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Spin Excitations in ^{48}Ca and ^{90}Zr with 319 MeV Protons

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TABLE OF CONTENTS

	page
List of Tables	vi
List of Figures	vii
Abstract	1
I. Introduction	3
A. M1 Transitions	6
B. Status of M1's	15
C. The Present Experiment	19
II. Experimental Procedure	24
A. The Proton Beam	24
B. The Spectrometer	26
C. Targets	31
D. Monitors	32
E. Standard Focal Plane	37
F. Focal Plane Polarimeter	41
G. Data Acquisition	44
H. Data Analysis Software	50
I. Small Angle Problems and Remedies	52
J. Summary of the Experiment	60
III. Data Analysis and Results	61
A. Analysis Software	61
B. Cross Section and Analyzing Power	63
C. Spin-Flip Probability	92
D. Background Correction	113
E. Spin-Flip Cross Sections	120

	page
IV. Theoretical Investigation	126
A. Formalism	126
B. Coulomb Excitation	146
C. Properties of Spin-Flip Probability	152
D. The 10.23 MeV state in ^{48}Ca	157
E. Study of ^{90}Zr	164
1. Low Lying States	165
2. The 9 MeV M1 Bump	172
3. Fine Structure in the M1 region	176
4. High Excitation Region	181
5. Multipoles of Continuum Spin Excitation	184
V. Conclusion	190
ACKNOWLEDGMENTS	194
REFERENCES	195
APPENDIX	203

LIST OF TABLES

	page
I. False asymmetry check of the Focal Plane Polarimeter with elastic scattering from ^{208}Pb	105
II. Background correction results of spin-flip probabilities for the M1 transitions in ^{48}Ca and ^{90}Zr	116
III. Effect of instrumental background on spin-flip probability	118
IV. The t-matrices of the Love-Franey interaction at 325 MeV	131
V. Strengths of the potentials in the Love-Franey interaction at 325 MeV	132
VI. Optical parameters used in the calculations for ^{48}Ca and ^{90}Zr	158
VII. Spectroscopic factors for the full f-p shell model calculations of McGrory and Wildenthal for the M1 transition in ^{48}Ca	160
VIII. The low lying natural parity states in ^{90}Zr	166
IX. The E1 states observed in Nuclear Resonance Fluorescence	168
X. Summary of renormalization of coulomb excitation calculations for the E1 states	171
XI. Comparison of fine structure in the M1 region with (e,e') results for ^{90}Zr	177
XII. Decomposition of E1 and M1 fractions in the M1 region of ^{90}Zr	180
XIII. Multipolar contribution to the spin excitation in the giant resonance continuum	185

LIST OF FIGURES

	page
1. Spin-flip transitions excitable in neighboring isobars in various reaction media	9
2. Spectra for : a) $^{90}\text{Zr}(p,n)^{90}\text{Nb}$, b) $^{90}\text{Zr}(e,e')^{90}\text{Zr}^*$, and c) $^{90}\text{Zr}(p,p')^{90}\text{Zr}^*$	16
3. S_{nn} and $\frac{d\sigma}{d\Omega} S_{nn}$ for $^{12}\text{C}(p,p')^{12}\text{C}$ at 400 MeV	21
4. Layout of the LAMPF switchyard and experimental areas	26
5. Schematic representation of Line C	29
6. Plan view of the High Resolution Spectrometer	30
7. Schematic view of Line C Polarimeter	34
8. Standard Focal Plane detector system	38
9. Typical anode plane delay line electronics setup	39
10. The Focal Plane Polarimeter layout	42
11. A two dimensional plot of θ_t vs ϕ_t	53
12. A two dimensional plot of θ_c and x_f	55
13. Low excitation background comparison for small angle scattering from ^{90}Zr	59
14. A spectrum of $^{12}\text{C}(p,p')^{12}\text{C}^*$ at 3.5°	70
15. Spectra of $^{48}\text{Ca}(p,p')^{48}\text{Ca}$ at 3.5° and 5°	72
16. Angular distributions for a) $\frac{d\sigma}{d\Omega}$ b) A_y for the 10.23 MeV state in the $^{48}\text{Ca}(p,p')^{48}\text{Ca}$ reaction at 319 MeV	73
17. a) $^{90}\text{Zr}(p,p')$ $E_p = 319$ MeV, $\theta_{lab} = 2.75^\circ$ spectrum b) the same spectrum at $\theta_{lab} = 4.25^\circ$	76
18. a-c) Cross section angular distributions for the states between 3 and 7 MeV.	77
19. Analyzing Power angular distributions for strongly excited states between 3 to 7 MeV.	80
20. a) Cross section and b) Analyzing power angular distributions for the total M1 peak for the $^{90}\text{Zr}(p,p')^{90}\text{Zr}$ reaction at 319 MeV.	82

	page
21. Expanded spectrum for the 9 MeV region in ^{90}Zr	85
22. a-c) Cross section angular distributions for the fine structure states in the M1 resonance region	88
23. Cross section and analyzing power for the region from 6 to 18 MeV in steps of 2 MeV	91
24. Schematic representation of Cone Test	98
25. Schematic representation of $\pi + \phi$ test	99
26. Inclusive analyzing Power of carbon analyzer	103
27. Spin-flip probability for the 10.23 MeV M1 transition in the $^{48}\text{Ca}(p,p')^{48}\text{Ca}$ reaction at 319 MeV	107
28. $^{90}\text{Zr}(p,p')$, $E_p = 319$ MeV, $\theta_{\text{lab}} = 3.5^\circ$ thick target spectrum	108
29. left and right scattered spectra accumulated by the FPP for normal and reverse beam polarizations for ^{90}Zr	109
30. Spin-flip probability for the 9 MeV M1 bump in ^{90}Zr .	110
31. Spin-flip probability and spin-flip cross section as a function of excitation energy for ^{90}Zr : a) $\theta_{\text{lab}} = 3.5^\circ$ and b) $\theta_{\text{lab}} = 5.0^\circ$	111
32. Spin-flip cross section angular distributions for 6 to 18 MeV in steps of 2 MeV	122
33. Typical DWIA calculations for simple particle-hole spin excitation cross sections	139
34. Spin-flip probability calculated with the central Love-Franey interaction for the M1 transition in ^{48}Ca .	154
35. Calculations of exchange effects and spin-orbit and tensor contributions to the spin-flip probability for the M1 transition in ^{48}Ca	155
36. Multipolar decomposition of the spin-flip cross section spectra for ^{90}Zr	187

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by

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ABSTRACT

Cross sections, analyzing powers, and spin-flip probabilities have been measured in the low momentum transfer region in the $^{90}\text{Zr}(p, p^-)^{90}\text{Zr}^*$ reaction at 319 MeV. A rich fine structure is observed for the first time in inelastic proton scattering in the previously proposed M1 giant resonance region. Angular distributions of most of these states are consistent with M1 excitation. The excitation energies of the fine structure states are in good agreement with similar structure seen in electron scattering; however, discrepancies in spin assignments remain. The measured cross section for the entire bump is about $37 \pm 10\%$ of the Distorted Wave Impulse Approximation (DWIA) prediction for the M1 strength in ^{90}Zr with simple wave functions. However, an analysis of the fine structure states reveals about 15% of the strength in the M1 region to be due to narrow E1 states; another 8% is attributed to M2 strength.

The spin-flip measurements for ^{90}Zr reveal a large spin-flip probability value for the M1 region; good agreement is obtained with DWIA calculations. However, a large cross section for spin excitations distributed uniformly over the excitation energy region from about

7 to 25 MeV is observed for the first time. The spin excitation strength in this giant resonance continuum is found to be about 0.80 mb/sr/MeV. Angular distributions for the spin-flip cross sections from 7 to 18 MeV in steps of 2 MeV have been analyzed with low multipole spin excitation calculations in the DWIA framework; the observed spin-flip strength in this region is found to be consistent with spin excitation involving angular momentum transfer of up to two.

Finally, cross section, analyzing power, and spin-flip probability data have also been obtained for the 10.23 MeV M1 transition in the $^{48}\text{Ca}(p,p')^{48}\text{Ca}^*$ reaction at 319 MeV. The quenching of M1 strength in ^{48}Ca relative to theoretical predictions is found to be consistent with previous work.

I. INTRODUCTION

Investigation of spin excitations in nuclei provides a unique insight into the magnetic properties of a nucleus. Following the development of various medium energy research facilities, a significant amount of theoretical as well as experimental endeavor over recent years has been devoted to the study of spin excitations. Although remarkable effort has been expended in this pursuit and progress has been achieved in its understanding, many of the features remain to be probed and understood. In particular, an in-depth understanding of the excitation of spin mode strength in the nuclear continuum of medium to heavy nuclei still remains elusive to researchers.

The primary focus of this dissertation will be the properties of spin excitation strength in ^{48}Ca and ^{90}Zr as revealed by the inelastic scattering of 319 MeV protons. Interpretation of the data obtained in the present work as well as comparison with available data with various other probes will also be discussed. In this introduction some of the fundamental aspects of spin excitations and their relevance to nuclear structure studies will be explained first. A brief historical account of this rapidly developing field and the motivation for the present measurements will follow.

The presence of an arbitrary distribution of charge and current, such as in a nucleus, gives rise to electric and magnetic moments which can be expressed as a series of multipole moments. Excitations of spin

operator and are designated as the magnetic multipole transitions. In sharp contrast to their counterparts, the electric multipole transitions, which are sensitive only to the effective charge distribution in the nucleus, the magnetic multipole transitions are sensitive to the effective "magnetization" within the nucleus. The nuclear magnetization arises primarily from: 1) the intrinsic magnetic moments of the neutrons and protons due to their internal degrees of freedom (spin part), 2) the magnetic moment produced by the orbital motion of the protons (current part), 3) exchange currents produced by the interaction of the nucleons determined by the internal dynamics of the collection of many nucleons; such effects include mesonic exchange currents (interaction part). Therefore, an insight into these magnetic degrees of freedom of a nucleus may be investigated through the study of magnetic multipole transitions.

The ground state of a nucleus may be thought of as the equilibrium state of a collection of nucleons. Perturbation of this state by external agents, such as scattering probes, may give rise to collective motion of the nucleons characterized by oscillations in the shape, density, or some other parameter associated with the equilibrium configuration. Such excitation mechanisms of the nucleus may be regarded as the excitation of multipole moments. Often, the excitation strength is concentrated in a small region of excitation energy forming a bump in the excitation spectrum; such a concentration is commonly referred to as a giant resonance. The well known Giant Dipole Resonance (GDR), observed systematically via the photonuclear reaction throughout the nuclear table, is a classic example of collective

nuclear motion. The excitation of the GDR is the consequence of a polarized vibration of the neutrons against the protons giving rise to an effective electric dipole moment. Occuring at an excitation energy of about $79 A^{-1/3}$ MeV and typically about 5 MeV in width, the GDR exhausts most of the ground state electric dipole strength. Since the discovery of the GDR in 1947, many other giant electric multipole resonances have been observed.

On the other hand, systematic observation of the analogous magnetic resonances, and the magnetic dipole (M1) resonance in particular, eluded experimenters, although a large concentration of M1 strength in medium to heavy nuclei had been predicted by theorists for many years.¹⁻⁷ Definitive discovery⁸⁻¹⁰ of the related Gamow-Teller resonance in medium and heavy nuclei came only in 1979, following which there has been a rapid growth in relevant experimental activities at many medium energy research facilities. Earlier attempts to find the analogue giant M1 resonances in medium and heavy nuclei had been unsuccessful.¹¹⁻¹⁴ The observation of concentrated M1 strength, in the Zr isotopes, was reported for the first time in 1981.¹⁵ This dissertation concerns the confirmation of the M1 resonance reported to have been observed in this region.

Why has it been so challenging to observe M1 strength in medium and heavy nuclei? The basic reason is that the spin dependent components of the nucleon-nucleus scattering amplitude are generally much weaker than the spin independent components. In addition, observation of M1 transitions requires detection of scattered particles at very low momentum transfer in order to selectively excite the M1

states over states of higher multipolarity. Inelastic proton scattering from a heavy nucleus at small momentum transfer has been plagued by enormous amounts of instrumental background related to the strong elastic scattering. This introduces significant ambiguity to the determination of M1 strength in (p,p') reactions. Furthermore, coulomb excitation of electric dipole (E1) states strongly excites such states at low momentum transfer, adding further difficulties in identifying M1 strengths.

A. M1 Transitions

Before going further into the historical background of M1 resonances, it is essential to briefly outline the selection rules and the isospin character of the M1 transition. As will be seen later, with some simplifying approximations, the M1 transition strength in the (p,p') reaction is very similar to the transition strength seen in electromagnetic interactions. Thus, it is necessary to consider the magnetic dipole operator μ , which may be expressed as:

$$\vec{\mu} = \sum_k g_s(k) \vec{s}(k) + g_l(k) \vec{l}(k) \quad (1)$$

where $\vec{s}$ and $\vec{l}$ are the spin and orbital angular momentum, respectively. The sum is over all nucleons. The spin and orbital g-factors for the k^{th} nucleon are denoted by $g_s(k)$ and $g_l(k)$, respectively, with their values in nuclear magnetons:

$$\begin{aligned}
g_s(p) &= 5.59; & g_d(p) &= 1; \\
g_s(n) &= -3.83; & g_d(n) &= 0.
\end{aligned}
\tag{2}$$

Here $p(n)$ denotes protons (neutrons). A more convenient way of expressing the magnetic dipole moment incorporates the isospin notation for protons and neutrons. If τ_z represents the Pauli matrix for the isospin projection, then Eq. (1) can be rewritten as:

$$\vec{\mu} = \sum_{k=1}^A \left\{ \left(g_s^0 - \frac{1}{2} \right) \vec{s}(k) - \left(g_s^1 \vec{s}(k) + \frac{1}{2} \vec{l}(k) \right) \tau_z(k) \right\} + \frac{1}{2} \vec{J}
\tag{3}$$

with

$$\begin{aligned}
g_s^0 &= \frac{1}{2}(g_s(p) + g_s(n)) = 0.88 \\
g_s^1 &= \frac{1}{2}(g_s(p) - g_s(n)) = 4.71.
\end{aligned}
\tag{4}$$

In Eq. (3), $\vec{J} = \sum (\vec{l} + \vec{s})$, is the total angular momentum. Both the initial and final states are eigenstates of $\vec{J}$, and so the operator $\vec{J}$ cannot cause transitions between states of different J . The reduced transition probability $B(M1)_{\uparrow}$ for an $M1$ transition between initial and final states with total angular momentum J_i and J_f , respectively, is proportional to the square of the reduced matrix element of the operator $\vec{\mu}$ between the two states. It is given by

$$B(M1)_{\uparrow} = \frac{1}{2J_i + 1} |\langle J_f T_f || \vec{\mu} || J_i T_i \rangle|^2. \quad (5)$$

The selection rules for the allowed quantum numbers for M1 transitions are:

$$\Delta J = \pm 1, 0 \quad 0 \rightarrow 0 \text{ forbidden}$$

no change in parity

$$\Delta T = \pm 1, 0 \quad 0 \rightarrow 0 \text{ forbidden} \quad (6)$$

In the independent particle shell model, M1 transitions are expected to occur between spin-orbit partners in a j-unsaturated shell; more specifically, the spin of one of the target nucleons must be flipped from $j = \ell + \frac{1}{2}$ to $j = \ell - \frac{1}{2}$ with $\Delta J = \pm 1$ without any change in parity. Equivalently, one then has orbital angular momentum transfer $\Delta L = 0, 2$ and spin angular momentum transfer $\Delta S = 1$ for the M1 transition. However, $\Delta L = 0$ dominates in the (p,p') reaction at low momentum transfer because of its lower multipolarity.

Isospin Properties

Examining Eq. (3), the $\sum_{\mathbf{k}} \vec{s}(\mathbf{k})$ term can be identified as a scalar in isospin space (isoscalar), and the term $\sum_{\mathbf{k}} \tau_z(\mathbf{k}) \vec{s}(\mathbf{k})$ as a vector in isospin space (isovector). However, it is easily seen that the M1 operator is predominantly an isovector operator. The transition probability due to a given operator is proportional to the square of the expectation value of the operator between the initial and final

Fig. 1: Levels excited in spin-flip transitions with various probes.

states; thus, the relative strength of the isoscalar and isovector parts of the M1 operator is determined by the square of the ratio of the isoscalar to isovector g-factors. From Eq. (4), this ratio $(g_S^0/g_S^1)^2$ is about 10^{-2} , which implies that for self conjugate ($N = Z$; $T_1 = T_0 = 0$) nuclei, isoscalar transitions ($T_f = T_0$) are strongly inhibited compared to their isovector ($T_f = T_0 \pm 1$) counterparts. This occurs because of the partial cancellation of the proton and neutron g-factors. This dominance of the isovector character of M1 the operator is known as Morpurgo's selection rule.¹⁶

An isovector transition in a nucleus with non-zero ground state isospin can reach excited states of different isospin; analogues of these isospin multiplets should be observable in the neighboring isobars. If we consider a target nucleus with neutron and proton numbers denoted by N and Z, respectively, with $N \neq Z$ and ground state isospin $T_0 = \frac{1}{2}(N - Z)$, isospin excitation of this nucleus will have analogues in the adjacent $N+1, Z-1$; $T_0 = T_0 + 1$ and $N-1, Z+1$; $T_0 = T_0 - 1$ isobars. Schematically represented in Fig. 1 are the possible states that can be excited by a spin-flip isospin-flip transition from the ground state in the target nucleus. The dashed lines in the figure connect the isobaric analogue states and the solid lines denote transitions that can be induced by various hadronic and electromagnetic interactions as indicated by arrows. In the target nucleus (N,Z), spin-flip, isospin-flip M1 transitions can be excited by the (p, p') , (e, e') , and (γ, γ') reactions. The isobaric analogues of the T_0 and $T_0 + 1$ states can be excited in the isobar (N-1, Z+1) by (p, n) , $(^3\text{He}, t)$, and (π^+, γ) charge exchange reactions whereas the $T_0 + 1$ analogue state

is excitable in the $(N+1, Z-1)$ isobar through the (n, p) , $(t, {}^3\text{He})$, and (π^-, γ) reactions. In addition, the $T_0 - 1$ and $T_0 + 1$ analogue states may decay to the ground state of the target nucleus via β^+ and β^- decay.

The M1 operator (see Eq. (3)) has both isoscalar and isovector components. The isoscalar part of the M1 transition operator cannot contribute to excitation of the $T_0 + 1$ state; therefore, a transition to the $T_0 + 1$ M1 state is purely isovector in nature. On the other hand, the isovector part can induce transitions to both the T_0 and $T_0 + 1$ states, implying that excitation of the T_0 , M1 state in the target nucleus is an admixture of both isoscalar and isovector transitions. Among the various possible interactions that can excite M1 states in the target nucleus, electromagnetic interactions such as (e, e') and (γ, γ') predominantly excite the isovector components of the M1 states because the reaction strength is proportional to $B(M1)$. The (p, p') reaction, however, may excite a significant isoscalar fraction in addition to the isovector components because the reaction strength depends on the proton-nucleus interaction.

The Gamow Teller Resonance

In the case of charge exchange processes leading to the $(N-1, Z+1)$ isobar as the product nucleus, all three components of the isospin multiplet are excitable. The operator responsible for such processes can be decomposed into its familiar spin independent (Fermi) and spin dependent (Gamow-Teller or GT) components, namely, $\sum_{\mathbf{k}} \tau_{\pm}(\mathbf{k})$ and $\sum_{\mathbf{k}} \sigma(\mathbf{k}) \tau_{\pm}(\mathbf{k})$, respectively. Both the Fermi and GT operators, being isovector operators, may produce isospin triplets of the excited

states, when the ground state isospin of the target nucleus T_0 is non-zero. However, the ground state isobaric analogue state of the product nucleus (IAS), characterized by $\Delta L = 0$ and $\Delta S = 0$, exhausts almost all of the Fermi strength. The GT strength is distributed among its isospin multiplets. The product states with isospin $T_0 - 1$, T_0 , and $T_0 + 1$ have transition probabilities proportional to nuclear structure factors and to the square of the relevant isospin Clebsch-Gordan coefficients: $(2T_0 - 1)/(2T_0 + 1)$, $1/(T_0 + 1)$, and $1/[2T_0 + 1)(T_0 + 1)]$, respectively. Therefore, for nuclei with large neutron excess, the $T_0 - 1$ state is the most strongly excited. The occurrence of a large concentration of $T_0 - 1$ strength (often referred to as the giant GT resonance although this term more properly refers to the entire multiplet) has been predicted by theorists for many years.¹⁷⁻¹⁸ The parent analogue corresponding to this $T_0 - 1$ GT state does not exist in the target nucleus since it is isospin forbidden. The strength contained in the T_0 and $T_0 + 1$ states, though much weaker, corresponds to the excited parent isobaric analogue states of the target nucleus. Since the selection rules for the GT states are same as those for M1's, the parent analogue of the T_0 and $T_0 + 1$ branches of the GT states should be observable as M1 transitions in the target nucleus.

Much of the theoretical development of the GT states, as briefly outlined above, was done as early as 1963. The existence of a giant GT resonance and its accessibility via the (p,n) reaction were postulated by Ikeda et al.¹⁷ However, only recently, following the development of high resolution neutron detection techniques at the Indiana University

Cyclotron Facility (IUCF), has the giant GT state been observed; it has been seen in a number of nuclei including medium and heavy nuclei.^{8-10,19} Observation of T_0 GT transitions have also been reported in two of these cases.²⁰

Comparison of Different Probes

Back angle electron scattering provides a very powerful means of observing M1 transitions because of the selectivity of such a technique to magnetic transitions.²¹ If we regard the electromagnetic interaction between the charge and current distribution in the nucleus and the electron in terms of the exchange of virtual photons, two components of the interaction may be defined in accordance with the polarization of the virtual photon: the longitudinal interaction, where the polarization of the photon is along the momentum transfer; and the transverse component, with the polarization perpendicular to the momentum transfer. Only the transverse interaction may contribute to spin-flip reactions. It can be shown,²¹ either from a simple plane wave formalism or from conservation of helicity of the electron, that at exactly 180° back angle scattering the longitudinal interaction vanishes. Furthermore, out of the two possible transverse interactions, transverse electric and transverse magnetic, the transverse magnetic components usually dominates for a given multipolarity.

Excitation of the M1 transition in (e,e') reactions occurs by means of electromagnetic interaction of projectile electrons with the nuclear magnetic dipole moment. Unlike the strong interaction in

hadron-nucleus scattering, the electromagnetic interaction, in principle, is well understood. Thus, although a closed form sum rule for the total M1 transition strength in the (e,e') reaction is not possible, a prediction for the strength dependent only on the nuclear structure factors can be made. In nucleon-nucleus scattering, such as (p,p') and (p,n) , due to the complex nature of the strong interaction involved in these reactions, similar predictions for the M1 strength cannot be made. However, for the GT strength in (p,n) reactions, because of the involvement of the same operator in the transition as in the case of β -decay, the β -decay matrix elements can be used⁹ to establish a closed form sum rule for the GT and Fermi strength observed in (p,n) . No such sum rules can be established for M1 transitions in (p,p') reactions.

In discussing spin excitations through different interaction mechanisms, it is imperative that justification for the comparison of the transition strengths be established. It has been pointed out earlier that the M1 operator is predominantly an isovector operator (Morpugo's selection rule); this implies that the electroexcitation of M1 transitions in (e,e') is dominated by isovector transitions. Similarly the GT operator, being an isospin lowering operator, is isovector in isospin space. Thus the M1 transition strength seen in (e,e') has the same isospin character as the strength observed in the analogue T_0 GT state in the (p,n) reaction. In the (p,p') reaction, the M1 transition can be mediated both by the isoscalar (τ_z) and isovector (τ_x, τ_y) parts of the proton-nucleus interaction. However, as described in Chapter IV, the isoscalar interaction is expected to be

small, and so the M1 strength in (p,p') under these conditions should be comparable to the respective strengths in (e,e') and (p,n) reactions. In addition to the spin and isospin terms of the M1 operator in Eq. (3), orbital current contributions due to the term $k_z(k)$ may also contribute to the transition probability and, thus, to the M1 transition strength in (e,e') reactions. On the other hand, both in (p,p') and (p,n) , orbital currents make negligible contributions to the M1 and GT transition strengths. However, orbital current contributions to the $B(M1)$ value, in the case of essentially pure neutron particle hole excitation, as in Ca and Zr, should be minimal. In this situation, (e,e') , (p,p') , and (p,n) reactions should yield directly comparable transition strengths.

B. Status of M1's

The experimental situation concerning M1 and GT states in medium and heavy ($A \geq 40$) nuclei remains highly controversial. A particular case of interest is Zr, for which there exists an abundance of experimental data as well as theoretical interest due to its relatively simple shell structure. Fig. 2a illustrates a spectrum (from Ref. 8) of the $Zr(p,n)$ Nb reaction with 120 MeV proton incident energy taken at 0.2° laboratory scattering angle. A large concentration of strength was observed at 5 MeV above the IAS and was identified as the $T_0 - 1$ branch of the GT resonance. The measured cross section angular distribution was described by a $L = 0$ theoretical prediction. The T_0 branch, the parent analogue of which should be observable as the M1

Fig. 2: Spectra of $^{90}\text{Zr}(p,n)^{90}\text{Nb}$ a), $^{90}\text{Zr}(p,p')^{90}\text{Zr}^*$ b), and $^{90}\text{Zr}(e,e')^{90}\text{Zr}$ c).

transition in (p,p') and (e,e') reactions, was also reported to have been observed ~ 9 MeV above the IAS. The energy resolution in the (p,n) data, about 1 MeV, is too poor to see fine structures in these peaks. These two states together exhaust about 60 % of the expected (p,n) sum rule strength for the GT resonance. The ratio of the measured cross section for the T_0 to the $T_0 - 1$ GT states was reported to be about 13; this ratio is expected to be 9 from the isospin coupling Clebsch-Gordan coefficients. More recently, the 200 MeV (p,n) reaction on ^{90}Zr by the same group has been reported with similar results.²² However, at 200 MeV the T_0 and $T_0 - 1$ branches were unresolved.

Prompted by the (p,n) work, inelastic proton scattering work by the Orsay group (Ref. 15) revealed a 1.5 MeV wide bump at 8.9 MeV of excitation energy in ^{90}Zr ; no fine structure was reported. Illustrated in Fig. 2b is the $^{90}\text{Zr}(p,p')^{90}\text{Zr}^*$ spectrum with $E_p = 201$ MeV taken at 4° laboratory scattering angle. The shape of the measured cross section angular distribution for the observed bump was adequately described by a $\Delta L = 0$ calculation with the distorted wave impulse approximation (DWIA). The experimental cross section of the reported bump accounts for 30 % of the M1 cross section predicted by the DWIA, and 70% of the strength observed in the analogue state in the (p,n) reaction.

The results reported previously with 165° electron scattering by the Darmstadt group²³ showed a significant discrepancy from the (p,n) results. Fig. 2c shows a spectrum of $^{90}\text{Zr}(e,e')^{90}\text{Zr}$ obtained with 35 MeV electrons at 165° scattering angle. In contradiction also to the (p,p') results, electron scattering reveals only a few isolated

narrow peaks ($E_x = 8.233, 9.000, \text{ and } 9.371 \text{ MeV}$) as definitely M1 transitions; several other possible M1 assignments in the excitation energy range of $7.7 - 9.5 \text{ MeV}$ were also reported. Such fine structure should have been observable in the (p, p') work of Ref. 15, if it was present. The upper limit for the total strength contained in these states was $B(M1)_{\uparrow} = 1.8 \pm 0.4 \mu_N^2$ which amounts to 11 % of the predicted M1 strength with bare g -factors. To complicate the matter further, preliminary results obtained through the resonance fluorescence work of Berg et al.²⁴ indicate the presence of a large number of narrow and resolved E1 states in the excitation energy region of 7 to 10 MeV of ^{90}Zr ; no definite M1 states were observed.

Despite the fact that considerable discrepancies exist concerning M1 transitions observed with different probes, there is a consensus with regard to one feature: the observed strengths are significantly "quenched" from the expected theoretical values. Various mechanisms to account for this quenching have been proposed, the most promising and widely pursued being the renormalization of nuclear magnetization due to $\Delta(1232)$ -particle and nucleon hole excitations.²⁵ In addition, effects of ground state correlations and meson exchange currents have also been considered^{26,27} as sources of quenching of M1 strength.

Even with the inclusion of various proposed quenching mechanisms, the measured M1 strength still remains significantly lower than theory. A question immediately arises: where is the missing M1 strength? Is the M1 strength in medium-heavy nuclei really "quenched" due to realistic nuclear structure effects or is it simply not being detected? Suggestions²⁸⁻³⁰ have been made that the missing strength is possibly

spread out into the higher excitation energy region; such a spreading could easily avert detection due to the strongly excited nuclear continuum in this region. These are some of the issues we will attempt to investigate in the present work.

C. The Present Experiment

Identification of M1 transitions in previous (p,p') work has been done by the measurement of cross section angular distributions. As pointed out earlier, the presence of overwhelming amounts of instrumental background at small scattering angles in such measurements introduces large systematic uncertainties in the measured cross section. Thus, the criterion for M1 assignment based on cross section measurements alone is unreliable. Furthermore, even if the uncertainty due to a choice of background can be eliminated, M1 strength identification in inelastic proton scattering from cross section angular distributions alone is not without ambiguity. Recent calculations of coulomb excitation of the GDR by Izumoto et al.³¹ show that, in heavier nuclei, coulomb excitation of E1 states yields cross sections peaking at forward angles very similar in shape to the predicted M1 cross section angular distribution. Therefore, the presence of significant amounts of E1 strength in the expected M1 resonance region would make it virtually impossible to isolate M1 states from neighboring E1's. As noted above, Berg et al. (Ref. 24) have suggested that there is indeed a large concentration of narrow E1 states in the proposed M1 resonance region of ^{90}Zr .

The foregoing discussion calls for the measurement of some other observable that enables the distinction between M1's and E1's to be made unambiguously. Since M1 transitions occur through spin transfer ($\Delta S = 1$), the measurement of spin-flip probability proves^{32,33} to be an excellent tool to make this distinction. Experimentally, one measures the polarization of the outgoing particles after scattering of an incident polarized beam. Existing experimental data³³ on the spin-flip probability for the well known $\Delta S = 1$ and $\Delta S = 0$ states in the $^{12}\text{C}(\vec{p}, \vec{p}')^{12}\text{C}$ reaction indicates that a high value for spin-flip probability at low momentum transfer is a unique signature of a $\Delta S = 1$ transition. This fact is qualitatively illustrated in the spectra of cross section and spin-flip cross section in Fig. 3 for the $^{12}\text{C}(\vec{p}, \vec{p}')^{12}\text{C}$ reaction; the natural parity $\Delta S = 0$ states in the lower excitation region (0^+ and 3^-) are strongly excited in the cross section (top) spectrum whereas they are nonexistent in the spin-flip cross section (bottom) spectrum. In contrast, the well known M1 states at 12.71 and 15.11 MeV are excited strongly in the spin-flip cross section spectrum. The measurement of the spin-flip cross section provides the ability to detect weak spin excitation strength even in the overwhelming presence of natural parity nuclear background such as in the giant resonance region of heavier nuclei.

The $^{90}\text{Zr}(p,n)^{90}\text{Nb}$ results of Ref. 8 indicate the presence of a wide $\Delta L = 1$ and $\Delta S = 1$ resonance at ~ 13 MeV above the IAS, providing evidence of giant spin-flip E1 or M2 type resonances. Measurement of spin-flip cross sections in this region of excitation, together with

Fig. 3: Spectra of spin-flip probability and spin-flip cross section for the $^{12}\text{C}(p, p')$ reaction at 397 MeV.

cross section angular distributions, may extend our knowledge of such spin-flip strength in the high excitation region.

The measurement of the analyzing power may provide information about the isospin character of the M1 states. Measured³⁴ analyzing powers for the 12.71 MeV $1^+, T = 0$ state in the $^{12}\text{C}(p, p')^{12}\text{C}$ reaction yield large negative values (~ -0.30); on the other hand, the 15.11 MeV $1^+, T = 1$ state in ^{12}C has small positive analyzing powers. Such an isospin dependence of the analyzing power has been observed systematically in ^{12}C throughout the energy range of $E_p = 400 - 800$ MeV;³⁴ pure isoscalar and isovector calculations of analyzing power with the DWIA also support this empirical observation. Therefore, measurements of the analyzing power may help establish the isospin transfer for the nuclei studied here.

The research described in this dissertation is primarily intended for the study of M1 transitions in ^{90}Zr through inelastic proton scattering at 319 MeV. The focus of the present work will be 1) to confirm the proposed M1 resonance in ^{90}Zr , 2) to understand the existing discrepancies among the various experimental data as outlined earlier, and 3) to search for the missing M1 strength in the high excitation region in ^{90}Zr . However, as a source of reference and calibration, the M1 transition in ^{48}Ca was also investigated. The ^{48}Ca nucleus is an excellent target because of the concentration of the observed M1 strength in a very narrow and isolated region of excitation which has been observed without major disagreement in both (p, p') and (e, e') reactions; its analogue has also been observed in the (p, n) reaction. In addition to the overwhelming availability of experimental

data, there is an abundance of theoretical understanding of the ^{48}Ca M1 transition due to its simple shell structure.

We have thus measured the differential cross section, the analyzing power, and the spin-flip probability in the reactions $^{90}\text{Zr}(\vec{p}, \vec{p}')^{90}\text{Zr}^*$ and $^{48}\text{Ca}(\vec{p}, \vec{p}')^{48}\text{Ca}^*$ at 319 MeV incident proton energy. As we shall show in the following chapters, a rich fine structure somewhat similar to the (e, e') spectrum is observed for the first time in inelastic proton scattering from ^{90}Zr in the proposed M1 giant resonance region. However, the angular distributions of most of these states are found to be consistent with M1 excitation. A considerable discrepancy with electron scattering thus persists. The spin-flip measurements reveal for the first time a large cross section for spin excitations distributed surprisingly uniformly over the excitation energy region from about 7 to 25 MeV; the M1 region does not stand out. The measured spin-flip cross section is found to be consistent with spin excitations involving angular momentum transfer of at least up to two.

A detailed description of the experimental procedure is presented in Chapter II. Chapter III deals with the analysis of the measured experimental observables and presentations of the final data. Theoretical interpretation of the data and comparison to previous experimental data with different probes are discussed in Chapter IV. Finally, concluding remarks are presented in Chapter V

II. EXPERIMENTAL PROCEDURE

The experiment was performed at the Clinton P. Anderson Meson Physics Facility (LAMPF) using the High Resolution Proton Spectrometer (HRS). The superiority of the HRS optics and data acquisition techniques made it an excellent choice for conducting small angle inelastic proton scattering experiments. In addition, the wide band capability of the newly developed Focal Plane Polarimeter (FPP) in measuring spin-transfer observables provided a unique tool for the study of spin excitation over a broad range of excitation energy.

The experimental work was conducted in two phases: the first phase was in June, 1982; and the second one was in October, 1982. In this chapter a brief description of the general features of the LAMPF linear accelerator (LINAC) and experimental facilities at HRS will be presented. Specific arrangements relevant to the present experiment along with problems encountered in the course of the experiment and solutions thereof will be discussed in more detail.

A. The Proton Beam

Beams of 319 MeV polarized protons provided by LAMPF were used in this experiment. The facility at LAMPF consists of a proton LINAC³⁵ capable of simultaneously accelerating both H^+ and H^- beams up to 800 MeV of proton energy. Three separate ion sources are used in

feeding the main LINAC with H^+ and either unpolarized H^- or polarized H^- (P^-) ions. In the present experiment, the Lamb-shift polarized ion source was used to obtain the desired spin orientation of the protons. The principle is based on the properties of the $2S_{1/2}$ metastable state of the hydrogen atom. The hyperfine interaction in the presence of an external magnetic field causes splitting of the doublets with spin projection of the proton as $+1/2$ and $-1/2$; thus, by choosing the correct magnetic field, the desired polarization can be obtained. In order to achieve this, 500 eV protons generated by a duoplasmatron are fed into a cell containing cesium vapor where about 43% of the protons are converted into hydrogen atoms in the $2S_{1/2}$ metastable state by electron pickup from the cesium. The mixture of ground state and metastable state hydrogen atoms is then passed through a spin filter which consists of a 25 cm RF cavity operating at 1608 MHz, a switchable axial magnetic field of 540 gauss, and a transverse electrostatic field of 20 V/cm. The electromagnetic coupling produced in the spin filter section causes all the hyperfine levels of the hydrogen atoms to decay to the ground state except for the component with the proton spin aligned along the direction of the magnetic field. The hydrogen atoms then pass through a cell of argon gas with a 10 gauss axial magnetic field; in the argon cell the metastable hydrogen atoms with their nuclear polarization along the magnetic field are selectively converted into polarized hydrogen ions (P^-) by electron pickup. The P^- beam is typically 75 - 85% polarized due to the presence of unpolarized ground state hydrogen atoms as the beam enters the argon cell. The polarization of the P^- beam can be chosen to be parallel or

Fig. 4: Layout of the LAMPF switchyard and experimental areas.

antiparallel to the beam momentum by switching the direction of the magnetic fields in the spin filter and the argon cell coherently. However, a spin precessor is used to obtain any desired direction of the beam polarization at the experimental sites.

The polarized beam is then accelerated to 750 keV and injected into an Alvarez drift tube section operating at 201 MHz which accelerates the beam to 100 MeV. The final stage of the acceleration is the 805 MHz side coupled cavity section consisting of 44 cavity modules. By selecting the desired number of modules, a final beam energy, stepwise variable, in the range of 197 to 800 MeV may be obtained. For the present experiment, 19 modules were used to provide the additional boost of 219 MeV to bring the hydrogen ions to their final energy of 318 MeV. The macro timing structure of the P^- beam consists of 40 pulses per second with each pulse 833 μ s long; this amounts to a duty factor of 3.3 %.

The final beam then enters the switchyard where a dipole magnet separates the H^+ and H^- components of the beam. Fig. 4 shows a schematic layout of the switchyard area. Subsequently, the H^+ beam proceeds to Area A, where it is utilized in producing various secondary meson beams for the experimental areas therein. The H^- beam enters Line X where a fraction of the beam is stripped of its electrons by means of a stripper section consisting of foil or wire strippers. Proper care is taken in choosing the stripper so as to control the spatial and momentum spread and the overall intensity of the stripped beam. The protons are then magnetically separated from the unstripped beam. The protons are then magnetically separated from the unstripped H^- ions by bending magnets LCBM01-03 on their entrance to Line C and

are fed to the HRS via Line C. The H^- beam proceeds to other experimental facilities.

Line C beam line optics has been described extensively elsewhere.³⁶ Briefly, Line C consists of a dispersion section, a twister section, and a matching section. A schematic of the Line C beam area is given in Fig. 5. First, the proton beam is horizontally dispersed by 18 cm/% and bent 57° by two bending magnets LCBM04,05. The dispersed beam is then twisted by 90° , without mixing of the phase space, by an array of quadrupole magnets, LCQM04-08. Following this, in the matching section, an additional set of quadrupole magnets, LCQM09-11, performs the task of dispersion matching of the beam on target with that of the spectrometer. Since the energy spread of the proton beam on target is typically two orders of magnitude higher than the desired energy resolution of the spectrometer, dispersion matching is critical in order to achieve the spectrometer resolution.

B. The Spectrometer

The HRS analyzes the momentum loss of the scattered protons after they undergo scattering in a two meter diameter scattering chamber. The HRS is a quadrupole-dipole-dipole (QDD) magnet configuration and operates in a vertical dispersion, horizontal scattering, and vertical analysis (VHV) mode.³⁷ To first order, the VHV mode of operation eliminates correlations between scattering angle and momentum loss of the scattered protons at the target, enabling independent measurement of the two. Fig. 6 shows a schematic of the HRS. The optics of the

Fig. 5: Schematic representation of Line C.

Fig. 6: SECTION HRS FACILITY

0 5 10
GRAPHIC SCALE
IN FEET

HRS magnets is designed to achieve a point to point focussing in the vertical ($\hat{x}$) direction and parallel to point focussing in the horizontal ($\hat{y}$) direction. The quadrupole magnet HSQM01, situated at the acceptance end of the HRS, focusses the scattered particles in the y-direction and defocusses them in the x-direction. The focussing in the y-direction allows for a larger angular acceptance with smaller gap in the bending magnets that follow the quadrupole. The two 3.5 meter radius bending magnets HSBM01,02 then bend the scattered particles by 150° in the vertical plane and deliver them to the focal plane detection assembly. Such a focussing mechanism implies that all scattered particles with the same momentum loss at the target get focussed at the same position in the x-direction whereas particles with the same horizontal scattering angle at the target end up at the same y-position at the focal plane. However, the above is valid only for heavier nuclei; for lighter nuclei, where the kinematic change of energy with scattering angle ($dE/d\theta$) is not negligible, the focal plane shows a correlation between scattering angle and momentum loss which is easily corrected in the analysis software.

C. Targets

The standard HRS scattering chamber and target holder assembly were used in the experiment. The ^{48}Ca target used in the experiment was an isotopically enriched ^{48}Ca foil with an areal density of 100 mg/cm^2 and a total area of 18.5 cm^2 . The isotopic content of the target was 94.47% ^{48}Ca , 5.32% ^{40}Ca , and the rest $^{42},^{43},^{44},^{46}\text{Ca}$. It was

stored in a can pressurized with argon gas; handling of the target was also performed in an atmosphere of argon. However, data taken with this target showed a significant amount of oxidation in the target.

Two natural zirconium foil targets of areal density 50 mg/cm² and 250 mg/cm² were used. Blank target frames similar in dimensions to the ones used for the above targets were also used to investigate possible background arising from beam halo effects. In addition to the two primary targets, the following targets were used for calibration and normalization purposes: a ¹²C (graphite) target of 44.6 mg/cm² areal density; a ²⁰⁸Pb target of 50 mg/cm² areal density; and two CH₂ targets of 6.4 and 200 mg/cm² areal densities. Finally, a 10 mil thick phosphor material foil was used for beam alignment purposes.

D. Monitors

To determine $\frac{d\sigma}{d\Omega}$, A_n , and S_{nn} , the polarization and intensity of the incident beam need be monitored. Several independent methods of measuring the polarization of the P⁻ beam are available at LAMPF. It is imperative that the incident beam polarization be measured precisely in order to determine A_n and D_{nn} accurately. Three independent methods, namely, the quench ratio method, the Line B polarimeter, and the Line C polarimeter were employed in measuring the beam polarization.

The polarization, P, of a projectile in a beam line polarimeter is measured utilizing the vector analyzing power (A_n) of a target nucleus. Scattering of a polarized projectile off a nucleus with non

zero A_n results in an asymmetric distribution of the scattered particles in the scattering plane; if I_0 is the incident flux of the projectile and I is the scattered flux then the scattering is governed by the following relation:

$$I = I_0 (1 + P \cdot \hat{n} A_n) \quad (7)$$

where $\hat{n}$ is the normal to the reaction plane. The flux of particles scattered to the left of the incident direction is given by $I_L = I_0(1 + P A_n)$, and that to the right is given by $I_R = I_0(1 - P A_n)$; hence an asymmetry, ϵ , may be defined as:

$$\epsilon = \frac{I_L - I_R}{I_L + I_R} = P A \quad (8)$$

Therefore, a knowledge of the A_n of the analyzing nucleus is all that is needed to determine P .

However, the presence of instrumental background introduces false asymmetry into the measurement of ϵ . From Eq. (7), it is important to note that a projectile polarized parallel to $\hat{n}$ (normal) and scattered to the right is equivalent to a projectile polarized antiparallel to $\hat{n}$ (reverse) and scattered to the left and vice versa. It has been shown in Ref. 38 that it is possible to eliminate false asymmetry by taking the geometric mean of right-normal and left-reverse and left-normal and right-reverse scattering yields.

Fig. 7: Schematic view of Line C polarimeter.

A schematic of the Line C polarimeter is shown in Fig. 7. A CH₂ target is used as the analyzer and two pairs of left-right detectors are placed at 17° and 67° to detect the scattered and recoil particles in coincidence. The choice of 17° lab scattering angle is based on the fact that the ¹H(p,p)¹H reaction has a broad maximum in A_n at this angle. A value of 0.380 ± 02 obtained from independent measurement of A_n for the ¹H(p,p)¹H reaction at 17° and corrected for the presence of carbon in the CH₂ target was used for the determination of P_b.

The P⁻ beam consists of a small fraction of unpolarized hydrogen ions (see section II.B). The quench ratio method³⁹ utilizes the relative intensities of the polarized to the unpolarized ions in the beam to determine the beam polarization. At regular intervals, the magnetic field in the spin filter is switched to a value where the polarized metastable hydrogen atoms are quenched to the ground state rendering an unpolarized beam to the argon cell. Thus the quenched beam intensity is the same as the intensity of the unpolarized fraction of the polarized beam. If I_Q and I_P denote the intensities of the quenched and polarized beams, respectively, then the intensity of the polarized ions is given by I_P - I_Q. Since P_b is equivalent to the fraction of the beam polarized, we have

$$P_b = \frac{I_P - I_Q}{I_P} = 1 - \frac{1}{Q} \quad (9)$$

where Q is referred to as the quench ratio.

Calculations⁴⁰ indicate that there is no significant depolarization of the beam during the transportation of the beam from the ion source to the experimental site. If we further assume that the spin vector of the hydrogen ions has precessed exactly 90° in a vertical plane between the source and Line C polarimeter, then the $\hat{n}$ -type beam polarization measured by the quench ratio method and with the Line C polarimeter should be the same.

The polarized ion source was operated in a normal-reverse quench cycle with a duration of 2 minutes each for the normal and reverse states followed by a 30 second quench state with a relaxation time of about 10 seconds between states. Signals indicating these states of the P^- beam were fed to the data acquisition electronics hardware which were used to gate the yields in various counters corresponding to these states. The beam polarization was calculated from equations 2.2 and 2.3 using the normal, reverse, and quench yields. In order to monitor the stability of the beam polarization and the consistency between the Line C polarimeter value with that of the quench method, the beam polarizations obtained with the two different methods were recorded at regular intervals. The Line C polarimeter consistently yielded a value about 7 % lower than the quench ratio measurements. Recollect that the quench ratio method measures the absolute beam polarization whereas the Line C polarimeter measures the $\hat{n}$ -component only. A consistent deviation of the two values indicates the presence of components of the beam polarization in the horizontal plane. To determine the correct set of values of the beam polarization, the A_n for the CH₂ target was measured with the spectrometer at 17° . By comparing the result with

the value of Ref. 37, the Line C polarimeter value was taken as the true $\hat{n}$ -beam polarization.

Determination of the absolute cross section of a scattering process requires that the incident beam flux be known precisely. This is best done by placement of a Faraday cup in the beam path. However, unavailability of a Faraday cup at HRS dictated a different method, comparison to previous p-p elastic scattering data, for absolute cross section measurements. In order to measure relative cross sections, measurement of relative beam flux between different data acquisition runs is called for. A set of ion chambers, ER02-04, placed along the beam axis inside the evacuated scattering chamber approximately 90 cm downstream from the target, were used to measure the relative beam intensity. The ion chambers contained a mixture of 90% Ar and 10% CO₂ at a pressure of 50 mm of Hg and operated at a potential difference of 90 volts. ORTEC model 439 model current digitizers were used to integrate the ER02-4 signals.

E. Standard Focal Plane

Situated at the focal plane of the HRS dipoles, an array of scintillators and multiwire drift chambers⁴¹ (MWDC) constitute the standard focal plane detection assembly. A schematic of the focal plane is illustrated in Fig. 8. The scintillators, S1-S4, provide pulse height and timing information for the particles emerging from HRS whereas the MWDC's, C1-C4, register their spatial coordinates.

Fig. 8

Fig. 9: Typical anode plane delay line electronics setup.

The scintillators S1 and S2 are approximately $76 \times 11 \times 0.64 \text{ cm}^3$ in size and are made of Pilot B and Pilot F material respectively; S3 and S4 are both made of NE110 and are $107 \times 14 \times 0.95 \text{ cm}^3$ each in size. The choice of scintillator sizes was made so as to define an acceptance smaller than the chamber sizes. All four are double ended scintillators, i.e., they are optically coupled to photomultiplier tubes at both ends. Passage of a particle through each scintillator thus produces two signals (top and bottom); both signals are then discriminated by LRS 621 leading edge discriminators and fed to an LRS 624 mean timer module. Hence, the output of the mean timer produces a signal independent of the position of the particle on the scintillator.

Each of the four MWDC's consists of a pair of planes for obtaining x- and y-position information. The four chambers have identical construction with an active area of $60 \times 10 \text{ cm}^2$ each. The chambers are grouped in pairs, C1,C2 and C3,C4; each pair is mounted on aluminum frames placed 50 cm apart. An array of anode wires with 8.13 mm relative spacing interlaced with a similar array of cathode wires constitute one detection plane. The wires are made of gold plated tungsten, 0.8 mil in diameter. The anode wires are connected to a delay line at regular intervals whereas the cathode wires are connected to a common ground bus. A high voltage power supply feeds 2100 volts to the anode planes. The two x- and y-planes within a given pair of chambers are offset exactly by half the wire spacing; this allows for improved spatial resolution. Each pair of chambers (four drift planes) is covered with 1 mil aluminized mylar. A mixture of argon, CO_2 , methylal, and isobutane is fed to the chambers to serve as

an ionizing medium. Passage of a charged particle through this gaseous mixture produces ion-electron pairs; the electrons are then collected by the nearest anode wire resulting in a signal in the delay line; these signals produced at each end of the delay line are then discriminated by LRS 821 discriminators and used as stop signals to obtain delay time information. Illustrated in Fig. 9 is the electronics set up for a typical anode plane. The time information obtained here is eventually converted to obtain information about the position of the particle (see section II.G).

F. The Focal Plane Polarimeter

The Focal Plane Polarimeter (FPP) is located at the downstream end of the HRS and is capable of measuring the transverse polarization of the scattered particles over a broad range of excitation energy. Since the front chambers in the standard focal plane detection system maintain the desired high energy resolution, no compromise between energy resolution and efficiency of the FPP need be made; this allows for the use of a very thick analyzing target to achieve high FPP efficiency. A detailed description of the FPP may be found in Ref. 42; however, for the sake of completeness, a brief description will be given here.

The FPP consists of a thick carbon analyzer followed by an array of scintillators and MWDC's. Fig. 10 shows a schematic representation of the FPP. In the FPP mode of operation, S3 and S4 of the standard focal plane configuration are removed and the two rear scintillators

Fig. 10: The Focal Plane Polarimeter layout.

SP1-SP2, rear MWDC's C5-8, and the desired number of carbon analyzing slabs are inserted into the detection system. All scintillators and chambers are mounted on rails using ball bearings and stops allowing for easy, reproducible transition between FPP and non-FPP modes of operation.

The carbon analyzer consists of up to 10 graphite slabs, $41 \times 41 \times 3 \text{ cm}^3$ each in size, mounted on individual slots allowing for independent insertion and retraction. Hence, analyzer thicknesses stepwise variable from 3 to 30 cm may be obtained. The actual analyzer thickness for a particular experimental need is determined by optimizing the desired FPP efficiency while keeping the multiple coulomb scattering in the graphite acceptably small. In the present experiment, 3 graphite slabs were used resulting in an analyzer thickness of 9 cm; this enabled us to achieve an overall FPP efficiency of better than 15 % while keeping the multiple coulomb scattered cone down to 1.8° .

Scintillators SP1-SP2 are a pair of double ended scintillators, $38 \times 76 \times 0.95 \text{ cm}^3$ each in size, overlapping each other laterally by 3 cm. The total area of the two SP scintillators allows for an acceptance of 25° in all directions for particles scattered from the center of the carbon analyzer. The anode signals from each end of SP1 and SP2 were processed in a manner similar to that described in section II.E; in addition, the SP1 and SP2 logic signals were ORed to give rise to one SP logic signal.

Chambers C5-C8 are similar to those described in section II.G except they have larger active areas of $60 \times 60 \text{ cm}^2$ each to comply with the 25° acceptance described above. They measure spatial coordinates of the particles scattered by the carbon analyzer enabling reconstruction of the trajectories of the analyzed particles.

G. Data Acquisition

The scintillator and drift chamber signals are transmitted via cables from the HRS focal plane to the Area C Counting House (CCH) where they are read into the CAMAC data acquisition modules;⁴³ the scintillator signals are fed to the analog to digital converters (ADC's) and the timing signals from the drift chambers are fed to the time to digital converters (TDC's) of the CAMAC system. The output of the CAMAC modules consists of 38 words of 8 bit length per event. A fast front-end processor, the Microprogrammed Branch Driver (MBD), controls the transfer of data from the CAMAC modules to the online PDP 11/45 computer in a buffered mode which in turn records them on magnetic tape for further analysis. Depending on the data arrival rate, a fraction of the events is analyzed online by the PDP 11/45 to help monitor the usefulness of the data being acquired.

Event Trigger

Passage of a particle, after being scattered at the target and analyzed by the HRS, with proper correlation in timing through a desired set of scintillators constitutes an "event" and generates an

"event trigger" signaling the CAMAC electronics modules to record information about the event. In order to ascertain that the event is not a product of background particles firing the scintillators, a definite coincidence logic for the scintillators was established. The event logic was set up as $(S1.S2).(SP1+SP2)$. Once a valid event signal arrives, it is gated by computer NOT BUSY, and the output, known as the event trigger, is used as the start signal for the ADC's and TDC's.

Particle Identification

The bending of trajectories in the HRS depends on the ratio of the momentum p , to the charge, z , of the particle; this makes it possible for any reaction product at the target with the same p/z to pass through the HRS and land at the focal plane detectors. It is then clear that a more definite method for particle identification (PID) is essential. The pulse height of a scintillator signal is proportional to $(zE/p)^2$, where E is the energy of the particle. On the other hand, the time of flight (TOF) is proportional to E/p . Therefore, a two dimensional plot of PH vs. TOF may be used to distinguish particles with different masses and charges. At 319 MeV protons travel at a speed of 0.196 m/ns; it takes protons 10.03 ns to traverse from S1 to SP1-2, which is much larger than the time resolution of the data acquisition electronics. The mean pulse height in S1 and the TOF between S1 and the SP scintillators were used for PID in the experiment.

Chamber Calibration

Timing information obtained from the fast delay line readout of the wire chambers was used to obtain the position of the particles. Signals produced at the positive end (P), and the negative end (N), of the delay line are used as stops for the TDC; the event trigger is used for the start. Thus, the outputs of the TDC yield the arrival times of the signal at each end with respect to the trigger.

If a charged particle traverses near the n^{th} anode wire, the arrival times of the signal at the P and N ends of the delay line

$$t_P = t_d + n \cdot \Delta t$$

$$t_N = t_d + (N - n) \cdot \Delta t \quad (10)$$

Here, t_d is the drift time of the electron in the gas; N is the total number of anode wires; and Δt is the delay time interval between two consecutive wires. From Eqs. (10) time sums and differences may be formed as:

$$t_+ = t_P + t_N = 2t_d + N\Delta t$$

$$t_- = t_P - t_N = 2n\Delta t - N\Delta t \quad (11)$$

It is clear from the expressions that t_- locates the position of the wire whereas t_+ gives the drift time in the gas, except for additional constant terms which can be corrected for in the software.

The TDC's convert t_0 to 9 bit data words representing wire position; the output, however, will be a function of the gain of the TDC's and the actual positions may be written as:

$$x = a + b \cdot t_0 \quad (12)$$

The chamber calibration process consists of the determination of the gain, b 's, of the TDC's for all chamber planes and adjustment of the constant offsets, a 's, such that all planes are spatially aligned. In order to achieve this, the spectrometer is set for quasi-elastic scattering so that there is a uniform distribution of events over the active area of the chambers. The spectra of crude anode wire positions are truncated to integral intervals of the known wire spacing by adjusting the b_1 coefficient; then offset values are added to each pairs of chambers to achieve spatial alignment.

Due to avalanche effects of the released electrons following the ionization of the medium, the drift position of the electrons is nonlinear with their drift time. Therefore, a lookup table of the drift positions for given drift times is used to obtain the drift positions.

Even though, from the above information, the wire and the distance of the particle from the wire may be obtained, it is not sufficient to locate whether the particle passed to the left or to the right of the wire. To resolve this ambiguity, the drift planes of the second chamber in the group are used. Independent position information from the two planes (which are offset physically by half a wire

spacing) gives rise to four possibilities for the position, two of which must be coincident, in principle, with the actual position. Therefore, out of the four possibilities, the mean of the two with the least difference is taken to be the actual position.

Trajectory reconstruction

For a complete description of the trajectories of the particles at the focal plane four quantities are needed: x_f , y_f , θ_f , and ϕ_f where x_f and y_f are defined to be the mean x and y position in chambers C1,2 respectively, and θ_f and ϕ_f are the vertical and horizontal angles of the trajectory with respect to the beam axis. From the difference in position information between C1,2 and C3,4, θ_f and ϕ_f are easily constructed. However, for the description of the scattering process, the focal plane trajectories need be traced back to the target; a knowledge of the dipole field settings and the transport optics of the HRS is, therefore, required. If we define δ as the deviation from the central momentum, p_c (determined from the dipole fields), $\delta = (p - p_c)/p_c$, where p is the momentum of the detected particle, then target quantities may be determined from the focal plane quantities using the inverse transport matrix⁴⁴ for the HRS.

Event Rejection

Bad events fall into two basic categories: 1) events that are a product of actual nuclear scattering but do not yield any information pertinent to the experiment and 2) instrumental background, which

arises from processes such as scattering from slits, flanges, and beam pipes, etc. In either case, these bad events should be rejected.

The measurement of the polarization of the scattered particles with the FPP is most efficient in an angular region where the product $\sigma \cdot A_n$ of the carbon analyzer is at a maximum. However, because of strong multiple coulomb scattering in the thick carbon analyzer, about 75 % of the analyzed events scatter between 0 and 2° where the inclusive A_n of carbon is close to zero. Therefore, such events yield very little polarization information and must be rejected.

Once the trajectories of the particles are reconstructed with the analysis software, the events that are identified as bad events can be rejected. However, when a significant fraction of the events falls into the category of such bad events, the analysis process becomes overly inefficient. In order to avoid this, the MBD was programmed to perform a crude construction of the event trajectories and reject unwanted events before they are passed on to the PDP 11/45 to be recorded on magnetic tape for future analysis. Since this method of data rejection is based on an approximate determination of trajectories, the limits of rejection were liberal enough to ensure non-rejection of good events. Over all, about 75% of the events were rejected this way in the present experiment.

H. Data Analysis Software

The data were analyzed using the standard "Q" data analysis software⁴⁵ developed at LAMPF. The Q system consists of three major components: the ANALYZER, the ALLTEST package, and the DISPLAY system. The ANALYZER and its associated programs are responsible for transferring raw data words from the MBD buffers to the data tapes for each event trigger. In addition, the ANALYZER is also capable of processing some or all of the raw events; in the "may-process" mode, used in online data acquisition, only a fraction of the events, depending on the available computer time between data transfers, is processed by the ANALYZER. It is clear then that the online analysis depends on the rate of data acquisition. In the present experiment, typically 20 to 30% of the data were analyzed online. The processing tasks performed by the ANALYZER consist of calculating various useful quantities, referred to as data words, from the raw data on tape (pulse heights and position information); the calculated data words are used in mapping particle trajectories and examining proper correlation between data words dictated by a desired good event. Furthermore, the calculated data words are also utilized by the ALLTEST package and the DISPLAY programs.

The ALLTEST package⁴⁶ enables the experimenter to conditionally accept events that fulfill criteria pertinent to the experiment. The heart of the package is the test file which contains a series of tests to be performed on every event that passes through the ANALYZER. These tests fall into two categories: tests performed on the data words

themselves, microtests; and tests performed on various logical combinations of the data words, macrotests. A total of 120 micro and macro tests may be defined in the test file; any of these tests, addressed by its sequence number in the test file, may be imposed upon the events in accumulating calculated data.

The DISPLAY system^{4/} facilitates interpretation of the data being accumulated through visual representation of the data. The DISPLAY software is capable of histogramming any of the data words (raw or calculated); in addition, it can produce two dimensional plots (dot-plots) for any pair of data words . The ability to accumulate histograms and create dot-plots as data are being accumulated enables the experimenter to monitor the progress of the experiment and check for desired correlations between any pair of data words ensuring the usefulness of the data online. The histograms are defined in the DISPLAY software by the experimenter by specifying the desired data word and its range. In addition, a test number defined in the ALLTEST package may be specified for the histogram; increment of the histogram will be subject to successful passing of an event through this test. A maximum of 40 histograms may be defined. The defined histograms may be saved on disk files for later examination and analysis.

Overall, the "Q" system allows the user to choose a software setup suitable to the needs of the experiment. Such flexibility enables the use of a standardized software package to fulfill the requirements of a variety of experimental setups.

I. Small Angle Problems and Remedies

Small angle inelastic proton scattering from a heavy nucleus is quite an experimental challenge. Many possible sources of background complicate the data analysis process. The sources of background in the scattered spectrum include inferior incident beam quality and undesired scattering processes that take place near the target, inside the HRS cavity, or in the detection system at the focal plane. We will consider each of these cases and remedial steps in the following.

A large spatial spread of the beam profile results in scattering from the beam pipes and seams. Furthermore, when collimator jaws are used in Line C to regulate the size of the beam spot and its intensity at the target, similar undesired scattering from the jaw faces may take place. In both cases, the momentum loss in these extraneous scattering processes produces lower momentum components in the beam. A nonstandard beam tuning procedure special to the small angle setup was used to reduce such beam related background. Recollect that the beam spot in Line C is an image of the stripped beam in Line X. Assuming that a clean beam is delivered to Line X from the LINAC, quality of the beam on target will depend on the beam-tune of Line X and Line C. To achieve this, all the collimating jaws in Line C were pulled out to their outer limit and the beam spot size and intensity were controlled with appropriate strippers in Line X. Several tests were done with thin aluminized mylar foils and tungsten wires as Line X strippers to optimize the beam quality on target; the horizontal and vertical beam profiles were monitored with a series of integrating secondary

Fig. 11: A two dimensional graph of θ_t vs ϕ_t .

ionization counters (ISIC) at different locations in Line C. Based on these tests, a tungsten wire was chosen for Line X stripper. The beam condition in Line X through Line C was monitored periodically throughout the experiment and retuning was performed as required to maintain a superior quality of the incident beam on target.

The straight-through beam, at a small angle setting of the HRS, is no longer able to travel down a beam pipe to the beam dump. Most of this beam hits a vertical steel post located at the entrance of the HRS cavity. This results in a very large room background accounting for a large counting rate in each individual detector at the focal plane. These particles do not produce the required coincidence to cause an event trigger. Thus, they do not contribute to background in the scattered spectrum; they merely limit the data acquisition rate. However, part of this beam enters the HRS cavity after undergoing glancing collisions with the hardware present near the acceptance of the HRS producing background events. This is illustrated in Fig. 11 in the two dimensional dot-plot of the horizontal and vertical scattering angles at the target, ϕ_t and θ_t , respectively; a large concentration of events at small ϕ_t independent of θ_t indicates scattering from objects located at the small angle side of the HRS acceptance. The ability to selectively accept events with BOX tests enables one to reject such instrumental background events. A box as shown in Fig. 11 was defined to reject these events.

The treatment of background events produced after the target by extraneous scatterings in the HRS cavity requires a thorough understanding of the undesired scattering processes. The presence of a

Fig. 12: A two dimensional plot of θ_c and x_f .

strong elastically scattered beam, typically three orders of magnitude higher than the inelastic peaks of interest, significantly amplifies instrumental background related directly or indirectly to the strong elastic peak in small angle scattering. Directly related to the elastic peak is the familiar background that falls off from low to high excitation energy, known as the elastic tail. The elastic tail is partly the product of straggling of the beam in the target. The shape and intensity of the elastic tail is a function of the energy of the incident protons and the density, thickness, and charge of the target. All other parameters remaining the same, the use of thinner targets helps in minimizing the elastic tail. On the other hand, too thin a target may result in an unacceptably low event rate. The choice of the target thickness thus was based on optimizing the elastic tail, counting rate, and energy resolution.

The possibility of reducing elastic related background through proper positioning of the elastic peak at the focal plane and changes in the hardware configuration and the detection system were also investigated. Note that there is about a 10 cm gap between the top edge of the drift chambers and the outer exit flange of the HRS; the HRS dipole fields can be set so as to position the elastically scattered peak anywhere in this gap. Tests with different positions of the elastic peak in this gap were conducted to isolate sources of elastic related background. An understanding of the origin of events that enter the focal plane detection system is provided by the two dimensional dot-plot of the x-position at the focal plane, x_f , and the corrected angle of incidence, θ_c (corrected for removing the known

correlation between x_f and θ_f); a typical dot-plot is illustrated in Fig. 12. Positioning of the elastic peak off the focal plane detectors implies that the elastically scattered protons travel very near the outer wall of the HRS dipoles; consequently, spurious rescattering of the elastic beam from the dipole walls, seams, and flanges etc. may occur. Such secondary scattered particles may no longer appear at or near the position of the elastic peak at the focal plane dictated by the magnet settings; they may end up in the inelastic regions of interest. However, these rescattered particles will have a downward angle of incidence at the focal plane compared to the good inelastic events. The rescattered elastic events are clearly indicated by intense bands with downward θ_c in the upper region of the dot-plot shown in Fig. 12. On the other hand, the somewhat spread out band of events in the bottom of Fig. 12 have an upward θ_c indicating possible rescattered particles from the inner walls of the HRS. This group of particles have lower momentum than that of the central ray; they are most likely to be a mixture rescattered particles from the lower momentum components of the beam and nuclear scattering into the high excitation energy region. The good inelastic events, with proper angles of incidence, lie in the rectangle marked between these two groups of spuriously scattered events. Thus, a BOX test was defined in the test file to designate the events that lie within the rectangular (solid) region of Fig. 12 to be the desired inelastic events. However, the boundary between the background and the good events is not so sharp as to clearly isolate the two kinds of events. Possible contamination

of the good events by the background events and their effects on the measured observables is discussed in Chapter. III.

Although positioning the elastic peak optimally above the edges of the drift chambers helped reduce the background, it did not eliminate all the rescattering problems. Note that the focal plane scintillators are larger than the active area of the drift chambers with phototubes at both ends; thus, rescattered particles from the elastic beam from the exit flange and the focal plane assembly holder still could produce background events in the low excitation region. In order to avoid this during our second phase of data acquisition, a specially designed single ended scintillator was used in place of the standard S1 to restrict the active area for the event trigger in the low excitation side of the focal plane (recall that S1 must produce a signal to get an event trigger). The new S1 was 0.1275" in thickness and had a phototube attached only at its lower end (high excitation end). The upper end was positioned in line with the active area of the drift chambers.

The dramatic improvement achieved in the background, particularly in the low excitation region, with all the measures discussed above included in the setup is shown in Fig. 13. The dashed spectrum was taken with the thick ^{90}Zr target and standard S1 used in the event trigger; the solid spectrum was the result of the use of the thin target and the special S1 arrangement as outlined above.

Fig. 13: Low excitation background with (solid) and without the special small angle setup.

J. Summary of the Experiment

The data were acquired for the present study under experiment 660 of the HRS. The spin-flip probability measurements were obtained during June, 1982. The data were obtained for the $^{90}\text{Zr}(\vec{p},\vec{p}')^{90}\text{Zr}^*$ and $^{48}\text{Ca}(\vec{p},\vec{p}')^{48}\text{Ca}^*$ reactions at 3.0° and 5.0° laboratory scattering angles. The thick (250 mg/cm^2) ^{90}Zr target was used in order to achieve high counting rates. The standard setup of the focal plane was using during these spin-flip probability measurements. The experiment was conducted over a period of ten days. Typically, about 10 nA of beam current was delivered onto the Line X stripper and about 80 pA was delivered on target, resulting in an event rate of about 400/sec. The polarization of the beam varied from 72 to 84%, but it was generally around 78%. About 30% of the total time was devoted to hardware setup, calibration, and systematic checks, and the remainder to data acquisition.

In the second phase of experiment 660, conducted for three days during October, 1982, emphasis was given to the measurements of cross sections and analyzing powers for $^{90}\text{Zr}(\vec{p},\vec{p}')^{90}\text{Zr}^*$ and $^{48}\text{Ca}(\vec{p},\vec{p}')^{48}\text{Ca}^*$ reactions. For these measurements, where high resolution and good peak to background ratio was desired, the thin (50 mg/cm^2) ^{90}Zr target and the specially designed S1 was used. The data were taken at 2.75° and 4.25° laboratory scattering angles. In addition, to investigate systematic effects, some of the earlier spin-flip measurements were repeated with this different focal plane setup and beam conditions.

III. DATA ANALYSIS AND RESULTS

During real time (online) data acquisition, fast event rates preclude the possibility of every event being analyzed. Priority is given to recording all events on tape and only a portion of the events are processed depending on the availability of computer processing time. A complete analysis of the data is performed offline where the recorded data are replayed and analyzed by the offline analysis software; this also allows the experimenter to thoroughly investigate and scrutinize the systematics and the quality of the data. In this chapter the principles and procedures of offline data reduction will be discussed. In addition, the final experimental results will be presented.

A. Analysis Software

The software package used for offline data analysis is very similar to the one described in Sec. II.H for online data acquisition. However, unlike online data acquisition, the ANALYZER is set for "must-process" mode which implies that every event recorded on tape must be processed by the ANALYZER. It is possible to define test file and histogram specifications entirely different from those used online; offline analysis, therefore, can include much more rigorous tests and so greatly improve the quality of the final data.

The standard HRS analyzer was used without any modifications for the offline analysis. However, the test package and the histogram specifications were altered significantly from those used online. For the analysis of the cross section and analyzing power data, tests were defined to divide the 2° -acceptance of the HRS into four angular bins of 0.5° width each; the histogram specifications were modified accordingly to accumulate spectra for all the bins for normal and reverse beam polarizations. The BOX tests imposed on the two dimensional plots of θ_c and x_f , and θ_{tgt} and ϕ_{tgt} were made more stringent to reject instrumental background events (see Sec. II.I). In addition, a BOX test was defined to restrict the computation of chamber efficiencies (see Sec. III.B) to the central region of the focal plane. For the analysis of the focal plane polarimeter data, however, the angular acceptance was kept at the standard 2° width in order to achieve the desired statistical accuracy. Additional tests were defined in the test file to accumulate sums (see Sec. III.C) for the computation of spin-flip probability in two-MeV bins of excitation energy. Histograms were defined to accumulate polarimeter spectra corresponding to left-right scattering in the carbon analyzer for normal and reverse beam polarizations. Analysis of the polarimeter data was performed using the program "DNC".⁴⁸

B. Cross Section and Analyzing Power

The differential scattering cross section at a scattering angle θ may be defined as the fraction of the incident flux scattered per unit solid angle per unit scattering center. This may be expressed as:

$$\frac{d\sigma}{d\Omega}(\theta) = \frac{dI}{I_0 \cdot n_t \cdot \Delta\Omega} \quad (13)$$

where dI is the number of particles scattered into the solid angle $\Delta\Omega$, n_t is the number of scattering nuclei per unit area, and I_0 is the incident beam flux. When n_t is expressed in cm^{-2} and $d\Omega$ in steradians (sr), $\frac{d\sigma}{d\Omega}$ has the dimension of cm^2/sr ; however, the differential cross section is traditionally expressed in the units of millibarns per steradian (mb/sr), which is 10^{-27} cm^2 .

It has been outlined in Sec. II.D that the scattering of a beam of polarized particles from a nucleus produces asymmetry of the scattered particle yields in the reaction plane. Utilizing Eq. (8), the analyzing power A_n of a nuclear reaction may be defined as the asymmetry ϵ produced per unit beam polarization P_b .

$$A_n = \frac{\epsilon}{P_b} \quad (14)$$

If the polarization state of the incident particle is allowed to switch between parallel (normal) and antiparallel (reverse) with respect to the normal to the reaction plane, it has been pointed out (Sec. II.D)

that the four possible yields, left-normal, I_L^N ; left-reverse, I_L^R ; right-normal, I_R^R ; and right-reverse, I_R^N ideally provide redundant information, namely $I_L^N = I_R^R$ and $I_L^R = I_R^N$. Thus, instead of counting yields by setting the spectrometer at the same angle on left and right (which is quite impractical and time consuming), it is adequate to keep the spectrometer fixed and gather yields N and R corresponding to normal and reverse polarization of the incident beam, respectively. This enables us to express the asymmetry as:

$$\epsilon = \frac{N - R}{N + R} \quad (15)$$

In the present experiment the yields corresponding to the HRS situated to the left of the incident beam were measured.

The purpose of the data analysis process is to reduce the crude experimental quantities into physical results. The crude quantities are often plagued by inefficiencies in the detection system and arbitrariness in their physical meaning. Thus, appropriate corrections to remove instrumental inefficiencies and normalization of the arbitrary quantities are needed. In the following sections, the methods used for inefficiency corrections and normalizations will be discussed.

The efficiency of the wire chambers in detecting particles traversing them depends on the particle flux. There is a relaxation time after detection of each particle and if a second particle arrives in this time interval the chambers fail to detect it. However, the

scintillators can handle much faster counting rate than the wire chambers; hence, events are written on tape whether or not they satisfied good chamber detection requirements. One can then compute chamber efficiency from the knowledge of how many particles produced good chamber events out of the total number of events triggered by the scintillators. However, in actuality the chamber efficiencies were calculated after demanding that the particles pass the PID test for protons and that they not lie at the edges of the focal plane; the latter requirement was imposed upon the discovery that in small angle scattering, events at the low excitation edge of the focal plane produce rescattered events with abnormal trajectories causing an artificial lowering of the chamber efficiency. Tests corresponding to the above requirements were defined in the test file and the chamber efficiency was calculated using the following:

$$CHMEFF = \frac{ALL\ CHAMBERS\ OK * PID * BOX\ 7}{PID * BOX\ 7} \quad (16)$$

where BOX 7 was a box test to exclude the outer edges of the focal plane.

Another kind of inefficiency arises from the fact that the computer is limited in speed to accumulate data; this inefficiency known as computer dead time can be computed from the scaler counts of total events and the events that produced triggers (recall that an event may produce a trigger only if the computer is not busy).

All measured yields may now be corrected for the inefficiency of the chambers and the computer to result in the corrected yields as follows:

$$\text{CORRECTED YIELD} = \frac{\text{MEASURED YIELD}}{\text{CHMEFF} * \text{COMPLT}} \quad (17)$$

where COMPLT is the computer live time.

During data acquisition, yields were measured corresponding to normal and reverse beam polarization with the polarization changing direction every minute. In order to compute asymmetry and cross sections the corrected normal and reverse yields must be normalized with respect to the beam flux corresponding to the polarization state. The counts registered in the beam line monitors, the ion chambers ERO2-4 (Sec. II.D), were used for this relative normalization process. ERO4 was chosen for normalization purposes; the selection of the specific ion chamber was made after monitoring the stability of all the ion chambers. In the case of very small angle setting of the HRS where it was impossible to put the ion chambers in the beam path without producing significant amount of background counts, the Line C polarimeter scintillators were used for obtaining normalized yields. We may write, therefore, the normalized yields as:

$$\text{NORMALIZED YIELD} = \frac{\text{CORRECTED YIELD}}{\text{ERO4}} \quad (18)$$

In addition to the overall inefficiencies outlined above, the wire chambers exhibited local inefficiencies at the edges of the wire chambers. The efficiency of the chambers at different focal plane positions in the x-direction was scanned with the help of the elastic scattering from ^{208}Pb and appropriate corrections were made to obtain uniform efficiency. Since the scattering angle was divided into four bins, the efficiency of the chambers in each of these bins was also measured and correction factors for each of the bins were used to yield the known shape of the angular distribution of the elastic scattering cross section from hydrogen.

Normalization

The normalized yield given by Eq. (18) represents the number of particles scattered into the solid angle of acceptance of the HRS (absolute yield), within a proportionality constant; this constant, however, is the same for normal and reverse beam polarization states and is a function of the gain of the beam-line monitors. As evident from Eq. (15), this constant cancels out in the calculation of the asymmetry, and hence, in the analyzing power. However, the determination of the differential cross section dictates the absolute measurement of the scattered fraction of the incident beam flux. Therefore, the proportionality constant between the normalized yield and the absolute yield must be determined. In the absence of a Faraday cup it is impossible to measure the incident beam flux absolutely. However, this inadequacy may be circumvented, if independently measured

and reliable cross section data are used as a reference as outlined in the following. If Y is the number of particles scattered into the solid angle $d\Omega$, substituting Y for dI in Eq. (13) yields

$$\frac{d\sigma}{d\Omega}(\theta) = \frac{Y}{N \cdot n_t \cdot \Delta\Omega}, \quad (19)$$

where N now becomes an arbitrary scale constant. If the yields Y and Y' for two reactions are measured with the same solid angle, and the differential cross section for one of the reactions, $\frac{d\sigma'}{d\Omega}(\theta')$, is known from previous independent measurements, then the unknown cross section $\frac{d\sigma}{d\Omega}(\theta)$ may be expressed as follows:

$$\frac{d\sigma}{d\Omega}(\theta) = \left(\frac{n_t'}{Y'}\right) \cdot \left(\frac{d\sigma'}{d\Omega}(\theta')\right) \cdot \frac{Y}{n_t} = N' \cdot \frac{Y}{n_t} \quad (20a)$$

where

$$N' = \frac{1}{N \cdot \Delta\Omega} = \left(\frac{n_t'}{Y'}\right) \cdot \left(\frac{d\sigma'}{d\Omega}(\theta')\right) \quad (20b)$$

Hence, the determination of the absolute normalization constant N' is all that remains to be determined to compute the absolute cross section for the unknown reaction. From Eq. (20b), note that it is adequate to measure the total yield Y' at only one scattering angle θ' for the reaction used as the reference.

The elastic scattering from hydrogen was used as the primary source of reference for the absolute normalization. The total yields for the elastic scattering from the hydrogen in two CH₂ targets of areal density 50 mg/cm² and 150 mg/cm² were measured at scattering angles of 3.5°, 13°, and 17°. The known cross section data were taken from Ref. 49. In addition, the p-p cross section predicted by the phase shift analysis of Arndt⁵⁰ were used for comparison purposes. Furthermore, as secondary sources of reference, elastic scattering data⁴⁹ at about 300 MeV for ²⁰⁸Pb were used for verification purposes.

Excitation Energy Calibration

Since the HRS operates in the energy loss mode, the calibration of the excitation energy of the reaction does not involve the determination of the incident beam energy precisely. If the position of a peak with a well known excitation energy, such as the elastic peak, is determined, the excitation energy as a function of the focal plane position can be calculated from magnetic field settings of the HRS bending magnets using the transport matrix of the HRS. In addition, this also determines the energy of the incident beam if the energy loss due to straggling in the target and the recoil energy from the kinematics of the reaction are known. Thus, when several targets are used in an experiment, once the beam energy is calibrated with a known peak in the spectrum of a reference target, the excitation energy can be determined accurately for any target even if no well known peaks are present in the spectrum. In the present experiment the 10.23 MeV

Fig. 14: A spectrum of $^{12}\text{C}(p,p^-)^{12}\text{C}^*$ at 3.5° .

M1 peak in ^{48}Ca was used as a primary reference for the excitation energy calibration. However, as a means of independent check, the two well known M1 states at 12.71 MeV and 15.11 MeV of ^{12}C were also used as references. The energy loss in the various targets used was calculated with the program LANDAU. Since there is a possibility that the incident beam energy changes when beam line bending magnet fields are tuned, calibration spectra were accumulated following every retune of the beam lines.

Cross Section and Analyzing Power Results

A spectrum of $^{12}\text{C}(p,p')^{12}\text{C}^*$ taken for excitation energy calibration purposes is presented in Fig. 14; the spectrum shows clearly the two M1 states mentioned above. Spectra of ^{48}Ca at 3.5° and 5.0° laboratory scattering angles are illustrated in Fig. 15; notice the insignificant amount of instrumental background in the small angle spectrum. The cutoffs in the low and high excitation region of the spectrum were introduced artificially via the MBD-cut data rejection method as outlined in Sec. II.G to reject unwanted portions of the spectrum. A very clear and isolated peak was observed at 10.23 MeV; the states at 6.92 and 8.88 MeV arise from an ^{16}O contaminant in the target. The energy resolution in these spectra with the 100 mg/cm² target was about 120 keV.

Normalized yields for the 10.23 MeV peak were extracted after subtracting a small linear background as shown in Fig. 15 and summing the yields in the spectra between 10.18 and 10.42 MeV. The angular

Fig. 15: Spectra of $^{48}\text{Ca}(p,p')^{48}\text{Ca}$ at 3.5° and 5° .

Fig. 16: Angular distributions for a) $\frac{d\sigma}{d\Omega}$ b) A_y for the 10.23 MeV state in the $^{48}\text{Ca}(p,p')^{48}\text{Ca}$ reaction at 319 MeV.

distribution for the absolute differential cross section and the analyzing power for this peak are shown in Fig. 16. The $\frac{d\sigma}{d\Omega}$ angular distribution shows a sharp fall off with increasing scattering angle, suggestive of a zero spin angular momentum transfer ($\Delta L = 0$) in the reaction. The values of A_n are mostly small with a hint of possible structure. A detailed interpretation of the results will be deferred until Chapter IV.

Spectra taken with the thin ^{90}Zr target at 2.75° and 4.25° laboratory scattering angles are illustrated in Figs. 17a-b. Necessary precautions have been taken in reducing the instrumental background as discussed earlier; a significant improvement in the quality of the spectra has been achieved compared to previous proton inelastic scattering from ^{90}Zr (Ref. 15). The energy resolution obtained was better than 70 keV. Qualitatively, the spectra show three distinguishable features: 1) on the low excitation side a group of narrow and well isolated states are clearly visible; 2) there is a wide bump present at the 9 MeV region and hints of fine structure within the bump; 3) the spectra then rise toward the high excitation region; when combined with spectra taken in the high excitation region alone (not shown in the figure), a peak which is several MeV wide and centered at about 16 MeV is clearly evident.

The excitation energies of several of the low lying states observed in the spectra were compared with previous measurements. A summary of the comparison is given in Table VIII (page 165). Excellent agreement was obtained for the strongly excited natural parity states. Good to fair agreement was obtained for the weakly excited states.

Fig. 17a

Fig. 17b

Fig. 18a

Fig. 18b

Fig. 18c

Fig. 19: Analyzing powers of the strongly excited low lying states of ^{90}Zr .

Some of these states are marked in the spectrum of Fig. 17b. The states listed in Table VIII were fitted after subtracting a background as shown in Fig. 17. A similar fitting procedure was adopted for these low lying states as the one for the more densely packed states in the 9 MeV region which is described later. The yields in the fitted peaks were then corrected and normalized and angular distributions for $\frac{d\sigma}{d\Omega}$ and A_n were obtained. The uncertainties in the A_n results for the weakly excited states were found to be too large to be useful. The angular distributions for $\frac{d\sigma}{d\Omega}$ of most of these low lying states and A_n of the strongly excited ones are shown in Fig. 18 and Fig. 19, respectively. The curves in these figures represent theoretical calculations that will be described in Chapter IV.

The 9 MeV region bump was first analyzed as the product of a single excitation mechanism. As evident from the spectra of Figs. 17, it is not straightforward to subtract a background from the bump in the fitting procedure as was the case for the low lying states; the large width of the bump and poor peak to background ratio introduce a large systematic uncertainty into the background subtraction procedure. There are no theoretical predictions for the shape of the nuclear continuum in this region. A large portion of the background underneath the bump most likely arises from the excitation of real states in the continuum, like the strongly excited giant resonances at higher excitation energy. However, the possibility of some instrumental background cannot be ruled out. In view of the uncertainty in the choice of a background curve, it is reasonable to assume that the background shape in this region is a smooth curve consistent with the

Fig. 20a

Fig. 20b

shape of the spectrum on the lower and higher excitation ends of the bump. Several such empirical background curves were tried; the uncertainty in the yield of the bump due to background variations was estimated at $\pm 20\%$. Yields for the bump were extracted by summing the yields measured in 20 keV bins from 7.4 to 10.2 MeV. In addition, the entire bump was fitted with a 1.8 MeV wide gaussian function with its centroid at 8.98 MeV; the fitted yield was found to be consistent with the summed yield. Angular distributions of $\frac{d\sigma}{d\Omega}$ and A_n for the entire bump were obtained after subtracting an empirical background as shown in Fig. 17; these angular distributions along with theoretical curves are shown in Fig. 20.

A more through examination of the bump reveals a rich fine structure; this is seen clearly in the spectra expanded around the 9 MeV region as illustrated in Fig. 21. It is important to note that such fine structure has not been observed previously in the reported inelastic proton scattering data (Ref. 15). Although the initial motivation of the experiment was to measure the spin-excitation strength in the proposed giant M1 resonance region of ^{90}Zr , observation of identifiable fine structure in the M1 bump, even in the spectra (see Fig. 28) taken with the thick ^{90}Zr target with a moderate resolution of 170 keV, prompted measurements with higher resolution and better statistics. The spectra shown in Figs. 17, obtained as a product of this, confirmed our original findings. The strengths and the excitation energies of the low lying states observed in both sets of spectra were consistent within the measured statistical accuracy. The rather high level density of states observed in the fine structure of

Fig. 21: Expanded spectrum for the 9 MeV region in ^{90}Zr .

the 9 MeV bump might make one sceptical as to the reality of the fine structure. However, it is important to note that the peak to valley yield ratio of the observed fine structure in this bump is well beyond three standard deviations of the measured statistics. In addition, the linearity of the wire chambers were also investigated to eliminate the possibility of spurious structure arising from such instrumental nonlinearities; the calibration of the wire chambers was fine tuned to eliminate misidentification of wires in the detection planes. Incorrect calibration of the drift positions (see Sec. II.G), is another possible source of spurious instrumental structure; however, such structure appears at a periodic interval of 4 mm in x-position of the focal plane, which approximately corresponds to 200 keV of excitation energy. The average separation of the fine structure peaks observed is about 100 keV, which rules out the possibility of bad drift position calibration as the cause. Nonetheless, to remove any uncertainty arising from the wire chamber non-linearities, measurements were repeated with the 9 MeV region placed at several different positions in the focal plane; the original findings were confirmed in these measurements. Most of the peaks observed in this region can be followed throughout the measured angular region.

The yields in the individual peaks were extracted after subtracting a background similar to the one chosen for the entire bump. The peak fitting program LOAF⁵¹ was used to fit these peaks. It was not possible to fit all the fine structure observed in this region; very weak structures with not well defined peak shapes were excluded from the fitting procedure. Two different line shapes were used as

reference for fitting the peaks: a 70 keV full width at half maximum gaussian function, which is consistent with the experimental energy resolution, and the line shape of the 3.84 MeV low lying state. The difference in the extracted yields with the two reference line shapes was found to be negligible. The peaks were fitted in groups of ten; fits were done with overlapping peaks in the groups to check for the consistency of the fits. Illustrated in Fig. 23 are the spectra expanded in the 9 MeV region; the solid line represents the fit obtained with a 70 keV width gaussian line shape. The uncertainty in the centroids of the peaks was estimated conservatively at 10 - 30 keV.

Angular distributions of $\frac{d\sigma}{d\Omega}$ obtained from the fitted yields for most of these fine structure states in the region 7 to 10 MeV are shown in Fig. 22a-c. In addition to the cross section obtained with the fitting methods, cross sections were also extracted in the region from 7 to 10 MeV by summing the yields in 100 keV intervals after background subtraction. The angular distributions of $\frac{d\sigma}{d\Omega}$ obtained in this manner was found to be consistent within errors with those obtained with the fitting procedure. In either case, the A_n results were found to involve large uncertainties and were excluded from our analysis. Many of the peaks may correspond to several unresolved states. The strength contained in the additional peaks too weak to be fitted was about 40 % of the strength contained in the entire 9 MeV bump. A complete table of the $\frac{d\sigma}{d\Omega}$ for all the states analyzed is presented in the Appendix.

The high excitation region of the spectra of Figs. 17 is dominated by the excitation of the natural parity giant E1 resonance which has been observed in photonuclear reactions^{52,53} at 16.4 MeV with

Fig. 22a

Fig. 22b

Fig. 22c

Fig. 23: Cross section and analyzing powers in steps of 2 MeV for ^{90}Zr .

a full width at half maximum of about 4 MeV. In addition, the nuclear continuum in this region is strongly excited. The observed continuum in this region is most likely due to the excitation of many broad resonances. In order to decompose the spin and parity of such resonances, angular distributions of the cross section and the analyzing power for this region were obtained in 2 MeV steps from 10 to 18 MeV without background subtraction. These are illustrated in Fig. 23.

C. Spin-Flip Probability

The spin-flip probability in a nuclear reaction is a measure of the transfer of polarization from the incident particle to the nucleus. The transfer of polarization can be determined by measuring the polarization of the incident and scattered particles. The incident particle polarization is measured with the help of the beam line polarimeters or the quench ratio method as outlined earlier. The scattered particle polarization is measured with the Focal Plane Polarimeter. In what follows we will briefly describe the principles of such measurements.

It is convenient to define coordinate systems for the incident and scattered particles: the incident particle system $(\hat{s}, \hat{n}, \hat{l})$, is the system defined by the scattering plane with $\hat{n} = (\vec{k}_i \times \vec{k}_f) / |\vec{k}_i \times \vec{k}_f|$ where $\vec{k}_i$ and $\vec{k}_f$ are the momenta of the incident and scattered particles respectively. The scattered particle system, $(\hat{s}', \hat{n}', \hat{l}')$, is defined with $\hat{l}'$ along the scattered particle momentum, $\vec{k}_f$, $\hat{n}' = \hat{n}$, and $\hat{s}'$ in

accordance with a righthanded frame. The scattered particle frame, therefore, may be arrived at by rotation of the incident particle frame about the $\hat{n}$ -axis by θ_s , where θ_s is the scattering angle. A complete description of the polarization transfer between the incident particle and the scattered particle in the scattering process is given by the five polarization transfer coefficients, commonly referred to as the Wolfenstein parameters,⁵⁴ D , R , A , R' , and A' ($D_{nn'}$, $D_{ss'}$, $D_{ls'}$, $D_{sl'}$, and $D_{ll'}$ respectively) along with the induced polarization $P(\theta)$, and the analyzing power A_n , as follows:

$$\vec{P}' = \begin{bmatrix} D & 0 & A \\ 0 & D & 0 \\ R' & 0 & A' \end{bmatrix} \cdot \vec{P} + \begin{bmatrix} 0 \\ P(\theta) \\ 0 \end{bmatrix} \quad (21)$$

where $\vec{P}$ and $\vec{P}'$ are the polarizations of the incident and scattered particles, respectively. The parameter D is known as the depolarization and the others are referred to as the spin rotation parameters.

In general the beam polarization can be oriented in any desired direction. However, when the incident beam is completely polarized perpendicular to the horizontal plane, as was the case in the present experiment, the particle polarization in the incident particle frame can be written as:

$$\vec{P} = \begin{bmatrix} \sin\phi \\ \cos\phi \\ 0 \end{bmatrix} \cdot P_n. \quad (22)$$

Here ϕ is the angle between the scattering plane and the horizontal plane and P_n is the n-type beam polarization as measured by the beam line polarimeter. For the sake of simplicity in the physical interpretation of D, let us assume $\phi = 0$. Then, substituting Eq. (22) in Eq. (21) and solving for $\vec{P}$ we get

$$P'_n = \frac{P_n D + P(\theta)}{1 + P_n A_n} \quad (23)$$

If the incident beam is completely unpolarized, substituting $P_n = 0$ in Eq. (23) one obtains the scattered particle polarization, $P'_n = P(\theta)$ and hence the name induced polarization. However, if $P_n \neq 0$, in addition to the induced polarization, the scattered particle polarization contains a fraction of the incident polarization determined by the magnitude of D. Since $|P'_n| < 1$, from Eq. (23) limits on the magnitude of D can be obtained:

$$-1 + 2|P(\theta)| \leq D \leq 1.$$

In the extremely simple case of $P(\theta) = A_n = 0$, Eq. (23) reduces to

$$P'_n = P_n; \quad \text{for } D = 1$$

$$P'_n = -P_n; \quad \text{for } D = -1$$

which implies that the scattered particle polarization remains the same (flips upside down) with respect to the incident particle polarization depending on whether $D = 1(D = -1)$.

Now we define the spin-flip probability in the $\hat{n}$ direction, $S_{nn'}$, as the probability that the scattered particle spin is flipped with respect to the incident particle spin as follows:

$$S_{nn'} = \frac{1}{2}(1 - D_{nn'}) \quad (24)$$

An $S_{nn'} = 1$ corresponds to $D_{nn'} = -1$ which implies complete flipping.

Spin-Flip Measurements

From the discussion outlined above, it is clear that the measurement of $S_{nn'}$ is synonymous with the measurement of $D_{nn'}$. Since the FPP is situated at the focal plane, the measured physical quantities are obtained in the focal plane (FP) coordinate system. However, the polarization measured at the focal plane can be easily transformed into the scattered particle frame if the angle of precession, χ , for the spin of the scattered particle in the HRS dipole magnets is known. It has been shown^{42,55} that with the inclusion of all the realistic effects neglected in the simplistic assumptions above, the full expressions for $D_{nn'}$ and $P(\theta)$ are

$$D_{nn'} = \frac{1}{P_n \langle \cos \phi \rangle^2 \langle \cos \chi \rangle} \left[\frac{P_n^+(FP) - P_n^-(FP)}{2} + P_n A_n \langle \cos \phi \rangle \frac{P_n^+(FP) + P_n^-(FP)}{2} \right]$$

and

$$P(\theta) = \frac{1}{P_n \langle \cos \phi \rangle^2 \langle \cos \chi \rangle} \left[\frac{P_n^+(FP) + P_n^-(FP)}{2} + P_n A_n \langle \cos \phi \rangle \frac{P_n^+(FP) - P_n^-(FP)}{2} \right] \quad (25)$$

Here $P_n^+(FP)$ and $P_n^-(FP)$ are the $\hat{n}$ -components of the polarization measured at the FPP in the focal plane coordinate system for normal and reverse beam polarization, respectively; $\chi = 269.7$ for protons, γ being the Lorentz factor for the scattered proton; and $\langle \rangle$ refers to the average over events.

In principle, $P_n^{\pm}(FP)$ can be measured from a left-right asymmetry measurement; however, since two dimensional wire chambers were used in the FPP for the detection of trajectories of the particles scattered by the carbon analyzer instead of fixed left-right counters, certain criteria for the acceptance of the scattered particles in these chambers were established in order to compute asymmetries. If θ_c and ϕ_c are the polar and the azimuthal scattering angles of particles scattered by the carbon analyzer, the distribution of the scattered particles may be written as

$$I(\theta_c, \phi_c) = I_0(\theta_c) \cdot (1 + \epsilon_n \cos \phi_c) f(\theta_c, \phi_c). \quad (26)$$

Here ϵ_n is the asymmetry in the $\hat{n}_{FP}$ -direction and $f(\theta_c, \phi_c)$ defines the criterion for the acceptance of an event by the FPP with a value

$$\begin{aligned} f(\theta_c, \phi_c) &= 1; && \text{accepted} \\ &= 0; && \text{rejected.} \end{aligned}$$

If the acceptance of events is restricted to comply with the requirement that the trajectories of the scattered particles for a given θ_c lie within the physical acceptance limits of the wire chambers for all possible values of ϕ_c , we have the so called "Cone Test" as illustrated in Fig. 24. The asymmetry ϵ_n for a given θ_c can be measured from the integrated yield of events passing the cone test criterion for θ_c . Further, if the θ_c is divided into narrow bins and the analyzing power of the carbon analyzer $A_c(\theta)$ is known for each bin, the weighted average of the polarization in the individual bins may be formed over the acceptance limits for θ_c to yield the final polarization of the scattered particle. It has been shown⁴² that for events satisfying the cone test, the $P_n(FP)$ is given by:

$$P_n(FP) = 2 \frac{\sum_{ev} I(\theta, \phi) \cos \phi A_c(\theta)}{\sum_{ev} I(\theta, \phi) A_c^2(\theta)} \quad (27)$$

Therefore, the measurement of $P_n(FP)^\pm$ with the cone acceptance test requires the accumulation of the following sums by the FPP for both normal and reverse beam polarization states

CONE-TEST

Fig. 24: Schematic representation of Cone Test.

$\phi + \pi - \text{TEST}$

Fig. 25: Schematic representation of $\pi + \phi$ test.

$$\sum_{ev} I(\theta, \phi) \cos \phi A_c(\theta) \quad \text{and} \quad \sum_{ev} I(\theta, \phi) A_c^2(\theta).$$

A second acceptance criterion with lower symmetry on the acceptance test is the $\pi + \phi$ test. This test requires that, for an event with its scattered trajectory given by (θ, ϕ) with respect to the unscattered trajectory to be accepted, the point corresponding to $(\theta, \phi + \pi)$ lie inside the acceptance limit of the wire chambers. Examples of accepted and rejected events under $\pi + \phi$ test are shown in Fig. 25. Due to the lower symmetry of this test a greater number of weighted sums is required to be accumulated in order to obtain polarization information. As shown elsewhere,⁴² this method requires the accumulation of the following five sums:

$$\begin{aligned} \sum_{ev} I(\theta, \phi) A_c(\theta) \cos(\phi), & \quad \sum_{ev} I(\theta, \phi) A_c(\theta) \sin(\phi), \\ \sum_{ev} I(\theta, \phi) A_c^2(\theta) \cos^2(\phi), & \quad \sum_{ev} I(\theta, \phi) A_c^2(\theta) \sin^2(\phi), \text{ and} \\ \sum_{ev} I(\theta, \phi) A_c^2(\theta) \cos(\phi) \sin(\phi). \end{aligned}$$

In comparing the two acceptance tests, note that the $\pi + \phi$ test is a subset of the cone test. For events that have their projected unscattered trajectories lying in the central region of the FPP wire chambers, the two tests become practically identical. It is only for the trajectories at the edges of these chambers that the difference in the selectivity of the two tests become apparent. Such events that fail the cone test may still be able to pass the $\pi + \phi$ test. For $D_{nn'}$ measurements, since the left-right asymmetry is the relevant quantity

to be measured, events that lie at the top and bottom ends of the focal plane (corresponding to lower and higher excitation, respectively) may still yield polarization information with the $\pi + \phi$ test whereas with the cone test requirement, these events may not provide any useful information. Thus in our present measurements of D_{nn} , the analysis was done with $\pi + \phi$ test requirements.

The histogramming system available in the "Q" system is not capable of incrementing histograms by fractional numbers; it can only increment the yield by unity each time an event passes the required tests for a given data variable. Since the sums for the FPP acceptance tests contain fractional numbers, histogramming proves to be unsuitable for accumulating these sums. However, these quantities in the FPP sums can be computed for all the events that lie in a given region of excitation energy and added to form the sum for that region. The limits of the excitation energy regions are specified by gate tests in the test descriptor file. However, the measurement of polarization by accumulating sums of events has the obvious pitfall that, once sums are accumulated for given regions of excitation energies and the data are replayed, the whole replay process must be repeated if the polarization measurement of an unspecified region of excitation is called for. Since in the present configuration of the "Q" system a maximum of thirteen such individual regions can be defined, an extensive study of polarization transfer properties over a large range of excitation energy is somewhat limited.

An alternate method of determining the polarization of the scattered particles for any region of excitation energy that lies within the good acceptance limits of the FPP is provided by the accumulation of the unweighted left-right histograms for normal and reverse beam polarization states. From these histograms unweighted $\hat{n}$ -type asymmetries for any region of excitation can be determined in a similar fashion as for the procedure outlined for A_n measurement earlier in this chapter. An effective inclusive analyzing power of the carbon analyzer is then determined by taking a weighted average of the analyzing powers for different θ_c bins. Thus, the polarization can be computed from the asymmetry and the effective analyzing power.

Calibration of Carbon Analyzer

To calibrate the analyzing power $A_c(\theta)$ of the carbon analyzer in the FPP, data were obtained with $\hat{n}$ -type polarized beam and the HRS set at 0° which allows the unscattered beam to enter the HRS directly. Thus, the polarization at the focal plane was the same as the incident beam polarization except for the precession in the HRS dipoles. Extremely low beam intensity was used to prevent saturation of the wire chambers. The polar angle was divided into 13 bins of 1° width from 4° to 14° and 2° widths from 14° to 20° . Since the Line C beam line polarimeter was not usable because of the very low beam intensity required for these measurements, the calibrated A_c were calculated from the accumulated sums for each bin using the beam polarization value obtained from the Line B polarimeter. In addition, the beam

Fig. 26: Inclusive analyzing Power of carbon analyzer.

polarization was monitored before and after these zero degree measurements with the Line C polarimeter; the two beam line polarimeters were found to be consistent in their results. To investigate systematic effects, data were taken with the straight through beam positioned at the upper ($\delta=-1\%$), central ($\delta=0\%$), and lower ($\delta=+1\%$) region of the focal plane. The A_c for each θ_c bin was measured and nonlinearities across the focal plane checked with a statistical accuracy of better than 0.01. Fig. 26 shows the results of the A_c calibration; the solid curve represents previous inclusive analyzing power data for ^{12}C at 300 MeV measured at SIN⁵⁶ and the dashed curve is the result of Ransome.⁵⁷ As evident from the figure our calibrated analyzing power differs from the previous results. Such a discrepancy may be indicative of the presence of false asymmetry in the polarimeter. However, when the calibrated analyzing power is used in acquiring the FPP sums, such false asymmetries are absorbed into the calibrated value.

In addition to a careful calibration of A_c , a check for instrumental asymmetries was also made. The elastic scattering from a spin zero target, from time reversal symmetry, has $P(\theta) = A_n$ and $D_{nn'} = 1$. These relations were investigated for the elastic scattering from a ^{208}Pb target at a laboratory scattering angle of 10° . The analyzing power of ^{208}Pb at this angle is large, 0.67, so the measurement is very sensitive to possible instrumental asymmetries. Measurements were carried out with the elastic peak positioned at three different positions of the focal plane; the above relations were

TABLE 1.

False asymmetry Checks for S_{nn} measurements with
elastic scattering from ^{208}Pb .

x_f (cm)	P	A_y^a	D_{nn}	S_{nn}
-15.0	0.69 ± 0.03	0.67 ± 0.05	1.01 ± 0.02	-0.01 ± 0.01
0.0	0.66 ± 0.03	0.67 ± 0.05	1.00 ± 0.02	0.00 ± 0.01
15.0	0.68 ± 0.03	0.67 ± 0.05	1.01 ± 0.02	-0.01 ± 0.01

a) From Ref. 49.

confirmed with a statistical accuracy of better than 0.01. Table I shows the results of the elastic D_{nn} and $P(\theta)$ measurements for ^{208}Pb .

Spin-Flip Results

The results of our S_{nn} measurements for the 10.23 MeV M1 state in ^{48}Ca are shown in Fig. 27. The values obtained, 0.44 ± 0.08 and 0.20 ± 0.07 at 3.5° and 5° laboratory scattering angles, respectively, are consistent with theoretical expectations as detailed in Chapter IV. The peak to background ratio for this peak is large enough for the background to introduce significant systematic uncertainty into these values. However, the data for ^{90}Zr were much more difficult to obtain. The reason is illustrated in Fig. 28 by the spectrum taken with the thick ($250\text{mg}/\text{cm}^2$) ^{90}Zr target at 3.5° laboratory scattering angle. Notice the marked difference in the instrumental background, particularly in the lower excitation region, from the spectra of Figs. 17 taken with the thin target; the deterioration of the background is attributed to the thick target. Shown in Fig. 29 are the left-right spectra accepted by the FPP for normal and reverse beam polarization states. The excitation energy region in these spectra were divided into 2-MeV bins and the FPP sums for these bins were accumulated for the measurement of spin-flip probability and the spin-flip cross section.

The values of the spin-flip probability and the spin-flip cross section for ^{90}Zr as function of excitation energy at 3.5° and 5°

Fig. 27: Spin-flip probability for the 10.23 MeV M1 transition in the $^{48}\text{Ca}(p,p')^{48}\text{Ca}$ reaction at 319 MeV.

Fig. 28: $^{90}\text{Zr}(p,p')$, $E_p = 319 \text{ MeV}$, $\theta_{\text{lab}} = 3.5^\circ$ thick target spectrum.

Fig. 29: left and right scattered spectra accumulated by the FPP for normal and reverse beam polarizations for ^{90}Zr .

Fig. 30: Spin-flip probability for the 9 MeV M1 bump in ^{90}Zr .

Fig. 31a: Spin-flip probability and spin-flip cross section as a function of excitation energy for the $^{90}\text{Zr}(\vec{p},\vec{p}')^{90}\text{Zr}^*$ reaction at $\theta_{lab} = 3.5^\circ$.

Fig. 31b: Spin-flip probability and spin-flip cross section S_{nn} as a function of excitation energy for the $^{90}\text{Zr}(p,p')$ reaction at $\theta_{\text{lab}} = 5^\circ$.

laboratory scattering angles are shown in Fig. 31a-b. These values were obtained without any subtracting any background.

When a reasonable background is subtracted from the 9 MeV bump, the values of S_{nn} obtained for the bump were 0.62 ± 0.20 and 0.20 ± 0.17 at 3.5° and 5° laboratory scattering angles, respectively. Note that these values are similar to the those for the 10.23 MeV peak in ^{48}Ca . Shown in Fig. 30 is the angular distribution of the spin-flip probability for the 9 MeV bump; theoretical calculations in the figure will be discussed in Chapter IV. In order to understand the significance of Figs. 32a-b and the systematic errors associated with the values for the 9 MeV bump, discussion of the background, in detail, is necessary. In particular, the polarization properties of the background and its effect on the measured polarization for a given region of excitation need to be well understood.

D. Background Correction

If the cross section and the polarization for a peak (or a given region of excitation energy) and the background associated with it are denoted by σ_p and P_p , and σ_b and P_b , respectively, the measured polarization for this region may be expressed as:

$$P_t^\pm = \frac{\sigma_p^\pm}{\sigma_t^\pm} \cdot P_p^\pm + \frac{\sigma_b^\pm}{\sigma_t^\pm} \cdot P_b^\pm \quad (28)$$

where $\sigma_t = \sigma_p + \sigma_b$, is the total cross section of the region. Furthermore, if we define a signal to noise ratio as the peak to background cross section

$$r^{\pm} = \frac{\sigma_p^{\pm}}{\sigma_b^{\pm}},$$

then, from Eq. (28), we get

$$P_p^{\pm} = \frac{1}{r^{\pm}} \{ (1 + r^{\pm}) P_t^{\pm} - P_b^{\pm} \}. \quad (29)$$

Hence, extraction of the polarization for a background contaminated peak requires the measurement of P_t , P_b , and r . But, experimentally, separate measurements of P_t and P_b are impossible since they occupy the same region of excitation energy. However, one can approximate the P_b for a given peak to be the same as the polarization in the immediate vicinity of the peak. To accomplish this, tests were defined in the test descriptor file to measure spin-flip probability for the left and right hand side of the peak of interest; P_b was taken to be the unweighted average of the two regions:

$$P_b^{\pm} = \frac{1}{2} (P_{b, \text{left}}^{\pm} + P_{b, \text{right}}^{\pm}). \quad (30)$$

In general the ratios $r^{\pm}$ are not equal to each other; thus, the effect of background correction for normal and reverse focal plane polarizations are different. However, for excitation regions with $A_n = 0$ one has $\sigma^+ = \sigma^-$ and therefore, $r^+ = r^-$. This implies that the focal plane polarizations $P_n^+(\text{FP})$ and $P_n^-(\text{FP})$ are corrected by the same amount. In such cases, the background correction technique outlined above can be directly applied to D_{nn} or S_{nn} . Thus, the background corrected spin-flip probability can be written as:

$$S_p = \frac{1}{r} \{ (1 + r) S_t - S_b \} \quad (31)$$

In the low momentum transfer region, the A_n for the background and the peaks of interest in the present study was found to be close to zero; this makes Eq. (31) quite suitable for the present work.

Such a method of background correction has little effect on strongly excited peaks ($r \gg 1$) such as for ^{48}Ca as is evident from Fig. 15. Here, for the 10.23 MeV M1 peak, the ratio r was 4.5 ± 0.1 and 3.0 ± 0.09 at 3.5° and 5.0° , respectively. The value of S_t was measured for the excitation energy region from 10.1 to 10.4 MeV; S_b was measured from 9 to 10 MeV and from 10.5 to 11.5 MeV to obtain the left and right background spin-flip probabilities, respectively. Summarized in Table II are the background corrected results for the M1 peak in ^{48}Ca .

TABLE II.

Summary of the background correction for the S_{nn} results
for ^{48}Ca and ^{90}Zr

	$(^{48}\text{Ca}(10.23 \text{ MeV}; \text{M1}))$		$^{90}\text{Zr}(8.98 \text{ MeV}; \text{"M1 Bump"})$	
	3.5°	5.0°	3.5°	5.0°
r	4.50 ± 0.20	3.00 ± 0.15	0.23 ± 0.05	0.24 ± 0.05
S_t	0.38 ± 0.07	0.17 ± 0.06	0.20 ± 0.02	0.14 ± 0.02
$S_{b,l}$	0.13 ± 0.10	0.08 ± 0.08	0.06 ± 0.02	0.08 ± 0.02
$S_{b,r}$	0.12 ± 0.10	0.08 ± 0.08	0.13 ± 0.02	0.16 ± 0.02
S_b	0.13 ± 0.10	0.08 ± 0.08	0.10 ± 0.02	0.12 ± 0.02
S_p	0.44 ± 0.08	0.20 ± 0.07	0.62 ± 0.10	0.20 ± 0.10

The determination of background polarization with the left-right averaging method is most applicable for narrow peaks with clearly defined background regions adjacent to the peak; the implicit assumption that the polarization of the background varies linearly from one side of the peak to the other is justifiable in such cases. However, for peaks that are several MeV wide as in ^{90}Zr , with a very large background, a more thorough understanding of the background spin-flip probability is required. In order to achieve this, the two dimensional plot of θ_k versus x_f was divided into twelve different region as shown in Fig. 12. The focal plane plane polarization for each of the regions was measured; the results obtained are listed in Table III. As expected, the $S_{nn'}$ for regions 1 - 4 was zero within the measured statistical accuracy; these regions are populated mostly by events arising from the elastic tail (see section II.I) and hence have zero spin-flip probability. Regions 5 - 8 showed mostly non zero $S_{nn'}$, except for x_f corresponding to the low excitation region of the focal plane. Note that these regions were the only regions accepted as good events; they correspond to events that have acceptable trajectories. The $S_{nn'}$ values obtained in regions 9 - 12 show a structure very similar to that in regions 5 - 8; however, the magnitudes of the $S_{nn'}$ values in these regions are diluted compared to those for the regions 5 - 8. Regions 9 - 12 contain many good inelastic events, but they are highly contaminated by background events having no spin-flip probability associated with them. Since the boundary between region 5 - 8 and 9 - 12 is rather unclear, it is quite possible that the so called good inelastic regions, 5 - 8, also contain such rescattered

TABLE III.

Effect of instrumental background on spin-flip probability. The focal plane was divided into 12 regions with BOX 4-6 and GATE 1-4; the values here are for the different regions for ^{90}Zr at 3.50_{lab} .

BOX		4		5		6	
GATE	Region	S_{nn}	Region	S_{nn}	Region	S_{nn}	
1	1	-0.003 ± 0.023	5	0.054 ± 0.043	9	0.072 ± 0.064	
2	2	0.045 ± 0.062	6	0.179 ± 0.041	10	0.138 ± 0.060	
3	3	0.043 ± 0.065	7	0.163 ± 0.041	11	0.154 ± 0.052	
4	4	0.062 ± 0.081	8	0.159 ± 0.052	12	0.152 ± 0.060	

elastic events as contaminants. These observations are consistent with the assumption that whatever instrumental background is still present in the spectrum of Fig. 28 derived from regions 5 - 8 is non-spin-flip in nature. Thus, substituting $S_b = 0$ in Eq. (30), we obtain

$$S_p = \left(1 + \frac{1}{r}\right) \cdot S_t = \frac{\sigma_t}{\sigma_p} \cdot S_p \quad (32)$$

Therefore, the presence of non-spin-flip instrumental background merely dilutes the measured value of spin-flip probability.

From our analysis that follows, we find that $S_{nn} = 0.62 \pm 0.20$ and 0.20 ± 0.17 for 3.5° and 5° , respectively, for the M1 bump in ^{90}Zr . The S_{nn} value for the M1 bump is largely affected by the correction for the background spin-flip probability. At 3.5° the ratio σ_p/σ_t for the 8 to 10 MeV region is $r = 0.23$ and the total spin-flip probability $S_t = 0.20 \pm 0.02$. If we choose $S_b = 0$, this implies $S_p = 1.0 \pm 0.08$. On the other hand, if S_b is assumed to be an unweighted average of the spin-flip probabilities of the regions 6 to 8 and 10 to 12 MeV, then $S_b = 0.10 \pm 0.05$. Thus, the S_{nn} value for the M1 bump using the left-right background correction method becomes 0.59 ± 0.10 . However, examination of the spectrum in Fig. 28 indicates that the contribution to the background in the 8 to 10 MeV is probably more due to the high excitation region than that of the lower excitation. Although, accurate weighing factors for the contribution to the background from the high and low excitation have to await theoretical calculations for this continuum region, from an empirical observation it is reasonable

to assume a contribution of about 2/3 from the high excitation side and about 1/3 from the lower side; thus a weighted average value for $S_b = 0.13 \pm 0.05$. This results in a value of $S_p = 0.47 \pm 0.10$.

Using the histogram method to compute the polarization, however, the background is treated entirely differently. Here the background subtraction is done from the yields of the left-right histograms of Fig. 29 for normal and reverse beam polarization; in this case the uncertainty shifts from determining S_b to choosing the proper background line in these spectra. After subtracting a background as shown in Fig. 29 the value at 3.5° is found to be $S_{nn} = 0.60 \pm 0.20$. Here the uncertainty due to the ambiguity in choosing a background has been included in the uncertainty of the quoted value; with variations in the background the S_{nn} value obtained from the background subtracted yields ranged from 0.49 to 0.72 with a statistical uncertainty of about 0.08. Thus, the S_{nn} value obtained with the FPP sums method and the histogram yield method are found to be consistent. The 5° results were arrived at with a similar background analysis.

E. Spin Flip Cross Section

In the presence of instrumental background whose origin can be conclusively traced to non-spin-flip processes, a better observable to measure is the "spin-flip cross section". We define the spin-flip cross section as the product of the differential cross section and the spin-flip probability, $\sigma.S$. Following the background correction

technique developed in the previous section, for zero background spin-flip probability, then, we rewrite Eq. (32) to obtain

$$\sigma_t S_t = \sigma_p S_p. \quad (33)$$

The measured spin-flip cross section is, therefore, the same as the spin-flip cross section of the peak. This prescription clearly demonstrates the quantity σS as an unambiguous means of examining the excitation of spin-flip strength in spectra contaminated with a significant amount of non-spin-flip background.

Angular distributions for the spin-flip cross section for the regions of excitation energy 6-8, 8-10, 10-12, 12-14, 14-16, and 16-18 MeV were also obtained; these are shown in Fig. 32a-c. The angular distributions were obtained from the cross section angular distributions obtained by summing the yields of the spectra taken for the spin-flip measurements (e.g. Fig. 28). Note that for cross section measurements the total angle of acceptance of the HRS was divided into four bins of 0.5 degree each whereas for the S_{nn} measurements the acceptance was divided into two bins. The spin-flip cross sections were obtained by assuming the spin-flip probability to be the same for the corresponding pair of bins in cross section. A conservative systematic error has been added to the data in Figs. 34 due to this assumption. Calculations were performed in order to decompose the components of the spin excitation multipoles in this continuum region. These calculations represented by the theoretical curves in these figures will be discussed in detail in Chapter V.

Fig. 32a

Fig. 32b

Fig. 32c

In addition to the spectra shown in Fig. 28, several other spectra were taken at 3.5° at a later time when the beam conditions were different, to investigate the possibility of systematic errors. At 319 MeV incident proton energy, the HRS is capable of obtaining polarimeter spectra of a 13 MeV wide range of excitation energy; the HRS dipole magnet fields were adjusted to obtain spectra with 9, 12, 15, and 20 MeV of excitation energy placed at the center of the focal plane. Thus, data were obtained for the total spin-flip probability and the spin-flip cross section for the range of excitation energy 5 to 25 MeV at 3.5° . The data obtained with all the different dipole field settings are included in Fig. 31a. Lack of time prevented a similar set of measurements at 5° ; only one spectrum with 12 MeV of excitation energy at the center of the focal plane was obtained.

IV. THEORETICAL ANALYSIS

This chapter will be devoted to the theoretical analysis and interpretation of the data for the experimental observables presented in the preceeding chapter. Calculations performed in the DWIA formalism with a microscopic model of the nucleus will be presented and compared to the data. In addition, calculation of coulomb excitation of E1 states in a coupled channels framework will be presented. Comparison to available data with other probes will also be discussed.

A. Formalism

The theoretical description of the scattering of a projectile nucleon from a system of nucleons such as in a nucleus involves the determination of the nucleon-nucleus scattering amplitude (T-Matrix); once the T-matrix is determined, construction of various observables from the T-matrix is straightforward. A formal representation of the T-matrix can be obtained from the solutions of the Schrödinger equation for the total hamiltonian of the system. If the hamiltonian H , is expressed as

$$H = H_0 + V$$

with

$$H_0 = U - \frac{\hbar^2}{2m} \nabla^2$$

such that U represents the nuclear potential energy (the elastic channel) and V is the interaction energy between the projectile and the nucleus, then V can be considered as a perturbation on H_0 . If the solutions for the unperturbed hamiltonian H_0 are known, then the scattering amplitude for the nuclear transition due to V can be obtained with the Distorted Wave Born Approximation (DWBA).

The nuclear potential energy can be expressed in the standard optical model of the nucleus as:

$$U(r) = Vf(r, r_v, a_v) + iWf(r, r_w, a_w) + \left(\frac{\hbar}{m_\pi c}\right)^2 [V_{so}g(r, r_{so}, a_{so}) + iW_{so}g(r, r_{wso}, a_{wso})] \quad (34)$$

Here $f(r)$ is the standard Woods-Saxon function and $g(r) = \frac{1}{r} \frac{d}{dr} f(r)$. If the parameters of $U(r)$ are chosen so that the solutions $\chi^{(+)}$ and $\chi^{(-)}$ of H_0 satisfying outgoing and incoming boundary conditions, respectively, describe the elastic scattering, then the DWBA T-matrix for the nuclear transition from an initial state $|i\rangle$ to $|f\rangle$ is expressed as:⁵⁸

$$T_{fi} = \int d\vec{r}_p D^*(\vec{r}) F(\vec{r}_p) \quad (35)$$

Here $\vec{r} = \vec{r}_i - \vec{r}_p$ is the target nucleon and projectile separation, and k_i and k_f are the initial and final wave numbers, respectively. The set

of quantum numbers $[s_1 m_1 t_1^z; s_f m_f t_f^z]$ and $[J_1 M_1 T_1^z; J_f M_f T_f^z]$ refer to the initial and final states of the projectile and target respectively. The function $D^*(\vec{r})$ is a distortion function given by the product of the distorted waves due to the optical potential

$$D^*(\vec{r}) = \chi^{(-)*}(\vec{k}_f', \vec{r}) \chi^{(+)}(\vec{k}_1, \vec{r}). \quad (36)$$

The function $F(\vec{r}_p)$ is referred to as the nuclear form factor and is given by the nuclear transition matrix element

$$F(\vec{r}_p) = \langle f | V | i \rangle. \quad (37)$$

The DWBA differential cross section is given by^{58,59}

$$\begin{aligned} \frac{d\sigma}{d\Omega} &= \left(\frac{\mu}{2\pi\hbar} \right)^2 \frac{k_f}{k_1} \frac{1}{(2J_1 + 1)(2s_1 + 1)} \sum_{M_1 M_f m_1 m_f} |T_{fi}|^2 \\ &= \sum_{m_1 m_f} \sigma_{m_1 m_f}. \end{aligned} \quad (38)$$

Here $\sigma_{m_1 m_f}$ is the partial cross section for scattering of the projectile in the initial state $|s_1 m_1\rangle$ to final state $|s_f m_f\rangle$. The spin-flip probability S_{nn} can be defined from the partial cross sections with the relation

$$S_{nn} = \frac{\sum_{m_1 m_f; m_1 \neq m_f} \sigma_{m_1 m_f}}{\sum_{m_1 m_f} \sigma_{m_1 m_f}} \quad (39)$$

Furthermore, denoting the spin projections of spin-up and spin-down projectiles as + and - respectively, the vector analyzing power A_n can be expressed as:

$$A_n = \frac{(\sigma_{++} + \sigma_{+-}) - (\sigma_{-+} + \sigma_{--})}{(\sigma_{++} + \sigma_{+-}) + (\sigma_{-+} + \sigma_{--})} \quad (40)$$

The form factor Eq. (37) is dependent on the form of the interaction and the structure of the particular nucleus considered. It may be viewed as a transition potential giving rise to the inelastic event. A convenient way of evaluating the integral in Eq. (34) is to expand this potential in terms of multipoles corresponding to definite orbital, spin, and total angular momentum transfer (LSJ). In addition, if O_T denotes a tensor of rank T (isospin transfer), such that $O_0 = 1$ for $T = 0$ isoscalar transitions and $O_1 = \tau_p$ for $T = 1$ isovector transitions, isospin dependence can be incorporated in the interaction by expressing V as:

$$V = \sum_T V^T O_T \quad (41)$$

with the static multipole expansion for V^T

$$V^T = \sum_{l,LSJM} (-)^{J-M} V_{LSJ,M}^T(r_p, r_1) T_{LSJ,-M}(\theta\phi). \quad (42)$$

Here T is the spin-angle tensor given by

$$T_{LSJ,M} = \sum_{m\lambda} \langle LSm\lambda | JM \rangle i^L Y_{Lm}(\theta\phi) S_{S,\lambda} \quad (43)$$

and S_S is a tensor of rank S with $S_0 = 1$ and $S_1 = \sigma_p$. Then the nuclear form factor can be expressed⁶⁰ as a sum of partial form factors

$$\begin{aligned} F(r) = \sum_{LSJT} i^{-L} G_{LSJ}^T(r) Y_{LM}^*(\theta\phi) (-)^{S_f - m_f} \langle s_i s_f m_i, -m_f | s\lambda \rangle \\ \times \langle J_i T_i M_i | J_f T_f \rangle \langle LSM\lambda | JM \rangle \\ \times \langle T_i T_i^z | T_f T_f^z \rangle \langle t_i T_i^z | t_i T_i^z \rangle \end{aligned} \quad (44)$$

The partial form factor G is given by

$$G_{LSJ}^T(r) = (-)^S \sqrt{2(2T+1)} \langle J_f T_f || V_{LSJ}^T(r) || J_i T_i \rangle \quad (45)$$

In general, the effective interaction V is a sum of two-body interactions between nucleon pairs; thus, the scattering amplitude may be expanded as a multiple scattering series of the projectile nucleon with the target nucleons. If the series is approximated as a single scattering only, the effective interaction reduces to:

$$V = \sum V_{ip}(1 - P_{ip}), \quad (46)$$

where V_{ip} is the effective interaction of the projectile with a nucleon inside the nucleus. Knock-on exchange contributions are included in the effective interaction with the help of the space exchange operator P_{ip} in Eq. (36) that exchanges the nucleons i and p . The effective

TABLE IV.

Strengths of the t-matrix in the Love-Franey interaction. The ranges are in fermi; the t 's are in MeV except for the tensor components which are in MeV/fm².

Real t-Matrix

Ranges: 0.15	0.25	0.40	0.55	0.70	1.40
t_0	1861.79	200.72	-277.58		
σ	-1945.77	893.67	-148.52		
t_T	985.28	-541.35	94.83		
$t_{\sigma T}$	-268.31	210.08	-21.63		3.50
t_{LS}	-4080.50	-153.39	11.62		
$t_{LS T}$	1824.80	-294.80	24.39		
t_T 20959.45	140.48	-89.40		2.42	
$t_{T T}$ -23981.85	3901.82	-95.18		20.21	

Imaginary t-Matrix

Ranges: 0.15	0.25	0.40	0.55	0.70	1.40
t_0	-807.72	-288.86	-28.76		
t_σ	-202.33	90.35	-27.90		
t_T	-1026.73	290.60	12.83		
$t_{\sigma T}$	-1188.22	260.99	-19.10		
t_{LS}	2184.13	-167.87	25.04		
$t_{LS T}$	-961.53	-8.62	-4.90		
t_T -37823.50	2492.65	-114.45		1.77	
$t_{T T}$ -43901.50	3292.38	-127.75		1.88	

TABLE V.

Absolute moments of the potentials in the Love-Franey interaction at 325 MeV obtained by the J_0 moment for the central force and J_2 moments for the spin-orbit and tensor forces.

force	moment (MeV-fm ³)	force	moment (MeV-fm ⁵)
V_0	259.2	V_{LS}	404.0
V_τ	18.2	$V_{LS\tau}$	86.7
V_σ	56.2	V_T	126.0
$V_{\sigma\tau}$	148.9	$V_{T\tau}$	2563.0

direct nucleon-nucleon interaction V_{ip} is usually expressed^{61,62,64} in the form:

$$V_{ip} = V^C(r) + V^{LS}(r)\vec{L}\cdot\vec{S} + V^T(r) S_{ip} \quad (47)$$

where

$$V^C = V_O^C(r) + V_\tau^C(r)\vec{\tau}_1\cdot\vec{\tau}_p + (V_\sigma^C(r) + V_{\sigma\tau}^C(r)\vec{\tau}_1\cdot\vec{\tau}_p)\vec{\sigma}_1\cdot\vec{\sigma}_p,$$

$$V^{LS} = V_O^{LS}(r) + V_\tau^{LS}(r)\vec{\tau}_1\cdot\vec{\tau}_p,$$

$$V^T = V_O^T(r) + V_\tau^T(r)\vec{\tau}_1\cdot\vec{\tau}_p,$$

and S_{ip} is the usual tensor operator given by

$$S_{ip} = 3(\vec{\sigma}_1\cdot\hat{r})(\vec{\sigma}_2\cdot\hat{r}) - \vec{\sigma}_1\cdot\vec{\sigma}_2.$$

In these equations, C, LS, and T denote central, spin-orbit, and tensor, r is the distance between the nucleons, $\vec{\sigma}$ and $\vec{\tau}$ are the Pauli spin and isospin operators, and $\vec{S} = (\vec{s}_1 + \vec{s}_p)$ is the total spin of the two interacting particles. Also, V_O , V_τ , and $V_{\sigma\tau}$ correspond to the transfer of one unit of spin, isospin, and both spin and isospin, respectively, in the direct transition amplitude.

If the kinetic energy of the projectile is much higher than the binding energy of the nucleon in the nucleus, V_{ip} may be reasonably approximated as the free nucleon-nucleon interaction (t_{NN}^0). Such an approximation, in this context, is known as the Distorted Wave Impulse Approximation (DWIA). In the DWIA calculations reported here, the effective interaction of Love and Franey⁶³ (LF) was used, in which a local finite range form of V_{ip} is assumed. They obtain individual

components of the interaction from empirically determined NN scattering amplitudes by parametrizing V_{ip} until its antisymmetrized matrix elements match those of the on-shell t_{NN} . The radial parts of the potentials were taken to be sums of Yukawa forms for central and spin-orbit parts; for the tensor component the sum consists of r^2 times the Yukawa terms. Table IV show the magnitudes of the t-matrices along with the range of corresponding Yukawa terms used in the calculations for $E_p = 319$ MeV in the $\sigma\tau$ representation. The J_0 moments for the central and the J_2 moments for the spin orbit and tensor forces for the V 's of Eq. (47) are given in Table V.

Spin Excitations

It is important to note that M1 transitions, or unnatural parity ($\Delta\pi = (-)^{J+1}$) transitions in general, involve one unit of spin transfer ($S = 1$) and can be mediated only by the spin dependent components of the effective interaction. In addition, for small q , the central part of the effective interaction dominates over other parts. Thus, for spin excitations in the low q region, we can approximate the effective interaction of Eq. (47) by the central spin-dependent pieces only:

$$V_{ip} = (V_\sigma + V_{\sigma\tau} \vec{\tau}_i \cdot \vec{\tau}_p) \vec{\sigma}_i \cdot \vec{\sigma}_p \quad (48)$$

In addition to the dependence on quantum numbers such as spin and isospin, the two-body interaction V_{ip} also depends upon the inter-nucleon separation $|\vec{r}_i - \vec{r}|$; following the treatment of Ref. 60, this dependence $v(|\vec{r}_i - \vec{r}|)$ may be expanded in terms of multipoles

$$v(|\vec{r}_1 - \vec{r}|) = \sum_{LM} v_L(r_1, r) Y_{LM}(\theta_1 \phi_1) Y_{LM}^*(\theta \phi). \quad (49)$$

If choose to define O_T and S_S such that $O_0 = 1$, $O_1 = \tau_1$, and $S_0 = 1$, $S_1 = \sigma_1$, then the multipole components (see Eq. 42) of the interaction for spin excitations becomes

$$v_{L1J,M}^T = \sum_i v_L(r_i, r) T_{L1J,M}(\theta_i \phi_i, \sigma_i) \times \begin{cases} v_\sigma & \text{for } T = 0 \\ v_{\sigma\tau\tau_1} & \text{for } T = 1 \end{cases} \quad (50)$$

and the partial form factor $G(r)$ reduces to

$$G_{L1J}^T = \langle J_f T_f | \sum_i [v_\sigma \delta_{T,0} + \sqrt{3} v_{\sigma\tau\tau_1} \delta_{T,1}] v_L T_{L1J} | J_i T_i \rangle. \quad (51)$$

Thus the determination of the T-matrix requires evaluating the reduced matrix elements of the isoscalar and isovector one body tensor operators of Eq. (51). Defining the transition densities for spin excitations as

$$\rho_{JL}^{ST}(r) = \langle J_f T_f | \sum_i v_L(r_i, r) T_{LSJ,M}(\theta_i \phi_i, \sigma_i) O_T | J_i T_i \rangle, \quad (52)$$

the form factor $G(q)$ in the momentum representation may be rewritten in terms of these transition densities

$$G_{L1J}^T(q) = v_\sigma \rho_{JL}^{S0}(q) + \sqrt{3} v_{\sigma\tau} \rho_{JL}^{S1}(q). \quad (53)$$

Particle-Hole Excitation

In a nucleus with a doubly closed shell structure excited states can be achieved by the excitation of a particle from a closed shell j_h to an unoccupied shell j_p , thus forming a particle-hole pair (j_p, j_h) . If we define $a_{j_p m_p t_p}^+$ as a creation operator for a particle in shell model orbital j, m from the closed shell ground state $|0\rangle$, the particle-hole state may then be defined as

$$|(j_p j_h^{-1})_{JT}\rangle = A_{JT}^+(j_p j_h) |0\rangle \quad (54)$$

where

$$A_{JT}^+(j_p j_h) = \sum_{m_p m_h t_p t_h} (-)^{j_p - m_p - t_p + 1/2} \langle j_p j_h m_p - m_h | J, -M \rangle \langle 1/2, 1/2, t_p, -t_h \rangle \\ \times a_{j_p m_p t_p}^+ a_{j_h m_h t_h}$$

Several such configurations may individually couple to result in the excited state $|JT\rangle$; the excited state may then be expressed as a coherent superposition of all such particle-hole configurations

$$|JT\rangle = \sum_{j_p j_h} Z_{JT}(j_p j_h) |(j_p j_h^{-1})_{JT}\rangle \quad (55)$$

where the spectroscopic coefficients Z_{JT} are proportional to the reduced matrix elements of the operator A^+ and are given by

$$Z_{JT}(j_p j_h) = \frac{\begin{Bmatrix} T_i T_f \\ T_i^z T_f^z \end{Bmatrix} \langle J_f T_f || A_{JT}^+ || J_i T_i \rangle}{\hat{J}_i \hat{J}_f} \quad (56)$$

where $\hat{x} = (2x + 1)^{1/2}$.

In the case of a closed shell nucleus with a 0^+ ground state, M1 transitions can occur to the excited 1^+ state by this particle hole excitation mechanism. The M1 selection rule (see Eq. 6) is satisfied when a nucleon in the closed shell orbital $j_h = \ell + \frac{1}{2}$ is transferred to its spin-orbit partner $j_p = \ell - \frac{1}{2}$ resulting in LSJ = 011 or 211. For such a transition the transition density can be separated into a radial part $I(r)$, and a spin-angle reduced matrix element. This matrix element, which represents a measure of the strength of the particle-hole excitation, can be shown to be a function of ℓ only; the transition density in such a case reduces to

$$\rho_{01}^{ST}(r) = -2 \left[\frac{\ell(\ell + 1)}{3\pi(2\ell + 1)} \right]^{1/2} I(r) \quad (57).$$

which is approximately $(8\ell/3)^{1/2}$ times the single particle matrix element,⁶⁰ for large values of ℓ . Thus, for heavy nuclei where the ℓ value of the spin unsaturated orbitals is usually large, the M1 transition is expected to exhibit a large cross section. Apart from statistical factors and sums over spin substates, a qualitative expression for the differential cross section in the (p,p') reaction for M1 transitions in the spin unsaturated shell may be written as:⁶⁴

$$\left(\frac{d\sigma}{d\Omega} \right)_{p,p'}(M1) \propto |D(q) [C_0 v_{\sigma} \rho_{10}^{s0}(q) + \sqrt{3} C_1 v_{\sigma\tau} \rho_{10}^{s1}(q)]|^2 \quad (58)$$

with C_0 and C_1 as the isospin coupling Clebsch-Gordan coefficient of Eq. (44) given by

$$C_T = \langle T_i T T_i^Z \lambda' | T_f T_f^Z \rangle \langle t_i T t_f^Z \lambda' | t_i t_i^Z \rangle. \quad (59)$$

Here the form factor and the transition densities are expressed in the momentum transfer representation which is obtained by the Bessel transform of the corresponding radial functions.

Typical DWIA cross section angular distributions for spin excitations to the 0^- , 1^+ , 2^- , and 3^+ states from the 0^+ ground state of ^{90}Zr for proton inelastic scattering at 319 MeV are illustrated in Fig. 33. Note that only the 1^+ angular distribution rises sharply as θ decreases toward 0° . The LF interaction outlined above was used (see Table IV). Simple particle-hole configurations, as indicated in Fig. 33, were used for these states. The exchange contributions were calculated exactly. The calculations were performed with the computer code DW81⁶⁵ which is a modified version of the original DWBA70 program. The radial integrations were carried out from the origin to 12 fm in steps of 0.06 fm. A total of 170 partial waves were used; for the exchange amplitudes 55 partial waves were used. The optical parameters used in the calculations are those of table VI and will be discussed later in this chapter.

Fig. 33: Typical DWIA calculations for simple particle-hole spin excitation cross sections.

Comparison with (e,e') and (p,n)

Inelastic scattering cross sections observed in various reactions are commonly characterized in terms of the reduced electromagnetic transition probability in order to facilitate comparison of the observed strengths. In the case of M1 transitions, the reduced transition probability $B(M1)$ is proportional to the square of the reduced matrix element of the magnetic dipole operator between the initial and final states given by Eq. (5).

The total excitation strength for a given multipole operator is often expressed as an energy weighted sum of the reduced transition probabilities over all possible final states. When the closure properties of a set of nuclear states can be used to express this total strength in a closed form in a model independent way, a sum rule is said to exist. For M1 transitions we can define an energy weighted sum $S(M1)$ as:

$$S(M1) = \sum_f (E_f - E_i) B(M1, i \rightarrow f). \quad (60)$$

Sum rules for different reactions and the implications thereof will be discussed later in this chapter.

Excitation of the nucleus by electromagnetic probes such as electrons provides a direct measurement of the electromagnetic transition probabilities of the nucleus. As discussed in the introduction, the interaction may be characterized by the exchange of virtual photons; hence, a longitudinal and a transverse component of the interaction may be defined in accordance with the polarization of

the virtual photon. The form factors associated with such interactions contain the nuclear structure information and are proportional to the corresponding reduced transition probabilities. The longitudinal form factor is sensitive only to the charge distribution of the nucleus and thus is referred to as the coulomb form factor. On the other hand the transverse form factor is a measure of the magnetic properties such as the spin and orbital current distribution of the nucleus. Only the transverse component may give rise to spin excitation. In addition, the transverse part can be separated into its electric and magnetic components. In plane wave Born approximation, the differential cross section for the J^{th} multipole in the (e, e') reaction is given by^{66,69}

$$\begin{aligned} \left(\frac{d\sigma}{d\Omega}\right)_{e,e'} = \alpha \left(\frac{q}{k_1}\right)^2 \frac{4\pi(J+1)}{J[(2J+1)!!]^2} \left\{ \frac{J}{J+1} B(CJ, q) v_L(\theta) \right. \\ \left. + [B(EJ, q) + B(MJ, q)] v_T(\theta) \right\}. \end{aligned} \quad (61)$$

Here $v_L(\theta)$ and $v_T(\theta)$ are kinematical factors for the longitudinal and transverse components, respectively. In back angle electron scattering i.e. $\theta = 180^\circ$, one has $v_L(\pi) = 0$; only the transverse form factors survive, thus selectively populating spin excitations. If momentum transfer is chosen to be small, lower multipolarity spin excitations such as M1 transitions are predominantly excited. In the small q region, assuming the $L = 2$ transition density is small, the back angle electron scattering cross section is then given by

$$\left(\frac{d\sigma}{d\Omega}\right)_{e,e'}(M1) = \frac{2\pi\alpha}{9} \left(\frac{q}{k_1}\right)^2 B(M1\uparrow; q). \quad (62)$$

If an orbital current transition density $\rho_{JL}^{\ell T}$ is defined analogous to the spin transition density in Eq. (52), the $B(M1\uparrow; q)$ in Eq. (62) can be shown⁶⁹ to be

$$B(M1, q)\uparrow = \sum_T (-)^T \langle T_1 T_1^Z T_0 | T_f T_f^Z \rangle \left[\frac{1}{2} g_s^T \rho_{10}^{sT}(q) + g_\ell^T \rho_{10}^{\ell T}(q) \right]^2 \quad (63)$$

The superscript T refers to isovector or isoscalar terms as usual with $g^0 = g^p + g^n$ and $g^1 = g^p - g^n$. From the above equations for the differential cross sections, it is clear that for electron scattering the cross section is a direct measure of the $B(M1)$ value. The evaluation of the sum rule (Eq. 60) for electron scattering of M1 states, however, is model dependent, since it depends on the nuclear structure of the states involved, and particularly on ground state correlations. For example, Auerbach's⁷⁰ calculations for ^{60}Ni and ^{90}Zr show a reduction of 20% in the sum rule prediction due to ground state correlations.

With nucleon probes, because of hadronic forces and exchange contributions due to indistinguishability from the target nucleons, the projectile-target interaction is complex and the excitation strength is no longer a fundamental property of the nucleus alone. The excitation strength of M1 states with nucleon probes will depend on the spin-spin and spin-isospin effective interaction between the projectile and the nucleus. Comparing Eq. 58 with Eq. 63, it is evident that the cross

section in the (p,p') reaction will be proportional to that in the (e,e') reaction provided 1) the V_0 component of the proton-nucleus effective interaction is negligible, and 2) the orbital current transition density for the (e,e') reaction is negligible. The orbital current transition density for a mostly neutron transition is expected to be small. The interpretation of the (p,p') cross section in terms of $B(M1)$ then depends entirely on the knowledge of the effective nucleon-nucleus interaction. In addition, the comments made above about nuclear structure effects in evaluating a sum rule are valid for the (p,p') reaction also.

In the case of (p,n) charge exchange reactions, the formalism developed earlier in this chapter still applies. However, spin-flip excitations in (p,n) reactions must necessarily be associated with isospin-flip. Thus, in Eq. (58) only the $V_{0\tau}$ piece of the effective interaction contributes, resulting in Gamow-Teller (GT) excitations discussed in chapter 1. Then the (p,n) cross section to the GT states can be expressed as:

$$\left(\frac{d\sigma}{d\Omega}\right)_{p,n}(GT) = 3 |D(q)C_1 V_{0\tau} \rho_{10}^{s1}(q)|^2. \quad (64)$$

Prediction of the (p,n) cross section to the GT states will, therefore, depend on the accuracy of the $V_{0\tau}$ effective interaction. Fortunately, however, model independent sum rules for zero degree (p,n) GT strength can be made by comparison with β -decay. The GT matrix element being the spin-isospin transition density, (Eq. 52) can be written as

$$\langle GT \rangle^2 = |\rho_{10}^{s1}(0)|^2. \quad (65)$$

The sums for the (p,n) and (n,p) GT transitions to all possible final states can be defined as:

$$S_{pn} = \sum_f |\langle GT; (p,n) \rangle|^2 \quad S_{np} = \sum_f |\langle GT; (n,p) \rangle|^2. \quad (66)$$

These are the same matrix elements that appear in the Gamow-Teller β^- and β^+ decays respectively.⁷¹ Using the commutator of the β -decay operator, a sum rule for (p,n) may be written as:^{22,72}

$$S_{pn} - S_{np} = 3(N - Z). \quad (67)$$

This sum rule is indeed model independent. To apply it to heavy nuclei, however, an additional assumption is made, that for nuclei with significant neutron excess, S_{np} is small, as expected.²² Thus, an approximate sum rule for such nuclei is given by

$$S_{pn} \gtrsim 3(N - Z). \quad (68)$$

Although similar sum rules for the (p,p') M1 strength cannot be established, charge independence of the nuclear forces suggests that the (p,p') cross section should be proportional to the cross sections of the T_0 and $T_0 + 1$ components of the GT resonance in (p,n). Recall that M1 transitions in (p,p') and the T_0 branch of the GT transition in (p,n) reactions are analogue states, where T_0 is the ground state

isospin of the target nucleus. For the purpose of comparison, a simple model is considered. Assume that the ground state consists of a completely filled $j = l + \frac{1}{2}$ shell. Furthermore, assume that the V_σ contribution to the M1 strength (Eq. 58) is negligible, that the M1 transition and the GT transitions are mediated only by the isovector component of the effective interaction, namely $V_{\sigma\tau}$, and that the distortion factors are the same for (p,p') and (p,n) reactions. A simple relation between the M1 and GT cross sections can then be obtained by comparing Eq. (58) and Eq. (64), with the ratios of the isospin coupling coefficients as a scaling factor:

$$\left(\frac{d\sigma}{d\Omega}\right)_{p,p'}(M1) = \left(\frac{C_1(p,n)}{C_1(p,p')}\right)^2 \left(\frac{d\sigma}{d\Omega}\right)_{p,n}(GT, T_0). \quad (69)$$

For (p,p') M1 transitions, $T_i = T_i^Z = T_0$; $T = 1$, $T^Z = 0$; and $T_f = T_f^Z = T_0$. The quantum numbers for GT transitions in the (p,n) reaction are given by $T_i = T_i^Z = T_0$; $T = 1$, $T^Z = -1$; and $T_f = T_0$, $T_f^Z = T_0 - 1$. Using Eq. (59), the proportionality constant in Eq. (69) reduces to

$$\begin{aligned} \left(\frac{C_1(p,n)}{C_1(p,p')}\right)^2 &= \frac{\langle T_0 1 \ T_0, -1 | T_0 - 1, T_0 - 1 \rangle \langle 1/2, 1, 1/2, -1 | 1/2, -1/2 \rangle}{\langle T_0 1 \ T_0 0 | T_0, T_0 \rangle \langle 1/2, 1, -1/2, 0 | 1/2, -1/2 \rangle} \\ &= \frac{T_0}{2} \end{aligned} \quad (70)$$

The simple model assumed here should be reasonably justifiable for the

nuclei and the incident energy chosen in the present study. Both ^{48}Ca and ^{90}Zr have neutron filled closed shells in the ground state. In addition, at 319 MeV, as seen from Table V, the ratio $|V_{\sigma\tau}/V_{\sigma}|^2 \sim 7$ in the LF effective interaction, which implies that neglecting the V_{σ} part is not a serious problem.

However, the T_0 and the $T_0 + 1$ branches of the GT states are often unresolved²² in the available (p,n) data. If the nuclear structure factors are the same, we may use the Clebsch-Gordan coefficients for the T_0 and $T_0 + 1$ GT states as a scaling factor between the M1 strength and the total GT strength:

$$\left(\frac{d\sigma}{d\Omega}\right)_{p,p'}(M1) = \frac{T_0}{2(2T_0 - 1)} \left(\frac{d\sigma}{d\Omega}\right)_{p,n}(GT; T_0, T_0 + 1). \quad (71)$$

B. Coulomb Excitation

The DWIA formalism discussed above describes excitation of the nucleus only through the nucleon-nucleus short range strong interaction. However, the electric charge of the proton also requires the long range coulomb interaction. If either the bombarding energy of the projectile is below the coulomb barrier of the nucleus, or the distance of closest approach of the projectile trajectory is greater than the strong absorption radius of the nucleus, the projectile and nucleus never approach sufficiently close to one another for a strong interaction to occur, but the target nucleus may still undergo a

transition due to the coulomb forces acting between the projectile and the nucleus. Such an excitation mechanism is referred to as "Coulomb Excitation".⁷³

Properties of Coulomb Excitation

Consider the classical picture of the scattering of a projectile of charge Z_p and velocity v from a target nucleus of charge Z_t and a strong absorption radius R . In addition to the usual optical potential of the nucleus, a repulsive coulomb potential $V_c = \frac{Z_p Z_t}{r}$, r being the relative projectile-target distance, now enters the scattering picture. The long range r^{-1} nature of the coulomb potential introduces an additional slowly varying distortion of the waves given by the coulomb phase shift $\eta \ln(2kr)$, where η is the dimensionless Sommerfeld parameter:

$$\eta = \frac{Z_p Z_t e^2}{\hbar v}$$

For a given target nucleus, at lower bombarding energies, $\eta \gg 1$. such a situation corresponds to a strong repulsive coulomb barrier that the projectile is unable to penetrate. Thus, coulomb excitation is the primary mechanism of nuclear transition and the effects of the coulomb interaction are significant over an appreciable range of scattering angle. However, at higher energies, where $\eta \ll 1$, the projectile is free to penetrate the coulomb barrier and undergo strong interaction. Two regions of interest may be defined: 1) the "nuclear region" denoted

by $d < R$, where the projectile-nuclear interaction is predominantly the short range nuclear strong interaction; and 2) the "coulomb region" characterized by $d > R$, where the effects of the nuclear forces are minimal and the long range coulomb forces dominate. Coulomb excitation becomes the dominant mechanism for nuclear transitions when the projectile is in the coulomb region; this corresponds to small scattering angles. The relation between d and the scattering angle θ may be expressed⁷⁴ as:

$$d = \frac{Z_p Z_t e^2}{2E} \left(1 + \operatorname{cosec} \frac{\theta}{2} \right)$$

$$\sim \frac{n}{k} \left(1 + \frac{2}{\theta} \right) \text{ for small } \theta \quad (72)$$

If one defines a critical angle θ_c corresponding to $d = R$, then trajectories with $\theta > \theta_c$ correspond to $d < R$. However, for $\theta < \theta_c$, $d > R$ which results in pure coulomb excitation of the nucleus. As an example proton inelastic scattering from ^{208}Pb at 200 MeV proton energy will experience coulomb excitation for scattering angles less than $\theta_c \sim 5^\circ$. In practice, however, the cut-off scattering angle between coulomb and nuclear excitation is not as well defined as the classical approximation for θ_c .

Nuclear structure effects on coulomb excitation have been ignored so far in the classical treatment presented here. With the usual static multipole moment expansion of the nuclear charge distribution, nuclear structure information can be folded into the treatment above.

The excitation of magnetic multipole transitions through coulomb excitation is forbidden due to the absence of spin coupling in the coulomb field. Consequently, one need be concerned only with electric transitions. For an electric multipole of order L , the coulomb excitation differential cross section may be written as⁷⁴

$$d\sigma_{EL} = \left(\frac{Z_p e}{\hbar v}\right)^2 \frac{1}{R(2L-1)} B(EL) \uparrow df_{EL}(\theta, \zeta) \quad (73)$$

where

$$\zeta = \frac{R E_x}{\hbar v}$$

and

$$df_{EL}(\theta, \zeta) = \frac{4\pi^2}{(2L+1)^3} \sum_M |Y_{LM}(\frac{\pi}{2}, 0)|^2 |I_{LM}(\theta, \zeta)|^2 \frac{d\Omega}{\sin^4 \frac{\theta}{2}}.$$

Thus, lower electric multipoles are more favorably excited via coulomb excitation. However, the monopole coulomb field merely deflects the trajectory of the projectile giving rise to elastic scattering; hence coulomb excitation cannot give rise to electric monopole transitions of the nucleus. All other parameters being the same, the electric dipole ($E1$) transitions thus dominate the coulomb excitation process.

Although the dependence of $d\sigma_{EL}$ on excitation energy is determined by the complicated function $f_{EL}(\theta, \zeta)$, a rough understanding of such dependence can be made from the asymptotic behaviour of $f_{EL}(\theta, \zeta)$. For large ζ , one has

$$f_{EL}(\theta, \zeta) \sim e^{-2\pi\zeta},$$

indicating an exponential decrease in the coulomb cross section with increasing excitation energy. Since the GDR, containing most of the electric dipole strength, occurs at relatively high excitation energy, coulomb excitation of the GDR, though not negligible, is considerably impeded. However, the weak isolated E1 states found scattered in the lower excitation region are strongly enhanced by coulomb excitation.

Since the present study concentrates on the small angle scattering region, in view of the foregoing discussion, estimation of the low lying E1 strength is important in order to facilitate extraction of spin-flip strength.

Method of Calculation

Calculation of coulomb excitation was performed with the coupled channels (CC) computer program ECIS.⁷⁵ A Detailed description of the calculational method may be found in Ref. 76.

A first order vibrational model of the nucleus with the ground state considered as the zero-phonon state and the excited E1 states as one-phonon states was used in the calculations. Computation of the coulomb integrals without any special techniques usually involves long range numerical integration and the inclusion of a large number of partial waves. To circumvent this rather cumbersome and slow process of evaluating integrals, ECIS uses the "local recurrence relations" of coulomb integrals to compute the coulomb functions outside the nuclear

region. If a matching radius is chosen so that the nuclear excitation has no effect beyond the chosen value, in this method, both nuclear and coulomb functions are numerically evaluated up to this radius after which the coulomb functions are evaluated using analytic forms and the recurrence relations. Such a technique enables the use of smaller limits for the radial integration than a full numerical evaluation of integrals up to infinite radius.

The absolute value of the cross section calculated by ECIS is determined by normalization of the calculated cross section to the reduced transition probability $B(E1)$ of the state under consideration. In the collective vibrational mode of the nucleus, the nuclear shape for the dipole vibration can be parametrized about the equilibrium shape as:

$$R(\theta, \phi) = R_0 \left(1 + \sum_{M=-1}^{+1} \beta_1 Y_{1M}(\theta, \phi) \right) \quad (74)$$

where β_1 is the deformation parameter. The reduced transition probability can be expressed⁷⁷ in terms of β_1 as:

$$B(E1)_{\uparrow} = \left(\frac{3}{4\pi} Z_t e R_0 \right)^2 \beta_1^2 \quad (75)$$

The deformation parameter and the $B(E1)$ strengths for E1 states are usually determined by comparing the calculated cross sections. Conversely, however, if the $B(E1)$ strengths of the E1 states are known from other independent measurements, they can be used to normalize the ECIS cross sections to predict an accurate cross section with coulomb

excitation assuming that the deformation of the charge distribution is same as the nuclear deformation.

In the nuclear resonance fluorescence (NRF) studies of Berg et al. (ref 24) the ground state decay widths Γ_0 of excited states can be measured from scattered γ -ray intensities. The radiative decay rate to the ground state from an excited E1 state can be shown⁷⁸ to be:

$$T(E1) = 1.59 \times 10^{15} E_x^3 B(E1) \text{ sec}^{-1}; \quad \Gamma_0(E1) = \hbar T(E1).$$

Using Eq.75 and Eq. 76, a relation between β_1 and Γ_0 can be established:

$$\beta_1 R_0 = \frac{7.08}{Z_t} \left[\frac{r_0}{E_x^3} \right]^{\frac{1}{2}} \quad (77)$$

where R_0 is in fm, Γ_0 is in ev, and E_x is in MeV. Values of β computed using Eq. 77 from known values of Γ_0 were used in the coulomb excitation calculations with ECIS.

C. Properties of Spin-Flip Probability

Although experimenters have been exploring nuclear structure through the measurement of σ and A_n since the advent of polarized beams, the observable S_{nn} , though a powerful tool in probing certain aspects of nuclear structure and forces, has remained relatively

unexplored due to lack of experimental techniques. The fundamental properties of S_{nn} need to be outlined in order to understand the significance of this observable. Assume a simple model for the nuclear scattering process, with only central forces and no knock-on exchange contributions or distortion of the projectile wavefunctions in both the entry and exit channels. Consider a nucleus with a 0^+ ground state, and a transition with transfer quantum numbers LSJ. Then it can be shown³² that the spin-flip probability reduces to a simple form dependent on the angular momentum coupling from the initial to the final state, but independent of the kinematics and dynamics of the reaction such as momentum transfer:

$$S_{nn}(L0J) = 0 \quad (78)$$

$$S_{nn}(L1J) = \frac{3J + 1}{2(2J + 1)} \quad J = L + 1$$

$$S_{nn}(L1J) = \frac{3J + 2}{2(J + 1)} \quad J = L - 1$$

$$S_{nn}(L1J) = \frac{1}{2} \quad J = L$$

The most striking feature that emerges immediately from the foregoing discussion is that for transitions involving no spin transfer ($S = 0$), $S_{nn} = 0$. On the other hand, for spin excitations ($S = 1$), S_{nn} is non-zero. More specifically, for the M1 transition $0^+ \rightarrow 1^+$; $L = 0$, the value of S_{nn} is $\frac{2}{3}$. This clearly establishes the credibility of S_{nn} as a signature of spin-flip transitions. However, the simplistic model assumed here (PWIA, V_C) is far from being a realistic representation of

Fig. 34: Spin-flip probability calculated with the central Love-Franey interaction for the M1 transition in ^{48}Ca .

Fig. 35: Calculations of exchange effects and spin-orbit and tensor contributions to the spin-flip probability for the M1 transition in ^{48}Ca .

the scattering process. Inclusion of non-central components of the force, distortion effects caused by the optical potential, exchange contributions, and coherent interference of higher L values (e.g. $L = 2$ for $M1$'s), tend to produce q dependent structure in the DWIA prediction of S_{nn} . This is illustrated in Figs. 34 and 35. The solid line in Fig. 34 is obtained with a PWIA calculation for a pure $L = 0$ $M1$ neutron transition in Ca whereas the dashed curve contains contributions from both $L = 0$ and $L = 2$ terms. Interference of the $L = 2$ component gives rise to an appreciable deviation in the S_{nn} value from $\frac{2}{3}$ at large q ; however, the small q region still remains reasonably unaffected. Next, the effects of distortion and inclusion of non-central forces in the interaction are investigated. These effects are summarized in the DWIA calculations shown in Fig. 35 for the same case as above. Comparison of the dotted DWIA curve of Fig. 35 with the corresponding FWIA one of Fig. 34 shows that for small q ($0 \leq q \leq 0.5$), deviations in S_{nn} from PWIA values due to distortion effects are minimal. Such a behaviour is well expected: in a single scattering approximation, small momentum transfer occurs when the projectile barely grazes the nuclear surface; thus, the effect of the optical potential is the least felt by the projectile resulting in almost plane wave scattering. Further, the inclusion of V_{LS} in the effective interaction in the small q region does not alter significantly the values for V_C alone. At larger q , however, the V_{LS} part contributes toward a significant lowering of the S_{nn} value. The V_T component of the interaction has a more pronounced effect even in the small q region than V_{LS} . By far the most significant deviations of S_{nn} from the

simple model can be attributed to the inclusion of knock-on exchange contributions of the tensor force; the spatial dependence of V_T makes it extremely sensitive to exchange effects.

Two very important deductions can be readily made from the response of S_{nn} to various components of the effective interaction:

- 1) measurement of S_{nn} as an identifying tool for M1 transitions must be made at the smallest possible laboratory scattering angle ensuring an almost model independent value of S_{nn} . Ideally, such measurements should be carried out at 0° scattering angle; unfortunately available experimental techniques preclude this possibility.
- 2) Measurement of S_{nn} angular distributions in the small q region provides a good test for the DWIA formalism. In addition, such measurements may provide a means of ascertaining the relative magnitudes of the central to non-central components of the effective interaction. In the case of a pure neutron transition, such as in ^{40}Ca , where the M1 transition is an equal admixture of isoscalar and isovector components, the angular distribution of S_{nn} does not provide sufficient information about individual pieces of the non-central t-matrix. However, a case study of S_{nn} measurements for pure isoscalar and pure isovector M1 transitions, such as in ^{12}C , may provide this additional information.

D. The 10.23 MeV State in ^{40}Ca

The M1 transition in ^{40}Ca proves to be an excellent test case for the study and understanding of such transitions. With eight excess neutrons filling up the $f_{7/2}$ shell in the pure Independent-Particle

TABLE VI.

Optical parameters used for DWIA calculations
for ^{48}Ca and ^{90}Zr at 319 MeV.

	V (MeV)	r_o (fm)	a_o (fm)	W (MeV)	r_w (fm)	a_w (fm)	r_c (fm)
^{48}Ca	7.56	1.22	0.86	15.33	1.38	0.57	1.24
^{90}Zr	8.65	1.28	0.72	18.56	1.26	0.66	1.26

	V_{so} (MeV)	r_{so} (fm)	a_{so} (fm)	W_{so} (MeV)	r_{wso} (fm)	a_{wso} (fm)
^{48}Ca	1.97	1.09	0.56	-1.76	1.09	0.56
^{90}Zr	1.34	1.10	0.56	-1.73	1.10	0.56

Model (IPM), Ca should exhibit an M1 transition with a pure neutron spin-flip from $f_{7/2}$ $f_{5/2}$. The IPM prediction for the reduced transition probability for a pure neutron $(f_{7/2}^{-1}, f_{5/2}^1)_1$ configuration is $B(M1) = 12.0 \mu_N^2$ using bare g-factors for neutrons.

A strong and narrow M1 state at 10.227 MeV of excitation energy with a strength $B(M1)_{\text{exp}} = 3.9 \pm 0.5 \mu_N^2$ was observed in the (e, e') work;⁷⁹ following this, seven additional weakly excited M1 states in the excitation energy range of 7 to 12 MeV were identified by the electron scattering group,⁸⁰ bringing the total M1 strength to $4.6 \pm 0.5 \mu_N^2$. The total observed strength is thus quenched by a factor γ ($\gamma = B(M1)_{\text{exp}}/B(M1)_{\text{theory}}$) of 0.38 from the IPM predicted value. The analogue of the M1 state ($T = 4$) was observed⁸¹ in (p, n) scattering at IUCF; 30% of the predicted T_0 , GT strength was reported to have been observed in this state. A series of inelastic proton scattering cross section measurements⁸²⁻⁸⁵ with proton energies $E_p = 45 - 200$ MeV have also reported the observation of the M1 state at $E_x \sim 10.2$ MeV. Comparison of the experimental cross section with an IPM DWIA calculation yields a quenching factor of approximately 30% for the reported (p, p') results.

DWIA calculations for $d\sigma/d\Omega$, A_n , and S_{nn} performed with the computer code DW81 with the IPM wave functions are shown as the solid curves in Figs. 16 and 30. The t-matrix of Love and Franey, as outlined in Sec. IV.A, was used in the calculations. The optical parameters used in the calculations were extrapolated from available data. Two different extrapolations were done: parameters were extrapolated from the parameters of ^{40}Ca by fitting the 300 MeV elastic

TABLE VII.

Reduced matrix elements of McGrory and Wildenthal for the 1^+ states in ^{48}Ca in a full f-p shell calculation.

E_x (MeV)	9.689		9.947	
(j_p, j_h)	$\Delta T=0$	$\Delta T=1$	$\Delta T=0$	$\Delta T=1$
$(f_{7/2}, f_{7/2})$	-0.0156	-0.0176	0.007	0.008
$(f_{7/2}, f_{5/2})$	-0.0310	-0.0340	0.165	0.184
$(p_{3/2}, p_{3/2})$	0.014	0.055	-0.006	-0.007
$(p_{3/2}, f_{5/2})$	0.023	0.026	0.020	0.022
$(p_{3/2}, p_{1/2})$	0.036	0.040	0.153	0.171
$(f_{5/2}, f_{7/2})$	-0.286	-0.319	1.261	1.410
$(f_{5/2}, p_{3/2})$	-0.023	-0.026	-0.014	-0.015
$(f_{5/2}, f_{5/2})$	0.010	0.011	-0.015	-0.017
$(p_{1/2}, p_{3/2})$	0.062	0.067	0.283	0.316
$(p_{1/2}, p_{1/2})$	0.041	0.045	0.039	0.044

scattering data⁴⁹ from TRIUMF with the optical parameters search program CUPID.⁸⁶ Secondly, an energy dependent extrapolation from 200 and 500 MeV⁸⁹ ⁴⁸Ca optical parameters was performed. The optical parameters obtained with each method were within 20% of each other; the average of the two, listed in Table VI, was used in the calculations. Knock-on exchange was calculated exactly in the computations.

Comparison of the experimental data in Figs. 16 and 30 reveals good agreement of the DWIA predictions with $d\sigma/d\Omega$ and S_{nn} ; however, agreement for A_n is only fair. The absolute cross section data is 28% of the IPM DWIA prediction, which is consistent with the quenching observed in (e,e') and the previous (p,p') results of Refs. 80,84.

Additional calculations with full f-p shell wave functions of McGrory and Wildenthal⁸⁷ (MW) were also performed and are shown as the dashed curves in Figs. 16 and 30. The MW wave functions include the $f_{7/2}$, $f_{5/2}$, $p_{3/2}$, and $p_{1/2}$ model space with ground state correlations; the spectroscopic factors for the various particle-hole pairs are shown in Table VII. The IPM and MW wave function curves in Figs. 16 and 30 show little difference in the shape of the angular distribution for all the observables. However, the magnitude of the cross section is $\sim 65\%$ of the IPM value. The observed $\frac{d\sigma}{d\Omega}$ data are thus 55% quenched relative to the MW prediction as opposed to the 72% relative to the IPM prediction. A major contribution to this reduction of the quenching with MW wave functions can be attributed to the 2p-2h ground state correlations.⁸⁷ Use of the MW wave functions results in a similar reduction in quenching of the observed M1 strength in (e,e') ; the $B(M1)$ value with the MW wave functions prediction is $7.3 \mu_N^2$ as opposed to a

value of $12 \mu_n^2$ for the IPM value. The observed quenching has been explained by non-nucleonic degrees of freedom.⁸⁰ In this calculation, the Δ -nucleon hole polarization was the major contributing factor, whereas contributions of core polarization effects accounted for the quenching to a lesser degree. The effects of meson exchange currents were found to be small. Such explanations with delta-hole mechanisms, however, remain controversial.

The shapes of A_n predictions with the MW wave functions are not appreciably different from those with the IPM wave functions. The experimental values for the 10.23 MeV state shown in Fig. 16 are mostly small and positive. Both calculations show agreement in the overall description of the data; however, neither show the possible structure in the shape of the angular distribution. These A_n data are similar to the values of A_n obtained³⁴ for the 15.11 MeV 1^+ , $T=1$ state in $^{12}\text{C}(p,p')$ with proton energies of 400 to 800 MeV. On the other hand, the 12.71 MeV 1^+ , $T=0$ state in ^{12}C yields values of A_n that are large and negative throughout the energy range. The pure neutron transition assumed here has equal amplitudes of isoscalar and isovector transition amplitudes. DWIA predictions assuming pure isoscalar transitions yield A_n values of around -0.3 in the small q region for both ^{12}C and ^{48}Ca ; this is also the case for the data for the 12.71 MeV isoscalar 1^+ transition between 400 and 800 MeV.³⁴ In view of this systematic trend of A_n for the isoscalar M1 transitions, the present A_n data for the 10.23 MeV state in ^{48}Ca indicate that the transition is almost pure isovector. Although the pure neutron transition assumed here implies equal contributions from isoscalar and isovector components to the

transition density, the results are not surprising given the weak strength of V_{σ} in the effective interaction.

Pion inelastic scattering provides a sensitive test of neutron-proton components of the transition density; for a pure neutron transition the ratio of π^- to π^+ cross sections for the M1 transition in ^{48}Ca should be: $\sigma(\pi^-)/\sigma(\pi^+) = 9$. Recent measurements⁸⁸ of $^{48}\text{Ca}(\pi^{\pm}, \pi^{\pm})$ cross sections to the 10.23 MeV state yield a $\sigma(\pi^-)$ to $\sigma(\pi^+)$ ratio consistent with a pure neutron transition; however, the uncertainty in the $\sigma(\pi^+)$ data are too large to arrive at any definite conclusion. Comparison of the more certain $\sigma(\pi^-)$ data with DWIA calculations reveals a quenching that is consistent with electron scattering results.

The S_{nn} data for the 10.23 MeV state in ^{48}Ca show good agreement with both the IPM and MW wave function calculations. Comparison of the two calculations indicates insensitivity of S_{nn} to the transition density. The very conspicuous structure in the shape of the DWIA S_{nn} predictions can be attributed mostly to exchange contributions of the tensor component of the effective interaction. The observed large value of S_{nn} at the smallest q measured validates the applicability of S_{nn} in identifying $\Delta S = 1$ transitions. The agreement of the angular distribution of the S_{nn} data with DWIA predictions provides a good test of the DWIA calculations and the t-matrix of Love-Franey.

Finally, no other M1 states of statistical significance, as suggested in Ref.80, were observed in the present study. The yield in the spectrum of Fig. 15 for the 8.88 MeV peak of ^{16}O is about 7% of that for the 10.23 MeV M1 state. Since the 8.88 MeV state is about at

the limit of detectability, this establishes a limit on the sensitivity of the present experiment in detecting any additional individual M1 states. From Ref. 80, these additional states have $B(M1)$ values between 2 and 3% of that of the 10.23 MeV states. Thus, our experimental sensitivity precludes the observation of such weak fragmentation.

E. Study of ^{90}Zr

The spectrum of ^{90}Zr is characteristic of proton scattering from a typical heavy nucleus. The lower excitation region, with narrow well resolved states, has been explored by scattering studies using different probes at various low to intermediate energies. However, except for natural parity giant resonances, the understanding of the medium and high excitation region still remains unclear. In the following sections, theoretical interpretation of the various regions of excitation of ^{90}Zr , with particular emphasis on spin excitations of lower multipolarity, will be presented. The $\frac{d\sigma}{d\Omega}$ and A_n data for the low-lying states observed in the present study and comparison to previous data will be discussed briefly in Section IV.E.1. In Sec. IV.E.2 the entire bump at about 9 MeV observed in ^{90}Zr will be analyzed as an M1 giant resonance. The fine structure states observed within the 9 MeV bump will be discussed individually in Sec. IV.E.3. Next, the high excitation region ($E_x = 10 - 17$ MeV) will be presented in Sec. IV.E.4. Finally, the S_{nn} results for the entire spectrum will be analyzed in Sec. IV.E.4.

In the present experiment, no cross section angular distribution data for the elastic scattering from ^{90}Zr were taken. The optical potential parameters were, therefore, extrapolated from nearby energies. The geometrical parameters were taken to be the same as those for 200 MeV of Ref. 15. The magnitudes of the real and imaginary central and spin orbit potentials were interpolated from the 200 MeV parameters of Ref. 15 and the 500 MeV parameters obtained by fitting the elastic data of Ref. 89 with the computer code CUPID. The parameters used for the 319 MeV DWIA calculations are listed in Table VI. The elastic cross section calculated with the extrapolated optical parameters of Table VI was compared with the predicted elastic cross section with the folding model calculation of Franey;⁹⁰ good agreement was obtained in the comparison.

1. The Low Lying States

Because of the nearly closed shell model structure of ^{90}Zr , low lying states are expected to arise from simple particle-hole configurations. The excitation of several natural parity states below ~ 4 MeV has been attributed to the distribution of the two valence protons among the $2p_{1/2}$ and $1g_{9/2}$ orbitals. Similarly, states in the region of 4-5 MeV are believed to arise from the multiplets of the $|1g_{9/2}, 2d_{5/2}\rangle$ neutron configurations. These states have been studied via transfer and pickup reactions.⁹¹ In addition, a significant amount of data from inelastic reactions with different probes is available on the low lying states. Although a detailed analysis of these states is

TABLE VIII

Comparison of the low lying states observed in the present work with literature. Errors on the present (p,p') excitation energies are ± 20 keV.

Ex(MeV)	L	β	Ex(MeV) ^a	J ^{π}
3.31			3.311	2 ⁺
3.84	2	0.051	3.846	2 ⁺
4.23	2	0.032	4.233	2 ⁺
4.49			4.480	(3 ⁻)
4.69	2	0.027	4.689	2 ⁺
5.11	3	0.057	5.119	3 ⁻
5.64	3	0.094	5.631	3 ⁻
5.78	3	0.014	5.780	3 ⁻

a) Ref. 92

outside the scope of the present study, it is essential to compare the present results with previous work as a consistency check. Furthermore, interpretation of the relatively well understood low lying states in the microscopic DWIA framework may provide better understanding of the more complex high excitation region.

Shown in Table VIII is a summary of the states observed in the region of 3.3 - 6.4 MeV and comparison to previous results obtained with high resolution inelastic scattering; the (p,p') results were taken from the 25 MeV proton scattering work of Van der Bijl.⁹²

The states observed at 3.31, 3.84, 4.23, and 4.69 MeV which have been assigned as 2^+ states in previous studies, are very well described by an $L = 2$ collective macroscopic description. The solid lines in Figs. 19a-c are the result of coupled channels calculations with the computer code ECIS. The deformation parameters β derived from the calculations are listed in Table VIII; these values are found to be in good agreement with previous low and high energy β values. In addition, microscopic DWIA calculations with the $\pi(g9/2)^2; 2^+$ configuration were performed for these states; these are shown as the dashed curves in Figs. 19a-c. The microscopic model shows overall reasonable agreement with the data; however, the agreement in the low q region is markedly poor. In passing, it should be mentioned that the configuration considered here is a simple one; calculations with increased model space configurations have been shown to produce better agreement with previous experimental data.⁹²

TABLE IX

Preliminary values of the excitation energies and the widths of more prominent E1 states observed in nuclear resonance fluorescence. Preliminary data were provided by U. Berg (Ref. 24). The values of β and $B(E1)^\dagger$ are calculated as described in the text

E_x (MeV)	Γ_0 (ev)	β ($\times 10^{-2}$)	$B(E1)^\dagger$ ($e^2 fm^2 \times 10^{-2}$)
6.295	2.20	0.2898	2.5180
6.424	1.40	0.2242	1.5080
6.762	2.50	0.2775	2.3090
6.876	1.00	0.1711	0.8783
7.249	0.52	0.1140	0.3898
7.478	0.80	0.1350	0.5462
7.806	0.70	0.1184	0.4202
8.113	0.90	0.1267	0.4812
8.237	1.25	0.1459	0.6386
8.413	1.80	0.1697	0.8631
8.440	3.60	0.2388	1.7100
9.140	4.70	0.2421	1.7580
9.330	1.00	0.1083	0.3516
9.487	1.80	0.1417	0.6019

Three of the 3^- states, at 5.11, 5.64, and 5.78 MeV, were also analyzed with macroscopic $L = 3$, coupled channels calculations. These are illustrated in Fig. 19 as the solid lines. The agreement of the data with the calculations was found to be very good. The β values of these states are also listed in Table VIII.

The analysis of the 1^- states, as discussed in Sec. IV.B, requires special consideration since the predominant mechanism for the excitation of these states at small q is coulomb excitation. Cecil et al.¹³ from their 25 MeV $(p,p'\gamma)$ work have identified three 1^- states at 5.51, 5.89, and 6.42 MeV in ^{90}Zr . The observation of these E1 states has also been reported in the results of Ref. 93. These studies extend only up to about 6.5 MeV of excitation energy. However, recent preliminary (γ,γ') data of Berg et al. (Ref. 24) reveal a number of E1 states in the region of excitation energy from 6.3 to 9.5 MeV; the excitation energy of the prominent states and their radiative decay widths are presented in Table IX

In this section we will restrict our discussion of the E1 states observed in (γ,γ') to those below 7 MeV of excitation energy; the states in the region 7 to 10 MeV will be discussed in the next section. In the present (p,p') study, states were observed at 6.30, 6.42, 6.75, and 6.88 MeV; although states were observed at 5.50 and 5.89 MeV, the yield in these states was inadequate for analysis of their cross section angular distribution. Cross section calculations with pure coulomb excitation were done in accordance with the formalism described earlier. The β values computed from the radiative widths of Ref. 24 (see Table IX) were used in the calculation. The calculations are

shown as solid lines in Fig. 18c. In addition, coulomb plus nuclear excitation calculation with an $L = 1$ collective form factor were also performed; these are represented by the dashed curves of Fig. 18c. The shape of the cross section angular distributions in the small q region for the states analyzed here is reasonably described by the pure coulomb excitation calculations whereas the coulomb-nuclear excitation curves agree with the data better over the range of measured q .

However, the absolute value of the cross sections is not reproduced by either of the calculations. In both cases the predicted cross section has been renormalized to fit the data. The renormalization factor required to describe the data is approximately the same ($N=2.5 \pm 0.4$) for the four states. The origin of such a constant is unclear. Few such calculations appear in the literature at high energies, and there are significant discrepancies among them.^{31,93,94}

In an attempt to investigate the ECIS calculations, we have also analyzed the E1 states observed in the $^{208}\text{Pb}(p,p')^{208}\text{Pb}^*$ reaction at $E_p = 212$ MeV.⁹⁵ The data were obtained at HRS in experiment 630 and a complete analysis of the results will be presented elsewhere. The decay widths for the relatively strong E1 states observed at 5.512, 7.063, and 7.332 MeV were taken from the NRF work of Chapuran et al.⁹⁶ In the present (p,p') work, with 70 keV of energy resolution, states were observed at 5.50, 7.06, 7.32 MeV. The shapes of the cross section angular distribution well described by the ECIS calculations. A similar renormalization factor of about 2.3 was required, however, to fit the absolute value of the cross section. A summary of the required

TABLE X.

Comparison of ECIS coulomb excitation calculations for ^{90}Zr with the data. The data and the calculation are at 2.2^oc.m. The ECIS calculations are as described in the text using the β values of Table IX.

E_x (MeV)	$(\frac{d\sigma}{d\Omega})_{\text{exp}}$ (mb/sr)	$(\frac{d\sigma}{d\Omega})_{\text{ECIS}}$ (mb/sr)	N^a
6.30	0.20 ± 0.03	0.10	2.0
6.42	0.17 ± 0.03	0.06	2.8
6.75	0.24 ± 0.03	0.09	2.7
6.88	0.13 ± 0.04	0.04	3.2

a) $N = (\frac{d\sigma}{d\Omega})_{\text{exp}} / (\frac{d\sigma}{d\Omega})_{\text{ECIS}}$

renormalization of the ECIS coulomb excitation calculations for ^{90}Zr is given in Table X. Thus, based on our empirical observation, we will include a renormalization factor of about 2.5 in the ECIS calculations in the subsequent analysis.

2. The 9 MeV M1 Bump

In the IPM model the M1 strength in ^{90}Zr is expected to reside exclusively in the pure neutron spin-flip transition $\nu g_{9/2} \rightarrow 8_{7/2}$. Model calculations^{9/} indicate this state to occur at an excitation energy of 8.8 MeV with a strength $B(M1)_{\uparrow}$ of $16\mu_K^2$.

Measurements with 165° high resolution inelastic electron scattering by the Darmstadt group (Ref. 23) reveal significantly fragmented M1 strengths. Although several tentative M1 assignments were attributed to states in the excitation energy range of 7.7 - 9.5 MeV, only three definite M1 states, at $E_x = 8.233, 9.000, 9.371$ were observed. The upper limit for the total observed strength in this (e,e') study was reported to be $1.8 \pm 0.4 \mu_K^2$ which accounts for about 11% of the IPM prediction. In (p,n) reactions, the analogue of the M1 transition was observed at 5 MeV above the IAS and the strength was found to be about 60% of the (p,n) sum rule. The (p,p') measurements at 200 MeV by the Orsay group (Ref. 15) observed a broad 1.5 MeV bump at about 9 MeV with a cross section at 4° of $2.8 \pm 0.4 \frac{\text{mb}}{\text{sr}}$. Inelastic proton scattering work at 200 MeV of Bertrand et al.⁹⁸ at TRIUMF reported observation of a giant resonance with a cross section

of $7.2 \pm 2.0 \frac{\text{mb}}{\text{sr}}$ at 4° ; however, these results do not appear definitive because of high backgrounds and poor energy resolution.

In the present (p,p') study a broad bump of 1.8 MeV width (full width at half maximum) was observed with the centroid of the bump at an excitation energy of 8.98 MeV. Although a very rich fine structure within the bump was observed here for the first time in (p,p') studies, an analysis of the entire bump as a single excitation will be presented here for comparison purposes (see Sec. IV.E.3 for the fine structure analysis). The cross section for the bump, at the smallest scattering angle measured (2°), obtained in a manner as outlined in Sec. III.B, is $5.6 \frac{\text{mb}}{\text{sr}}$. The momentum transfer at a scattering angle of 4° at $E_p = 200 \text{ MeV}$ is 0.24 fm^{-1} . The present datum, interpolated at the same momentum transfer, is $3.4 \pm 0.5 \frac{\text{mb}}{\text{sr}}$, which is about 20% higher than the 200 MeV Orsay data. A conservative estimate of the systematic errors due to various choices of empirical backgrounds is included in the cross section value given here.

The DWIA prediction with the $\nu(g_{9/2}, g_{7/2})_1^+$ IPM wave functions and LF t-matrix of Table IV are shown as the solid curves in Figs. 20 and 31. The optical parameters used were the extrapolated parameters as shown in Table VI. Exchange contributions were calculated exactly. The predictions have been normalized to fit the low q region of the data. Qualitatively, the DWIA, IPM prediction is in very good agreement with the angular distribution of the $\frac{d\sigma}{d\Omega}$ data in the low q region. However, at larger q the cross section data show a systematic upward deviation away from the DWIA curve. This indicates that the bump is predominantly $\Delta L = 0$ in the small q region, but there

appears to be a significant contribution due to transitions of higher multipolarity at larger q . In any event, the predicted DWIA cross section for the M1 state at 319 MeV and $q = 0.24 \text{ fm}^{-1}$ is $9.2 \frac{\text{mb}}{\text{sr}}$ which is about 25% smaller than the DWIA prediction at 200 MeV at the same q . Uncertainty in the 319 MeV optical parameters due to extrapolation is estimated to introduce a systematic uncertainty of 20% to the DWIA calculations. Comparison of the present data to the DWIA cross section, thus, yields only $37 \pm 10\%$ of the predicted cross section. From a comparison of the measured angular distribution with similar DWIA calculation, the latest analysis of the 200 MeV data yields 30%. Thus, the present measurements are in agreement with Orsay results in the observed quenching of the M1 bump.

In the $^{90}\text{Zr}(p,n)^{90}\text{Nb}$ reaction spectrum of Ref. 22 at 200 MeV, the $T_0 - 1$ and T_0 branches of the GT states are not resolved; the experimental cross section for the total GT strength was reported to be $50 \pm 12 \frac{\text{mb}}{\text{sr}}$. In order to compare the (p,n) results with the present (p,p') work, extrapolation of the data to 0° scattering angle and 200 MeV is required; the DWIA predictions for the cross section were used for the purpose of extrapolation. Extrapolating the cross section data of the present study at 2.0° to 0° with the help of the DWIA cross section angular distribution, the (p,p') cross section at 0° is about $7.3 \frac{\text{mb}}{\text{sr}}$ at 319 MeV. Assuming that the ratio of the 0° DWIA cross sections for (p,p') at 200 MeV and 319 MeV can be used as a scaling factor, our 319 MeV data should correspond to a 0° (p,p') cross section of $8.4 \frac{\text{mb}}{\text{sr}}$ at 200 MeV.

With the simple model for comparison discussed in Sec. IV.A, using Eq. 71 and the $T_0 = 5$ for the ground state of ^{90}Zr , for (p,p') reactions a transition to the analogue T_0 , M1 state should have a 0^0 cross section of $12.5 \pm 3.1 \frac{\text{mb}}{\text{sr}}$ at $E_p = 200$ MeV. Considering the simplicity of the model for comparison assumed here, the present (p,p') value of $8.4 \frac{\text{mb}}{\text{sr}}$ is adequately consistent with the (p,n) results of 200 MeV.

The assumptions made in order to simplify the process of comparison with (p,n) can be corrected to provide a more realistic picture. For ^{90}Zr with $T_0 = 5$, the ratio of the $T_0 - 1$ to T_0 GT states in (p,n) reaction is 9; however, lower energy (p,n) studies at $E_p = 120$ MeV, where the T_0 and the $T_0 - 1$ GT states are well resolved, have measured this ratio to be 13. This deviation has been explained by Bertsch et al.⁹⁹ as a result of configuration mixing of the T_0 GT state with a low lying $T = 4$ state with the configuration $(8g_{7/2}, 8g_{9/2}^{-1})_1^+$. Taking this experimental ratio to be valid for the unresolved GT states at 200 MeV, the scaled cross section for (p,p') becomes $\sim 9 \frac{\text{mb}}{\text{sr}}$, which makes the agreement with the present data better.

The effect of neglecting the V_σ contribution in (p,p') can be calculated in a straightforward manner; DWIA calculations made with and without the V_σ part in the LF interaction indicate that at 319 MeV interference between V_σ and $V_{\sigma\tau}$ interactions results in a net cancellation of the predicted cross section making it $\sim 15\%$ lower than that with $V_{\sigma\tau}$ alone. Furthermore, the assumption that the distortion factors for (p,p') and (p,n) are the same is not exact. A distortion factor for (p,p') at 319 MeV was calculated from the ratio of the PWIA

to DWIA cross section. This ratio was found to be ~ 10.5 whereas the ratio for (p,n) is 12. Thus, the combined effect of V_0 and distortion will further reduce the predicted (p,p') cross section by $\sim 30\%$. The net result is that the present (p,p') and (p,n) results are in good agreement within errors about the overall strength of the peak.

3. Fine Structure in the M1 Region

The expanded spectra of Fig. 21 clearly show rich fine structure in the 9 MeV M1 transition region. No such fine structure was observed in the previous proton work of Ref. 15. In an attempt to isolate M1, M2, and E1 transition strengths in this region, the cross section angular distribution of these peaks were compared with DWIA calculations for M1 and M2 transitions assuming $\nu(g_{9/2}^{-1}, g_{7/2})$ and $\nu(f_{5/2}^{-1}, g_{7/2})$ configurations, respectively. These calculations are illustrated by the solid curves in Fig. 22; both curves have been normalized to agree with the data. In addition, the shape as well as the magnitude of the cross sections were compared with coulomb excited E1 calculations. It is now possible to make a detailed comparison between the (e,e') and (p,p') results. The centroids of peaks identified in the present (p,p') study from their cross section angular distributions as consistent with mostly M1 excitation are listed in Table XI. Peaks fitted at 7.53, 7.94, and 8.62 MeV appear to have M2 angular distributions. A number of states in Table XI have angular distributions that indicate significant M2 strength as well; their M2 contribution at the most forward angles, however, is very small. The

TABLE XI

Comparison of (p,p') and (e,e') results. The peak numbers for (p,p') correspond to the labels in Fig. 21; only states with angular distributions consistent with 1^+ are listed. Errors on (p,p') excitation energies are ± 20 keV; errors on (e,e') excitation energies are ± 10 keV. Definite assignments from (e,e') are underlined.

No.	(p,p') Ex(MeV)*	$(e,e')^a$	
		Ex(MeV)	J^π
		7.77	$(1^+, \quad)$
		7.87	$(1^+, 2^-)$
3	8.13	8.14	$1^+, (2^-)$
4	8.26	<u>8.23</u>	<u>1^+</u>
5	8.37	8.37	(1^+)
		8.60	(1^+)
7	8.73		
8	8.98	<u>9.00</u>	<u>1^+</u>
9	9.14		
10	9.28		
11	9.38	<u>9.37</u>	<u>1^+</u>
		9.44	$1^+(2^-)$
12	9.48		
		9.52	$(1^+, 2^-)$
13	9.60		
14	9.77		

DWIA calculations indicate at 2.2° the M2 cross section is only about 7% of that of the M1. Also listed in Table XI are the definite and possible M1 assignments from electron scattering. Table XI reveals better agreement between the (p,p') and (e,e') results than previously apparent³⁻⁵, but it is clear that substantial problems remain. At the positions where definite 1^+ states are located in (e,e') , strength consistent with 1^+ is observed also in (p,p') . This is true also of several states assigned as possibly 1^+ in (e,e') , especially if the 9.44 and 9.52 MeV states in (e,e') are seen as the 9.48 MeV state in (p,p') . However, several equally strong states, consistent with 1^+ in (p,p') , lie in regions where no M1 strength is identified in (e,e') . For example, the strongest peak in the present (p,p') spectrum, at 9.28 MeV, is not seen as a 1^+ in (e,e') .

Several of the E1 states reported in the preliminary findings of Berg et.al (see Table IX) appear at the same excitation energies as states listed in Table XI. The radiative widths Γ_0 for the proposed E1 states in this region are, with some exceptions, comparable to values of Γ_0 for E1 states around 6.5 MeV; these low lying E1 states have been discussed earlier. Now, Γ_0 is proportional to $E_X^3 B(E1)$; the Coulomb excitation cross section is approximately proportional to $E_X^{-1} B(E1)$. This indicates that, around 9 MeV, only E1 states with widths much larger than those around 6.5 MeV should be identifiable in the (p,p') experiment. A similar argument can be based on the observed strength of the giant dipole resonance.

In the present work, five of the states listed in Table XI as M1's have excitation energies consistent with the proposed E1 states of Table IX; these states are at 8.13, 8.26, 9.14, 9.38, and 9.48 MeV. Although the (γ, γ') indicates a strong E1 state at 8.44 MeV, no statistically significant state was observed at this energy. In (p, p') reactions, however, it is not possible to confirm E1 characteristics of these peaks from cross section angular distributions. Furthermore, the uncertainty in the present spin-flip probability data for such narrow peaks is too large to be useful in isolating E1's from M1's in this region with such a large density of states. However, it is conceivable that the states labeled as M1's in the present work are a mixture of E1 and M1 strengths. To help decompose the E1 and M1 strengths in this excitation region, coulomb excitation cross sections were calculated for these five states using the preliminary Γ_0 values in Ref. 24. These predictions were compared with the experimental cross sections of these peaks at 2.2° ; the result is summarized in Table XII. The ratio of the renormalized ECIS prediction to the data reveals that, except for the peak at 9.14 MeV, the experimental cross sections for the other four states are well above the predicted E1 cross sections.

Note here that the calculations of E1 cross sections are based on preliminary Γ_0 values. Clearly some discrepancies exist between the preliminary NRF results and the present (p, p') data. For example, the 8.41 and 8.44 MeV states, strongly excited in the NRF work, are not seen in the present study. In addition, the excitation energies of the states that are seen in both studies do not match precisely. Thus, only a tentative E1 strength decomposition of the 9 MeV fine structure

TABLE XII.

Decomposition of E1 and M1 fractions in the M1 region of ^{90}Zr . The data and the calculation are at 2.2°c.m. The ECIS calculations are E1 coulomb excitation calculations as described in the text using the β values of Table X.

E_x (MeV)	E_x^a (MeV)	$(\frac{d\sigma}{d\Omega})_{\text{exp}}$ (mb/sr)	$(\frac{d\sigma}{d\Omega})_{\text{ECIS}}^b$ (mb/sr)	R^c
8.13	8.11	0.12 ± 0.04	0.042	0.33
8.26	8.24	0.18 ± 0.04	0.054	0.30
9.14	9.14	0.15 ± 0.04	0.135	0.90
9.38	9.33	0.13 ± 0.04	0.026	0.20
9.48	9.48	0.20 ± 0.04	0.045	0.22

a) Ref. 24

b) Renormalization factor as discussed in the text is included.

c) $R = (\frac{d\sigma}{d\Omega})_{\text{ECIS}} / (\frac{d\sigma}{d\Omega})_{\text{exp}}$

can be made. Furthermore, the uncertainty in the coulomb excitation calculations due to the need for a renormalization factor and the assumed collective form factor for the nuclear excitation of these proposed E1 states make it more difficult to reach a definitive conclusion.

From the foregoing discussion, it is possible to subtract likely E1 and M2 strengths from the total cross section contained in all the fitted peaks to establish a better estimate of the observed M1 strength. The total cross section for the fitted peaks at 2.2° amounts to $2.3 \frac{\text{mb}}{\text{sr}}$. The estimated coulomb excitation cross section of the five proposed E1 states is $0.32 \frac{\text{mb}}{\text{sr}}$, about 15% of the fitted strength. Allowing for 7% M2 strengths as discussed before, we conclude that the best estimate for the M1 strengths contained in the peaks of Table XI is about 78% of the total cross section. Recall that only 60% of the total "M1 bump" was fitted to obtain individual states. If we further assume that the M1, E1, and M2 decomposition of the transition strength in the unfitted region exhibits a behaviour similar to the fitted region, this amounts to a reduction of the total M1 strengths to about $30 \pm 8\%$ compared to the DWIA prediction.

4. High Excitation Region

The natural parity giant electric multipole resonances, namely, E0, E1, and E2 at 16.3, 16.4, and 14.3 MeV in ^{90}Zr , are expected⁹⁹ to dominate the high excitation energy region. These are expected to be $\Delta S = 0$ transitions. Comparison of the cross section angular

distribution of the 14-16 MeV region without any background subtraction with E1 coulomb excitation calculations (see Fig. 23) shows very good agreement in the low q region. However, the large background underneath these states giving rise to a strongly excited nuclear continuum above about 7 MeV is very little understood. The absence of well defined structure in the shape of the spectra of ^{90}Zr in the high excitation energy region presented in the previous chapter is suggestive of the fact that there are possibly many broad resonances situated close to one another resulting in such a structureless spectrum.

In addition, contributions from single-collision quasifree N-N scattering is also a possible factor in this continuum. However, the response function with a Fermi-gas model for quasifree scattering is expected¹⁰⁰ to be very small in the low q region due to Pauli blocking. Calculations reported¹⁰¹ for the continuum region in the $^{116}\text{Sn}(p,p')^{116}\text{Sn}^*$ reaction at 800 MeV with the semi-infinite slab (SIS) approximation of Bertsch and Scholten¹⁰⁰, where Pauli blocking at low q is partially suppressed, indicate the quasifree cross section decreases rapidly with smaller q . No such calculations for ^{90}Zr at about 300 MeV exist at this time. Finally, note that the S_{nn} value predicted from Arndt's N-N phase shift analysis at 3.5^0 is 0.37 and 0.20 for free p-p and p-n scattering. However, as mentioned earlier, contributions to the spin-flip cross section due to quasifree scattering are expected to be small.

The most striking feature of the present spin-flip measurements is the unexpectedly large spin excitation strength observed for the first time in this high excitation region. The measured spin-flip cross section for ^{90}Zr is about 0.80 mb/sr/MeV throughout the region of excitation energy from 8 to 25 MeV at 3.5° laboratory scattering angle; the spin-flip probability in the same region is about 0.16. If we take the value of $S_{nn} = 0.44$ for the 10.23 MeV state in ^{48}Ca as a rough indicator of expected S_{nn} for spin-excitations (this is also consistent with DWIA predictions for low multipole spin-excitations), then about one-third of the cross section observed in the continuum is due to spin excitations. This appears very large considering that the dominant natural parity resonances here should have $\Delta S = 0$.

Note here that no spin excitation in ^{90}Zr have previously been observed in ^{90}Zr above 10 MeV. The results reported by Moss et.al¹⁰¹, however, show an apparent discrepancy with the present results; they have found that, the spin-excitation strength in $^{90}\text{Zr}(p,p')^{90}\text{Zr}$ at 500 MeV and 12° scattering angle averaged over 12 to 29 MeV excitation energy, and for $^{208}\text{Pb}(p,p')^{208}\text{Pb}$ at 400 MeV and 4.5° scattering angle averaged over 7 to 22 MeV, is consistent with zero. At the 500 MeV proton energy used in their measurements the spin excitation strength is weaker relative to those of natural parity in the LF interaction than at 300 MeV. They have measured an average S_{nn} for a wide range of excitation energy extending beyond the present study. In addition, at the large q value of their measurement, the S_{nn} values for spin excitations as indicated by present DWIA calculations are expected to

be small. In light of this, this apparent discrepancy may not be surprising.

Multipoles of Continuum Spin Excitations

A qualitative examination of the cross section angular distributions in 2 MeV steps for the high excitation region (see Fig. 23) suggests the contributions of higher multipole excitations of the nucleus as indicated by the rise of the angular distributions with q . In order to make a preliminary effort to decompose the spin and parity of such excitations, we consider here possible transitions with angular momentum transfer up to $L = 2$; since our data are limited to the low q region, such a restriction is reasonable. This limitation also makes conclusions about the multipole distribution of the strength very tentative. It is discussed here only to indicate possibilities. The possible transitions are those with the transfer quantum numbers J^π given by 0^+ and 1^+ for $L = 0$; 0^- , 1^- , and 2^- for $L = 1$; and 1^+ , 2^+ , and 3^+ for $L = 2$. Our DWIA calculation indicates that the angular distributions for $S = 0$ and $S = 1$ transitions for the same L have a very similar shape. Furthermore, the possible contribution of instrumental backgrounds to the cross sections in this region adds difficulties to the interpretation of the cross section data. Thus from the cross section angular distributions of Fig. 23, a multipole decomposition proves to be not without ambiguity.

DWIA calculations for the spin-flip cross section of these states at 319 MeV were performed with the LF interaction. Consideration of the spin-flip cross section eliminates 0^+ states, as well as problems

TABLE XIII.

Relative multipolar contributions to the spin excitation continuum. The spin-flip cross section values are obtained from the fitting procedure described in the text.

E_x (MeV)	$(S.\sigma)_{exp}$ (mb/sr)	$(S.\sigma)^{1+}$ (mb/sr)	$(S.\sigma)^{2-}$ (mb/sr)	$(S.\sigma)^{3+}$ (mb/sr)
<u>$\theta_{lab} = 3.5^\circ$</u>				
8 - 10	1.90 ± 0.18	1.261	0.370	0.361
10 - 12	1.20 ± 0.20	0.297	0.137	0.795
12 - 14	2.03 ± 0.28	0.593	0.219	1.087
14 - 16	2.27 ± 0.30	0.586	0.411	1.121
16 - 18	1.62 ± 0.32	0.126	0.520	0.991
<u>$\theta_{lab} = 5.0^\circ$</u>				
8 - 10	0.99 ± 0.12	0.164	0.230	0.652
10 - 12	1.49 ± 0.12	0.038	0.085	1.433
12 - 14	2.14 ± 0.13	0.077	0.136	1.955
14 - 16	2.30 ± 0.13	0.076	0.255	2.019
16 - 18	2.13 ± 0.13	0.016	0.322	1.785

with instrumental background. Simple particle-hole configurations, $\nu(g_{9/2}^{-1}, g_{7/2})_{1+}$, $\nu(f_{5/2}^{-1}, g_{7/2})_{2-}$, and $\nu(g_{9/2}^{-1}, d_{5/2})_{3+}$, were used in the calculations. The observed spin-flip cross section for a given region of excitation in the continuum was expressed as a linear combination of the 1^+ , 2^- , and 3^+ excitations

$$(S.\sigma)_{\text{exp}} = a(S.\sigma)^{1^+} + b(S.\sigma)^{2^-} + c(S.\sigma)^{3^+}. \quad (79)$$

Here $(S.\sigma)$ is used in place of $S_{nn} \frac{d\sigma}{d\Omega}$ for brevity. The solid curves in Fig. 32 were obtained by determining the coefficients a , b , and c in a χ^2 minimization search to fit the measured spin-flip cross section for each of the regions presented in the figure. The relative contributions of $L = 0$ (1^+), $L = 1$ (2^-), and $L = 2$ (3^+) transitions are indicated by dashed, dotted, and dashed-dotted curves in the figures. Comparison of the resultant curves with the data in Fig. 32 indicates that the high excitation spin-flip continuum at these low q regions is well explained by excitation of spin multipole vibration modes of up to $L = 2$ (3^+). The relative contribution of the different multipoles to $(S.\sigma)_{\text{exp}}$ is summarized in Table XIII. As expected, these DWIA curves (Fig. 32) indicate the dominance of the $L = 0(1^+)$ transition in the 8-10 MeV region at 3.5° . However, the high excitation region at larger q appears to be mostly $L = 2$ excitations. The relative contribution of $L = 0$ and $L = 1$ strength is not well determined. Finally, the strength distribution of the different multipoles obtained from this limited analysis is represented by histograms along with the data in Fig. 36. Examination of the strength distribution histogram at 3.5° reveals a

Fig. 36: Multipolar decomposition of the spin-flip cross section spectra for ^{90}Zr .

large concentration of M1 ($L = 0$) strength in the 9 MeV region; however, these preliminary calculations indicate a significant spreading of the M1 strength into higher excitation energy which bears further analysis. The M1 strength contained in 10 to 18 MeV in this histogram is about 60% of the DWIA prediction; recall that the 9 MeV bump contains only about 37% of the DWIA prediction for the cross section. Thus, it may be that the missing M1 strength is spread out into the high excitation region; as discussed before, uncertainty in the spin-flip data and the limited range of angular distributions makes it difficult for a more definitive conclusion. In addition, more detailed calculations with realistic form factors are necessary.

Such calculations would also be necessary to confirm the other striking feature observed in the analysis of the present data, the apparent strong contribution of $L = 2$ spin excitation to the giant resonance continuum of ^{90}Zr . It is interesting to note that Osterfeld¹⁰² in a recent article has predicted a strong $L = 2$ (3^+) resonance dominating the high excitation continuum in the $^{48}\text{Ca}(p,n)^{48}\text{Sc}$ reaction; this resonance is expected to be around 25 MeV with a full width at half maximum of about 4.5 MeV. No such resonance is indicated in ^{40}Ca . The appearance of this 3^+ resonance in ^{48}Ca is attributed to a more attractive single particle potential for the $g_{9/2}$ and $g_{7/2}$ orbitals. Note that similar effects should be observable in ^{90}Zr . In addition, calculations of transition strengths by Auerbach for proton scattering from ^{90}Zr indicate the presence of an $L = 2$ (3^+) excitation at about 24 MeV. Recent results of (p,n) scattering by

Rapaport et al.¹⁰³ have also revealed broad $L = 1$ and $L = 2$ resonances in Ni isotopes.

V. CONCLUSION

The present (p,p') measurements have yielded some expected and some not so expected results. First of all, the spin-flip measurements on ^{48}Ca carried out for the first time have revealed, as expected, a large value for the spin-flip probability of the 10.23 MeV M1 state in the low q region. The DWIA prediction for the spin-flip probability is in good agreement with the data, providing credibility to the DWIA calculations. The present cross section data are quenched relative to the DWIA prediction; the quenching is similar to that observed in previous work.

Secondly, the present measurements have revealed considerable fine structure in the M1 region of ^{90}Zr that has not been observed in previous intermediate energy proton scattering work. Angular distributions of most of the fine structure states are consistent with DWIA calculations for M1 transitions. The excitation energies of many of these states are in good agreement with electron scattering results, but discrepancies in spin assignments and strengths for some of the states continue to be a problem. The present study has provided a good first step in reconciling some of the existing disagreements between (p,p') and (e,e') results. However, more measurements with higher resolution and better statistical accuracy are needed to help resolve the remaining discrepancies; such measurements are currently under way at LAMPF.

The measured cross section for the entire 9 MeV bump in ^{90}Zr reveals only about 30% of the M1 cross section predicted by the DWIA using the LF interaction and a simple particle-hole configuration in agreement with the previous (p,p') results of Ref. 15. The spin-flip probability for the bump was found to be 0.60 ± 0.20 at 3.5° . The angular distributions for the cross section, analyzing power, and spin-flip probability are in good agreement with the DWIA; thus, an M1 assignment for this bump is consistent with our measurements. However, about 15% of the strength in the bump in the low q region is estimated to be due to weakly excited E1 states; this estimate is based on the preliminary data from the nuclear resonance fluorescence data of Ref. 24. Another 7% is attributed to M2 contributions.

Concerning the discrepancies between the (p,p') and (e,e') reactions, it is probably most important is to note that the (p,p') data extend to much smaller q values (0.16 fm^{-1}) than the (e,e') data (0.35 fm^{-1}), so that the relative sensitivity of (p,p') to M1 states is considerably enhanced. Possible current contributions to electron scattering are unlikely to be important for the mostly neutron M1 transitions, but they may enhance the M2 states which predominate in electron scattering. While $\Delta T=0$ transitions contribute relatively more to proton scattering, these amplitudes are weak in the Love-Franey interaction; large $\Delta T=0$ contributions seem ruled out by our A_y data. In addition, the discrepancy observed here with respect to the preliminary studies, where no M1 states were observed, can probably be attributed to the relative insensitivity of the (γ,γ') reaction to M1 excitations. Note that in (γ,γ') reactions, radiative decay widths are

measured. A similar analysis as that for the E1 states (see Sec. IV.B) reveals that the widths of M1 states in the 9 MeV region typically are about four times smaller than the widths of typical E1 states in this region, while they have comparable cross sections in the (p,p') reaction. Thus, the dominance of the E1 states in the (γ,γ') work may not be surprising.

More definitive conclusions regarding the spin assignments of the fine structure observed here, however, must await future measurements. The lack of elastic scattering data at about 300 MeV introduces a systematic uncertainty into the theoretical calculations presented here due to uncertainty in the optical parameters. Although the analyzing power and spin-flip probability are relatively unaffected by the optical parameters, the cross section calculations were found sensitive enough to introduce an estimated 20% error in the strengths quoted above. This makes the measurement of elastic scattering highly desirable. In principle, spin-flip measurements can be used to isolate non-spin-flip E1 states, but these remain very difficult for fine structure peaks with large background. Present DWIA calculations at somewhat lower energies indicate larger spin-flip probability values for the M1 peak. In addition, preliminary experimental results at LAMPF indicate a significant improvement in the continuum background in medium and heavy nuclei at lower proton energies. Thus, spin-flip probability measurements with improved polarimetry techniques at about 200 MeV proton energies might facilitate better understanding of this region of excitation in ^{90}Zr .

Finally, contrary to expectations, the present spin-flip measurements for ^{90}Zr have revealed a large spin excitation cross section for the giant resonance continuum distributed almost uniformly over the excitation energy region from 7 to 25 MeV. The measured spin-flip cross section in this region is approximately 0.80 mb/sr-MeV. No spin excitations have been identified previously in this region of ^{90}Zr ; only natural parity giant resonances have been observed. Preliminary analysis of the angular distribution of spin-flip cross sections from 8 to 18 MeV in steps of 2 MeV indicates that the strongly excited spin-flip continuum is consistent with spin excitations involving considerable $L = 2$ components. Once again, more measurements are clearly needed to explore further this fundamental aspect of nuclear structure. Measurements extending the present data base to higher excitation energy and larger scattering angles will permit a more extensive analysis of the multipole decomposition in the continuum. In addition, measurements of other Wolfenstein parameters as suggested by Bertsch et al.¹⁰⁰ might help resolve the spin assignments for these proposed spin excitations. But it is clear that the present measurements of spin-flip cross sections have opened a new frontier in the understanding of the nuclear continuum.

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REFERENCES

1. A. Arima and H. Horie, Prog. Theor. Phys. 12, 623 (1954).
2. R. Blin-Stoyle and A. Perks, Proc. Roy. Soc. 67A, 885 (1954).
3. V. Gillet, A. Green, and E. Sanderson, Nucl. Phys. 88, 321 (1966).
4. H. Mavromatis, L. Zamick, and G. Brown, Nucl. Phys. 80, 545 (1966).
5. T. T. S. Kuo, Nucl. Phys. A112, 325 (1968).
6. J. D. Vergados, Phys. Lett. 36B, 12 (1971).
7. J. S. Dehasa, J. Speth, and A. Faessler, Phys. Rev. C11, 607 (1977).
8. D. E. Bainum, J. Rapaport, C. Goodman, D. Horen, C. Foster, M. Greenfield, and C. Goulding, Phys. Rev. Lett. 44, 1751 (1980).
9. C. D. Goodman, C. Goulding, M. Greenfield, J. Rapaport, D. Bainum, C. Foster, W. G. Love, and F. Petrovich, Phys. Rev. Lett. 44, 1755 (1980).
10. B. D. Anderson, J. Knudson, P. Tandy, J. Watson, R. Madey, and C. Foster, Phys. Rev. Lett. 45, 699 (1980).
11. W. Knüpfer, R. F. Ly, A. Friebe, W. Mettner, D. Meuer, A. Richter, E. Spamer, and O. Titze, Phys. Lett. 77B, 367 (1978).
12. A. Richter, in Nuclear Physics with Electromagnetic Interactions, edited by H. Arenhovel and D. Dreschel, Lecture Notes in Physics, Vol. 108 (Springer, Berlin, 1979) p. 19.

13. F. E. Cecil, G. T. Garvey, and W. J. Braithwhite, Nucl. Phys. A232, 22 (1974).
14. H. Ikegami, T. Yamazaki, S. Morinobu, I. Katayama, M. Fujiwara, Y. Fujita, T. Itahashi, and S. I. Hayakawa, in Proceedings of the Symposium on Highly Excited States in Nuclear Reactions, Osaka, edited by H. Ikegami and M. Muraoka (Osaka University, Osaka 1980)
15. N. Anantaraman, G. M. Crawley, A. Galonsky, C. Djalali, N. Marty, M. Morlet, A. Willis, and J. Jourdain, Phys. Rev. Lett. 46, 1318 (1981).
16. G. Morpugo, Phys. Rev. 110, 721 (1958).
17. K. Ikeda, S. Fujii, and J. I. Fujita, Phys. Lett. 3, 271 (1963); S. Fujii, J. I. Fujita, and K. Ikeda, Phys. Rev. 133, B549 (1964); K. Ikeda, Prog. Theor. Phys. 31, 434 (1964); J. I. Fujita and K. Ikeda, Nucl. Phys. 67, 145 (1964).
18. G. R. Satchler, Nucl. Phys. 55, 1 (1964).
19. C. Goodman, "The (p,n) Reaction and the Nucleon-Nucleon Force," edited by C. Goodman, S. Austin, S. Bloom, J. Rapaport, and G. Satchler, (Plenum, New York 1980) p. 149, and references therein.
20. B. D. Anderson, J. Knudson, P. Tandey, J. Watson, R. Madey, and C. Foster, Phys. Rev. Lett. 45, 699 (1980); C. D. Goodman, Nucl. Phys. A314, 241 (1982).
21. L. W. Fagg, Rev. Mod. Phys. 47, 683 (1975).
22. C. Gaarde, J. Rapaport, T. Taddeucci, C. Goodman, C. Foster, D. Bainum, C. Goulding, M. Greenfield, D. Horen, and E. Sugarbaker, Nucl. Phys. A369, 258 (1981).

23. D. Meuer, R. Frey, D. Hoffmann, A. Richter, E. Spamer, O. Titze, and W. Knapfer, Nucl. Phys. A349, 309 (1980).
24. U. E. P. Berg, private communication.
25. A. Bohr and B. Mottelson, Phys. Lett. 100B, 10 (1981); see for example, Proceedings of the Symposium on Delta-Nucleus Dynamics, edited by T.-S. Lee, D. F. Geesman, J. P. Schiffer, Argonne National Lab. Report ANL-PHY-83-1, CONF-830588.
26. M. Kohno and D. W. L. Sprung, Phys. Rev. C26, 297 (1982).
27. A. Arima, Proc. Int. Conf. on Spin Excitations in Nuclei, Telluride, Colorado (1982); Plenum, in press.
28. G. E. Brown, J. S. Dehsea, and J. Speth, Nucl. Phys. A330, 290 (1979)
29. G. F. Bertsch and I. Hamamoto, Phys. Rev. C26, 1323 (1982).
30. O. Sholten, G. F. Bertsch, and H. Toki, Phys. Rev. C27, 2975 (1983).
31. T. Izumoto, Y.-W. Lui, and D. H. Youngblood, Phys. Rev. C24, 2179 (1981).
32. W. D. Cornelius, Ph. D. Thesis, Texas A & M University, (~~1979~~), unpublished.
33. S. J. Seestrom-Morris, J. M. Moss, J. B. McClelland, W. D. Cornelius, T. A. Carey, N. M. Hintz, M. Gazzaly, M. A. Franey, S. Nanda, and B. Aas, Phys. Rev. C25 3035 1982.
34. K. W. Jones, Ph. D. Thesis, Rutgers University, (1982), unpublished.

35. M. S. Livingston, "LAMPF - A Nuclear Research Facility," LANL Report LA-6878-MS (1977).
36. G. S. Blanpied, Ph. D. thesis, LANL Report LA-7262-T (1978).
37. B. Zeidman, "The HRS spectrometer system: Concepts and Design," LANL Report LA-4773-MS (1971).
38. A. Rahbar, Ph. D Thesis, LANL Report LA-9505-T (1982).
39. M. McNaughton, Phys. Rev. C23, 1128 (1981).
40. R. M. Mobley, BNL Report, BNL-50120, 430 (1968).
41. C. L. Morris, H. A. Thiessen, and G. W. Hoffmann, IEEE Trans. Nucl. Sc. NS-25, 141 (1978).
42. J. B. McClelland, J. F. Amann, W. D. Cornelius, H. A. Thiessen, and B. Aas, to be published.
43. B. Zacharov, "CAMAC Systems: A Pedestrian's Guide", Daresbury Nucl. Phys. Lab., Daresbury, England (1972).
44. J. McGill, Ph. D. Thesis, LANL Report LA-8937-T (1981).
45. M. Kellogg, M. Minor, S. Schaler, W. Spencer, R. Thomas Jr., and H. van der Beken, "Introduction to Q," LANL Report LA-7001-M (1978).
46. J. F. Amann, "ALLTEST - A Data Testing Package," unpublished (1979).
47. R. J. Ridge, "DSP - A Histogram Display System," unpublished (1979).
48. S. K. Nanda, "DNC - Calculation of Polarization Transfer Observables from Polarimeter Sums and Histograms", unpublished (1982).

49. D. Hutcheon, private communication.
50. R. Arndt, WI82 solution, Scattering analysis Interactive Dial-in program.
51. G. Smith, "LOAF - A peak fitting program," unpublished.
52. B. L. Berman, J. T. Caldwell, R. R. Harvey, and F. C. Flutz, Phys. Rev. 162, 1098 (1967).
53. A. Lepretre, H. Beil, R. Bergere, P. Carlos, A. Veysiere, and M. Sugawara, Nucl. Phys. A175, 609 (1971).
54. L. Wolfenstein, Phys. Rev. 96, 1654 (1954).
55. W. D. Cornelius, "Mixing of Final State Spin Components in the HRS," LANL Report, LA-9130-MS.
56. D. Besset et al., Nucl. Inst. and Meth. 166, 379 (1979).
57. R. D. Ransome, Ph. D. Thesis, LANL Report LA-8919-T (1981).
58. F. Petrovich, Nucl. Phys. A251, 143 (1975).
59. G. R. Satchler, Nucl. Phys. 55, 1 (1964).
60. G. R. Satchler, Nucl. Phys. 77, 481 (1966).
61. W. G. Love, Nucl. Phys. A127, 129 (1969).
62. W. G. Love, "The (p,n) Reaction and the Nucleon-Nucleon Force," edited by C. Goodman, S. Austin, S. Bloom, J. Rapaport, and G. Satchler, (Plenum, New York 1980) p. 23.
63. W. G. Love and M. A. Franey, Phys. Rev. C24, 1073 (1981).
64. F. Petrovich, "The (p,n) Reaction and the Nucleon-Nucleon Force," edited by C. Goodman, S. Austin, S. Bloom, J. Rapaport, and G. Satchler, (Plenum, New York 1980) p. 115.

65. J. R. Comfort, DW81: modified version of the computer program DWBA70 by R. Shaeffer and J. Raynal, unpublished.
66. T. deForest Jr. and J. D. Walecka, Adv. Phys. 15, 1 (1966).
67. L. W. Fagg, W. L. Bendel, R. A. Tobin, and H. T. Kaiser, Phys. Rev. 171, 1250 (1968)
68. L. W. Fagg, W. L. Bendel, E. C. Jones Jr., and S. Numrich, Phys. Rev. 187, 1378 (1969).
69. F. Petrovich, W. G. Love, and R. J. McCarthy, Phys. Rev. C21, 1718 (1980).
70. N. Auerbach, private communication.
71. A. Bohr and B. Mottelson, Nuclear Structure (Benjamin, New York, 1969) Vol. II, p. 411.
72. W. G. Love, private communication.
73. See, for example, J. M. Eisenberg and W. Greiner, Excitation Mechanisms of the Nucleus, (North-Holland, Amsterdam, 1970).
74. G. R. Satchler, Introduction to Nuclear Reactions, (John Wiley & Sons, New York, 1980).
75. J. Raynal, computer program ECIS, unpublished.
76. J. Raynal, Phys. Rev. C23, 2671 (1981).
77. A. Bohr and B. Mottelson, Nuclear Structure (Benjamin, New York, 1969) Vol. II, p. 345.
78. A. Bohr and B. Mottelson, Nuclear Structure (Benjamin, New York, 1969) Vol. I, p. 382.
79. W. Steffen, H. Gräf, W. Gross, D. Meuer, A. Richter, E. Spamer, and O. Titze, Phys. Lett. 95B, 23 (1980).

80. W. Steffen, H. Gräf, A. Richter, A. Harding, W. Weise, U. Deutschmann, G. Lahm, and R. Neuhausen, Nucl. Phys. A404, 413 (1983).
81. B. D. Anderson et al., Phys. Rev. Lett. 45, 699 (1980); C. D. Goodman, Nucl. Phys. A374, 241 (1982).
82. G. P. A. Berg, W. Hurlimann, I. Katayama, S. Martin, J. Meissburger, J. Romer, B. Styczen, F. Osterfeld, G. Gaul, R. Santo, and G. Sondermann, Phys. Rev. C25, 2100 (1982).
83. K. E. Rehm, P. Kienle, D. Miller, R. Segal, and J. Comfort, Phys. Lett. 114B, 15 (1982).
84. G. M. Crawley, N. Anantaraman, A. Galonsky, C. Djalali, N. Marty, M. Morlet, A. Willis, and J. Jourdain, Phys. Lett. 127B, 322 (1983).
85. Y. Fujita, N. Fujiwara, S. Morinobu, T. Yamazaki, T. Itahashi, S. Imanishi, and S. I. Hayakawa, Phys. Rev. C25, 678 (1982).
86. J. R. Comfort, Optical model search code CUPID, unpublished (1982).
87. J. B. McGrory and B. H. Wildenthal, Phys. Lett. 103B, 173 (1981).
88. D. Dehnhard, private communication.
89. G. W. Hoffmann, private communication.
90. M. Franey, private communication.
91. J. Ball and C. Fulmer, Phys. Rev. 172, 1199 (1968); H. Fann, J. Schiffer, and U. Strobusch, Phys. Lett. 44B, 19 (1973).
92. L. T. van der Bijl, Ph. D. Thesis, University of Amsterdam (1982), unpublished.

93. C. Djalali, N. Marty, M. Morlet, A. Willis, V. Comparat, and R. Frascaria, Z. Phys. 298, 79 (1980).
94. C. Djalali et al., Nucl. Phys. A338, 1 (1982).
95. S. K. Nanda et al., to be published.
96. T. Chapuran, R. Vodhanel, and M. K. Brussei, Bull. Am. Phys. Soc. 22, 1022 (1977).
97. J. Speth, E. Werner, and W. Wild, Phys. Rep. C33, 127 (1977).
98. F. E. Bertrand et al., Phys. Lett. 103B, 326 (1981).
99. G. Bertsch, D. Cha, and H. Toki, Phys. Rev. C24, 533 (1981).
100. G. F. Bertsch and O. Scholten, Phys. Rev. C25, 804 (1982).
101. J. Moss et al., Phys. Rev. Lett. 48, 789 (1982).
102. F. Osterfeld, Phys. Rev. C26, 762 (1982).
103. J. Rapaport et al., Nucl. Phys. A410, 371 (1983).

APPENDIX

Cross sections, analyzing powers, and spin-flip probabilities for the entire 9 MeV M1 bump in the $^{90}\text{Zr}(p,p^-)^{90}\text{Zr}^*$ reaction at 319 MeV.

$\theta_{\text{c.m.}}$ (deg.)	$d\sigma/d\Omega$ (mb/sr)	Error (mb/sr)	$\theta_{\text{c.m.}}$ (deg.)	A_n	Error
2.03	5.609	0.148	2.03	-.024	0.034
2.54	4.057	0.186	2.54	-.027	0.032
3.04	3.747	0.183	3.04	0.044	0.036
3.55	2.581	0.004	3.55	-.024	0.037
3.81	1.743	0.143	3.81	-.047	0.054
4.06	1.860	0.084	4.06	0.014	0.032
4.32	1.352	0.073	4.31	0.056	0.001
4.57	1.434	0.085	4.57	0.057	0.043
4.83	1.072	0.074	4.83	0.083	0.056
5.07	1.051	0.085	5.07	0.075	0.060
5.34	0.937	0.076	5.34	0.220	0.067
5.85	0.905	0.075	5.85	0.182	0.070

$\theta_{\text{c.m.}}$ (deg.)	S_{nn}	Error
3.60	0.62	0.19
5.16	0.20	0.16

Cross sections, analyzing powers, and spin-flip probabilities for the 10.23 MeV M1 state in the $^{48}\text{Ca}(p,p^-)^{48}\text{Ca}^*$ reaction at 319 MeV.

$\theta_{\text{c.m.}}$ (deg.)	$d\sigma/d\Omega$ (mb/sr)	Error (mb/sr)	$\theta_{\text{c.m.}}$ (deg.)	A_n	Error
2.83	3.156	0.069	2.83	0.073	0.044
3.34	2.844	0.067	3.34	0.101	0.044
3.85	2.404	0.061	3.85	0.118	0.050
4.37	1.883	0.057	4.37	0.063	0.038
4.88	1.479	0.036	4.88	-.036	0.044
5.40	1.184	0.034	5.40	0.047	0.052
5.91	0.965	0.032	5.91	0.139	0.058

$\theta_{\text{c.m.}}$ (deg.)	S_{nn}	Error
3.60	0.44	0.07
5.14	0.21	0.06

Cross sections data for the low lying states in the $^{90}\text{Zr}(p,p')^{90}\text{Zr}^*$ reaction at 319 MeV. The systematic error in the absolute normalization of the data is about 10%.

$\theta_{\text{c.m.}}$ (deg.)	$d\sigma/d\Omega$ (mb/sr)	Error (mb/sr)	$\theta_{\text{c.m.}}$ (deg.)	$d\sigma/d\Omega$ (mb/sr)	Error (mb/sr)
$E_x=3.31$ MeV			$E_x=3.84$ MeV		
2.03	0.129	0.015	2.03	0.722	0.032
2.54	0.158	0.019	2.54	0.691	0.037
3.04	0.173	0.019	3.04	0.657	0.035
3.55	0.183	0.012	3.55	0.706	0.021
4.06	0.181	0.014	4.06	0.825	0.026
4.57	0.255	0.016	4.57	1.134	0.031
5.07	0.335	0.017	5.07	1.536	0.034
$E_x=4.23$ MeV			$E_x=4.49$ MeV		
2.03	0.277	0.023	2.03	0.154	0.020
2.54	0.245	0.025	2.54	0.105	0.020
3.04	0.241	0.023	3.04	0.113	0.018
3.55	0.300	0.015	3.55	0.120	0.011
4.06	0.360	0.019	4.06	0.122	0.013
4.57	0.451	0.021	4.57	0.129	0.013
5.07	0.611	0.022	5.07	0.168	0.013
$E_x=4.69$ MeV			$E_x=5.11$ MeV		
2.03	0.179	0.021	2.03	0.087	0.020
2.54	0.193	0.024	2.54	0.100	0.020
3.04	0.202	0.022	3.04	0.113	0.021
3.55	0.202	0.013	3.55	0.133	0.012
4.06	0.253	0.017	4.06	0.163	0.014
4.57	0.334	0.019	4.57	0.182	0.015
5.07	0.439	0.020	5.07	0.206	0.014
$E_x=5.64$ MeV			$E_x=5.78$ MeV		
2.03	0.093	0.020	2.03	0.086	0.018
2.54	0.120	0.022	2.54	0.087	0.021
3.04	0.150	0.022	3.04	0.079	0.019
3.55	0.147	0.012	3.55	0.063	0.021
4.06	0.135	0.015	4.06	0.058	0.012
4.57	0.184	0.016	4.57	0.087	0.013
5.07	0.244	0.016	5.07	0.103	0.013

Analyzing powers for the strongly excited low lying states in the $^{90}\text{Zr}(p,p^-)^{90}\text{Zr}^*$ reaction at 319 MeV.

$\theta_{\text{c.m.}}$ (deg.)	A_n	Error	$\theta_{\text{c.m.}}$ (deg.)	A_n	Error
$E_x=3.84$ MeV			$E_x=4.23$ MeV		
2.03	0.002	0.045	2.03	-.018	0.082
2.54	-.012	0.038	2.54	-.042	0.072
3.04	0.091	0.039	3.04	0.035	0.070
3.55	0.174	0.029	3.55	0.322	0.048
4.06	0.257	0.022	4.06	0.285	0.035
4.57	0.296	0.019	4.57	0.346	0.031
5.07	0.358	0.016	5.07	0.348	0.027
$E_x=4.69$ MeV			$E_x=5.64$ MeV		
2.03	-.108	0.116	2.03	-.263	0.202
2.54	-.112	0.087	2.54	0.070	0.125
3.04	-.036	0.081	3.04	0.208	0.102
3.55	0.205	0.062	3.55	0.329	0.070
4.06	0.271	0.045	4.06	0.376	0.069
4.57	0.387	0.037	4.57	0.391	0.056
5.07	0.380	0.033	5.07	0.304	0.051

Cross sections for the weakly excited E1 states in the $^{90}\text{Zr}(p,p^{-})^{90}\text{Zr}^{*}$ reaction at 319 MeV. The systematic error in the absolute normalization of the data is about 10%.

$\theta_{\text{c.m.}}$ (deg.)	$d\sigma/d\Omega$ (mb/sr)	Error (mb/sr)	$\theta_{\text{c.m.}}$ (deg.)	$d\sigma/d\Omega$ (mb/sr)	Error (mb/sr)
$E_x=6.30$ MeV			$E_x=6.42$ MeV		
2.18	0.200	0.023	2.18	0.171	0.022
2.59	0.140	0.022	2.59	0.140	0.025
2.99	0.135	0.023	2.99	0.126	0.023
3.40	0.064	0.020	3.40	0.089	0.021
3.70	0.057	0.010	3.70	0.084	0.011
4.11	0.059	0.012	4.11	0.071	0.014
4.52	0.042	0.014	4.52	0.054	0.011
4.92	0.047	0.011	4.92	0.036	0.010
$E_x=6.75$ MeV			$E_x=6.88$ MeV		
2.18	0.244	0.029	2.18	0.135	0.033
2.59	0.209	0.027	2.59	0.129	0.024
2.99	0.165	0.024	2.99	0.081	0.020
3.40	0.115	0.019	3.40	0.053	0.019
3.70	0.095	0.014	3.70	0.043	0.010
4.11	0.088	0.013	4.11	0.039	0.013
4.52	0.064	0.012	4.52	0.039	0.011
4.92	0.070	0.013	4.92	0.026	0.012

Cross sections for the fine structure states in the M1 transition region in the $^{90}\text{Zr}(p,p')^{90}\text{Zr}^*$ reaction at 319 MeV. The systematic error in the absolute normalization of the data is about 10%.

θ (deg.)	$d\sigma/d\Omega$ (mb/sr)	Error (mb/sr)	θ (deg.)	$d\sigma/d\Omega$ (mb/sr)	Error (mb/sr)
$E_x=7.53$ MeV			$E_x=7.94$ MeV		
2.18	0.078	0.038	2.18	0.082	0.026
2.59	0.087	0.033	2.59	0.057	0.029
2.99	0.074	0.024	2.99	0.096	0.027
3.40	0.092	0.026	3.40	0.046	0.024
3.70	0.082	0.015	3.70	0.076	0.015
4.11	0.093	0.016	4.11	0.072	0.014
4.52	0.102	0.015	4.52	0.076	0.015
4.92	0.113	0.015	4.92	0.095	0.015
$E_x=8.26$ MeV			$E_x=8.37$ MeV		
2.18	0.177	0.033	2.18	0.108	0.032
2.59	0.129	0.031	2.59	0.100	0.031
2.99	0.126	0.031	2.99	0.099	0.029
3.40	0.115	0.030	3.40	0.079	0.030
3.70	0.078	0.016	3.70	0.052	0.015
4.11	0.057	0.015	4.11	0.045	0.017
4.52	0.072	0.019	4.52	0.035	0.016
4.92	0.114	0.017	4.92	0.021	0.011
$E_x=8.73$ MeV			$E_x=8.98$ MeV		
2.18	0.202	0.033	2.18	0.202	0.035
2.59	0.139	0.033	2.59	0.162	0.034
2.99	0.160	0.033	2.99	0.141	0.035
3.40	0.110	0.034	3.40	0.098	0.029
3.70	0.063	0.017	3.70	0.084	0.016
4.11	0.033	0.016	4.11	0.086	0.017
4.52	0.039	0.015	4.52	0.075	0.017
4.92	0.048	0.016	4.92	0.118	0.018
$E_x=9.14$ MeV			$E_x=9.28$ MeV		
2.18	0.147	0.033	2.18	0.239	0.034
2.59	0.101	0.036	2.59	0.228	0.038
2.99	0.077	0.035	2.99	0.185	0.034
3.40	0.068	0.038	3.40	0.118	0.036
3.70	0.041	0.018	3.70	0.071	0.016
4.11	0.021	0.018	4.11	0.065	0.016
4.52	0.019	0.015	4.52	0.036	0.016
4.92	0.007	0.043	4.92	0.054	0.017

[continued]

θ (deg.)	$d\sigma/d\Omega$ (mb/sr)	Error (mb/sr)	θ (deg.)	$d\sigma/d\Omega$ (mb/sr)	Error (mb/sr)
$E_x=9.38$ MeV			$E_x=9.48$ MeV		
2.18	0.130	0.037	2.18	0.202	0.039
2.59	0.115	0.048	2.59	0.219	0.069
2.99	0.142	0.035	2.99	0.177	0.036
3.40	0.082	0.050	3.40	0.148	0.034
3.70	0.100	0.037			
4.11	0.029	0.019	4.11	0.050	0.040
4.52	0.025	0.017	4.52	0.049	0.019
4.92	0.030	0.020	4.92	0.028	0.021
$E_x=9.60$ MeV			$E_x=9.77$ MeV		
2.18	0.113	0.041	2.18	0.088	0.032
2.59	0.234	0.037	2.59	0.164	0.035
2.99	0.120	0.034	2.99	0.101	0.038
3.40	0.091	0.042	3.40	0.083	0.031
3.70	0.072	0.017			
4.11	0.060	0.017	4.11	0.028	0.016
4.52	0.039	0.020	4.52	0.019	0.016
4.92	0.023	0.016	4.92	0.045	0.017

Cross sections and analyzing powers in steps of 2 MeV without background subtraction for the high excitation region in the $^{90}\text{Zr}(p,p^-)^{90}\text{Zr}^*$ reaction at 319 MeV.

θ (deg.)	$d\sigma/d\Omega$ (mb/sr)	Error (mb/sr)	A_n	Error	$d\sigma/d\Omega$ (mb/sr)	Error (mb/sr)	A_n	Error
$E_x=8-10$ MeV					$E_x=10-12$ MeV			
2.18	10.44	0.11	0.040	0.011	8.61	0.10	0.003	0.012
2.59	9.40	0.12	0.024	0.010	8.80	0.11	0.017	0.011
2.99	8.42	0.11	0.035	0.011	8.54	0.11	0.001	0.011
3.40	6.95	0.10	0.043	0.012	7.42	0.11	0.034	0.011
3.70	5.84	0.06	0.087	0.010	5.87	0.06	0.086	0.010
4.11	4.88	0.06	0.149	0.009	5.30	0.06	0.133	0.009
4.52	4.96	0.06	0.173	0.009	5.65	0.06	0.160	0.009
4.92	4.79	0.06	0.223	0.009	5.82	0.06	0.197	0.008
$E_x=12-14$ MeV					$E_x=14-16$ MeV			
2.18	11.38	0.12	-0.017	0.010	17.43	0.15	-0.022	0.008
2.59	11.37	0.13	0.014	0.009	15.69	0.15	0.001	0.008
2.99	10.96	0.13	0.022	0.010	14.70	0.15	0.012	0.008
3.40	9.53	0.12	0.038	0.010	12.21	0.14	0.022	0.009
3.70	7.26	0.06	0.097	0.009	8.87	0.07	0.089	0.008
4.11	6.81	0.07	0.122	0.008	7.73	0.07	0.089	0.008
4.52	7.15	0.07	0.165	0.008	7.63	0.07	0.123	0.008
4.92	7.35	0.07	0.213	0.008	7.75	0.07	0.159	0.007
$E_x=16-18$ MeV								
2.79	15.48	0.17	0.020	0.007				
3.30	14.06	0.16	0.029	0.007				
3.81	12.21	0.15	0.037	0.008				
4.31	10.67	0.13	0.064	0.008				
4.32	9.03	0.07	0.041	0.006				
4.83	8.70	0.07	0.054	0.006				
5.34	8.73	0.07	0.076	0.006				
5.85	8.67	0.07	0.112	0.007				

Spin-flip probabilities and spin-flip cross sections without background subtraction for the $^{90}\text{Zr}(p,p')^{90}\text{Zr}^*$ reaction at 319 MeV.

E_x (MeV)	$\theta_{\text{lab}}=3.5^\circ$				$\theta_{\text{lab}}=5^\circ$			
	S_{nn}	Error	$d\sigma/d\Omega \cdot S_{\text{nn}}$ (mb/sr)	Error (mb/sr)	S_{nn}	Error	$d\sigma/d\Omega \cdot S_{\text{nn}}$ (mb/sr)	Error (mb/sr)
3.0	0.012	0.027	0.110	0.170				
5.0	0.021	0.020	0.105	0.170	0.034	0.028	0.131	0.110
7.0	0.060	0.020	0.358	0.130	0.100	0.022	0.475	0.110
9.0	0.200	0.015	1.610	0.150	0.137	0.015	1.020	0.110
11.0	0.130	0.020	1.200	0.150	0.176	0.014	1.500	0.120
12.0	0.170	0.030	1.900	0.220				
13.0	0.170	0.020	2.350	0.200	0.209	0.013	2.100	0.130
14.0	0.140	0.030	1.330	0.200				
15.0	0.165	0.020	2.200	0.710	0.215	0.012	2.330	0.140
16.0	0.150	0.030	1.790	0.630				
17.0	0.105	0.020	1.390	0.550	0.251	0.015	2.180	0.130
18.0	0.140	0.030	1.340	0.640				
19.0	0.170	0.040	1.410	0.490				
20.0	0.130	0.040	1.520	0.450				
21.0	0.140	0.040	1.460	0.440				
23.0	0.125	0.040	1.740	0.500				
25.0	0.178	0.060	1.520	0.450				