}Rh have also been observed in the $^{100}\text{Ru}(^3\text{He}, d)^{101}\text{Rh}$ reaction (Phelps and Sarantites 1970). This would appear to substantiate the fact that considerable mixing occurs between the single-particle ($p_{3/2}$ and $f_{5/2}$) and core-excited states.

The experimental evidence for the niobium, rhodium and silver isotopes is consistent with an interpretation which involves the configuration mixing of single-particle and core-excited states of like spin and parity. In view of this fact, it has been assumed that this mixed structure is

also present in the technetium isotopes.

(c) Negative Parity States in ^{101}Tc

Due to the large decay energy of 2.84 MeV, many negative parity levels in ^{101}Tc have been populated in the decay of ^{101}Mo . It is expected that all of the $1/2^-$, $3/2^-$ or $5/2^-$ states observed in ^{97}Tc or ^{99}Tc will also be found in ^{101}Tc and therefore most of the following discussion will concern the latter nucleus. Negative parity levels can arise from holes in the $p_{1/2}$, $p_{3/2}$ and $f_{5/2}$ proton subshells or from the coupling of these hole states to the one- or two-phonon excitations of the even-even core. Table 4.1 indicates the negative parity states which are expected at low excitation.

The $1/2^-$ state at 207.53 keV is the $p_{1/2}$ hole state in ^{101}Tc . The J^π values and excitation energies of the 288.45 ($3/2^-$) and 394.40 ($5/2^-$) keV states indicate that these are the most likely candidates for the states which can be ascribed to coupling of the $p_{1/2}$ proton hole to the first 2^+ state of the even-even core. The center of gravity of these two states is $\sim 60\%$ of the average energy (590 keV) of the first 2^+ state in ^{100}Mo (700 keV) and ^{102}Ru (475 keV). This value is in reasonable agreement with the value of 65% obtained in the silver and rhodium nuclei. No $B(E2)$ values are known for the 288.45 and 394.40 keV states and the $5/2^- \rightarrow 3/2^-$ transition has not been observed (upper limit of intensity < 0.1). Nevertheless, by comparing the ratio of the $5/2^- \rightarrow 3/2^-$ and $5/2^- \rightarrow 1/2^-$ reduced transition probabilities in ^{107}Ag and ^{101}Tc

Table 4.1

Negative Parity States in the Technetium Isotopes

<u>Description</u>	<u>Configuration</u>	<u>J^π value</u>	<u>Designation</u>
p _{1/2} hole	$(\pi p_{1/2})^{-1} (\pi g_{9/2})^4_0 \equiv (1/2^-)$	1/2	1/2
p _{1/2} hole + 1-phonon	$[(1/2^-) (2^+)]_J$	3/2 ⁻ , 5/2 ⁻	3/2 ⁻ , 5/2 ⁻
	$[(1/2^-) (0^+)]_J$	1/2 ⁻	1/2 ^{-'}
p _{1/2} hole + 2-phonon	$[(1/2^-) (2^{+'})]_J$	3/2 ⁻ , 5/2 ⁻	3/2 ^{-'} , 5/2 ^{-'}
	$[(1/2^-) (4^+)]_J$	7/2 ⁻ , 9/2 ⁻	7/2 ^{-'} , 9/2 ^{-'}
p _{3/2} hole	$(\pi p_{3/2})^{-1} (\pi p_{1/2})^2_0$ $(\pi g_{9/2})^4_0$	3/2 ⁻	3/2 ^{-*}
f _{5/2} hole	$(\pi f_{5/2})^{-1} (\pi p_{3/2})^4_0$ $(\pi p_{1/2})^2_0 (\pi g_{9/2})^4_0$	5/2 ⁻	5/2 ^{-*}

it appears that the $5/2^- \rightarrow 1/2^-$ transition is at least an order of magnitude weaker than that observed in ^{107}Ag . Thus, the core-coupling model alone does not appear to be completely adequate for a description of the 288.45 and 394.40 keV states. It is likely that these states contain significant single-particle admixtures from the $2p_{3/2}$ and $1f_{5/2}$ hole states.

If it is assumed that the predominant configuration of the 288.45 and 394.40 keV states is $[(p_{1/2} \text{ hole}) (2^+)]_J$, then the discussion can be extended to the other negative parity states. It is expected that the states $(3/2^-$ and $5/2^-)$ attributed to the coupling of the $p_{1/2}$ hole state to the 2^+ two-phonon state of the core will occur at ~ 645 keV (60% of 1075 keV, the average energy of the 2^+ state in ^{100}Mo (1050 keV) and ^{102}Ru (1103 keV)). Therefore, one of the $5/2^-$ levels at 669.45 and 742.22 keV can be identified with the $5/2^-$ state. The other is probably the $f_{5/2}$ hole state ($5/2^{-*}$). The 742.22 keV state decays only to the $7/2^+$ state at 9.317 keV. In addition to decaying to the $7/2^+$ state, the 669.45 keV state decays to the $3/2^-$ and $5/2^-$ levels at 288.45 and 394.40 keV respectively, and to the 500.45 keV state. In view of the decay patterns, it is most likely that the $f_{5/2}$ hole state forms the predominant configuration of the 742.22 keV state while the 669.45 keV state has the $p_{1/2}$ proton hole coupled to the 2^+ core state as its predominant configuration.

As shown below, the $7/2^+$ first excited state is primarily

composed of the $(\pi g_{9/2})^3_{7/2+}$ configuration. The $f_{5/2}$ hole state, which has the configuration $(\pi f_{5/2})^{-1}(\pi p_{3/2})^4_0(\pi p_{1/2})^2_0(g_{9/2})^4_0$, can decay to the $7/2^+$ state if a $g_{9/2}$ proton ($\ell=4$) is changed into an $f_{5/2}$ proton ($\ell=3$) and if the remaining three $g_{9/2}$ protons are recoupled to give a spin of $7/2$. Although this is not a simple single-particle transition, it is still an E1 transition (somewhat hindered). The core-excited states (e.g. $[(\pi f_{5/2})^6_0(\pi p_{3/2})^4_0(\pi p_{1/2})^2_0(\pi g_{9/2})^4_{2+}]_{5/2-}$) can only decay to the $7/2^+$ state by changing a $g_{9/2}$ proton ($\ell=4$) into a $p_{1/2}$ proton ($\ell=1$) and then recoupling the remaining $g_{9/2}$ protons. This is an ℓ -forbidden ($\Delta\ell=3$) transition and E1 transitions would be expected to be highly hindered. Since all three $5/2^-$ states decay to the $7/2^+$ first excited state, it is clear that considerable mixing of the $5/2^-$ states takes place and that the $f_{5/2}$ hole strength is distributed among the three levels at 394.90, 669.45 and 742.22 keV. It is not clear why the 394.40 keV state decays more strongly to the $5/2^+$ state at 15.606 keV than to the $7/2^+$ state at 9.317 keV. The other $5/2^-$ states do not appear to decay to the $5/2^+$ state.

There are three states in Table 4.1 yet to be accounted for; these can be associated with the negative parity levels at 500.45, 616.20 and 622.15 keV. The 616.20 keV level decays very strongly to the 207.53 keV state and also to the 288.45, 394.40 and 500.45 keV states. The 500.45

and 622.15 keV levels decay only to the 288.45 keV state. The experimental results for the rhodium and silver isotopes showed that the decays from the $3/2^-$ and $5/2^-$ states, which were attributed to coupling of the $p_{1/2}$ hole state to the 2^+ core state, to the $p_{1/2}$ state were very weak. Therefore, the 616.20 keV level is most likely the $p_{3/2}$ hole state. The transition to the $p_{1/2}$ state would be a $(p_{3/2})^{-1}(p_{1/2})_0^2 \rightarrow (p_{3/2})_0^4(p_{1/2})$ transition; such a transition would not be hindered.

Since the spins of the 500.45 and 622.15 keV levels are not known, it is not clear which of these is the $3/2^-$ state or the $1/2^-$ state. The existence of transitions from levels at 1103.5 and 1141.5 keV would tend to support a $3/2^-$ assignment for the 500.45 keV state. The former levels are believed to be $[(p_{3/2} \text{ hole})(2^+)]_J$ states (see below). However, until the spins are measured, no positive assignments can be made.

As in the case of the $5/2^-$ levels, there is probably considerable mixing of the configurations forming the $3/2^-$ states and therefore all three $3/2^-$ states will contain some $p_{3/2}$ hole strength.

Five other negative parity levels at 1103.5, 1141.5, 1188.04, 1197.3 and 1231.3 keV have been identified in ^{101}Tc . Since the excitation energy of these levels is approximately equal to the $p_{3/2}$ or $f_{5/2}$ hole excitation energy plus the 2^+ excitation energy, they can probably be attributed to the

coupling of the $p_{3/2}$ or $f_{5/2}$ hole states with the 2^+ core state. The former coupling gives J^π values of $1/2^-$, $3/2^-$, $5/2^-$ and $7/2^-$; the latter gives J^π values of $1/2^-$, $3/2^-$, $5/2^-$, $7/2^-$ and $9/2^-$. The two lowest energy states at 1103.5 and 1141.5 keV decay to the $3/2^-$ state ($p_{1/2}$ coupled to the 2^+) at 288.45 keV. In addition, the 1103.5 and 1188.04 keV states decay to the $p_{1/2}$ state at 207.53 keV and the 1188.04 keV state decays to the predominantly $p_{3/2}$ state at 616.20 keV. The decay patterns of these states are consistent with an interpretation of a $p_{3/2}$ hole state coupled to the 2^+ core state. These states would be the $1/2^-$, $3/2^-$ and $5/2^-$ members of the multiplet. Until further spin values are assigned, it is impossible to make any further arguments.

The preceding discussion of the negative parity levels in ^{101}Tc has accounted reasonably well for the character of these states. The mixing of the single-particle and core-excitation configurations probably explains the absence of the $5/2^- \rightarrow 1/2^-$ transitions. The E2 transition is not able to compete with the unhindered M1 transitions which can proceed through the single-particle admixtures existing in each state.

d) Negative Parity States in ^{99}Tc and ^{97}Tc

The low-energy level structures of ^{98}Mo and ^{100}Ru , the core nuclei for ^{99}Tc , are similar to those of ^{100}Mo and ^{102}Ru respectively, the core nuclei for ^{101}Tc . Therefore,

the low-energy negative parity states in ^{99}Tc should have a structure which corresponds to that of ^{101}Tc . Thus, the center of gravity of the multiplet formed by the coupling of the $p_{1/2}$ hole state to the one-phonon state might be expected to be at ~ 400 keV (60% of 664 keV, the average energy of the 2^+ state in ^{98}Mo (788) keV) and ^{100}Ru (540)), if experience in silver, rhodium and ^{101}Tc is any guide. The center of gravity of the observed $3/2^-$ (509.13 keV) and $5/2^-$ (671.6 keV) states is at ~ 605 keV, a value which is much larger than expected. This suggests that the core-excitation interpretation is incorrect for these states and that the core-excited states are not populated in the decay of ^{99}Mo with its rather low Q-value. The 509.13 keV state ($3/2^-$ or $1/2^-$), which is strongly populated in beta decay, deexcites only to the level at 142.68 keV ($p_{1/2}$ hole state). The data suggest that the $p_{3/2}$ hole state is the predominant configuration of the 509.13 keV level. Similarly, the 671.6 keV level ($5/2^-$) can be associated with the $f_{5/2}$ hole state. If these assignments are correct, a ($^3\text{He}, d$) experiment on ^{98}Mo should find a considerable fraction of the $2p_{3/2}$ and $1f_{5/2}$ strengths in the 509.13 and 671.6 keV levels, respectively.

At higher excitation in ^{99}Tc , negative parity levels are found at 1004.0, 1073.0, 1129.3 and 1199.4 keV. These levels occur at about the same energy as the group of negative parity levels in ^{101}Tc . Therefore, they can reasonably be attributed to the coupling of the $p_{3/2}$ or $f_{5/2}$

hole states to the one-phonon excitations of the core. The 1003.9 keV level decays very strongly to the $5/2^+$ state at 181.07 keV (predominant configuration $(\pi g_{9/2})^3_{5/2^+}$). Therefore, the 1003.9 keV level probably contains considerable $f_{5/2}$ strength and can be associated with the $3/2^-$ member of the $(f_{5/2^+} + 2^+)$ core-excited multiplet. A unique spin assignment has not been made for the two levels at 1129.3 and 1199.4 keV. However, both levels decay to the $p_{1/2}$ state at 142.68 keV and to the $p_{3/2}$ state at 509.13 keV. Thus, they are probably two members of the $(p_{3/2^+} + 2^+)$ core-excited multiplet. There is insufficient data available on the weakly populated 534.4 and 1073.0 keV levels to make a meaningful assignment. One can conjecture that the 534.4 keV level may be the $(p_{1/2^+} + 2^+)_{5/2^-}$ core-excited state.

The 574.2 ($3/2^-$) and 656.8 ($5/2^-$) keV levels are the only negative parity states identified in ^{97}Tc . As in ^{99}Tc , these states appear to be too high in excitation energy to be associated with the $(p_{1/2^+} + 2^+)$ core-excited multiplet. They are assumed to have considerable $p_{3/2}$ and $f_{5/2}$ hole strength, respectively.

e) The Low-Lying Positive Parity States

The previous discussion has concentrated on the negative parity states found in the technetium isotopes, the remainder of the discussion will be confined to the positive parity states. Much of the discussion will concentrate on ^{101}Tc due to the larger number of levels observed experimentally.

As shown in Fig. 4.1, the splitting of the low-lying $5/2^+$ and $7/2^+$ states and their separation from the $9/2^+$ ground state decrease with increasing neutron number. Since the filling of the $1g_{9/2}$ proton subshell results in a major shell closure at $Z=50$, it is expected that other positive parity shell model orbitals, such as the $1g_{7/2}$ and $2d_{5/2}$, will be found at higher excitation energies. Therefore, these orbitals cannot be identified with the low-lying $5/2^+$ and $7/2^+$ states found in the odd mass technetium isotopes.

A simple weak-coupling interpretation of these $5/2^+$ and $7/2^+$ levels is inappropriate in the case of ^{101}Tc since the average energy of the one-phonon vibrational state in the neighbouring even-even nuclei is ~ 630 keV. Therefore, very strong perturbations would be required to displace the one-phonon states down to within ~ 15 keV of the ground state. The situation is not so clear in the other technetium isotopes because the $5/2^+$ and $7/2^+$ states are found at higher excitation energies than in ^{101}Tc . Although these states may contain some phonon admixture, this will not be the major component since they are still found at much lower excitation energies than expected from weak-coupling predictions.

These low-lying states could be described in terms of the quasi-particle-phonon coupling (QPC) model of Kisslinger and Sorensen (1963). For the $g_{9/2}$ shell, this intermediate coupling model predicts positive-parity levels with spins

$j, j \pm 1, j \pm 2$ (where j is the quasi-particle of spin $9/2$) at an excitation corresponding to the vibrational phonon energy $\hbar\omega$. For the technetium isotopes, this model predicts that the excitation of the $9/2^+$ state increases rapidly with neutron number so that the $1/2^-$ state becomes the ground state for $A > 95$: this prediction does not agree with the experimental data, since the ground state is $9/2^+$ for these isotopes. The QPC calculations also fail to account for the existence of low-lying $5/2^+$ and $7/2^+$ states in the technetium isotopes. Several attempts have been made to improve the QPC coupling scheme. Sherwood and Goswami (1966) have extended this theory through the inclusion of correlation effects ("backward-coupling") in the core states and have been able to reduce the energy separations between phonon and quasi-particle levels. The calculated level energies for the technetium isotopes (Goswami and Sherwood 1967) are in reasonable agreement with the $9/2^+ - 1/2^-$ level spacing observed experimentally for the lighter technetium nuclei, but predict a $1/2^-$ ground state for ^{101}Tc in disagreement with experiment. The $5/2^+, 7/2^+, 11/2^+$ and $13/2^+$ phonon levels are degenerate in their calculation; a lowering of the energy of the phonon levels with increasing neutron number is predicted. In a more recent calculation, Goswami and Nalcioglu (1968) have used the measured quadrupole moments of ^{104}Ru and ^{102}Pd to evaluate the quadrupole interaction matrix element in the EQPC theory of Sherwood and Goswami (1966). This cal-

culatation removed the phonon degeneracy of the EQPC and was able to account for the increased splitting of the $5/2^+$ and $7/2^+$ states in going from the Tc isotopes to the Ag isotopes. Since low-lying $13/2^+$ and $11/2^+$ states are predicted for all isotopes the failure to observe these states experimentally (Bond et al. 1971) casts some doubt on the applicability of this model to the odd-mass technetium nuclei.

Kisslinger (1966) has shown that the $(g_{9/2})^3$ three-quasi-particle state with spin $7/2^+$ can be significantly lowered in energy by the quadrupole interaction. There is some evidence for a $5/2^+$ state being lowered from the unperturbed three-quasi-particle energy as well. An exact diagonalization of the pairing plus quadrupole-force Hamiltonian for the three-particle states of the $g_{9/2}$ shell model configuration by Drozdov (1962) has shown that the splitting of the $7/2^+$ and $9/2^+$ levels is particularly sensitive to the relative strengths of the pairing and quadrupole interactions. For a strong quadrupole interaction, the $7/2^+$ and $9/2^+$ states can become degenerate.

An important test of these core-coupling approaches would result from comparison of the E2 transitions strengths between the $5/2$, $7/2$ and $9/2$ states. It is noted here that electronic measurement of the mean lives of the 9.317 and 15.601 keV levels in ^{101}Tc would provide this valuable information since the M1/E2 branching ratios are now known.

(f) Postive Parity States in ^{101}Tc

Of the various models proposed, the three-quasi-particle picture seems to give the best agreement with the

experimental data for the technetium isotopes. This model accounts for the presence of low-lying $7/2^+$ states and provides a mechanism for explaining the observed $7/2^+ - 9/2^+$ splitting as a function of increasing neutron number. In addition to the $5/2^+$ and $7/2^+$ seniority-three states, the $(\nu g_{9/2})_J^3$ configuration gives rise to other seniority-three states with J^π values of $3/2^+$, $9/2^+$, $11/2^+$, $15/2^+$, $17/2^+$ and $21/2^+$; of these, only the $3/2^+$ member is likely to be populated in the decay of ^{101}Mo . As indicated in section 4.3(a) the technetium ground states also contain some admixture from the $(\pi g_{9/2})_J^5$ configuration which consists of 20 states including single $1/2^+$ and $3/2^+$ states, two $5/2^+$ and two $7/2^+$ states. The 606.41 ($1/2^+$ or $3/2^+$) and 711.12 ($3/2^+$) keV levels could be associated with two of the three low-spin levels expected from the $(\pi g_{9/2})_J^3$ and $(\pi g_{9/2})_J^5$ configurations. This suggestion is supported by the strong decays of the 606.41 and 711.12 keV levels to the low energy states at 15.601 and 9.317 keV.

The observation of many $1/2^+$ or $3/2^+$ states between 1 and 2.5 MeV in the decay of ^{101}Mo make it clear that the above description is much too simple to explain the level structure of ^{101}Tc . Two additional facts support this conclusion. First of all, the 606.41 keV state decays quite strongly to the $1/2^-$ state at 207.53 keV and to the $[p_{1/2} + 2^+]_{3/2^-}$ state at 288.45 keV. The 711.12 keV state also decays to the latter state. If the $3/2^+$ states were pure $(\pi g_{9/2})_{3/2^+}^3$ or

$(\pi g_{9/2})^5_{3/2^+}$ configurations, these decays would be highly forbidden. In addition, these levels are populated with allowed transitions in the decay of ^{101}Mo . If the ground-state of ^{101}Mo ($1/2^+$) is predominantly a $\nu s_{1/2}$ single-particle state (Hjorth and Cohen 1964), allowed beta transitions could not proceed to states consisting of pure $g_{9/2}$ proton configurations. Therefore, it is clear that the $1/2^+$ and $3/2^+$ states in ^{101}Tc must contain admixtures from other types of configurations.

Consideration of the following example may serve to illustrate the complexity of the problem. Two possible neutron components in the $1/2^+$ ground state of ^{101}Mo are

$$\begin{aligned} & \left[(\nu d_{5/2})^6_0 \left[(\nu g_{7/2})^2_0 + (\nu d_{3/2})^2_0 \right] (\nu s_{1/2}) \right]_{1/2^+} + \left[(\nu d_{5/2})^{-2}_0 \left[(\nu g_{7/2})^4_0 + (\nu d_{3/2})^2_0 \right] (\nu s_{1/2}) \right]_{1/2^+} \\ & + (\nu g_{7/2})^2_0 (\nu d_{3/2})^2_0 \left[(\nu s_{1/2}) \right]_{1/2^+} \end{aligned}$$

while the proton configuration is $(\pi p_{1/2})^2_0 (\pi g_{9/2})^2_0 + (\pi g_{9/2})^4_0$.

An allowed beta decay from the ground state of ^{101}Mo to a state involving the $(\pi g_{9/2})^3$ and/or the $(\pi g_{9/2})^5$ configuration could proceed through the transition $(\nu g_{7/2}) \rightarrow (\pi g_{9/2})$ to form the configuration

$$\begin{aligned} & \left[(\pi p_{1/2})^2_0 (\pi g_{9/2})^3_{J'} + (\pi g_{9/2})^5_{J'} \right] \left[(\nu d_{5/2})^6_0 (\nu g_{7/2}) (\nu s_{1/2}) \right]_K \\ & + \left[(\nu d_{5/2})^{-2}_0 (\nu g_{7/2})^3_I (\nu s_{1/2}) \right]_{K'}. \end{aligned}$$

in ^{101}Tc . If $I = 7/2^+$, then $K' = K = 3^+$ or 4^+ . Under these assumptions, the configuration will contain levels with spin given by $\vec{J} = \vec{J}' + \vec{K}$. For the case where $J' = 9/2^+$,

the configuration will contain levels with spins varying from $1/2^+$ to $17/2^+$ and will include two $3/2^+$ states. The general configuration gives rise to ~ 105 states with J^π values from $1/2^+$ to $29/2^+$.

In the above example, the beta decay populates these states through the neutron part of their wavefunction while the subsequent decay to the ground-state $(\pi g_{9/2})^3_J$ multiplet occurs through the proton part of the wavefunction.

It should be noted that at excitation energies of between 1 and 3 MeV, proton states from the next major shell ($2d_{5/2}$, $2d_{3/2}$, $1g_{7/2}$ and $3s_{1/2}$) will also occur. Proton states with $\ell=0$ and $\ell=2$ strength have been observed in ^{93}Tc between 2.5 and 6.2 MeV with the $(^3\text{He}, d)$ reaction (Kozub and Youngblood 1971); these states are expected to occur at lower excitation energies as neutron pairs are added. Allowed beta decays will be able to populate these states through the various neutron components present in the ^{101}Mo ground state. One such decay, the $\nu s_{1/2} \rightarrow \pi s_{1/2}$ transition, should be particularly favoured since the initial and final states involved in the decay are identical except for the odd $s_{1/2}$ nucleon. Two possible candidates for such a decay are the 1962.18 and 2047.58 keV states in ^{101}Tc . These levels are populated with very low $\log f_0 t$ values (4.90 and 4.65, respectively). Other decays, such as a $\nu g_{7/2} \rightarrow \pi g_{7/2}$, will leave unpaired neutrons in several orbitals and will form a group of states similar to the configuration discussed above; the neutron configuration is the same, but the proton

configuration will be different.

It is clear that considerable mixing of the various configurations will take place and that many of states will have both proton and neutron character.

(g) A Comparison of $5/2^+$ and $7/2^+$ States in ^{97}Tc , ^{99}Tc and ^{101}Tc

Several high spin states in ^{101}Tc have been observed at 515.19 ($5/2^+$ or $7/2^+$), 533.55 ($5/2^+$ or $7/2^+$), 886.65 ($5/2^+$ or $7/2^+$) and 1598.94 ($5/2^+$ or $7/2^+$) keV. All of these states populate the ground-state multiplet with considerable strength indicating that they contain $(\pi g_{9/2})^3$ and/or $(\pi g_{9/2})^5$ components in their wavefunctions.

Experiments on $^{93}_{41}\text{Nb}_{52}$ (Stelson et al. 1971) and $^{115}_{49}\text{In}_{66}$ (Dietrich et al. 1970), which have a $g_{9/2}$ particle and a $g_{9/2}$ hole respectively, have found a multiplet of levels (J^π values - $5/2^+$, $7/2^+$, $9/2^+$, $11/2^+$, $13/2^+$) at low excitation energy that corresponds to the states expected for the coupling of a 2^+ core state to the $g_{9/2}$ particle (hole). However, the $B(M1)$ values for transitions from the $7/2^+$ and $11/2^+$ states to the ground state are found to be quite large, in disagreement with the simple model. Moreover, in the case of ^{115}In , it is necessary to invoke strong coupling between the hole and the vibrational motion in order to explain the large multiplet splittings observed. In the technetium isotopes, it is very likely that the states belonging to the core-excitation multiplet are admixed with the three- and five-

quasi-particle states involving $g_{9/2}$ protons.

The high spin levels in ^{99}Tc can be compared with those in ^{101}Tc . For example, the 761.6 keV ($5/2^+$) and 762.3 keV ($7/2^+$) levels in ^{99}Tc have similar properties to the 515.19 keV and 533.55 keV levels, respectively. This would tend to favour $5/2^+$ and $7/2^+$ assignments for the 515.19 and 533.55 keV levels respectively, in ^{101}Tc . Coulomb excitation experiments on ^{99}Tc (Bond et al. 1971) have found that the $B(E2)$ strength is concentrated in four (possibly five) states (the 140.51, 181.07 and 726.3 keV levels and the 762 keV doublet). None of these levels have a spin greater than $9/2$. This data is in conflict with the weak-coupling model which predicts that five states with spins of $5/2$ to $13/2$ should have been found. The observation of core-excitation strength in the 140.51 and 181.07 keV levels definitely indicates that these levels contain admixtures from both core-excitation and quasi-particle states. The core-excitation strength may be fragmented over many levels. This would explain the failure to observe the higher spin members of the core-excitation multiplet. This is also true for the levels at 761.6, 762.3 and 726.3 keV. These data would tend to support the view that the corresponding levels in ^{101}Tc also contain such admixtures. Since the 920.5 keV level in ^{99}Tc decays strongly to the $1/2^-$ state, it can probably be associated with the 606.41 keV state in ^{101}Tc .

The 785.07 and 855.41 keV levels in ^{97}Tc have J^π values of $5/2^+$ or $7/2^+$ and decay patterns which resemble those of the 515.19 and 533.55 keV levels respectively, in ^{101}Tc and the 761.6 and 762.3 keV levels, respectively in ^{99}Tc . Therefore, the most likely spin and parity values for this pair of levels are $5/2^+$ and $7/2^+$, respectively.

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APPENDIX A

A brief description of the use of JAGSPOT follows.

A number of regions (called windows) are chosen from the gamma-ray spectrum to be analyzed. In addition to the channels containing the peak, a number of adjacent channels are included, so that a meaningful background can be defined. For each window, the starting channel, the number of channels (≤ 100), the number of peaks (≤ 6), and the approximate channel position of each peak are supplied as input to the program. Initial values for the background parameters α and β , and for the Gaussian intensities γ , are calculated from the data. Starting values for δ and ϵ are supplied by the program. A non-linear least-squares fit to the data, in which all parameters are varied simultaneously is then performed. After each iteration, the values of the parameters are revised and a new iteration begins. The iteration procedure is terminated when the changes in the parameters are less than the associated errors, or when 20 iterations have been performed. When attempting to fit weak peaks or peaks with very bad statistics, it is desirable to reduce the number of fitting parameters as much as possible.

It was found at Chalk River that $1/\delta$ and $1/\epsilon$ were linear functions of energy. Therefore, an option was added to the program to allow it to define calibration curves of $1/\delta$ and $1/\epsilon$ versus energy. The values employed to generate these curves are taken from δ and ϵ parameters calculated from previous fits to a number of the stronger peaks in the spectrum. The user can

specify that the fitting of the weak peaks be done with δ and ϵ fixed at values determined from these calibration curves while all the other parameters are varied in the normal fashion. A considerable improvement in the ability of the program to accurately fit weak peaks is attained with this method. It should be emphasized that δ and ϵ are independent parameters. The user can specify that δ be varied during the least-squares fit in the normal fashion, but that ϵ be taken from a calibration curve of $1/\epsilon$, or vice versa.

A gamma-ray spectrum may contain some peaks whose energies are known. These may be from standard sources introduced to the "unknown" spectrum being studied, as in the "mixed source" technique explained in section 2.6, or from gamma rays in the "unknown" spectrum whose energies have been previously determined. Such peaks may be specified by entering the known energy and associated error on a data card. After peak fitting has been performed, the user may ask the program to supply a calibration curve (linear, quadratic, cubic, quartic) of channel number versus energy which may be used in the subsequent analysis of other parts of the same spectrum.

A number of comments concerning the operation of the program JAGSPOT may be useful.

(a) Although the peak positions calculated by the program are relatively insensitive to the background underlying the peaks, the areas calculated are not. Normally the fits will be best if fairly wide windows are selected. The larger number of channels leads to a better definition of the background.

Since the fitting function represents the background by a linear term, care must be taken to avoid choosing windows in which the linear approximation is invalid. For example, when a peak occurs at or near the Compton edge of another peak, the linear background can be a rather poor approximation. In such cases, a manual fitting procedure is often required.

(b) The "shape" standards required to obtain the $1/\delta$ and $1/\epsilon$ curves should be distributed fairly uniformly throughout the region of the spectrum in which the calibration curves are to be used. If this is done, the curves will be representative of the variation of these parameters over the entire region, and not just a particular part of that region. If the curves are properly defined, the fits are found to be only moderately sensitive to the values of δ and ϵ , since they are slowly varying quantities. Poor fits can result at low channel numbers (<1000), where the intercept term can dominate the slope term in the calibration curves, unless a sufficient number of "shape" standards are chosen from the low-channel region. The precautions described above are particularly applicable to those experiments in which the amplifier gain is adjusted to a low value, so that a large energy range can be covered. Under these circumstances, the combination of high resolution and low gain leads to a situation where each peak is defined by very few channels. For this reason, the intensity measurements should be checked by repeating the experiments at several gain settings.

(c) The distortions caused by high count rates usually result in unacceptable JAGSPOT fits. As explained in section 2.3, one of the effects of high rates is pulse pile-up. The high energy bulge on the peak, caused by this pile-up, tends to compensate for the tailing on the low energy side. The fitting program will try to approximate the peak by a Gaussian shape. Unfortunately, the bulge is not a smooth addition to the high side of the peak, and therefore, the fit deviates considerably from the true shape. This is particularly true for multiplets or unresolved peaks. Since the pile-up bulge from the lowest energy peak will be added to the next highest energy peak and so on, large errors in areas and peak positions can arise. In addition, the apparent shape of these peaks will be different from that of single peaks in the spectrum; therefore, any comparison between the two groups is clearly impossible. If correct fits are to be consistently obtained, moderate rates should be used or else some form of pile-up rejection or compensation should be provided so that pulse pile-up will not be a contributing factor to peak distortion.

(d) Since each window has only one value of δ and ϵ , the peaks selected for "shape" standards need not be isolated peaks in the spectrum. Multiplets in which one strong peak is resolved from the other components can also be used. The δ and ϵ values are then taken from this strong peak.

(e) Often there is some doubt as to whether a particular peak has a singlet, doublet or multiplet structure. Using the δ and ϵ parameters from the standard curves, the suspected multiplet

can be fit for each of the possible choices. The best fit will then be indicative of the number of peaks that can be placed in the region being investigated.

(f) Difficulties sometimes arise when attempting to fit weak peaks found on the low or high energy sides of very strong peaks. In the former case, part or all of the weak peak may be included in the tail of the strong peak. In the later case, the program may try to account for the presence of the weak peak (all or part of it) by fitting the larger peak with a Gaussian shape. In either case, the peak position and area calculated for the weak peak may be seriously in error. Often, it is wise to check some of these situations by fitting the weaker peaks by hand. Similarly, JAGSPOT will sometimes have trouble giving correct peak positions and areas when attempts are made to fit unresolved doublets. The program will often predict two equal intensity peaks.

APPENDIX B

The purpose of this appendix is to document some of the properties of $^{88}_{38}\text{Sr}_{50}$, $^{90}_{40}\text{Zr}_{50}$, some of the even-even and odd-even isotopes of zirconium and molybdenum and some of the odd-odd and odd-even isotopes of niobium. In most cases, no attempt is made to detail accurate amplitudes of the components contributing to the wavefunctions of the states which are discussed. Only the presence of these components is important to the discussion of Chapter 4. This survey has proved invaluable in providing an interpretation for the nuclei investigated in this thesis.

B.1 The Nucleus $^{88}_{38}\text{Sr}_{50}$

According to the simple shell model, the ground state of ^{88}Sr is described as having a closed neutron shell at $N = 50$ and a semiclosed proton shell at $Z = 38$. Since more energy would be required to promote a neutron over the major shell gap than to excite a proton, it is expected that the lowest-lying states in ^{88}Sr will be comprised largely of proton core-excited configurations. The neutron core-excited configurations should begin to appear only about 4 or 5 MeV in ^{88}Sr .

This simple picture has been fairly well substantiated

by the $^{87}\text{Sr}(d,p)^{88}\text{Sr}$ data of Cosman and Slater (1968) and by the $^{89}\text{Y}(d,^3\text{He})^{88}\text{Sr}$ experiments of Kavaloski et al. (1967). The (d,p) work showed that the states below 4 MeV were very weakly populated with the majority of the transition strength populating the configurations $[(\nu d_{5/2})(\nu g_{9/2})^{-1}]_J$, $[(\nu d_{3/2})(\nu g_{9/2})^{-1}]_J$ and $[(\nu s_{1/2})(\nu g_{9/2})^{-1}]_J$. In addition, there were many states ($\sim 40\%$) which had non-stripping character. The $(d,^3\text{He})$ data showed that several of the low-lying states (the 1.84 and 3.21 MeV) had the expected proton particle-hole character, with the $[(\pi p_{1/2})(\pi p_{3/2})^{-1}]_{2+}$ and $[(\pi p_{1/2})(\pi f_{5/2})^{-1}]_{2+}$ configurations respectively, dominating. Since these states were also weakly populated in the (d,p) reaction, this provided clear evidence that neutron configurations such as $[(\nu d_{5/2})(\nu g_{9/2})^{-1}]_{2+}$ are weakly mixed with the above proton configurations.

Thus, although ^{88}Sr appears to possess many of the properties of a good core nucleus, configuration mixing has to be invoked to give a satisfactory explanation of all the experimental details.

B.2 The Zirconium Isotopes

The zirconium isotopes are somewhat similar to the molybdenum isotopes which are the core nuclei for the technetium isotopes of concern in this thesis. The structure of the zirconium nuclei is of course expected to be much simpler because they have fewer particles outside the closed shell.

(a) The Proton Configurations of the Even-Even Zirconium Ground States

In the simple shell-model description of ${}^{90}_{40}\text{Zr}_{50}$, the subshells up to and including the $1g_{9/2}$ are completely filled by neutrons to form a major shell closure at $N = 50$. The $p_{3/2}$ proton subshell closure is expected to occur at $Z = 38$ (strontium) and the $p_{1/2}$ subshell closure to occur at $Z = 40$ (zirconium). However, because of the difference in pairing energy, the configurations $(\pi p_{1/2})^2$ and $(\pi g_{9/2})^2$ are nearly degenerate. Therefore, the two protons beyond the closed $2p_{3/2}$ shell ($Z = 38$) are expected to be distributed between the two orbitals, a situation which will produce configuration mixing in the ${}^{90}\text{Zr}$ ground state.

Experimental support for this view is provided by a number of proton pickup and stripping reactions on even zirconium targets. In all of these, the relative population of the proton states in the final nucleus is a measure of the ground-state configuration of the zirconium isotope being studied. The results are summarized in the table below. It is evident that as more neutrons are added, the $2p_{1/2}$ orbital is seen to move deeper in the potential than the $1g_{9/2}$. For ${}^{96}\text{Zr}$, the $2p_{1/2}$ level is nearly full and the $1g_{9/2}$ almost empty. The data presented below are consistent with a $65\%(\pi p_{1/2})_0^2 + 35\%(\pi g_{9/2})_0^2$ description of the 0^+ ground state of ${}^{90}\text{Zr}$ and a $35\%(\pi p_{1/2})_0^2 - 65\%(\pi g_{9/2})_0^2$ description for the 0^+ excited state at 1.75 MeV in ${}^{90}\text{Zr}$.

Structure of the 0^+ Ground State of Even-Even
Zirconium Nuclei

<u>Isotope</u>	<u>Configuration</u>		<u>Reaction</u>	<u>Reference</u>
	$(\pi p_{1/2})_0^2$	$(\pi g_{9/2})_0^2$		
^{90}Zr	70%	30%	$^{90}\text{Zr}(p, \alpha)^{87}\text{Y}$	Fulmer and Ball (1965)
	64%	36%	$^{90}\text{Zr}(d, ^3\text{He})^{89}\text{Y}$	Preedom et al. (1968)
	71%	29%	$^{89}\text{Y}(^3\text{He}, d)^{90}\text{Zr}$	Day et al. (1964)
	56%	44%	$^{90}\text{Zr}(^3\text{He}, d)^{91}\text{Nb}$	Cates et al. (1969)
^{92}Zr	55%	45%	$^{92}\text{Zr}(d, ^3\text{He})^{91}\text{Y}$	Preedom et al. (1968)
	44%	56%	$^{92}\text{Zr}(^3\text{He}, d)^{93}\text{Nb}$	Cates et al. (1969)
^{94}Zr	66%	34%	$^{94}\text{Zr}(d, ^3\text{He})^{93}\text{Y}$	Preedom et al. (1968)
	63%	37%	$^{94}\text{Zr}(^3\text{He}, d)^{95}\text{Nb}$	Cates et al. (1969)
^{96}Zr	85%	15%	$^{96}\text{Zr}(d, ^3\text{He})^{95}\text{Y}$	Preedom et al. (1968)
	91%	9%	$^{96}\text{Zr}(^3\text{He}, d)^{97}\text{Nb}$	Cates et al. (1969)

(b) The Closed $N = 50$ Shell in $^{90}_{40}\text{Zr}_{50}$

Evidence of a very good neutron shell closure at $N = 50$ is provided by the $^{90}\text{Zr}(p, d)^{89}\text{Zr}$ reaction data of Ball and Fulmer (1968). The four states in ^{89}Zr at $0.0(9/2^+)$, $0.588(1/2^-)$, $1.094(3/2^-)$ and $1.450(5/2^-)$ MeV exhaust a major portion of the single-particle strength expected for creating neutron holes in the filled levels between $N = 28$ and $N = 50$. The weakly excited state at 1.626 MeV has been identified as arising

from the pickup of one of a pair of $2d_{5/2}$ neutrons present in the ground state of ^{90}Zr in the form of a two-particle, two-hole (2p-2h) excitation. The probability of a neutron pair being present in the $2d_{5/2}$ level was about 3%.

(c) The Nucleus $^{88}_{40}\text{Zr}_{48}$

Except for the presence of two holes in the $1g_{9/2}$ neutron shell, the ground state of ^{88}Zr is expected to be somewhat similar to the ground state of ^{90}Zr . However, a much more complex structure was indicated by the large enhancement in the population of the ^{88}Zr ground state in the $^{90}\text{Zr}(p,t)^{88}\text{Zr}$ reaction (Ball et al. 1969). Since the $N = 50$ shell closure at ^{90}Zr is quite good, this enhancement could not be ascribed to 2p-2h impurity components of the ^{90}Zr ground state. Instead, it implies significant admixtures of the $(\nu p_{1/2})^2$ and $(\nu p_{3/2})^2$ configurations in the ^{88}Zr ground state. These correlated pairs lead to coherent addition of the transition strength in the two-neutron transfer reaction exciting the ground state. The lack of any appreciable $\ell = 0$ strength to excited states support this interpretation. The observed strengths to the 2^+ and 4^+ states indicates that these states are formed by mixtures of the form: $(p_{3/2})^{-2}_{2^+} + [(p_{1/2})^{-1}(p_{3/2})^{-1}]_{2^+}$ and $(f_{5/2})^{-2}_{4^+} + [(p_{3/2})^{-1}(f_{5/2})^{-1}]_{4^+}$.

(d) The Ground States of ^{92}Zr , ^{94}Zr and ^{96}Zr

The simplest shell-model picture for the even-even zirconium isotopes with $N > 50$ is to assume that the additional neutrons outside the $N = 50$ core fill the $2d_{5/2}$ orbital. However, the experimental transition strengths for the (p,t) reactions (Ball et al. 1971) are found to be significantly larger than the predictions of this simple model. The enhancements increase with target mass indicating the increased participation of other orbitals as neutrons are added. The experimental transition intensities for the $^{92,94,96}\text{Zr}$ ground-state reactions seem to be intermediate between the predictions of the pure $(d_{5/2})^n$ picture and the pairing-vibrational model calculations of Sorensen (1971). Although the $(d_{5/2})^n$ model is inadequate for a complete description, this neutron orbital forms a dominant component of the neutron portion of the ground-state wavefunctions of these zirconium isotopes.

(e) Single-Particle Strength in the Odd-Mass Zirconium Isotopes

Single-neutron transfer reactions on ^{90}Zr can be used to locate the single-particle states in ^{91}Zr . In addition, the $^{96}\text{Zr}(d,p)^{97}\text{Zr}$ reaction can be used to investigate the fullness of the $2d_{5/2}$ orbital in ^{96}Zr . The (d,p) and (d,t) reaction studies on the zirconium isotopes by Cohen and Chubinsky (1963) have determined the distribution of the $2d_{5/2}$, $3s_{1/2}$, $2d_{3/2}$ and $1g_{7/2}$ single-particle strengths in these nuclei. For example, the 0.0, 1.21, 2.07 and 2.19

MeV levels of ^{91}Zr contain major fractions of these neutron states, respectively. Beyond ^{90}Zr , the added neutrons are primarily filling the $2d_{5/2}$ shell-model orbital. This orbital is effectively empty for ^{90}Zr and essentially full for ^{96}Zr . The absence of low-lying $d_{5/2}$ states in ^{97}Zr confirms the latter result. The uncertainties in their data allow up to about 10% of the $3s_{1/2}$ and $2d_{3/2}$ states to be filled in ^{96}Zr .

(f) Configuration Mixing in ^{91}Zr

There is a considerable evidence for configuration mixing in the low-lying states of the odd-mass zirconium nuclei. For example, the (p,d) and (d,t) reactions on ^{91}Zr (Ball and Fulmer (1968) and Cohen and Chubinsky (1963), respectively) show that the ratio of the $(\pi p_{1/2})_0^2$ and $(\pi g_{9/2})_0^2$ configurations in the ground states of ^{90}Zr and ^{91}Zr is almost identical. The 2.184 MeV 2^+ state in ^{90}Zr is usually interpreted as a recoupling of the $g_{9/2}$ proton pair. These reactions are believed to populate this state through a small amount ($\sim 1.3\%$) of the $[(\pi g_{9/2})_2^2 (\nu d_{5/2})]_{5/2^+}$ configuration in the ^{91}Zr ground state. Thus, although the ^{91}Zr ground state is predominantly a $d_{5/2}$ neutron coupled to the ^{90}Zr ground state, it contains small admixtures of other components such as the one above and $[(\pi g_{9/2})_4^2 (\nu d_{5/2})]_{5/2^+}$. However, there was no evidence for any significant mixture of neutron-excited configurations such as $[(\pi g_{9/2})_2^2 (\nu s_{1/2})]_{5/2^+}$ in the ^{91}Zr ground

state. These conclusions were confirmed by the (d,p) studies of Bingham and Halbert (1970). The two $1/2^+$ states (1.21 and 2.58 MeV) in ^{91}Zr were interpreted as arising from a mixing of the $s_{1/2}$ single-particle state with the $[(^{90}\text{Zr})_{2+} + \nu d_{5/2}]_{1/2^+}$ configuration.

(g) The Nucleus $^{92}_{40}\text{Zr}_{52}$

The structure of ^{92}Zr is also more complex than the simple shell model would indicate. The excitation of the 1.204 ($1/2^+$), 2.04 ($3/2^+$) and 2.18 ($7/2^+$) MeV states in the (p,d) and (d,t) reactions on ^{92}Zr (Ball and Fulmer (1968) and Cohen and Chubinsky (1963), respectively) shows that the ^{92}Zr ground state cannot be a pure $(d_{5/2})^2$ configuration. These states can be populated in the pickup reaction if the ground state of ^{92}Zr contains an admixture of two-particle excitation of the neutron pair from the $2d_{5/2}$ to other single-particle levels ($s_{1/2}$, $d_{3/2}$, $g_{7/2}$). Most of the amplitude ($\sim 90\%$) is still associated with the $(d_{5/2})^2_0$ configuration.

(h) Mixing of Single-Particle and Core-Excited States

Core excitations also play a role in the structures of levels in the zirconium isotopes. Moore et al. (1967) found three resonances in the inelastic excitation curves for the $^{90}\text{Zr}(p,p')^{90}\text{Zr}$ reaction. These resonances had the same spacing as the $2d_{5/2}$, $3s_{1/2}$ and $2d_{3/2}$ analogue states in the elastic excitation curve, but shifted upwards by 1.75 MeV, the excita-

tion energy of the first excited state in ^{90}Zr . These states have the structure expected from the coupling of single-particle states with the excited ^{90}Zr core. Similar results were obtained by DuBard and Sheline (1969). These experiments show that there is considerable mixing between the single-particle and core-excited states. If the $1.46(5/2^+)$ and $2.58(1/2^+)$ MeV states in ^{91}Zr were pure core-coupled states, the 1.46 MeV level should have been populated by an $\ell = 0$ transfer and the 2.58 MeV level should not have been seen. However, the transition to the 1.46 MeV state had a mixed $\ell = 0 + \ell = 2$ character, and a transition to the $s_{1/2}$ state was observed. The data also showed that the $d_{5/2}$ neutron was coupled to other core-excited states (e.g. the 2^+ state at 2.18 MeV in ^{90}Zr).

(i) The Nucleus $^{98}_{40}\text{Zr}_{58}$

The low energy structure of ^{98}Zr is considerably more complicated than that of ^{96}Zr . Blair et al. (1969) have used the (t,p) reaction to show that the three lowest excited states are all 2^+ . The observed cross sections suggest that they arise from the $(vd_{3/2})^2_{2^+}$, $[(vd_{3/2})(vg_{7/2})]_{2^+}$ and $[(vs_{1/2})(vd_{3/2})]_{2^+}$ configurations which are expected to show considerable mixing. The observed 4^+ states were interpreted as arising from the $[(vs_{1/2})(vg_{7/2})]_{4^+}$ and $[(vd_{3/2})(vg_{7/2})]_{4^+}$ configurations. No $\ell = 0$ strength $> 2\%$ of the ground-state strength was found up to 3.2 MeV. This is in marked contrast to the (t,p)

results for $^{92,94,96}\text{Zr}$ where in each case several $\ell = 0$ transitions were observed. Their results suggest that the ground state of ^{98}Zr should be described with a wavefunction comprised of an in phase sum of the $(\nu s_{1/2})_{0+}^2$, $(\nu d_{3/2})_{0+}^2$ and $(\nu g_{7/2})_0^2$ configurations coupled to the ^{96}Zr ground state (almost pure $(d_{5/2})_{0+}^6$). Transitions to the excited 0^+ state would then have been inhibited by destructive interference.

B.3 The Molybdenum Isotopes

A comparison of the molybdenum and zirconium isotopes is useful in determining the influence of the two extra protons on the level structures found for the latter. In particular, it is expected that the presence of the two $1g_{9/2}$ protons will cause a downward displacement of the $1g_{7/2}$ neutron orbital due to the long-range part of the neutron-proton interaction.

(a) Proton and Neutron Configurations in $^{92}_{42}\text{Mo}_{50}$

The $^{92}_{42}\text{Mo}_{50}$ nucleus is characterized by four protons outside the $^{88}_{38}\text{Sr}_{50}$ core. Therefore, the low-energy spectrum of ^{92}Mo should be described in terms of configurations based on the four valence protons in the $2p_{1/2}$ and $1g_{9/2}$ shells. Calculations of the wavefunctions for the lowest energy positive parity states at $0.0(0^+)$, $1.51(2^+)$, $2.29(4^+)$ and $2.52(0^+)$ MeV have been carried out by Auerbach and Talmi (1965). These wavefunctions are found to be strong admixtures of the configurations shown in the table below.

Configurations in $^{92}_{42}\text{Mo}_{50}$

Energy (MeV)	J^π	Description
0.0	0^+	$(\pi p_{1/2})_0^2 (\pi g_{9/2})_0^2 + (\pi g_{9/2})_0^4$
1.51	2^+	$(\pi p_{1/2})_0^2 (\pi g_{9/2})_2^2 + (\pi g_{9/2})_2^4$
2.29	4^+	$(\pi p_{1/2})_0^2 (\pi g_{9/2})_4^2 + (\pi g_{9/2})_4^4$
2.52	0^+	$(\pi p_{1/2})_0^2 (\pi g_{9/2})_0^2 - (\pi g_{9/2})_0^4$

Ohnuma and Yntema (1969) have used the $^{92}\text{Mo}(d,t)^{91}\text{Mo}$ reaction to investigate the low energy levels of ^{91}Mo . The results are similar to those obtained for ^{89}Zr . The first four excited states are populated with neutron pickup from the $1g_{9/2}$, $2p_{1/2}$, $2p_{3/2}$ and $1f_{5/2}$ orbitals. In addition, there is evidence for a weak $d_{5/2}$ state. The small cross section for this state (the analogous state in ^{89}Zr is thought to be a 2p-2h excitation) is evidence for good closure of the $N = 50$ neutron shell in ^{92}Mo . The results are in good agreement with those obtained for ^{90}Zr (Ball and Falmer 1968).

(b) The Ground State of $^{94}_{42}\text{Mo}_{52}$

The levels in ^{92}Mo have been investigated with the $^{94}\text{Mo}(p,t)^{92}\text{Mo}$ reaction (Moalem et al. 1971). If the ground state of ^{94}Mo were composed only of seniority $\nu = 0$ proton and neutron configurations, it would be impossible to excite the low-lying 2^+ and 4^+ levels in ^{92}Mo by means of a direct

(p,t) reaction. Since these states are observed, the ground state of ^{94}Mo must have a configuration given by

$$\begin{aligned} & \left[\left[(\pi p_{1/2})_0^2 (\pi g_{9/2})_0^2 + (\pi g_{9/2})_0^4 \right] (v d_{5/2})_0^2 \right]_{0+} \\ & + \left[\left[(\pi p_{1/2})_0^2 (\pi g_{9/2})_2^2 + (\pi g_{9/2})_2^4 \right] (v d_{5/2})_2^2 \right]_{0+} \\ & + \left[\left[(\pi p_{1/2})_0^2 (\pi g_{9/2})_4^2 + (\pi g_{9/2})_4^4 \right] (v d_{5/2})_4^2 \right]_{0+}. \end{aligned}$$

The ^{92}Mo ground state is strongly excited via the first component which accounts for about 93% of the ^{94}Mo ground-state wavefunction. The 2^+ and 4^+ states are weakly excited through the second and third terms, respectively. In this model, the 2.52(0 $^+$) MeV state in ^{92}Mo , which is orthogonal to the ^{92}Mo ground state, is excited because the proton configuration of the ^{94}Mo ground state contains the second and third components. For the states above 3.85 MeV in ^{92}Mo , the most important contributions are expected to arise from neutron lp-lh

$$\left[(v g_{9/2})^{-1} (v d_{5/2}) \right]_J \text{ and } 2p-2h \left[(v g_{9/2})^{-2} (v d_{5/2})^2 \right]_J \text{ excitations.}$$

In fact, the strength of the transition to the 0 $^+$ state at 3.85 MeV indicates that it cannot be explained by pickup from a pure $(g_{9/2})^{10}$ configuration. Therefore, it must be a 2p-2h state.

(c) Configurations in ^{92}Mo , ^{94}Mo and ^{96}Mo

The level schemes of the even-even molybdenum isotopes have also been investigated with the $(\alpha, 2n)$ reaction (Lederer et al. 1971). The energies, deexcitation patterns and transition probabilities for the low-lying states populated in ^{92}Mo were

shown to be in good agreement with the model based on the proton configurations $(\pi g_{9/2})^4$, $(\pi p_{1/2})^2 (\pi g_{9/2})^2$ and $(\pi p_{1/2}) (\pi g_{9/2})^3$. For isotopes with $N > 50$, other excitations are possible. The $(\nu d_{5/2})_J^2$ excitations ($J = 2$ or 4) are expected to be the lowest of these. Such excitations can also couple to proton excitations to form the configurations: $[(\pi d_{5/2})^2 (\pi g_{9/2})^2]_J$ and $[(\pi d_{5/2})^2 (\pi p_{1/2}) (\pi g_{9/2})]_J$. Other states can arise from the excitation of neutrons to the higher orbitals: $s_{1/2}$, $g_{7/2}$, $d_{3/2}$, $h_{11/2}$. The experimental level schemes of ^{94}Mo and ^{96}Mo were found to be consistent with this interpretation. For example, the 5^- and 4^- states in ^{94}Mo are known to be predominantly $[(\pi p_{1/2}) (\pi g_{9/2})^3]_J$ excitations from the $^{93}\text{Nb}(^3\text{He}, d)^{94}\text{Mo}$ studies of Cates et al. (1969). These data also indicate that the lowest 6^+ , 8^+ and 10^+ states have the predominant configurations $[(\pi p_{1/2})^2 (\pi g_{9/2})^2]_J$. Because of the low $\log f_0 t$ value for electron-capture decay to the 6^+ state, it has been suggested (Lederer et al. 1971) that this state is a $[(\nu d_{5/2}) (\nu g_{7/2})]_J$ excitation. Thus, there would appear to be mixing of the two configurations. Further evidence of configuration mixing has been suggested for the 7^- and 9^- states.

After decreasing between $N = 50$ and $N = 52$, the energies of the lowest 2^+ and 4^+ states in the molybdenum isotopes remain relatively constant. This feature has been interpreted (Lederer et al. 1971) as due to a change from a nearly pure $(\pi g_{9/2})^2$ configuration to one which is predominantly $(\nu d_{5/2})^2$.

It is expected that considerable mixing will take place between the two configurations.

(d) Configurations in $^{98}_{42}\text{Mo}_{56}$

If the $\nu d_{5/2}$ orbital were totally filled in ^{98}Mo , only proton excitation would be possible at the lower excitation energies and therefore the level scheme would resemble that of ^{92}Mo . For example, the 0^+ first excited state in ^{98}Mo (Hübenthal et al. 1967) might have the proton structure $[(\pi p_{1/2})^{-2}(\pi g_{9/2})^4]_{0^+}$. The analogue resonance data of Courtney and Moore (1970) showed resonances corresponding to neutron single-particle states built on the 0^+ ground state. Consequently, the ground state must contain a large number of neutron configurations other than just $(\nu d_{5/2})^6_{0^+}$. Their (p,p') data has shown that the 0^+ excitation cross section for ^{98}Mo is very much larger than for ^{96}Zr . Therefore, the addition of two protons has altered considerably the configuration of the 0^+ first excited state; it has a much more complex structure than the corresponding states in zirconium. In addition, the experimental level structure of ^{98}Mo (Lederer et al. 1971, Heck et al. 1971) is much more complex at low excitation than that for ^{92}Mo . Therefore, it appears that the simple model considered above is inadequate and that there exists considerable configuration mixing in the ^{98}Mo ground state. This conclusion is in good agreement with the (d,p) and (d,t) results.

(e) The 0^+ and 2^+ Excited States in ^{94}Mo to ^{106}Mo

The splitting involved for the 0^+ , 2^+ , 4^+ seems too great for an explanation in terms of the vibrational model. Heck et al. (1971) have suggested that the nucleus ^{98}Mo may be regarded as a representative nucleus for the transitional region between vibrating spherical nuclei and rotating deformed nuclei. The (p,t) and (t,p) reactions (Taketani et al. (1972) and Casten et al. (1972), respectively) have been used to investigate the level schemes of the even-even molybdenum isotopes from mass 94 to mass 102. No excited 0^+ states were found in ^{94}Mo or ^{96}Mo but 0^+ states at ~ 700 keV were populated with large cross sections in all three of the heavier nuclei.

The heavier isotopes ^{104}Mo and ^{106}Mo have been produced using primary fission of ^{252}Cf (Cheifetz et al. 1970). They have identified excited states with J^π values of 2^+ , 4^+ and 6^+ . A comparison of the excitation energy of the first 2^+ state indicates that it remains approximately constant for mass 94 to 98, decreases slightly for mass 100 and 102 and then decreases sharply at mass 104. The energy decreases from 787 keV in ^{98}Mo to 172 keV in ^{106}Mo . At the same time, the $B(E2; 2 \rightarrow 0)$ values increase from ~ 10 to ~ 150 single-particle units in going from ^{98}Mo to ^{106}Mo . These observations suggest that the isotopes $^{98}, ^{100}, ^{102}\text{Mo}$ may form a series of transitional nuclei whose shapes change gradually from spherical to deformed. It should be noted that the E_{4^+}/E_{2^+} ratio of 3.15 is still somewhat smaller than the limiting value of 3.33 for a perfect

rotor. The large cross sections to the excited 0^+ states in the (t,p) and (p,t) reactions also support the transitional interpretation.

(f) Single-Particle Strength in the Odd-A Molybdenum Isotopes

These nuclei have some importance to the present study because of their similarity to the ruthenium isotopes. In particular, the properties of ^{97}Mo and ^{99}Mo should be relevant to the subsequent discussion of ^{101}Ru .

The nuclear structure of the odd-mass molybdenum isotopes has been investigated with the (d,p) and (d,t) reaction studies of Hjorth and Cohen (1964), the (d,p) studies of Moorhead and Moyer (1969), the (d,t) studies of both Ohnuma and Yntema (1969) and Diehl et al. (1970). These experiments have shown that in addition to the $2d_{5/2}$ orbital, the $3s_{1/2}$, $2d_{3/2}$, $1g_{7/2}$ and $1h_{11/2}$ single-particle orbitals are being filled. These data also indicate that the centers of gravity for the $2d_{3/2}$, $1g_{7/2}$ and $1h_{11/2}$ single-particle states are considerably lower than the corresponding centers in the zirconium isotopes. Table B.1 presents a comparison of the centers of gravity and emptiness factors for the zirconium and molybdenum isotopes. It is clear from an inspection of the table that the addition of neutrons in each isotopic series brings the positions of the $s_{1/2}$, $d_{3/2}$, $g_{7/2}$ and $h_{11/2}$ orbitals down with respect to the $d_{5/2}$ orbital and also decreases the "emptiness factor". The ground-state spin of both ^{99}Mo and ^{97}Zr is $1/2^+$ and

TABLE A.1

Comparison of Centers of Gravity and Emptiness Factors
for Zirconium and Molybdenum Nuclei

Nucleus	Center of gravity (MeV)					Emptiness Factor				
	$s_{1/2}$	$d_{3/2}$	$d_{5/2}$	$g_{7/2}$	$h_{11/2}$	$U_{1/2}^2$	$U_{3/2}^2$	$U_{5/2}^2$	$U_{7/2}^2$	$U_{11/2}^2$
$^{91}_{40}\text{Zr}_{51}$ a)	1.66	2.76	0.04	2.77	2.68	1.37	1.51	1.07	1.07	0.58
$^{93}_{40}\text{Zr}_{53}$ b)	1.15	2.40	0.0	2.4		1.13	1.01	0.54	0.92	
$^{95}_{40}\text{Zr}_{55}$ b)	1.43	2.20	0.0	(2.6)		1.09	1.00	0.30	0.40	
$^{97}_{40}\text{Zr}_{57}$ b)	0.0	1.37		1.64		0.98	0.83		0.85	
$^{93}_{42}\text{Mo}_{51}$ c)	1.63	1.97	0.081	1.58	2.30	1.06	0.98	0.89	0.45	0.33
$^{95}_{42}\text{Mo}_{53}$ c)	1.28	1.98	0.222	1.19	1.93	0.80	0.95	0.76	0.35	0.26
$^{97}_{42}\text{Mo}_{55}$ d)	0.88	0.93	0.0	0.79		0.81	0.91	0.44	0.54	
$^{99}_{42}\text{Mo}_{57}$ c)	0.22	0.64	0.132	0.41	0.68	0.80	0.70	0.23	0.64	0.14
$^{101}_{42}\text{Mo}_{59}$ d)	0.35	0.20	0.31	0.26		0.69	0.34	0.28	0.52	

a) Bingham and Halbert (1970)

b) Cohen and Chubinsky (1964)

c) Moorhead and Moyer (1969)

d) Hjorth and Cohen (1964)

therefore it might be expected that $^{98}_{42}\text{Mo}_{56}$ would show the same good closure of the $2d_{5/2}$ neutron subshell as $^{96}_{40}\text{Zr}_{56}$. However, the presence of a strongly excited $2d_{5/2}$ level at 97 keV in ^{99}Mo shows that there is an incomplete closure of the $2d_{5/2}$ subshell in ^{98}Mo . In fact, it appears from the data that the $\nu d_{5/2}$ orbital is only half full for ^{98}Mo . In contrast to the lower mass number molybdenum isotopes, ^{99}Mo shows considerable $1g_{7/2}$ single-particle strength. The isobaric analogue data of Moore et al. (1966) is in good agreement with the (d,p) results. The (d,t) data are consistent with the fact that ^{94}Mo , ^{96}Mo and ^{98}Mo are not pure $(\nu d_{5/2})^n$ configurations, but contain considerable components from other neutron orbitals. For example, the excitation of the $0.95(1/2^+)$ and $1.5(3/2^+)$ MeV states in ^{93}Mo indicate that the ^{94}Mo ground state wavefunction must include configurations of the form: $(\nu s_{1/2})^2_{0^+}$ and $(\nu d_{3/2})^2_{0^+}$.

(g) Configurations in ^{95}Mo and ^{97}Mo

Levels in ^{95}Mo and ^{97}Mo have been populated with (α, xn) reactions by Lederer et al. (1971) and Mesko et al. (1972). In addition to populating many of the levels found in previous reaction and decay studies, they have observed a number of high spin (up to $23/2^+$) states. Both groups have measured the angular distributions of many of the deexcitation gamma rays.

The $3/2^+$ first excited state and the lowest $9/2^+$ state in each isotope can be associated with the seniority-three states $(\nu d_{5/2})^3_{3/2^+}$ and $(\nu d_{5/2})^3_{9/2^+}$, respectively. This assignment is supported by the lifetime of the $3/2^+$ state (in ^{95}Mo), the weak excitation in the transfer reactions, and the agreement (^{95}Mo) with the calculations of Auerbach and Talmi (1965). In order to account for the low $\log f_0 t$ value for electron capture to the $9/2^+$ state in ^{95}Mo (Chilosì et al. 1969), it was necessary to postulate admixtures such as $[(\pi d_{5/2})^2_0 (\nu g_{7/2})]_{9/2^+}$ and $[(\pi g_{9/2})^2_2 (\nu g_{7/2})]_{9/2^+}$.

In contrast to the states described above, the lowest $1/2^+$ and $7/2^+$ states are strongly populated in the transfer reactions. The description of the $7/2^+$ state as the $[(\nu d_{5/2})^2_0 (\nu g_{7/2})]_{7/2^+}$ configuration is supported by the presence of a positive parity "band" which decays by a series of E2 transitions ($15/2^+ \rightarrow 11/2^+ \rightarrow 7/2^+$). These levels are assumed to have the predominant configurations $[(\nu d_{5/2})^2_{0^+, 2^+, 4^+} (\nu g_{7/2})]_J$.

The $(\alpha, 3n)$ reaction populates positive parity states up to spin $23/2$ in these nuclei. They have been explained in terms of the configurations $[(\pi g_{9/2})^2_6 (\nu d_{5/2})^3_{5/2}]_{17/2^+}$, $[(\pi g_{9/2})^2_8 (\nu d_{5/2})^3_{5/2^+}]_{19/2^+}$ and $[(\pi g_{9/2})^2_8 (\nu d_{5/2})^2_0 (\nu g_{7/2})]_{23/2^+}$ for ^{95}Mo , with similar configurations for ^{97}Mo .

The existence of the seniority-three $(\nu d_{5/2})^3$ states in

^{97}Mo is further evidence that the $\nu d_{5/2}$ shell is not closed at $N = 56$ in molybdenum. The strong similarity of the ^{95}Mo and ^{97}Mo level structures argues against the existence of a $(d_{5/2})^{-1}$ hole-state configuration in ^{97}Mo .

B.4 The Niobium Isotopes

Niobium ($Z = 41$) has two fewer protons than technetium ($Z = 43$). Therefore, knowledge of the level structures and configurations of the niobium isotopes will be a valuable asset in interpreting the level schemes of the technetium isotopes which are the subject of investigation in this thesis.

(a) The Nucleus $^{91}_{41}\text{Nb}_{50}$

In agreement with the simple shell model, ^{91}Nb has a $9/2^+$ ground state and a $1/2^-$ first excited state. These correspond to the configurations $[(\pi p_{1/2})^2_0 (\pi g_{9/2}) + (\pi g_{9/2})^3_{9/2}]_{9/2^+}$ and $[(\pi p_{1/2}) (\pi g_{9/2})^2_0]_{1/2^-}$, respectively. In addition to these, the $(d, ^3\text{He})$ and $(^3\text{He}, d)$ reactions (Ohnuma and Yntema (1968); Cates et al. (1969), Bassini and Picard (1969) and Knöpfle et al. (1970)) populate $3/2^-$ and $5/2^-$ levels which contain appreciable strength from the $p_{3/2}$ and $f_{5/2}$ hole states. Instead of the two hole states expected, there are actually two $3/2^-$ (1310 and 1615 keV) and one $5/2^-$ (1850 keV) state. Moreover, the (p, n) reaction (Kim et al. (1970, 1971), Matsuki et al. (1971) and Matthews et al. (1971)) excites many levels including a $5/2^-$ (1188 keV) state. Other $3/2^-$ and $5/2^-$ states can arise through

a weak coupling of the $p_{1/2}$ proton to the 2^+ excitation of the even-even ^{90}Zr core. Since the excitation energy of this 2^+ state is at 2.18 MeV, the energy of the core-excited states would be much higher than the observed states. Therefore, a better explanation is to attribute them to the $[\pi p_{1/2})(\pi g_{9/2})^2]_J$ configuration which would not be expected to be populated in the single-particle transfer reactions. The unperturbed energies of the two $3/2^-$ states are close enough to allow significant mixing, but the two $5/2^-$ states are too far apart. This would explain why both $3/2^-$ states are populated in the transfer reactions, but only one of the $5/2^-$ states.

Above 3 MeV, the deuteron spectrum in the $(^3\text{He}, d)$ reaction (Knöpfle et al. 1970) becomes very complex with many states strongly populated with $\ell = 0, \ell = 2$ or $\ell = 4$ transfer. Presumably, these states are due to the excitation of the proton to states in the next major shell $50 \leq Z \leq 82$ such as $2d_{5/2}$, $3s_{1/2}$, $2d_{3/2}$ and $1g_{7/2}$. They observe considerable fragmentation of these states. It is expected that this complex array of states will occur at lower excitation in the technetium isotopes and may explain some of the complexity found in the decay studies of the present thesis.

(b) The Nucleus $^{93}_{41}\text{Nb}_{52}$

The low energy level structure of ^{93}Nb is similar to that of ^{91}Nb . The $9/2^+$ ground state, the $1/2^-$ first excited

state and several $3/2^-$ states are populated in the ($^3\text{He}, d$) and ($d, ^3\text{He}$) experiments (Cates et al. (1970) and Ohnuma and Yntema (1968), respectively). Again there is evidence of considerable mixing between single-particle ($p_{3/2}$) and core-excited states ($p_{1/2}$ coupled to $^{92}\text{Zr } 2^+$).

Since ^{93}Nb is stable, it has been possible to investigate its level structure with ($n, n'\gamma$) and ($p, p'\gamma$) experiments (Rogers et al. 1970) and Coulomb excitation (Stelson et al. (1971); Kregar and Seaman (1972)). A multiplet of positive parity levels (742($7/2^+$), 809($5/2^+$), 952($13/2^+$), 980($11/2^+$) and 1080($9/2^+$)) were identified; these levels have the spin values expected for the coupling of a $g_{9/2}$ single particle to a 2^+ core state. The sum of all the $B(E2)$ s was observed to be approximately equal to the $B(E2)$ for excitation of the first 2^+ state in ^{92}Zr . In disagreement with the single-particle core weak-coupling model, some of the ground state transitions had large $B(M1)$ values. In addition, the $5/2^+$, $9/2^+$ and $13/2^+$ states have been populated in transfer reactions such as ($^3\text{He}, d$), (α, t) and (α, d) (Zisman and Harvey 1971). Thus, it would appear that there is considerable mixing between single-particle and core-excited states.

As in ^{91}Nb , a large number of positive parity states, which can be attributed to proton transfer to the next major shell, were observed. These states are found at much lower excitation (1.6-4.5 MeV) than in ^{91}Nb .

(c) The Odd-Odd Niobium Isotopes $^{92}_{41}\text{Nb}_{51}$ and $^{94}_{41}\text{Nb}_{53}$

The simple shell model predicts that the ^{92}Nb nucleus will have a single $d_{5/2}$ neutron and a single $g_{9/2}$ proton in addition to the filled orbitals corresponding to the ^{90}Zr core. The lowest configuration should give rise to six states with J^π values ranging from 2^+ to 7^+ . The (d,t) and (p,d) pickup reactions (Sheline et al. (1964) and Bhatin et al. (1971); Sweet et al. (1964)) on ^{93}Nb (assumed $(\pi g_{9/2})(\nu d_{5/2})^2_0$) populate this multiplet with considerable strength. The spin sequence is 7^+ , 2^+ , 3^+ , 5^+ , 4^+ and 6^+ . This multiplet is also populated in the ($^3\text{He},d$) reaction (Cates et al. 1969). They also observe the two states of the $(\pi p_{1/2})(\nu d_{5/2})$ multiplet.

The nucleus ^{94}Nb has three protons and three neutrons outside the ^{88}Sr core; therefore, the low energy spectrum of ^{94}Nb should contain configurations based on the $\pi g_{9/2}$, $\pi p_{1/2}$ and $\nu d_{5/2}$ orbitals. The simplest of these is expected to be $(\pi g_{9/2})(\nu d_{5/2})^3$. Since the three $2d_{5/2}$ neutrons may couple to partial configurations with $J = 5/2, 3/2$ or $9/2$, there are 20 possible states. It is expected that there should be considerable mixing of these states. Low-lying negative parity states should arise from configurations such as

$$(\pi p_{1/2})^{-1}(\pi g_{9/2})^2(\nu d_{5/2}).$$

Studies of the levels in ^{94}Nb with (d,p) and (n, γ) measurements (Jurney et al. 1968) have identified a number of multiplets

discussed above. The $[(\pi g_{9/2})(\nu d_{5/2})^3_{5/2+}]$ configuration is observed at low excitation; the spin sequence is 6^+ , 3^+ , 4^+ , 7^+ , 5^+ and 2^+ . At slightly higher excitation is the $[(\pi g_{9/2})(\nu d_{5/2})^3_{3/2+}]$ configuration; the spin sequence is 5^+ , 4^+ , 6^+ and 3^+ . It should be noted that the 2^+ state from the former configuration is ~ 200 keV higher than the other states of the configuration. This is probably due to the fact that there is no 2^+ state from the latter configuration to mix strongly with it. It would appear that there is considerable mixing of the 3^+ , 4^+ , 5^+ and 6^+ states. They have also identified states belonging to the $[(\pi p_{1/2})^{-1}(\pi g_{9/2})^2_0(\nu d_{5/2})^3_{5/2}]$, $[(\pi p_{1/2})^{-1}(\pi g_{9/2})^2_0(\nu d_{5/2})^3_{9/2}]$ and $[(\pi g_{9/2})(\nu d_{5/2})^2_0(\nu s_{1/2})]$ configurations.

It would appear that the low-energy spectrum of ^{94}Nb can be reasonably well interpreted in terms of the simple model proposed above. From the point of view of this thesis, the interest in the odd-odd niobium isotopes lies in the fact that the observed configurations represent a coupling of a $g_{9/2}$ proton to a $(d_{5/2})^3$ neutron configuration. In $^{101}_{44}\text{Ru}_{57}$ only the neutron configurations play an important role in the low-energy structure. The ground state of ^{101}Ru is $(\nu d_{5/2})^{-1}(\nu g_{7/2})^2_0$ and one may expect to find the $(\nu d_{5/2})^3(\nu g_{7/2})^4_0$ configuration at somewhat higher excitation. Thus many of the levels in ^{101}Ru at ~ 1 MeV may well be due to a $(d_{5/2})^3$ multiplet. As the discussion above has shown, con-

figuration mixing is very important in these nuclei. It is thus not surprising that we have not been able to uniquely associate any levels with one of the $(d_{5/2})^3$ multiplets.