Scattering of polarized protons by yttrium, iron and nickel nuclei

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SCATTERING OF POLARIZED PROTONS BY YTTRIUM, IRON AND NICKEL NUCLEI

PROEFSCHRIFT

TER VERKRIJGING VAN DE GRAAD VAN DOCTOR IN DE TECHNISCHE WETENSCHAPPEN AAN DE TECHNISCHE HOGESCHOOL EINDHOVEN, OP GEZAG VAN DE RECTOR MAGNIFICUS, PROF. DR. P. VAN DER LEEDEN, VOOR EEN COMMISSIE AANGEWEZEN DOOR HET COLLEGE VAN DEKANEN IN HET OPENBAAR TE VERDEIDIGEN OP DIENSDAG 16 MEI 1978 TE 16.00 UUR

DOOR

JOSEPH PETRUS MARIA GERARDUS MEISSLSEN

geboren te Brunssum

(Druk. Willink Heerlen)
DIT PROEFSCHRIFT IS GOEDGEKEURD
DOOR DE PROMOTOREN

PROF. DR. O.J. POPEMIA
EN
DR. B.J. VERHAAR.
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CHAPTER I

INTRODUCTION AND SUMMARY

The study of nuclear scattering is an important source of information on nuclear structure. Especially the scattering of protons and neutrons should be sensitive to details of the nuclear wavefunctions because of the weak absorption of nucleons in the nucleus. A more stringent test of the spin dependent parts of the interaction involved in the scattering process is provided by experiments with polarized beams. In this thesis we present results of scattering experiments performed on yttrium and some iron and nickel isotopes with polarized proton beams at energies around 20 MeV.

In a scattering experiment a beam of particles is incident on a target foil. An incident particle may be scattered by a target nucleus. If the internal energy of the target nucleus remains unchanged the scattering is called elastic, whereas in an inelastic process a definite amount of the energy of the incident particle is used to excite the nucleus. A schematic diagram of a scattering process is given in figure 1.1, where the scattering angle $\theta$ and the normal to the scattering plane

![Schematic diagram of a scattering process.](image)

are defined; $\vec{k}_i$ is the momentum of the incident particle and $\vec{k}_f$ the momentum of the scattered particle. As shown in figure 1.2 the energy spectrum of the scattered particles at a fixed scattering angle reveals peaks which correspond to the energies of the excited states of the nucleus. The excitation strength of such a state is characterized by the differential cross section $d\sigma(\theta)/d\Omega$ given by the number of counts in a peak, normalized with respect to the number of incident particles, the number of target nuclei (atoms/cm$^2$) and the solid angle in which the scattered protons are detected. An additional quantity, the analysing power $A(\theta)$ can be measured
in experiments with polarized beams. In these experiments the differential cross section is given by:

$$\frac{d\sigma(\theta)}{d\Omega} = \left[1 + \vec{P} \cdot \vec{n} \ A(\theta)\right] \frac{d\sigma(\theta)}{d\Omega}_{\text{unpol}}$$  \hspace{1cm} (1.1)$$

where $\vec{P}$ is the polarization of the incident proton beam. Thus the analysing power is given by (1.2) as the difference in cross sections for the opposite spin directions, normalized to their sum and to the degree of polarization of the beam:

$$A(\theta) = \frac{1}{P} \times \frac{\left(\frac{d\sigma}{d\Omega^-} - \frac{d\sigma}{d\Omega^+}\right)}{\left(\frac{d\sigma}{d\Omega^-} + \frac{d\sigma}{d\Omega^+}\right)}$$ \hspace{1cm} (1.2)$$

The experimental arrangement used in our experiments to measure the angular distributions of the differential cross sections and analysing powers is described in chapter 2. In chapter 3 we report upon our experimental procedure and data analysis.

Information about an excited nuclear state like spin, parity and details of the wavefunction is obtained by comparing the angular distribution of the experimental differential cross section and analysing power with predictions from theoretical models. Usually the Distorted Wave Born Approximation (DWBA) is used to describe the reaction at the bombarding
energies and for the target nuclei investigated in our experiments. In this approximation it is assumed that the elastic scattering is far out the most important process that occurs, so the inelastic processes can be treated as perturbations. The relative motion of the proton before and after the inelastic event is then described by elastic scattering wavefunctions that are calculated in the optical model. In this model the interaction of the projectile with the target nucleus is approximated by a spherical potential which consists of a central complex part and a spin-orbit part.

Two different models are currently used for the description of the nuclear states. The collective model treats the nucleus as a whole and the excited states are considered either as collective vibrations of the nuclear surface or as collective rotations of a permanently deformed nucleus. Then the optical potential contains nonspherical parts which give rise to the excitation of the nucleus. With this macroscopic DWBA description only information about the collective behaviour of the nucleus is obtained. A more detailed picture of the nuclear states is provided by the shell model, where the nuclear states are described in terms of the motion of the individual nucleons. In the microscopic DWBA, based on this model, the interaction which is responsible for the excitation is then composed of effective two-body interactions between the projectile nucleon and the target nucleons. The formalism of the optical model and the DWBA is summarized in chapter 4.

The macroscopic DWBA has been shown to reproduce the shape of the angular distributions reasonably well for an extensive number of cases, (Sa70). Generally it is used as a reliable tool to determine the orbital angular momentum transferred to the target nucleus. Moreover deformation parameters obtained in this way agree well with the values determined from other types of experiments. However, it is also known that the analysing powers for the first excited \(2^+\) states of some closed neutron shell nuclei \((N=28, N=50)\), measured with protons of about 20 MeV, are found to be much larger as compared to neighbouring nuclei with an open neutron shell or as compared to the next excited state with the same target spin \((G167, H669, G169)\). This distinct behaviour can be described by the macroscopic DWBA only if the spin-orbit term in the optical model is deformed two to three times as much as the central part, whereas

\[ \text{A list of references is given at the end of each chapter.} \]
for neighbouring nuclei or higher excited states no significant higher spin-orbit deformation is needed. There are indications that this effect might depend on the projectile energy. Experiments at 30 MeV and at 40 MeV show that the analysing power of the first excited $2^+$ state in $^{54}\text{Fe}$ differs not significantly from analysing powers of other $2^+$ states in this mass region (Ka70, Fr67); also in a recent experiment at 30 MeV on $^{90}\text{Zr}$ and $^{92}\text{Zr}$ no differences are found between the first excited $2^+$ states of these nuclei (Sw76).

To investigate a possible energy dependence we performed a series of experiments on $^{54}\text{Fe}$, $^{56}\text{Fe}$, $^{58}\text{Ni}$, $^{60}\text{Ni}$ and $^{62}\text{Ni}$ with proton energies around 20 MeV. The results of these experiments are presented in chapter 5 together with a macroscopic DWBA-analysis. In the data on the first excited state in $^{54}\text{Fe}$ the anomalous behaviour shows up indeed and also its energy dependence is clearly ascertained. With ratios of the deformation of the spin-orbit to the central part ranging from 1.5 at 24.6 MeV to 3.0 at 17.2 MeV good fits are obtained for the measured analysing powers. However, at an energy of 15.3 MeV this parametrization in terms of a larger spin-orbit deformation fails and we can not achieve a reasonable fit to our data.

One can argue whether full microscopic calculations would be more successful in describing the above mentioned effects. Investigations of Raynal (Ra71) in this direction emphasized the structural differences in the excited states, which for the closed neutron shell nuclei are expected to be dominated by proton excitations. Since the two-body spin-orbit interaction is strongest between like nucleons indeed larger spin-orbit deformations, in the macroscopic DWBA, are expected for protons exciting protons than for protons exciting neutrons. So far the microscopic analyses do not succeed very well in reproducing the data. This might be partly due to the use of inadequate expressions for the nuclear wavefunctions. Also the information about the non-central parts of the effective nucleon-nucleon interaction is scarce at present.

To study the microscopic DWBA and to acquire information about the effective two-body spin-orbit force we performed an experiment on $^{89}\text{Y}$. The ground and first excited states of $^{89}\text{Y}$ are generally believed to be well described by simple zero order shell model wave functions. The elastic scattering to these states is thus regarded as a good test for the effective interaction. The experimental results and the
macroscopic and microscopic DWBA analyses of this experiment are
given in chapter 6.

The microscopic calculations reveal that only a reasonable fit to both
the differential cross section and the analysing power is obtained if
along with the contribution of a single-particle excitation also a
macroscopic core-polarization component is taken into account. In
general the influence of the core-polarization contribution is so
large that it obscures the information on the effective nucleon-nucleon
interaction. Definite conclusions should therefore await a microscopic
description of the core-polarization contribution.

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CHAPTER II

THE EXPERIMENTAL ARRANGEMENT FOR MEASUREMENTS WITH A POLARIZED PROTON BEAM

2.1 Introduction

The experiments described in this thesis have been performed with the polarized proton beam of the Eindhoven University of Technology cyclotron (Ve62). In the sections 2.2, 2.3 and 2.4 the polarized proton source, the injection of the polarized beam into the cyclotron and the beamguiding system are described respectively. With this set up we have succeeded to obtain a proton beam of, on the average, 20 nA on the target with an energy resolution \Delta E/E of about 4 \times 10^{-3} and a polarization of 75\%. For our experiments we used a scattering chamber with a multi-detector system, which is described in section 2.5. The accessory electronics will be discussed in section 2.6.

2.2 The polarized proton source

For the production of polarized protons an ion-source of the atomic beam type is used. The method employed in such a source stems from an idea by R. Fleischmann (Cl56). A general review about polarized ion sources has been given, among others, by Haebeler (Ha67), whereas progress reports on the construction of these sources were given by Fleischmann (Fr65), Glaerisch (Gl70), Donally (Do70) and Clegg (C175) at the five-annual symposia on polarization phenomena. So here we shall only briefly review the general principles of the atomic beam method and at the same time report the essential parameters of the Eindhoven source. The basic concepts of this particular source originated from the work of Van der Reide (Ve72).

A schematic diagram of the ion source is shown in figure 2.1. On the basis of this diagram we will outline the principle of this type of source. In the dissociator hydrogen atoms are produced by dissociating molecular hydrogen in a high frequency discharge. This discharge is excited inside a water cooled quartz tube. Through an orifice at the end of this tube hydrogen atoms can escape from the
discharge forming, after collimation by a small diaphragm (the skimmer), an atomic beam that enters into a sextupole magnet.

The electron and proton spin of a hydrogen atom decouple in a strong magnetic field in the way schematically shown in figure 2.2. In the inhomogeneous field of the sextupole magnet the atom experiences a force of which the direction depends on the orientation of the electron spin. The atoms with electron spin up are focussed towards the axis of the magnet, whereas the atoms with electron spin down are deflected from this axis. Thus a Stern-Gerlach separation between the hyperfine structure states 1+2 and 3+4 is obtained. With an appropriate diaphragm behind the sextupole magnet the components 1+2 are selected, yielding an atomic beam which is polarized in the electron spins. Next the atomic beam is polarized in the proton spins in a weak field transition unit where hyperfine transitions from substate 1 to substate 3 are effectuated following a method suggested by Abragam and Winter (Ab62). These transitions are induced by means of an alternating magnetic field directed along the beam axis and a static field perpendicular to the axis. After the beam has passed the transition unit it practically only contains particles in substate 2 and 3. The atomic beam is then ionized by electron bombardment in the strong magnetic field of a solenoid and thus a longitudinal polarized proton beam is obtained. By a system of electrodes this beam is extracted from the ionizer and accelerated to 5 keV since the ionizer is operated at +5kV against earth. Finally the spin direction is rotated by 90° in crossed elec-
Fig. 8.2
Energy-level diagram of the hydrogen atom in a magnetic field. The energy is given in units $E = 5.82 \times 10^{-6}$ eV and the magnetic field in units $B = 5.07 \times 10^{-2}$ T.

Electric and magnetic fields, a so-called Wien filter. By switching both the electric and magnetic fields, the spin direction of the protons can be reversed. The essential parameters of our source are listed in Table 2.1. The source delivered about 3 μA on a cup just behind the ionizer.

<table>
<thead>
<tr>
<th>Table 2.1 Parameters of the atomic beam source.</th>
</tr>
</thead>
<tbody>
<tr>
<td>Dissociator:</td>
</tr>
<tr>
<td>warm-resist quartz tube</td>
</tr>
<tr>
<td>inner diameter 6 mm</td>
</tr>
<tr>
<td>&quot; length 275 mm</td>
</tr>
<tr>
<td>inner diameter nozzle 0.6 - 1.8 mm</td>
</tr>
<tr>
<td>r.f. discharge</td>
</tr>
<tr>
<td>frequency 80 MHz</td>
</tr>
<tr>
<td>power dissipated by oscillator 240 W</td>
</tr>
<tr>
<td>pressure in the discharge 2 torr</td>
</tr>
<tr>
<td>Separator:</td>
</tr>
<tr>
<td>conical shaped</td>
</tr>
<tr>
<td>hole diameter aperture 3 mm</td>
</tr>
<tr>
<td>distance to dissociator 12 mm</td>
</tr>
<tr>
<td>&quot; to sextapole 24 mm</td>
</tr>
<tr>
<td>Sextapole:</td>
</tr>
<tr>
<td>sextapole magnet</td>
</tr>
<tr>
<td>distance opposite pole tips 10 mm</td>
</tr>
<tr>
<td>field at pole tips 0.6 T</td>
</tr>
<tr>
<td>length of sextapole 740 mm</td>
</tr>
<tr>
<td>Weak field transition unit:</td>
</tr>
<tr>
<td>static field 0.5 T</td>
</tr>
<tr>
<td>alternating magnetic field amplitude 0.1 mT</td>
</tr>
<tr>
<td>frequency 17 MHz</td>
</tr>
<tr>
<td>Ionizer:</td>
</tr>
<tr>
<td>magnetic field 0.125 T</td>
</tr>
<tr>
<td>length ionizer 190 μm</td>
</tr>
<tr>
<td>operated at 5 kV against earth</td>
</tr>
</tbody>
</table>
2.3 The injection and the acceleration of the beam

Through a system of focusing and steering elements the polarized proton beam is transported to the entrance of a trochoidal injection system, by which the beam is radially injected into the cyclotron. This system is an exact copy of a second version of the Saclay injector, the principles of which have been described by Beurtey and Durand (Be67). The main part of the injector, which is shown in figure 2.3, consists of two pair of copper electrodes which are oppositely charged. During operation of the injector the voltages on the two pairs of electrodes are adjusted in such a way that the Lorentz force acting on the protons is balanced electrostatically. By the construction of the injector an electric field results that together with the magnetic field possesses strong focusing properties. With this system we were able to transmit about 70% of the incoming beam to the centre of the cyclotron.

Notwithstanding this high transmission it first seemed impossible to extract any accelerated beam from the cyclotron. One should realize that the electrodes of the injector form a vertical aperture of 8 mm through which the beam accelerated by the cyclotron has to pass about 300 times. A typical diagram of the accelerated beam current against the estimated distance from the centre is shown in figure 2.4. These currents were measured with a target of 10 mm height and 3 mm length having an insulating strip of about 1 mm in the middle which divided the target in an upper and lower part. This target was attached on the "weather-side" of the injector system and could be moved outwards.
In the first revolutions the beam intensity strongly diminishes and up to about 50 mm a vertical oscillatory pattern of the beam shows up. It is also seen that the middle of the beam does not coincide with the median plane of the injector and that the current measured on the lower part seems to cut off at about 60 mm. From that distance the total intensity diminishes slowly and at 140 mm a current of approximately 4 nA is left. At radii larger than 200 mm it was impossible to find any measurable current. Similar results were obtained with a diaphragm attached to the conventional (unpolarized) ion source. This diaphragm had a vertical aperture similar to that of the injector. It was only possible to accelerate the beam further outward if one or both of the inner two pair of correction coils were excited asymmetrically, as suggested by Poussard (Po74). In this way, finally, mean values during our experiments of 50 nA of internal beam and 35 nA of extracted beam were obtained.

2.4 Beam transport system

The beam extracted from the cyclotron was transported by means of bending magnets, steering magnets and quadrupole lenses (Ha70), as shown in figure 2.5, to the scattering chamber located at experiment station IV. At various positions beamstops were used to measure the
beam current. The two $45^\circ$ bending magnets MB4 and MC1 and the quadrupole lenses between these were operated in a doubly achromatic mode (Sc73). The slits between the quadrupole lenses QE2 and QE3 were adjusted to form an aperture of 3 by 3 mm. With appropriate settings of the beam transport elements this aperture was imaged on a target foil in the scattering chamber, giving a beamspot of about 1.5 mm in diameter.

2.5 The scattering chamber and polarization monitor

In figure 2.6 an outline of the scattering chamber and polarization monitor is shown. The scattering chamber has an inside diameter of 560 mm and an inside height of 90 mm. In the centre a target of 10 mm in diameter is mounted on a nickel frame of 14 mm by 26 mm. To suppress slit-scattered protons we placed at the entrance of this chamber two diaphragms of successively 6 and 8 mm in diameter.

The protons scattered by the target foil are detected by two arrays of four 3 mm Si (Li) detectors, each mounted in the median plane of the scattering chamber and adjustable at different angle settings. The detectors are placed obliquely, under an angle of $45^\circ$ with respect to the incoming scattered
protons. The array used for the detection of protons scattered in the forward region (20°-90°) -the forward block- is placed at the right side of the beam. The detectors in this array are 5° apart at a distance of 250 mm from the target. By a diaphragm in front of each detector a horizontal angular acceptance of about 1° results. The second array, which is used for the more backward angles (60°-165°) -the backward block- is placed at the left side of the beam at a distance of 125 mm from the target; here the detectors are 10° apart and an identical diaphragm as used in the forward block defines a horizontal angular acceptance of about 2°. Channels in the detector blocks, in front of the angle defining slits, restrict the view of the detectors to the direct environment of the target. Permanent magnets are placed at the top and bottom of these channels to sweep away the secondary electrons arriving from the target.

In the vertical plane through the beam two 500 μm Si detectors with aluminium degraders in front are placed at a forward scattering angle of 45°. The total number of elastically scattered protons detected in these "out of plane" detectors is used for normalization purposes.

The degree of polarization of the beam is continuously monitored in the polarization monitor: a separate small scattering chamber located downstream the main chamber. Here the proton energy is degraded by means of an aluminium sheet to an energy of 15.5 MeV. After collimation these protons are scattered from a thick carbon target (about 0.3 mm). The
elastically scattered protons are detected in the horizontal plane on both sides of the beam at a scattering angle of 52.5°. From the well-known analysing power of ¹²C in this energy region (Me76) the polarization of the beam can be deduced.

2.6 The electronic equipment

The detection system, which is the same for each of the eight detectors, is shown schematically in figure 2.7. In a semi-conductor detector an incident proton creates by ionization electron hole pairs. Due to the electric field across the detector these charge carriers are collected, which results in a charge pulse proportional to the energy loss of the proton. After preamplification and conversion into a voltage pulse the signal is fed into a pulse-shaping circuit and a fast logic circuit. In this fast circuit a discriminator generates a logical pulse whenever the filter-amplifier output exceeds a minimum pulse height, corresponding to an energy of about 1 MeV. This logical pulse is fed into a routing
unit, which may enable a linear gate. Next the amplified and stretched analog signal is sent through a mixer to an analog to digital converter (ADC) where the pulse is converted into a channel-address proportional to the input pulse amplitude. In all our experiments a conversion gain of 1024 channels is used of which by subtraction of 512 channels only the upper half of a spectrum is stored in the memory of a PDP-9 computer.

Instead of the fast logic circuit, described above, we used in our first experiments at 17.2 MeV, 20.4 MeV and 24.6 MeV a discriminator with a relatively large dead time. Since in the following experiments larger beam currents were obtained a pile up inspection was included, for which the present much faster discriminator was needed.

On the basis of Figure 2.8 we will point out some further features of our electronic set up, the main functions of which are processed by the routing unit and the control unit.

The routing unit handles the following five tasks:

a) After the routing unit has been triggered by a discriminator pulse a gate signal is sent to the corresponding linear gate and stretcher which enables the input.

b) Within 20 ns after a discriminator pulse is received by the routing unit any other discriminator pulse will be disabled by means of internal input gates. Moreover a check is performed on accidental coincidences between two detector channels. If such a coincidence is detected the gate signals to the linear gates are suppressed.

c) A dead time signal is given as soon as a discriminator pulse is registered. This signal lasts until a ready signal, indicating that the event is stored in the memory of the computer, is received. The routing unit is reset by this ready signal or by an internally generated signal, if in a preset time the ADC has not started a conversion.

d) The pile up protection works as follows. Whenever a pulse is accepted for analysis a start signal is fed into the start input of a time to pulse height converter (TPHC). During the analysis of a pulse the discriminator pulses, which correspond with an event in the same channel, are transmitted through the routing unit to the stop input of the TPHC. When a time interval smaller than 2μs is measured a prohibitive signal
is fed into the ADC. We established that for a longer interval no distortion of the spectrum occurred by pile up effects.

e) The detector which has triggered the routing unit is identified by three routing bits.

The control unit handles the following four tasks:

a) After a preset number of counts is detected by the two "out of plane" detectors a signal is sent to the polarized ion source where the direction of the polarization is changed by reversing the magnetic and electric fields in the Wien-filter.

b) This reversal is checked together with a number of essential operating conditions of the ion source and the injection system. An experimental run is stopped whenever one of these conditions is not fulfilled.

c) A routing bit is generated which indicates the direction of the polarization.

d) After a preset number of spin direction reversals an interrupt is sent to the routing unit and the PDP-9 computer by which an experimental run is stopped.

The four routing bits, three for the detector identification and one for the spin direction, together with nine bits from the ADC output give a division of the PDP memory into 16 x 512 channels. The number of elastically scattered protons detected by each of the two "out of plane" detectors and each of the two detectors in the polarization monitor are registered for the two spin orientations in scalers. After an experimental run the total number of counts in each scaler together with the accumulated spectrum of 16 x 512 channels are stored on DEC-tape. For further analysis the spectra may be transmitted to a B7700 computer.

In addition to the above mentioned actions several others are possible which can give us some insight in the status of the experiment by means of a plot or a display of the spectrum and a first estimate of the peak contents in a rough peak scan. All these actions are governed by the experiment monitor system KORAM (Ra77).
Fig. 2.8 Block schema indicating the function of the routing unit and the control unit, as explained in the text.
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CHAPTER III

PROCEDURE FOR MEASUREMENTS WITH THE POLARIZED PROTON BEAM

3.1 Introduction

In this chapter we will describe and discuss the procedure we followed in measuring the analysing powers and differential cross sections. Our general procedure is outlined in section 3.2. In section 3.3 the influence of possible sources of systematic errors caused by a non-proper spin flip, uncertainties in the actual scattering angles and uncertainties in the beam polarization are estimated. At least with respect to the inelastic scattering data, these errors are found to be negligible as compared to the statistical errors. In section 3.4 we give our results on the analysing powers in the elastic scattering from $^{12}$C, which we collected alternately with the data on the iron, nickel and yttrium nuclei. Generally these data are found to be consistent with those from the literature. Finally in section 3.5 a description of the analysis of the energy spectra is given.

3.2 Experimental procedure

3.2.1 The measurement of the analysing power and differential cross section

From the definition of the analysing power, given in chapter 1, it is easily shown that one can determine the analysing power from the number of particles recorded by two detectors placed on opposite sides of the beam at the same scattering angle, or from the number of particles recorded by one detector at a fixed scattering angle in two separate runs with opposite spin directions. As we intended to measure angular distributions over a wide angular region we decided upon the eight single detector set up described in section 2.5.

With this set up measurements with forward and backward positioned detectors were done simultaneously. Generally the following measuring scheme was carried out:
Forward block, detectors at backward block, detectors at

<p>| | | | | | |</p>
<table>
<thead>
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<th></th>
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</tr>
</thead>
<tbody>
<tr>
<td>a.</td>
<td>20°, 25°, 30°, 35°;</td>
<td>b.</td>
<td>90°, 100°, 110°, 120°;</td>
<td></td>
<td></td>
</tr>
<tr>
<td>b.</td>
<td>40°</td>
<td>c.</td>
<td>130°</td>
<td></td>
<td></td>
</tr>
<tr>
<td>c.</td>
<td>60°</td>
<td>d.</td>
<td>135°</td>
<td></td>
<td></td>
</tr>
<tr>
<td>d.</td>
<td>80°</td>
<td>e.</td>
<td>165°</td>
<td></td>
<td></td>
</tr>
<tr>
<td>e.</td>
<td>80°</td>
<td>f.</td>
<td>90°</td>
<td></td>
<td></td>
</tr>
<tr>
<td>f.</td>
<td>80°</td>
<td></td>
<td>65°</td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

In the first four runs the target was positioned at an angle of 45° with the beam direction, by which the forward scattered protons were detected in transmission mode and the backward scattered protons in reflection. In the last two runs this situation was reversed by having the target rotated over 90°.

The differences in counting rate between the forward and backward positioned detectors, due to the decrease of the cross sections with increasing scattering angle, was partly compensated by the four times larger solid angles for the detectors in the backward block as compared to those in the forward block.

Usually we made several runs at the same angular setting. This precaution was taken, to avoid that a temporary decrease in the quality of an energy spectrum resulted in a bad or even a worthless measurement. In each run the energy spectra were taken alternately for the opposite spin directions with periods determined by a preset number of elastically scattered protons detected by the two "out of plane" detectors. So a measure of the number of incident protons times the number of target nuclei was found. To obtain the angular distributions of the differential cross section the total amount of detected protons in one run in both "out of plane" detectors together was taken to normalize the measurements at different angular settings. This normalization procedure was preferred over a normalization based on the integrated beam current because in that way the accuracy might have been affected seriously by inhomogeneities of the target. Since the elastic scattering cross section decreases rapidly with increasing scattering angle the counts registered by one out of plane detector should sensitively depend on the precise scattering angle. For this reason we used two detectors situated at the same scattering angles above and below the beam. Thus to first order changes in the beam position or beam direction did not affect the total number of counts detected. In several experiments the ratio of the
normalization factors for the differential cross section resulting from the integrated beam current and that from the out of plane detectors have been compared. Deviations smaller than a few per cent have been observed for measurements with a fixed target position, whereas larger deviations, up to 10%, occurred between measurements with a different position of the target.

An additional check on the normalization was provided by the measuring scheme, by which at various stages of the experiment the backward detectors comprised an angular region which had been measured before. From inspection of these measurements no systematic discrepancies emerged.

3.2.2. Beam tuning and determination of the energy of the beam.

By the beam transport system, as described in section 2.4, the proton beam was guided to the scattering chamber. An aperture of 3 mm in diameter, at the target position, was used to steer the beam through the scattering chamber and polarization monitor. To that end, quadrupole and steering magnets were adjusted in such a way that the current measured on this aperture was minimized while the current on the beam stop behind the polarization monitor was kept at a maximum. In all our experiments the condition was imposed that this aperture intercepts one per cent of the beam intensity at most.

Because of the uncertainty about the precise radius at which the beam is extracted from the cyclotron, the real beam energy can be different from the energy at which the cyclotron is set. We therefore determined, at the beginning of each series of experiments, the energy of the beam by the "cross-over" method (Ba64), (Sm64). In this method the "cross-over" angle is determined at which the protons scattered from the hydrogen in a polyethylene target have the same energy as those scattered inelastically from $^{12}$C, thereby producing the first excited state of this nucleus at an excitation energy of 4.433 MeV. Since this angle depends strongly on the incident proton energy it is an accurate measure for this energy. The measurements were performed with one of the detectors of the forward block. To eliminate errors due to a slight misalignment of the scattering chamber with respect to the beam axis, the cross-over angle was determined from measurements on both sides of the beam. Both angles generally were found to be equal within 0.1°. The beam energy
was deduced from the average of these values. Normally an accuracy of about 0.1 MeV was attained.

3.2.3 Targets

As targets self-supporting foils were used, which were obtained from A.E.R.E. Harwell. The thicknesses of these targets and their purities are given in table 3.1.

Table 3.1. Thickness and purities of the targets

<table>
<thead>
<tr>
<th>nuclide</th>
<th>thickness in mg/cm²</th>
<th>purity in %</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{54}$Fe</td>
<td>1</td>
<td>97.48</td>
</tr>
<tr>
<td>$^{56}$Fe</td>
<td>1</td>
<td>99.97</td>
</tr>
<tr>
<td>$^{58}$Ni</td>
<td>1</td>
<td>99.91</td>
</tr>
<tr>
<td>$^{60}$Ni</td>
<td>1</td>
<td>99.85</td>
</tr>
<tr>
<td>$^{62}$Ni</td>
<td>1</td>
<td>99.39</td>
</tr>
<tr>
<td>$^{89}$Y</td>
<td>3</td>
<td>99.0</td>
</tr>
</tbody>
</table>

3.2.4 Scattering by a polyethylene and a mylar target

At each angular setting the measurements on a specific target were "sandwiched" between similar measurements with a polyethylene and mylar target of 2 mg/cm² and 1 mg/cm² thickness respectively.

Carbon and oxygen contaminants, which showed up in all our spectra more or less, were easily corrected for with the data from the polyethylene and mylar targets. From the position of the peaks corresponding to the elastic and inelastic scattering by $^{12}$C we calibrated the energy scale. Moreover for scattering angles smaller than 45°, also an accurate determination of the scattering angle could be made from the kinematical shift of the $^{12}$C-peak. A further check on our experimental procedure was performed, by comparing our results on the elastic scattering from $^{12}$C with results from the literature.
3.3 Sources of systematic errors

3.3.1 Non proper spin flip

The direction of the spin is reversed by reversing the electric and magnetic field in the Wien filter. However, if the Lorentz-force acting on the protons is not exactly cancelled by the electric force, this Wien filter also acts as a steering element in the beam transport system between the ionizer and the injector. In reversing the spin direction it is thus possible that the position and direction of the proton beam at the entrance of the injector is affected slightly. If such changes in position and direction are conveyed to the scattering chamber on the target one introduces a so-called non-proper spin flip. These non-proper spin flip components may give rise to considerable systematic deviations (Oh75).

For this reason our first experiments, on the iron and nickel nuclei, have been done by switching the hyper fine transition unit on and off. Thus we measured alternately with a polarized and unpolarized beam. However, measuring the asymmetries between the number of counts detected in two runs with opposite spin direction is a more efficient method. In a two times shorter measuring period about the same accuracy is obtained as compared to an experiment in which the polarization is switched on and off. Therefore it was checked thoroughly that by a spin reversal by means of the Wien-filter no noticeable variations in beam position or direction occurred at the target, not even when the Wien filter was deliberately used as a steering element in such a way that the total beam current was diminished by about a factor of two. So all our other experiments which we will report on were done by switching the polarization direction of the beam with the Wien filter.

In effect these checks on beam variations were repeated in each experiment. Small in plane variations should strongly affect the positions of the protons scattered by the hydrogen in the polyethylene and mylar energy spectra. Variations perpendicular to the scattering plane were checked by the ratios of the counts collected by each of the two "out of plane" detectors in the spin up and spin down states. From both checks we estimated upper values for a variation of the beam direction of 0.1° or for a beam displacement of 0.1 mm, at the target position, presupposing that the direction and position variations did not cancel out.
3.3.2 The actual scattering angle

The detectors and the diaphragms in front of them always were placed at fixed positions in the detector blocks. The angular settings of these blocks were read, with an estimated error of 0.1°, from a graduation on the wheels, on which these blocks were attached. The angular distance between the detectors in each block was verified by the observed kinematical shift of the hydrogen peak in the polyethylene spectra measured for scattering angles between 20° and 60°. Thus the actual scattering angles were found to be correct within 0.1° if the direction of the beam axis coincided with the zero line of the scattering chamber.

From the determination of the energy of the beam (Cf. section 3.2.2.) slight misalignments appeared of at most 0.1°. However, instable elements in the ion-source, the injector, the cyclotron or the beam guiding system could result in random variations or a slow drift in the position or direction of the beam. Since the spin direction was changed about every twenty seconds such variations should unlikely affect the resulting differential cross sections or the analysing powers, though the actual scattering angle to which these measurements corresponded might deviate from the nominal settings. Until about 40° the scattering angles were calculated in each experiment from the positions of the 12C-elastic, the 12C-inelastic and the 1H peak in the energy spectra obtained for the scattering from the mylar and polyethylene targets. These calculations revealed deviations from the nominal settings of 0.2° at most.

Deviations of this order will only affect the elastic scattering cross section and analysing power in those region where rapid changes in the angular distributions are observed, such as for the cross section at forward angles.

3.3.3 Uncertaintly in the polarization of the beam

The beam polarization was continuously monitored in the polarimeter (Cf. section 2.5.). Fluctuations in the beam polarization were automatically corrected for since we determined the mean beam polarization in an experimental run from the total number of elastically scattered protons registered by each of the polarization monitor detectors for each spin direction.

It is a disadvantage, however, of a single detector system that the measured analysing power and differential cross section depend in first
order on the difference in the beam polarizations for the up and down
direction.

If we define the polarization $\mathbf{P}^+$ in the spin up and $\mathbf{P}^-$ in the spin
down state by:

\[
\begin{align*}
\mathbf{P}^+ &= \mathbf{P} + \delta\mathbf{P} \\
\mathbf{P}^- &= \mathbf{P} - \delta\mathbf{P}
\end{align*}
\]  
(3.1)

it is easily shown that to first order in $\delta\mathbf{P}$ the analysing power $A(\theta)$
and the differential cross section $\sigma(\theta)$ are given by:

\[
\begin{align*}
A(\theta) &= (1 + A_m \delta\mathbf{P}) A_m \\
\sigma(\theta) &= (1 - A_m \delta\mathbf{P}) \sigma_m
\end{align*}
\]  
(3.2)

where the analysing power $A_m$ and the differential cross section $\sigma_m$
are measured, assuming $\delta\mathbf{P}$ to be zero; the upper sign and lower sign
respectively refer to scattering to the left and to the right with
respect to the beam direction.

The polarization of the beam has to be determined from the following
relations:

\[
\begin{align*}
L^+ &= n N' \Omega_L \sigma_o (1 + P^+ A) \\
R^+ &= n N' \Omega_R \sigma_o (1 - P^+ A) \\
L^- &= n' N' \Omega_L \sigma_o (1 - P^- A) \\
R^- &= n' N' \Omega_R \sigma_o (1 + P^- A)
\end{align*}
\]  
(3.3)

where $L$ and $R$ are the number of elastically scattered protons detected
by the left and right detector, $n$ is the number of incident protons,
$N$ the target thickness (atoms/cm$^2$), and $\Omega_L$, $\Omega_R$, $E_L$, $E_R$ are the solid
angles and efficiencies for the left and right detector, respectively.
Assuming that the effective target thickness $nN$ was the same for the
two counting periods we found values for $\delta\mathbf{P}$ ranging from $-3\%$ till $5\%$.
However, variations in the effective target thickness could readily
account for the supposed differences in the beam polarization or for
the fluctuations observed in it. We remark that differences in the
effective target thickness for opposite spin directions might be caused
by a small non-proper spin flip component as the measuring periods for
both spin directions were determined by the effective target thickness
of the target placed in the scattering chamber.
Since we could not deduce the exact value of $\delta P$ from our measurements, we omitted the correction (3.2.) and took the mean polarization derived from the geometrical means of (3.3), which is to first order correct in $\delta P$. This correction might only become significant for the elastic scattering data, where it could give rise to systematical errors of about the same order of magnitude as the statistical ones.

3.4 The analysing power for $^{12}\text{C}(p,p)$

In figure 3.1. we give the results for the analysing powers for the elastic scattering from $^{12}\text{C}$, which we have measured alternately with the data on the iron, nickel and yttrium nuclei. These data may be compared to the numerous data from the literature. The data of Meyer (Ma76), Martin (Ma76) and Craig (Cr66) cover the energy region of 11.5 MeV up till 28 MeV. From these we learn that a consistency check over the whole angular region is only incidently possible because of the strong energy-dependence in this reaction. For this reason knowledge of the exact incident proton energy and the energy resolution of the beam becomes of major importance if comparisons are to be made. The analysing powers, given in figure 3.1., all agree qualitatively with the above mentioned data. At those scattering angles where the energy dependence is not too strong the absolute values are also quantitatively consistent. From these values the absolute calibration of the beam polarization is found to be correct within 5%.

3.5 Analysis of the energy spectra

On the average a total of twenty spectra of 8192 channels were stored in one experiment on one target at one energy. These spectra were analysed off-line with the aid of a PDP-9 computer. This analysis is briefly described below.

First of all the spectra measured with a specific detector at a specific scattering angle were compared to each other, either by a visual inspection on a display or by a chi-square criterion. In the latter case the chi-square value of each spectrum with respect to that of each other spectrum has been evaluated. From these values we decided which spectra could be added. Eight new data blocks were
Fig. 3.1 Analysing power for the elastic scattering of protons by $^{12}\text{C}$. 
formed, out of the added spectra and the corresponding polyethylene and mylar spectra, that contained all the information collected with a specific detector at a specific scattering angle. These data blocks were used in our further analysis which was completed in two stages.

In the first stage the polyethylene and mylar spectra were considered by means of a display. The peaks owing to the reactions $^{12}\text{C}(p,p)$, $^{12}\text{C}(p,p')$ ($E_x = 4.433$ MeV), $^1\text{H}(p,p)$ and $^{16}\text{O}(p,p)$ were localized successively. From the estimated width of a peak two regions around the peak were defined to which a rectilinear background was fitted. These regions were taken the same in the spectra measured for the two spin states. The peak region and the background region could be adjusted in an interactive way. After the background had been subtracted, peak properties like the position and content were calculated together with the total content and the asymmetry for the different spin states. From the position of the two carbon peaks we determined the energy calibration, whereas from the position of the hydrogen peak the scattering angle could be verified. Both calculations were performed with a relativistic description of the reaction kinematics. Next in the energy spectrum obtained from the target studied the elastic peak and the peaks due to the scattering from $^{12}\text{C}$ and $^{16}\text{O}$ impurities were analysed.

In a second stage the energy calibration obtained from the analysis of the polyethylene spectrum was used to calculate the position of the peaks corresponding to the inelastic scattering process. These peaks were analysed in a similar fashion as described for the first stage.

No peak fitting was provided in our programmes for the analysis of the spectra. Generally this was found to be redundant because our study concerned only a few strong excited levels, at low excitation energies, which were well separated. However, for the experiment on $^{54}\text{Fe}$ at 15.3 MeV, in which also reasonable statistics for higher excited states had been obtained, and for the experiment on $^{89}\text{Y}$, which revealed also other than carbon and oxygen contaminants, a peak fitting programme had to be used for a reliable analysis. In these cases we made use of the automatic peak search and fitting programme Poespas. This programme, which is a version of Spapas (B175), was run on a B7700-computer.
REFERENCES

Ma76  P. Martin, private communication.
Sm64  R. Smythe, Rev. Sci. Instr. 35 (1964) 1197.
CHAPTER IV

SOME ASPECTS OF THE THEORY OF ELASTIC AND INELASTIC SCATTERING

4.1 Introduction

In the description of nuclear reactions one usually distinguishes between transitions involving only a few degrees of freedom of the target nucleus and those with excitations of many degrees of freedom. In the first class of transitions the initial and final states differ only slightly from each other and it is therefore expected that the excitation proceeds via a direct reaction, whereas the second class may proceed via a complicated compound nucleus mode. Of course this distinction is not so strict and in every reaction both modes will be present. However, if the energy of the incident particle is sufficiently high, there are so many states to which the compound nucleus can decay that the probability of decay to any one of them will be small. It is therefore likely that for the bombarding energies used in our experiments, which are well above the \((p,n)\) threshold, the compound nucleus contributions are sufficiently small to be neglected in our analysis.

In this chapter we describe the direct reaction models employed in the theoretical analyses of our experiments. We mainly intended to test their applicability with regard to the analysing powers. So no refinements are suggested and therefore we confine ourselves to a discussion of some of the characteristic features. A detailed description of the use of these models and a justification of the basic approximations are given for instance by Mc. Carthy (Ca68), Jackson (Ja70), Austern (Au70) and Hodgson (Ho71).

In section 4.2. we give a short outline of the general description of a scattering process. It is indicated how the angular distributions of the differential cross section and the analysing power can be calculated from transition amplitudes. Simplifying assumptions have to be made which lead to manageable expressions for these amplitudes.

The elastic scattering can be described by the optical model. In our analyses we used the phenomenological optical model, in which the interaction between the incident particle and the nucleus is represented
by a complex potential the parameters of which are determined empirically. This potential is described in section 4.3. In this section also a short outline of a somewhat more fundamental model is given, which is known as the reformulated optical model. Though we did not use this model, its basic concepts may be helpful in the interpretation of the parameters of the phenomenological potential.

The Distorted Wave Born Approximation (DWBA) has been used to analyse the inelastic scattering. The transition amplitude in this approximation is evaluated in section 4.4. This amplitude can be reduced to a form in which the basic information about the excitation of the nucleus is contained in a radial formfactor or transition density. Expressions for the formfactor, both for a macroscopic and a microscopic DWBA description, are discussed.

4.2 Description of the scattering process

Here we will give a short description of a scattering process in which an incident proton is scattered by a nucleus, which thereby may be excited. The system of the projectile and the target nucleus is described by a Hamiltonian \( H \):

\[
H = H_0 + T + V
\]  

(4.1.)

which comprises the Hamiltonian for the internal motion of the target nucleus \( H_0 \), the kinetic energy operator for the relative motion of the projectile and the target nucleus \( T = (-\hbar^2/2\mu)\nabla \) and the interaction \( V \) between the projectile and the nucleus. The cross section for a scattering process can be obtained by solving the time-independent Schrödinger equation

\[
(H - E) \psi^{(+)} = 0
\]

(4.2.)

Evidently such a solution \( \psi^{(+)} \) of this equation is sought that, in the region where the interaction has fallen to zero, it consists of an incoming plane wave in the initial channel \( i \) and outgoing scattered waves in all the possible final channels \( f \):
\[
\psi^{(+)}_{i} \rightarrow e^{i \frac{\mathbf{k}_i \cdot \mathbf{r}_i}{\hbar}} \phi_i \eta_i = \frac{\mu}{2\pi \hbar^2} \sum_{f} T_{f, i} \langle \mathbf{k}_f, m_f, M_f; \mathbf{k}_i, m_i, M_i \rangle \phi_f \eta_f \frac{e^{i \mathbf{k}_f \cdot \mathbf{r}}}{r}
\]

(4.3.)

This equation defines the transition amplitude \( T_{f, i} \). The internal wavefunctions of the target nucleus and the spin wavefunctions of the proton are denoted respectively by \( \phi \) and \( \eta \); the spin projections of the target spin and proton spin are referred to explicitly by respectively \( M \) and \( m \).

We mention that the distortion of the incoming and outgoing waves due to the Coulomb interaction is not indicated.

The differential cross section for a particular transition \( i, m_i, M_i \rightarrow f, m_f, M_f \) can be expressed as:

\[
\frac{d\sigma}{d\Omega} = \left( \frac{\mu}{2\pi \hbar^2} \right)^2 \frac{k_f}{k_i} |T_{f, i}(\mathbf{k}_f, m_f, M_f; \mathbf{k}_i, m_i, M_i)|^2
\]

(4.4.)

If the initial states are unpolarized and the polarizations of the final states are not measured the differential cross section is obtained by summing over the final states \( m_f \) and \( M_f \) and averaging over the initial states \( m_i \) and \( M_i \). This results in:

\[
\left\{ \frac{d\sigma}{d\Omega} \right\}_{\text{unpol.}} = \frac{1}{(2s+1)(2I_i+1)} \left( \frac{\mu}{2\pi \hbar^2} \right)^2 \frac{k_f}{k_i} \sum_{m_f, M_f} |T_{f, i}(\mathbf{k}_f, m_f, M_f; \mathbf{k}_i, m_i, M_i)|^2
\]

(4.5.)

Here the spin of the incident proton and the target nucleus are \( s \) and \( I \) respectively.

Obviously more detailed information from a scattering experiment can be obtained by measuring the separate cross sections in case one or even more of the initial or final spin states have been selected. This for instance is achieved in experiments with a polarized proton beam in which the cross sections for the two possible spin orientations of the incident proton are measured separately. Generally the beam is only partly polarized, so in calculating the differential cross section one ought to take a weighted average over the initial spin orientations. As shown by Wolfenstein (Wo59) the expression (1.1.), given in chapter 1,
is obtained with the cross section for an unpolarized beam given by (4.5.) and the analyzing power by:

\[
\Lambda = \frac{\sum_{m_f m_i M_i} \left| T_{f,i} (k_f, m_f, M_f) \right|^2 - \left| T_{f,i} (k_f, m_f, M_f) \right|^2}{\sum_{m_f m_i M_i} |T_{f,i}|^2}
\]

(4.6.)

The quantization axis is taken here along the normal to the scattering plane.

The analyzing power and differential cross section can be calculated straightforward once a suitable expression for the transition amplitude is obtained. To that end the total wavefunction \( \psi^{(+)} \) is expanded in terms of the complete set of eigenstates \( \phi_\alpha (\vec{x}) \) of the internal Hamiltonian \( H_0 \).

If we assume here for brevity the projectile to be spinless and moreover neglect the exchange of the proton with one in the target nucleus this expansion becomes:

\[
\psi^{(+)} (\vec{r}, \vec{\xi}) = \sum_\alpha \psi^{(+)} (\vec{r}) \phi_\alpha (\vec{\xi})
\]

(4.7.)

After substitution of (4.7.) into (4.2.), multiplication with \( \phi_\alpha \) and integration over the internal coordinates \( \vec{\xi} \) one obtains a system of coupled equations for the functions \( \psi_\alpha \) which for each channel \( \alpha \) has the form:

\[
(T - E_\alpha) \psi_\alpha^{(+)} (\vec{r}) = \sum_\alpha \langle \phi_\alpha | V | \phi_\alpha \rangle \psi_\alpha^{(+)} (\vec{r})
\]

(4.8.)

with \( E_\alpha = \hbar^2 k_\alpha^2 / 2m \). Often in the discussion of scattering processes it is convenient to split the interaction \( V \) into two parts \( U \) and \( W \). Usually \( U \) is taken to represent an average interaction potential which is independent of the internal coordinates of the target nucleus. Then the coupled equations can be written as:

\[
(T - E_\alpha) + \langle \phi_\alpha | U (\vec{r}, \vec{\xi}) | \phi_\alpha \rangle \psi_\alpha^{(+)} (\vec{r}) =
\]

\[
- \sum_{\alpha' \neq \alpha} \langle \phi_\alpha | U (\vec{r}, \vec{\xi}) | \phi_{\alpha'} \rangle \psi_{\alpha'}^{(+)} (\vec{r})
\]

(4.9.)
Of course such an infinite system of coupled equations can not be solved exactly. In the DWBA it is assumed that the elastic scattering is far out the most dominant process. Thus only those terms at the right hand side are taken into account which describe the coupling to the ground state in first order in \( W \). Then for any of the inelastic channels of the following set of equations results:

\[
(T - E_f + \langle \psi_f | U | \phi_i \rangle) \chi_{i}^{(+)}(\vec{r}) = 0
\]

\[
(T - E_f + \langle \psi_f | U | \phi_i \rangle) \psi_{f}^{(+)}(\vec{r}) = - \langle \phi_i | W | \phi_i \rangle \chi_{i}^{(+)}(\vec{r})
\]  

(4.10.)

Here the ground state is denoted by \( \phi_i \) and an excited state by \( \phi_f \). The wavefunction \( \chi_i \) describes the elastic scattering by the potential \( \langle \phi_i | U | \phi_i \rangle \). Regarding the physical condition given in equation (4.3.), \( \chi_i \) should comprise both an incoming and an outgoing scattered wave, whereas \( \psi_f^{(+)} \) stands for an outgoing wave only. The amplitude of this outgoing wave contains the transition amplitude for the inelastic scattering. This amplitude is given by (To61):

\[
T_{f,i}^{\text{DWBA}} = \int \chi_{i}^{(-)}(\vec{r}_i, \vec{r}) \langle \phi_i | W | \phi_i \rangle \chi_{i}^{(+)}(\vec{r}, \vec{r}_i) \, d\vec{r}
\]  

(4.11.)

The superscripts on the elastic scattering wavefunctions, solutions of the homogeneous equations, denote whether the scattered wave is outgoing (+) or incoming (−).

As a consequence of the approximation made in the derivation of the transition amplitude only those transitions are considered which lead directly from the ground state of the nucleus to an excited state. This means that in cases where the direct transition from the ground to the excited state is hindered or forbidden i.e. \( \langle \phi_f | W | \phi_i \rangle = 0 \), or in cases where one of the excited states is populated too strongly, the DWBA will fail and one has to include more transitions by means of a coupled channel calculation (Ta65). In practice the DWBA has been found to give a reliable description if the differential cross section for the inelastic scattering is much smaller than the elastic scattering cross section. Regarding the calculations of Percy (Pe63) this condition is fulfilled for all the reactions considered in this thesis.
4.3 The optical potential

4.3.1 The phenomenological optical potential

The elastic scattering can be described with the optical model of the nucleus. This model has been successfully applied in describing the huge amount of data obtained in the last fifteen years for the elastic scattering of nucleons and composite particles. A comprehensive review of the analyses with this model has been given by Hodgson (Wa71).

In the optical model the interaction between the incident particle and the target nucleus is represented by a central one-body potential, which is usually assumed to be local. The expression for this potential is well-established now. For the sake of completeness the optical potential used in our analyses is given here:

\[ U(r) = V_c(r) - V_o \frac{f(r,r_o,a_o)}{a} + i(W_v f(r,r_i,a_i) + 4a_i W_d \frac{d}{dr} f(r,r_i,a_i)) \]

\[ + \left( \frac{\hbar}{2m_c} \right)^2 V_{s.o.} \frac{1}{dr} f(r,r_{s.o.},a_{s.o.}) \frac{\hat{J} \cdot \hat{l}}{a^2} \]  \hspace{1cm} (4.12)

Here \( r \) is the distance between the centre of the nucleus and the incident proton and \( V_c, W_v, W_d, \) and \( V_{s.o.} \) are the strength parameters. For the radial formfactors the generally adopted Saxon-Woods form is taken, which is defined by:

\[ f(r,r_x,a_x) = \left( 1 + \exp \left( \frac{(r-r_{s.o.})}{a_{s.o.}} \right) \right)^{-1} \]  \hspace{1cm} (4.13)

with \( r_x \) being the usual radius parameter. The diffuseness \( a_x \) is a measure of the surface thickness and \( \Lambda \) is the mass number of the target nucleus.

The main features of the elastic scattering are described by the real potential \( V_o(r) \) together with the Coulomb potential \( V_c(r) \), which is calculated as due to a uniform charge distribution of the nucleus with radius \( r_c \Lambda^{1/3} \). The spin orbit potential has been introduced primarily to describe the polarization phenomena (Ad54). By analogy with the shell model a Thomas form is chosen. The factor \( \frac{\hbar}{m_c} \) is the reduced Compton wave length of the pion; \( \hat{J} \) denotes the angular momentum and \( \hat{l} \) the Pauli spin operator of the incident proton.

Since the optical model considers only the wavefunctions in the elastic channel explicitly, an imaginary term is needed in the potential to describe the removal of particles by non-elastic scattering processes.
This imaginary term consists of a volume part with strength $W_v$ and a surface-peaked part with strength $4\alpha_i W_i$. For energies below 20 MeV the surface absorption has been found the most important of these two terms. However, with increasing energy the volume absorption must also be included to reproduce the measurements.

The expression for the optical potential in equation 4.12. contains 14 parameters in all. These are chosen to describe as well as possible the experimental angular distributions of the differential cross section and the analysing power. To that end the Schrödinger equation is solved numerically (Me66). To provide an objective measure of the goodness of the fit a chi-square value is calculated which is defined by:

$$
\chi^2 = \sum_i \left( \frac{\sigma_{\text{exp}}(\theta_i) - \sigma_{\text{O.M.}}(\theta_i)}{\Delta \sigma_{\text{exp}}(\theta_i)} \right)^2 + \sum_i \left( \frac{A_{\text{exp}}(\theta_i) - A_{\text{O.M.}}(\theta_i)}{\Delta A_{\text{exp}}(\theta_i)} \right)^2
$$

(4.14.)

where $\sigma_{\text{exp}}$ and $A_{\text{exp}}$ are the experimental differential cross sections and the analysing powers with their corresponding errors $\Delta \sigma$ and $\Delta A$ respectively. The theoretical values calculated with the optical potential are $\sigma_{\text{O.M.}}$ and $A_{\text{O.M.}}$. The summation runs over all the experimental points taken at different scattering angles $\theta$. We used the programme OPTIMO (Vo72) to find a set of parameters which fits the experimental distributions. In this programme all or some of the parameters are varied in an iterative way to search for a minimum of the chi-square value.

From many analyses it has been observed that several different sets of parameters resulted in equally good fits to the experimental data measured for a specific nuclide at a fixed value of the bombarding energy. Well known, for instance, is the $V_0 r^n$ ambiguity with $n = 2$ (Ca73). Other ambiguities have been reported for the imaginary potential (Ag75). Although such ambiguities may be caused by incompletely measured angular distributions or by insufficient statistical accuracy it is more likely that they arise from that fact that the measured angular distributions only depend on the asymptotic wavefunction. Due to such ambiguities a systematic trend of the parameters to be expected from separate searches for several nuclides or at several bombarding energies might be obscured. This problem can be overcome by a so-called global analysis in which measurements on many nuclides and energies are fitted simultaneously. Such analysis has been made for instance by Bacchetti and Greenles (Be69), who performed a search on an extensive number
of experimental data for $A>40$ and $E <50$ MeV. Parameters from their search are listed in Table 4.1. To avoid excessive numerical work and to reduce the risk of obtaining unphysical solutions these BG parameters were taken to initiate the search process.

Table 4.1. Global parameter set of the proton optical potential as determined by Beauchet and Greenlees (strength parameters in MeV and geometry parameters in fm).

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>$V_0$</td>
<td>$54 - 0.32 E_p + 24 \frac{N-Z}{A} + 0.4 \frac{Z}{A^{1/3}}$</td>
</tr>
<tr>
<td>$a_0$</td>
<td>0.75; $r_0 = 1.17$</td>
</tr>
<tr>
<td>$W_d$</td>
<td>$11.8 - 0.25 E_p + 12 \frac{N-Z}{A} + 0.4 A^{1/3}$</td>
</tr>
<tr>
<td>$W_v$</td>
<td>$0.22 E_p - 2.7$</td>
</tr>
<tr>
<td>$a_i$</td>
<td>0.51 + 0.7 $\frac{N-Z}{A}$; $r_i = 1.32$</td>
</tr>
<tr>
<td>$V_{s.o.}$</td>
<td>6.2</td>
</tr>
<tr>
<td>$a_{s,o.}$</td>
<td>0.75; $r_{s,o.} = 1.01$</td>
</tr>
</tbody>
</table>

4.3.2 The reformulated optical potential

A successful attempt to calculate the optical potential using a less phenomenological approach has been made, about ten years ago, by Greenlees, Pyle and Tang (Gr68). In their reformulated optical model the real part of the optical potential is taken as the sum of the interactions of the incoming particle with the individual nucleons in the target nucleus. If the density distribution of nuclear matter is $\rho_m(\mathbf{r'})$ and the nucleon-nucleon interaction $v(|\mathbf{r}-\mathbf{r'}|)$, this summation results in:

$$v(\mathbf{r}) = \int \rho_m(\mathbf{r'}) v(|\mathbf{r}-\mathbf{r'}|) \, d^3 \mathbf{r'}$$  \hspace{1cm} (4.15)

The real central potential and the spin orbit potential were derived by considering the appropriate components of a nucleon-nucleon interaction. Greenlees and co-workers took as variable parameters the radius and the surface diffuseness of the nucleonic distribution and also the
strengths of the central part and the spin orbit part to take into account
deficiencies in the model such as the omission of the exchange process.
Moreover a phenomenological imaginary potential was added. Thus, with
two parameters less, they achieved fits comparable to the fits resulting
from phenomenological analyses.

From this study of Greenlees et al. two conclusions were drawn,
which are also important for the phenomenological analyses. Firstly
the root mean square radius of the real central potential is a well-
deﬁned geometrical quantity and secondly the volume integral of this
potential per target nucleon contains the information on the strength
of the interaction. For the Saxon–Woods form these quantities are given
approximately by:

$$<r^2>^\frac{1}{2} = \left\{ \frac{3}{5} \frac{R_e^5}{R_0} \left( 1 + \frac{7}{3} \frac{\pi a}{R_0^2} \right)^2 \right\}^\frac{1}{2} \quad (4.16,.)$$

and

$$J_0/A = \frac{4\pi V}{3 R_0^3} \left( 1 + \frac{\pi a}{R_0^2} \right)^2 \quad (4.17,.)$$

with $R_o = r_o A^{1/3}$.

In recent years more elaborate analyses have been performed with
considerable success. For instance, it has been found that by the
inclusion of the exchange process, brought about by the antisymmetrization
of the complete wavefunction for the incident particle and those of the
target nucleons, the energy dependence of the real central potential
could be explained for a large part. These developments have been re-
viewed by Sinha (S175).

Since we used the optical potential mainly as a tool to derive suitable
elastic scattering wavefunctions for our DWBA analyses we did not carry
out a microscopic analysis.

4.4 The Distorted Wave Born Approximation

4.4.1 General expressions

In the expression for the DWBA transition amplitude given in section
4.2. the dependence on the spin state of the projectile has been suppressed.
Because of the presence of the spin–orbit term in the optical potential
the elastic scattering wavefunctions $\chi_k^\alpha (r)$ become matrices in the projectile
spin space. Referring to expression (4.3) \( \chi^{(+)} \), for instance, can be written as:

\[
\chi^{(+)}(k_i, r) = \hat{\gamma}_i^{m_i} |s, m_i^{(+)} \rangle \chi^{-(+)}_m^{(+) m_i}(k_i, r)
\]  

(4.18)

where \( |s, m_i^{(+)} \rangle \) is a spin eigenfunction of the proton. Substitution in (4.11) results in a transition amplitude:

\[
T_{f,i} = \sum_{m_f^{(-)}} \chi^{-(-)}_i^{m_i m_f}(k_i, r) < \alpha f I_f M_f, m_f^{(-)} | W | \alpha_i I_i M_i, m_i^{(+)} \rangle \chi^{(+)}_m^{(+) m_i}(k_i, r) \, dr
\]  

(4.19)

Here the target state is denoted by the total angular momentum \( I \), its projection \( M \) and additional quantum numbers \( \alpha \). The matrix element between the elastic scattering wavefunctions, for which the bracket notation implies integration over all the coordinates except \( r \), describes the excitation of the target\( \alpha_i I_i M_i \rightarrow \alpha_f I_f M_f \) and a possible change in the projectile spin. It is clear that a change in the projectile spin may also occur through the influence of the optical potential for the initial and final channel.

To evaluate the transition amplitude it is convenient to make a multipole expansion of the nuclear matrix element \( \langle f | W | i \rangle \) into terms which correspond to the transfer to the nucleus of a definite angular momentum \( j \) composed of an orbital part \( L \) and a spin part \( S \). These transferred angular moments are defined by:

\[
\vec{j} = \vec{l} - \vec{s}; \quad \vec{s} = \vec{s}_L - \vec{s}_f \quad S = 0, 1; \quad \vec{l} = \vec{j} - \vec{s}
\]  

(4.20)

In case of a central interaction \( W \) the transition amplitude then can be written in the form (8a64):

\[
T_{f,i} = \sum_{LSJ} \sum_{m_f^{(-)}} \chi^{-(-)}_i^{m_i m_f}(k_i, r) f_{LSJ}(r) Y^{LSJ}(\hat{r}) \chi^{(+)}_m^{(+) m_i}(k_f, r) \, dr
\]  

(4.21)

where we have omitted several vector coupling coefficients corresponding to the coupling in (4.20.) and \( M' = M_f - M_i + m_f - m_i \). The parity change in the transition is given by \( \eta_i f M_f = (-1)^L \). Often only one combination of LSJ values is important in a given transition. Thus the information
on the target wavefunctions and the interaction $W$ is contained in only one radial formfactor $F_{LSJ}(r)$. The macroscopic and microscopic description of these formfactors are given in the following subsections.

We mention that if non central terms appear in the interaction $W$, such as a tensor or spin-orbit interaction, in general the simple expression (4.21.) can no longer be obtained. The same remark holds for the inclusion of exchange effects, which we have neglected throughout this section. Consequently a simple comparison of different nuclear models by means of the resulting formfactors is then no longer possible.

4.4.2 Macroscopic description

In the macroscopic description of the nucleus the collective model is used (Bo53). In this model the nuclear states are considered as vibrations of the nucleus about a mean spherical shape or as rotations of a permanently deformed nucleus. In both cases the surface of the nucleus can be described by:

$$R(\theta,\phi) = R_0 \left(1 + \alpha(\theta,\phi)\right)$$ \hspace{1cm} (4.22.)

where $R_0$ is the radius of the equivalent spherical surface and $\alpha(\theta,\phi)$ describes the deformation. It is assumed now that the interaction potential follows the deformation of the nuclear surface. This assumption leads to an extension of the optical potential to include non-spherical parts. Then the potential is given to first order in the deformation $\alpha$ by:

$$V(r_0,R(\theta,\phi)) = V(r_0,R_0) + \alpha(\theta,\phi) R_0 \frac{dV(r_0,R_0)}{dR_0}$$ \hspace{1cm} (4.23.)

The first term is now associated with the spherical optical potential $U_{O.M.}$, which does not change the nuclear state. The second term gives rise to the transitions to excited states. Generally the deformation is expanded into multipoles.

$$\alpha(\theta,\phi) = \sum_{LM} \alpha_{LM} Y_{LM}(\theta,\phi)$$ \hspace{1cm} (4.24.)
For the excitation of a state with a specific $l^\pi$ only one term of this expansion has to be considered, which gives as a form factor (Ba62):

$$F_{\text{LOL}}(r) = (2L + 1)^{-\frac{1}{2}} \beta_L \frac{dU_{O.M.}}{dR_{o}}$$  \hspace{1cm} (4.25.)

where $\beta_L$ is the deformation parameter. Thus the macroscopic description leads to an universal formfactor which is given by the derivative of the optical potential. As the differential cross section is proportional to the square of the transition amplitude the value of $\beta_L^2$ is determined directly from the absolute magnitude of the differential cross section. Actually it is the only parameter as the parameters of the optical potential are determined from the elastic scattering.

In the first applications of the macroscopic description only the real part of the optical potential was deformed. Later extensive experimental data showed definitely that also the imaginary as well as the spin-orbit part should be deformed (Sa66, Sa70). From the several possible prescriptions for calculating the deformed spin-orbit potential (C168), the full Thomas form is found to improve the fit to the experimental angular distributions of the analysing power substantially. This full Thomas term was derived by Sherif and Blair (Sh68, Sh69, Sh70) by rewriting the spin-orbit term in the optical potential in the form $\hat{V}_{s.o.}(r_{s.o.}, R_{s.o.}) \times \frac{1}{i} \hat{\nabla}$ for which the non-spherical part, then again to first order in $\alpha(\theta, \phi)$ is given by:

$$W_{s.o.}(r) = \left[ \frac{\hbar}{m c} \right]^2 \langle \hat{V}_{s.o.} \rangle \left[ \hat{\nabla} \alpha(\theta, \phi) \frac{df_{s.o.}}{d\Omega_{s.o.}} \times \frac{1}{i} \hat{\nabla} \right] \hat{\sigma}$$  \hspace{1cm} (4.26.)

The inclusion of this term into DWBA calculations has been discussed in detail by Sherif (Sh69), Raynal (Ra70) and Verhaar (Ve72, Ve74).

Theoretical predictions for the several types of deformed potentials are compared in figure 4.1. It is seen that the differential cross section is raised if an imaginary term is included but the shape does not change significantly at the forward angles. However, larger effects are seen at the backward angles, especially for the analysing powers. The introduction of the full Thomas term has only a slight effect on the cross section, while the analysing power is shifted to more positive
values at forward scattering angles.

The macroscopic description outlined above, refers to low-lying excited states of even-even nuclei. If the target nucleus can be described by an even-even core plus an extra nucleon the weak coupling model may be applied (Sh61, B159). In this model excited states occur which are described by the extra nucleon coupled to an excited state of the core. Thus a multiplet is formed with spin I ranging from $|j_p - I_c|$ till $j_p + I_c$, where $p$ refers to the extra nucleon and $c$ to the core. The angular distributions for the cross sections and analysing powers are then the same as those for the excitation of the core. For the absolute magnitude of the cross sections a simple relation holds:

$$\beta_L^2(I_i = j_p \rightarrow I_f) = \frac{(2I_f+1)}{(2I_i+1)(2L+1)} \beta_L^2(0 \rightarrow I_c)$$

(4.27)

For convenience often $\beta'_L$ is used for the deformation of the odd nucleus. From (4.27) it is obvious that the values for $\beta_L^2$ should show a $2I_f+1$ dependence and that their sum over the weak coupling multiplet should equal the $\beta_L^2$ for the corresponding state in the even-even nucleus.
4.4.3 Microscopic description

In comparison with the macroscopic description, the microscopic description is a more detailed and fundamental one. The ground and excited states of the nucleus are described in terms of the motion of the individual nucleons. The interaction is assumed to be given by the sum of effective interactions between the projectile o and the nucleons j of the target nucleus:

$$W = \sum_j U_{oj}$$  \hspace{1cm} (4.28.)

A general expression for the central part of $U_{oj}$ is:

$$U_{oj} = - \left( V_{00}(r_{oj}) + V_{10}(r_{oj})(\vec{\sigma}_o \cdot \vec{\sigma}_j) \right) +$$

$$V_{01}(r_{oj})(\vec{\sigma}_o \cdot \vec{\tau}_j) + V_{11}(r_{oj})(\vec{\sigma}_o \cdot \vec{\sigma}_j)(\vec{\tau}_o \cdot \vec{\tau}_j)$$  \hspace{1cm} (4.29.)

where $\vec{\sigma}$ and $\vec{\tau}$ are the spin and iso-spin operators. If the projectile enters and leaves as a proton only the z-components of $\vec{\tau}$ contribute and the interactions effectively reduce to:

$$U_{oj} = - \left( V_0 + V_1 (\vec{\sigma}_o \cdot \vec{\tau}_j) \right) g(r_{oj})$$  \hspace{1cm} (4.30.)

Here for convenience the radial factors of the different parts of the interaction are taken equal. To arrive now at the expression for the DWBA transition amplitude, given in (4.21.), the function $g(r_{oj})$ is expanded into multipoles:

$$g(r_{oj}) = \sum_{LM} g_L(r_{oj}) Y_{LM} (\vec{r}_j) Y^*_{LM} (\vec{r}_o)$$  \hspace{1cm} (4.31.)

By inserting this expansion into the matrix element $<I|W|i>$ the following radial form factors are obtained (Sa64): 

$$F_{LSJ}(r_j) = -4\pi^{1/2} V_S (2I_{j} + 1)^{-1}$$

$$\times \sum_j e^{|I_f|} g_L(r_{oj}) T_{LSJ}(\vec{r}_j, \phi_j) \quad \text{for} \quad (S=0,1)$$

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where the tensor $I_{LSJ}$ is constructed from the spherical harmonics and the components of the spin operator. Since the operator appearing in (4.32.) is a one-body operator in the nuclear coordinates, it will couple initial and final states which differ only in the state of one nucleon. So in the most simple case one target nucleon is excited and the final and initial target states may be represented by single particle wavefunctions. As shown by Johnson et al. (Jo66) the reduced matrix element can then be written as follows:

$$<I_f | g_{LSJ} | I_i> = (M_L \delta_{S0} + N_{SL} \delta_{S1}) I_L (r_0)$$  (4.33.)

where the multipole coefficients $M_L$ and $N_{SL}$ depend, apart from $L$ and $J$, on the nuclear wavefunctions only. The formfactor $I_L (r_0)$ is given by:

$$I_L (r_0) = \int u_{i i j j} (r_0) g_L (r_0, r_j) u_{i i j j} r_j^2 dr_j$$  (4.34.)

where the $u_{i i j j}$ are the radial parts of the single particle wave functions. In the limit of a zero-range interaction the shape of the formfactor is given by the product of the wavefunctions $u_f$ and $u_i$. This illustrates that formfactors with a shape different from the surface peaked macroscopic model formfactors may occur. However, in the description of the more collective states a large number of single particle wavefunctions is involved and hence the formfactor becomes a combination of single particle formfactors. Such combinations add constructively, resulting in a total formfactor which resembles in strength and in shape the real part of the macroscopic formfactor (Gl66).

Sometimes the ground and low excited states of open shell nuclei near magic or semi-magic nuclei can reasonably be described by the motion of a few valence nucleons. Which means that rather simple shell-model configurations can be used in the description of these states. However, the effective charges needed to explain the observed electromagnetic transition strength indicate that such a description is not complete. So contributions of the other nucleons, which are referred to as the "core", have to be added in the description of the nuclear states. If the target system is divided into two parts, one comprising
the valence nucleons and the other the core, these core-polarization contributions arise from the interaction between the two parts. A simple macroscopic approximation has been proposed by Love and Satchler (Lo67) to include the core polarization terms into the microscopic description of inelastic scattering. By treating the coupling between the core and the valence nucleons as well as that between the core and the projectile according to the macroscopic description, they derived that the core contributions are accounted for by a renormalized interaction. This results in an additional formfactor which is proportional to the macroscopic formfactor:

$$\Gamma_L(r_0) = 4\pi V_0^L \cdot \gamma_L \cdot \frac{dU_{S.M.}}{dR} \cdot \gamma_i \cdot \gamma_j \cdot \frac{dU_{O.M.}}{dR} \quad (4.35)$$

Here $U_{S.M.}$ is the shell-model potential, by which the valence nucleons are bound and $U_{O.M.}$ the optical model potential. The core coupling parameter $\gamma_L$ can be obtained from the effective charge needed to describe the observed electromagnetic transition strength, thus in principle no additional parameter is introduced.

In early studies of the microscopic description of the inelastic scattering (Gr66, Sa67) the effective nucleon-nucleon interaction was adjusted to reproduce the absolute magnitude of the measured cross sections. The resulting potential strengths then were found to be significantly larger than those derived from the free nucleon-nucleon scattering. This could be partly understood from the neglect of the core coupling contributions. Moreover, it was noticed that the effective interaction depended on the bombarding energy and on the angular momentum transferred in the reaction. In several studies (Am67, At70, Lo70) this dependency could be removed by the inclusion of knock-on exchange effects. These exchange effects follow from the antisymmetrization between the wavefunctions of the projectile and the target nucleons (Au70). By this antisymmetrization an amplitude additional to the amplitude (4.11.) is introduced which for a single particle excitation is of the form:

$$T_{f,i} = -\int \frac{d\Gamma_{f,i}}{d\Gamma_{o,j}} \cdot \frac{d\Gamma_{o,j}}{d\Gamma_{f,i}} \cdot \hat{\gamma}_i \cdot \hat{\gamma}_j \cdot \frac{d\Gamma_{o,j}}{d\Gamma_{f,i}} \cdot \gamma_j$$

(4.36.)
This knock-on contribution may be visualized by a process in which the incoming particle is captured into the nucleus whereas a target nucleon is ejected. The exchange amplitude cannot be factorized like the direct amplitude, in an integral over the internal coordinates and an integral over the coordinates of the distorted waves. This greatly complicates the calculation of the exchange contributions (Ge71). From calculations in which these knock-on contributions were treated exactly, it has been found that they increase strongly with increasing bombarding energy up till about 25 MeV whereafter their influence decreases slowly. A strong enhancement with angular momentum transfer appears at all energies (Cf. (Lo70)). Considering the region around 20 MeV, which is the subject of our study, the ratio of the total cross sections with and without exchange is found to be approximately a factor of 1.2 for L=0 and a factor 5 for L=5.

The past decade encouraging results have been obtained with the microscopic description applied to inelastic scattering of protons (Cf. (Sc69, Ha71, Ha73, Wa73, Sa73, Sc77)). In these studies extensive particle-hole wavefunctions have been used to describe low-lying collective states, whereas less collective states were described with single-particle excitations and the macroscopic inclusion of core contributions. The long-range part of the even-state Hamada-Johnston potential gave an adequate description of the central part of the effective interaction. This Hamada Johnston potential (Ha62), which is derived from free nucleon-nucleon scattering also possesses rather strong non-central parts such as a spin-orbit force and a tensor force, which are given by:

\[ V_{LS}(r_{12}) = \frac{V_L S(r_{12})}{2m} (\hat{r}_1 \times \hat{r}_2) \times (\hat{p}_1 \times \hat{p}_2) \times (\sigma_1 \times \sigma_2) \]

\[ r_{12}^2 V_T(r_{12}) S_{12} = r_{12}^2 V_T(r_{12}) \left( \frac{3(\sigma_1 \times \hat{r}_{12})(\sigma_2 \times \hat{r}_{12})}{r_{12}^2} - (\sigma_1 \times \sigma_2) \right) \]  

The inclusion of these terms in the microscopic description of scattering processes has been discussed in detail by Love (Lo72) and Love and Parish (Lo70a). Although analyses of inelastic scattering incorporating these non-central parts are scarce it has been found in some particular cases that it might be important to include them in the effective interaction (Au72, Ra70, Ra71, Ra75, Wa73).
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CHAPTER V

THE SCATTERING OF POLARIZED PROTONS BY IRON AND NICKEL NUCLEI.

5.1 Introduction

In this chapter the scattering of polarized protons by the nuclides $^{54}\text{Fe}$, $^{56}\text{Fe}$, $^{58}\text{Ni}$, $^{60}\text{Ni}$ and $^{62}\text{Ni}$ will be discussed. For the iron isotopes experiments have been performed at bombarding energies of 15.3 MeV, 17.2 MeV, 20.4 MeV and 24.6 MeV, whereas for the nickel isotopes data were taken at energies of 20.4 MeV and 24.6 MeV. In all these experiments we measured the differential cross sections and analysing powers at scattering angles from $20^\circ$ to $165^\circ$ in $5^\circ$ steps, for the elastic scattering and for the inelastic scattering leading to several excited states.

The elastic scattering data are presented in section 5.2. From these data we deduced parameters for the optical potentials which are used in the analysis of the inelastic scattering results. This analysis is described in section 5.3. As mentioned in chapter 1 in the present analysis special attention is paid to a possible energy dependence of the analysing powers for the first $2^+$ state in $^{54}\text{Fe}$. That is why we performed DWBA calculations with a macroscopic formfactor only. These calculations may also indicate to what extent energy dependent effects can be described by more elaborate descriptions, presupposing a direct reaction process.

The main results from our analysis are discussed at the end of section 5.2 and 5.3. We notice that the analysis of the experiments performed at 17.2, 20.4 and 24.6 MeV has been published recently (Hæ77).

5.2 Elastic scattering

5.2.1 Optical model analysis

Optical model fits to the elastic scattering data for each nucleus and energy were made in a search on the parameters of the optical potential. The form of this potential was given in
section 4.3.1. For all the calculations we used the search programme OPTIMO (Vo72). Relative errors of 3% were assigned to the cross sections of the iron isotopes, whereas an error of 2% was taken for the nickel isotopes. These errors were estimated from the uncertainty in the relative solid angle calibrations of the eight detector system and from the uncertainty in the normalization of the measurements at different angular settings. The statistical errors were always found to be less. The errors assigned to the analysing powers were the statistical uncertainties or at least an absolute error of 0.01. With the exception of the experiment at 20.4 MeV the differential cross sections and analysing powers were measured simultaneously.

The parameters of the global potential of Becchetti and Greenlees (Be69) (cf. chapter 4, section 4.3.1.) were used as starting values in our searches. We found that these parameters fitted the analysing powers of the elastic and inelastic scattering rather well (Ha75). To ensure that the different ways along which a search can be performed, results in the same set of parameters, we applied in our analysis of the 20.4 and 24.6 MeV data the following procedure. First we fitted the analysing power varying all the strength parameters and of the geometry parameters only those of the spin orbit potential. Next, with the resulting parameters of the first search as startvalues, the differential cross section was fitted by varying all the parameters but those of the spin orbit geometry. Finally both the analysing power and the cross section was fitted simultaneously by varying all the parameters. An alternative procedure was applied, starting now with the cross section data and fitting subsequently the analysing power and both the cross section and the analysing power. We found that these search procedures converged more or less to the same final parameter set. Once this had been established the data at 17.2 MeV and 15.3 MeV were fitted by varying all the parameters directly. In all these searches the final parameterset was assumed to be obtained if the reduction in the total chi-square value was found to be less than 0.1. The best fit (BP) optical model parameters, resulting from searches on all the parameters, are listed in table 5.1.
Table 5.1 Best-fitted optical model parameters obtained in a search on all the parameters.\(^a\)

<table>
<thead>
<tr>
<th>Nucleus</th>
<th>Energy (MeV)</th>
<th>(v_0) (MeV)</th>
<th>(v_\infty) (MeV)</th>
<th>(a_0) (MeV)</th>
<th>(W_\infty) (MeV)</th>
<th>(W_0) (MeV)</th>
<th>(r_1) (fm)</th>
<th>(a_1) (MeV)</th>
<th>(V_{\infty}) (MeV)</th>
<th>(r_{\infty}) (fm)</th>
<th>(a_{\infty}) (fm)</th>
<th>(x_0)</th>
<th>(x_\infty)</th>
</tr>
</thead>
<tbody>
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<td>(^{15}Fe)</td>
<td>15.3</td>
<td>55.3</td>
<td>154.2</td>
<td>3.21</td>
<td>6.00</td>
<td>1.32</td>
<td>0.57</td>
<td>8.48</td>
<td>0.97</td>
<td>0.66</td>
<td>49</td>
<td>431</td>
<td></td>
</tr>
<tr>
<td>17.2</td>
<td>55.2</td>
<td>154.2</td>
<td>3.21</td>
<td>6.00</td>
<td>1.32</td>
<td>0.57</td>
<td>8.48</td>
<td>0.97</td>
<td>0.66</td>
<td>49</td>
<td>431</td>
<td></td>
<td></td>
</tr>
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<td>20.4</td>
<td>52.6</td>
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<td>154.2</td>
<td>3.21</td>
<td>6.00</td>
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\(^a\) Coulomb radius parameter \(r_0 = 1.23\) fm for all nuclei and energies.

\(^{**}\) Search process terminated because of too slow convergence.

Table 5.2 Near-circular geometry optical model parameters obtained in a search on the depth parameters only.\(^a\)

<table>
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<tr>
<th>Nucleus</th>
<th>Energy (MeV)</th>
<th>(v_0) (MeV)</th>
<th>(r_\infty) (fm)</th>
<th>(a_\infty) (MeV)</th>
<th>(W_\infty) (MeV)</th>
<th>(W_0) (MeV)</th>
<th>(r_1) (fm)</th>
<th>(a_1) (MeV)</th>
<th>(V_{\infty}) (MeV)</th>
<th>(r_{\infty}) (fm)</th>
<th>(a_{\infty}) (fm)</th>
<th>(x_0)</th>
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<td>559</td>
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<td>1.38</td>
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<td>307</td>
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</table>
In optical model analyses it is customary to perform searches in which the geometry parameters are kept fixed on suitable values. This procedure has the advantage that ambiguities between the strength and geometry parameters do not obscure the general trends, which are expected for the strength parameters. For this reason such a procedure was followed in our analysis too. From the BF parameters at 20.4 MeV and 24.6 MeV we deduced average geometry parameters for the real and spin-orbit potential while the geometry of the imaginary potential was allowed to vary from nucleus to nucleus. Keeping the geometry parameters fixed the strength parameters were varied. Thus we obtained the fixed geometry (FG) potentials listed in table 5.2.

The results of our experiments together with the optical model calculations are given in figure 5.1 to figure 5.6. As the results of the BF potentials and the FG potentials are hardly distinguishable by eye we give only the FG curves. For $^{54}$Fe at 15.3 MeV (fig. 5.1) theoretical curves calculated with the parameters from an experiment at 15.13 MeV of Lombardi et al. (Lo72) are also given. For $^{54}$Fe at 17.2 MeV, 20.4 MeV and 24.6 MeV (figure 5.2) and for $^{62}$Ni at 20.4 MeV and 24.6 MeV (figure 5.6) the FG results are compared to the results from the Becchetti and Greenlees (BG) parameters.

It can be seen that generally the optical model descriptions fit the experimental data quite well and that the quality of the fits to the analysing powers is comparable to that of the fits to the cross sections. Difficulties arise in describing the data at 15.3 MeV and 17.2 MeV; especially at forward scattering angles where the cross sections are reproduced not so well. It is likely that the fit at these forward angles is influenced by the behaviour at the more backward angles, where effects not taken into account in the optical model description are expected to become more important. We therefore tried to improve the description at the forward angles by imposing larger relative errors (about 10%), and thus less statistical weight, upon the experimental points for scattering angles larger than 95°. In this way only minor improvements were found at forward angles whereas the parameters did not change significantly.
Fig. 5.1 Differential cross section and analysing power for elastic scattering by $^{54}$Fe and $^{56}$Fe at 15.3 MeV. Solid curves, FC potentials; dashed curves, potential parameters from Lombardi et al. (Lo72). As customary the cross section is divided by the Rutherford cross section $\sigma_R$ (scattering from a point charge 2e).

3.2.2 Normalization of the differential cross section

As we measured the differential cross sections relatively only, the normalization was used as an additional parameter in our searches. For the analysis of the data at 20.4 MeV and 24.6 MeV the normalization was determined together with the other parameters in the BP search and it was kept fixed in the FC search. This procedure raised some problems at the lower bombarding energies because there the norm was not so well determined, as will be shown further on in this section. At 17.2 MeV the norm was deduced by fitting only the data up to 100° in a FC search.
Fig. 5.3 Differential cross section and analysing power for elastic scattering of protons \( E_p = 17.2, 20.4, 24.6 \text{ MeV} \) by \(^{54}\text{Fe}\). Solid curves, PG potentials; dashed curves, BG potentials.
Fig. 5.3 Differential cross section and analysing power for elastic scattering of protons ($E = 17.2, 20.1, 24.6$ MeV) by $^{56}$Fe. The curves through the points are optical model (FG) fits.
Fig. 5.4 Differential cross section and analysing power for elastic scattering of protons \( E_p = 20.4, 24.6 \text{ MeV} \) by \(^{58}\text{Ni}\).

The curves through the points are optical model (FG) fits.

At 15.3 MeV the normalization was varied as well as in the FG search; the normalization factors for the experimental cross sections were then taken from the FG search, which yielded values 5% and 10% smaller respectively for \(^{54}\text{Fe}\) and \(^{56}\text{Fe}\) as compared to the BF normalization.

To investigate to what extent the uncertainty in the absolute normalization procedure affects the optical model parameters deduced from our experiments we performed a series of searches on the \(^{56}\text{Fe}\) data at 17.2 MeV and 24.6 MeV with normalization factors \( N_C \), fixed at 0.9, 0.95, 1.05 and 1.10 times the originally determined value. Although the parameters were found to vary systematically with \( N_C \), in both cases only minor differences of a few percent or even less appeared between the extreme values for the parameters of the real and spin-orbit part and the radius of the imaginary potential. For the experiment at 24.6 MeV the strength of the volume absorption \( N_v \) was found to
Fig. 5.5 Differential cross section and analysing power for elastic scattering of protons \( E_p = 20.1, 24.6 \text{ MeV} \) by \( ^{60}\text{Ni} \).

The curves through the points are optical model (FO) fits

decrease whereas \( a_{1d} \) increased with approximately the same amount (8%) with increasing \( N_0 \). This effect was strongly enlarged for the experimental data at 17.2 MeV for which even fits without almost any volume absorption were obtained. A similar ambiguity between the volume and surface absorption term has been reported a few years ago by Agrawal and Sood (Ag75) who also demonstrated that the total volume integral of the absorptive part of the proton optical potential is a much better determined quantity. Also in our searches this volume integral varies only a few percent, yielding smaller values for larger \( N_0 \) whereas the volume integral of the real potential, as expected, shows an opposite tendency.

If an increase of 50% in the chi-square value of the cross section is regarded as a guide to determine the accuracy of our normalization procedure, we found that for the 24.6 MeV data the normalization factor is rather well defined within 10%. However, larger errors of about
Fig. 5.6 Differential cross section and analysing power for elastic scattering of protons ($E_p = 20.4, 24.6$ MeV) by $^{62}$Ni. Solid curves, FG potentials; dashed curves, BG potentials.

15% are estimated for the 17.2 MeV data. These larger errors are most probably brought about by the difficulties encountered in fitting the data at the forward scattering angles.

5.2.3 Discussion

The chi-square values obtained in our BF and FG searches are comparable to those reported from other studies in this energy and mass region (Sa67, Ko67, Lo72). From tables 5.1 and 5.2 it is observed that the chi-square values for $^{54}$Fe at 15.3 MeV and 17.2 MeV, $^{56}$Fe at 15.3 MeV and $^{58}$Ni at 20.4 MeV are significantly larger than the other values. These larger values reflect the difficulties encountered in the search process, which for $^{54}$Fe and $^{56}$Fe resulted in an inferior
fit at the forward scattering angles and for $^{58}$Ni in a somewhat worse
description at the backward angles. Much larger chi-square values for
$^{54}$Fe as compared to other nuclei in this mass region have been reported
previously by Kossanyi-Denay et al. for 18.6 MeV (Ro67) and by Lombardi
et al. for 15.13 MeV (Lo72). We notice that in the experiments of Lombardi
similar discrepancies occur, as found in our experiments, for the scattering
of protons from $^{50}$Ti, $^{51}$V, $^{52}$Cr and $^{54}$Fe around bombarding energies of
15 MeV. The dashed curve in figure 5.1 has been calculated from the
optical model parameters given by Lombardi, only the normalization was
adjusted to our experimental data, yielding a value merely a few percent
larger than that obtained in our PC search. Although a somewhat better
description is obtained at the forward scattering angles the fit over the
entire angular region remains poor, whereas the total chi-square value is
about a factor of two larger than that obtained in our PC search.

It may be obvious to ascribe the inferior fits at the lower bombarding
energies to compound nucleus contributions. However, one should realize
that these bombarding energies are amply above the $(p,n)$ threshold (9 MeV
for $^{54}$Fe and 5.4 MeV for $^{56}$Fe) and that it has been shown that beyond this
threshold the compound nucleus contributions rapidly decrease (Re69).
Moreover, Lombardi et al. reported that in the excitation functions they
measured no prominent resonances were found.

Having indicated the serious difficulties we found in describing
the experimental results at the lower bombarding energies, we now will
consider the general trends which nevertheless can be observed in the
optical potential parameters. Since it can be expected that particularly
the parameters of the spin-orbit potential are determined from
polarization experiments, we will discuss the results for these
parameters first. We learn from table 5.1 that generally the
parameters $V_{so}$ and $a_{so}$ scatter largely. Contrary to this the
parameter values for the radius $r_{so}$ show remarkably small fluctuations.
Moreover, each individual parameter set clearly demonstrates the need
for a 10-15% smaller radius for the spin-orbit potential as compared
to that for the real potential. In the reformulated optical model such
differences are readily explained by the much shorter range of the
two-body spin-orbit force as compared to the range of the central two-
body force (Gr68). The mean values of the spin-orbit strength $V_{so}$ and
the diffuseness $a_{so}$ in table 5.1, are somewhat smaller than the BC values
which are $V_{so} = 6.2$ MeV and $a_{so} = 0.75$ MeV. However, a striking correlation
between the strength and diffuseness clearly shows up. A larger, respectively smaller, value for \( a \) is always accompanied by a larger, respectively smaller, value for \( V_{so} \). To investigate such an ambiguity we performed several searches on the \( ^{56}\text{Fe} \) data at 24.6 MeV in which the diffuseness \( a \) was fixed on values ranging in steps of 0.05 fm from 0.40 fm to 0.65 fm.

A search on all the strengths parameters and the radius parameter \( r_{so} \) was done, whereas \( a_o, r_o, a_i \) and \( r_i \) were fixed on the FO-values. We found that all the varied parameters except \( V_{so} \) were not significantly affected by the imposed value of \( a \). The spin-orbit strength varied in such a manner that \( V_{so} \sim a_n \) (\( n \approx 0.5 \)) was found to be constant. This result is surprising, if one realizes that the maximum value of \( V_{so} (r) \) is proportional to \( V_{so} / r_{so} a_{so} \). Aside from this unexpected result we conclude that if an increase of 50\% in the chi-square value for the analysing power is regarded as a significant difference in the quality of the fit, the spin-orbit strength and the diffuseness are not so well determined giving values of 0.50 < \( a \) < 0.62 fm and 5.4 < \( V_{so} < 6.2 \) MeV.

Considering the values of \( V_o \) in table 5.1 and 5.2 we found in general a dependence on the energy close to that given by Bacchetti and Greenlees (Be69). We did not find a clear dependence on \( N, Z \) and \( A \). The parameters of the imaginary potential largely deviate from the BG parameters. At 15.3 MeV and 17.2 MeV the volume absorption decreases and the surface absorption increases with energy. The opposite behaviour, which is also expected from the BG potential, is observed at 20.4 MeV and 24.6 MeV.

Regarding the already mentioned ambiguity between \( W_v \) and \( W_d \) it seems more reasonable to compare the volume integrals of the imaginary potential. In order to facilitate this comparison we give in table 5.3 these volume integrals per nucleon \( J_v/\Lambda \) together with the volume integrals per nucleon \( J_o/\Lambda \) and the root mean square radius \( \langle r_o^2 \rangle^{1/2} \) of the real central potential. Contrary to the marked energy dependence of \( J_o/\Lambda \) it is seen that for \( J_v/\Lambda \) such a behaviour does not appear.

Though the values for \( J_v/\Lambda \) scatter they are all found to be consistent with the value of 115 \pm 15 MeV fm\(^3\) reported by Agrawal and Sood (Ag75) for proton energies between 10 MeV and 60 MeV and target nuclei with mass numbers between 40 and 208. We notice that the several FG values for \( J_o/\Lambda \) at a specific energy show a remarkable agreement. Thus no dependence on \((N-Z)/2\) is found, whereas one may get the impression that \( J_v/\Lambda \) slightly increases with increasing neutron excess.
Table 5.3 Volume integrals and root mean square radii derived from the best fits (BF) and fixed geometry (FG) potentials.

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<th>Nucleus</th>
<th>Energy (MeV)</th>
<th>( J^g/h ) (MeV fm(^3))</th>
<th>( J_z/A ) (MeV fm(^3))</th>
<th>( \langle r^2 \rangle^{1/2} (\text{fm}) )</th>
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<td>FG</td>
<td>BF</td>
<td>FG</td>
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5.3 Inelastic scattering

5.3.1 DWBA analysis

The results of our experiments are shown in figure 5.7 to 5.17 together with theoretical curves, which have been calculated using the DWBA programme written by L.D. Tolsma and adapted by B.J. Verhaar et al. (Ve72) to include the full Thomas spin-orbit coupling. The coupling between the ground and excited states is described with a complex macroscopic formfactor, which includes a full Thomas term. In our calculations we took for the optical potentials the fixed geometry potentials as given in section 5.2.1. For \( ^{62}\text{Ni} \) we give, for comparison, also results of calculations with the Becchetti Greenlees potential. Non-locality corrections were not included. On the whole we did not find our analysis to be sensitive for the several suitable potentials we tried.
By fitting the differential cross sections at the first maximum to the calculated curves we obtained the deformation parameters \( \beta_1 \) listed in Table 5.4. Due to our normalization procedure, described in section 5.2.1, errors must be assigned to these values of approximately 5%. In view of these errors the agreement between the parameters deduced for the several bombarding energies is quite good. An exception from the deformation parameters for the second \(^2\) state in \(^{54}\)Fe of which that derived at 15.3 MeV is found to be larger. In general our results agree well with previous results (Ve70, Au77, Ko76, Ki75, Ve74).

Table 5.4 Deformation parameters for the central potential \( \beta_1 \) from the present analysis.

<table>
<thead>
<tr>
<th>Nucleus</th>
<th>State</th>
<th>15.3 MeV</th>
<th>17.2 MeV</th>
<th>20.4 MeV</th>
<th>24.6 MeV</th>
</tr>
</thead>
<tbody>
<tr>
<td>(^{54})Fe</td>
<td>(^2)</td>
<td>1.41 MeV</td>
<td>0.15</td>
<td>0.15</td>
<td>0.16</td>
</tr>
<tr>
<td>(^{54})Fe</td>
<td>(^2)</td>
<td>2.95 MeV</td>
<td>0.17</td>
<td>0.14</td>
<td>0.13</td>
</tr>
<tr>
<td>(^{54})Fe</td>
<td>(^2)</td>
<td>3.16 MeV</td>
<td>0.09</td>
<td>0.11</td>
<td>0.11</td>
</tr>
<tr>
<td>(^{54})Fe</td>
<td>(^4)</td>
<td>2.95 MeV</td>
<td>0.17</td>
<td>0.15</td>
<td>0.15</td>
</tr>
<tr>
<td>(^{54})Fe</td>
<td>(^4)</td>
<td>3.84 MeV</td>
<td>0.17</td>
<td>0.15</td>
<td>0.15</td>
</tr>
<tr>
<td>(^{54})Fe</td>
<td>(^4)</td>
<td>4.28 MeV</td>
<td>0.15</td>
<td>0.15</td>
<td>0.15</td>
</tr>
<tr>
<td>(^{56})Fe</td>
<td>(^2)</td>
<td>0.88 MeV</td>
<td>0.23</td>
<td>0.26</td>
<td>0.26</td>
</tr>
<tr>
<td>(^{56})Fe</td>
<td>(^2)</td>
<td>2.65 MeV</td>
<td>0.10</td>
<td>0.10</td>
<td>0.10</td>
</tr>
<tr>
<td>(^{58})Ni</td>
<td>(^2)</td>
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<td>0.24</td>
<td>0.24</td>
</tr>
<tr>
<td>(^{58})Ni</td>
<td>(^2)</td>
<td>4.08 MeV</td>
<td>0.24</td>
<td>0.24</td>
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<tr>
<td>(^{60})Ni</td>
<td>(^2)</td>
<td>1.33 MeV</td>
<td>0.27</td>
<td>0.27</td>
<td>0.27</td>
</tr>
<tr>
<td>(^{60})Ni</td>
<td>(^2)</td>
<td>4.04 MeV</td>
<td>0.27</td>
<td>0.27</td>
<td>0.27</td>
</tr>
<tr>
<td>(^{62})Ni</td>
<td>(^2)</td>
<td>1.17 MeV</td>
<td>0.22</td>
<td>0.22</td>
<td>0.22</td>
</tr>
<tr>
<td>(^{62})Ni</td>
<td>(^3)</td>
<td>3.75 MeV</td>
<td>0.19</td>
<td>0.19</td>
<td>0.19</td>
</tr>
</tbody>
</table>

The \(^2\) and \(^3\) state in \(^{58}\)Ni, \(^{60}\)Ni and \(^{62}\)Ni (fig. 5.7 to 5.9).

From these figures one sees that the experimental data on these three isotopes are much alike for each bombarding energy. No large systematic changes with mass number are observed. The analysing powers are rather well reproduced by the theoretical curves. For the \(^2\) states as well as for the \(^3\) states these curves are calculated with the same deformation parameter for the central and spin-orbit part. We mention that no improvements are obtained if the ratio of these parameters is changed.
Fig. 5.7 Differential cross section and analysing power for inelastic scattering of protons ($E_p=20.4, 24.6$ MeV) by $^{58}$Ni. The curves are obtained from macroscopic DWBA calculations using the FG potentials.
Fig. 5.8 Differential cross section and analysing power for inelastic scattering of protons ($E_p = 20.4, 24.6$ MeV) by $^{60}$Ni. The curves are obtained from macroscopic DWBA calculations using the FG potentials.
Fig. 5.9 Differential cross section and analysing power for inelastic scattering of protons ($\gamma_p=20.4, 34.6$ MeV) by $^{62}$Ni. The curves are obtained from macroscopic DWBA calculations using the FG potentials (solid curves) and the BG potentials (dashed curves).
Generally it is seen that the analysing powers agree much better with the theoretical predictions than the cross sections do. It is especially remarkable that the slopes of the experimental cross sections are definitely steeper than those resulting from the theory. One can observe that this discrepancy becomes less serious at the higher bombarding energy. This agrees with the results from the experiments of Karban et al. at 30 MeV (Ka70) and from the experiment of Fricke et al. at 40 MeV (Fr67) who found that both the analysing power and the differential cross section were reproduced fairly well over a large angular region. It may be questionable whether in our experiment the theoretical curves fail to predict the cross section at the forward or at the backward scattering angles, as the absolute scale of the theoretical cross sections, determined by the deformation parameter, is not exactly known. However, since our deformation parameters are only 10-15% larger than those reported by Fricke et al., it is near to assume that the main deficiency occurs at the backward angles.

We did not succeed to obtain a better fit at these backward angles by trying the BP or BG parameter sets. Also no significant improvements resulted if the deformation parameters for the imaginary or the spin-orbit potential were varied.

The $2^+$ (0.85 MeV) state in $^{56}$Fe (fig. 5.10 and 5.11).

From figure 5.10 we see that difficulties similar to that encountered for the Ni isotopes arise in the description of the differential cross sections at the backward angles. It is clear that this discrepancy is not so large here. The analysing powers are well described by the DWBA calculations over the entire angular region. At forward angles the experimental data are found to be slightly more positive than for the Ni isotopes. This can readily be described by a somewhat larger spin-orbit deformation of about 1.25 times the deformation of the central potential. One may get the impression that this factor should be larger for a better fit to the analysing power at forward angles at 15.3 MeV.
Fig. 6.10 Differential cross section and analysing power for inelastic scattering of protons ($E_p = 17.2, 20.4, 24.8$ MeV) by $^{56}$Fe. The curves are obtained from macroscopic DWBA calculations using the FG potentials.
Fig. 5.11 Differential cross section and analysing power for inelastic scattering of protons \(E_p=15.3\) MeV by \(^{56}\text{Fe}\). The curves are macroscopic DWBA calculations using the FG potential.

The \(2^+\) state (1.41 MeV) in \(^{54}\text{Fe}\) (fig. 5.12 and 5.13).

If one compares the data on the analysing powers for \(^{56}\text{Fe}\) with those measured at the same bombarding energy for \(^{56}\text{Fe}\) or the \(\text{Ni}\) isotopes, one immediately observes significant differences, mainly at the forward scattering angles. The larger analysing powers for \(^{54}\text{Fe}\) around 30\(^\circ\) were reported about ten years ago by Glashausser et al. (G167) and Hendrie et al. (He69) for bombarding energies of 18.6 MeV and 19.6 MeV respectively. However, at higher bombarding energies Karban et al. (30 MeV)(Ka70) and Fricke et al. (40 MeV) (Fr67) found no evidence that the analysing power for this \(2^+\) state in \(^{54}\text{Fe}\) behaves different from analysing powers for other \(2^+\) states in this mass region. In our data we observe a clear energy dependence in the maximum value for the analysing powers around 30\(^\circ\). This energy
Fig. 5.18 Differential cross section and analysing power for inelastic scattering of protons (E_p = 17.2, 20.4, 24.6 MeV) by $^{56}$Fe. The curves are obtained from macroscopic DWBA calculations using the FG potentials.
Fig. 5.13 Differential cross section and analysing power for inelastic scattering of protons ($E_p = 15.3$ MeV) by $^{54}$Fe. The curves are obtained from macroscopic DWBA calculations using the FG potential.
Pig. 5.14 Chi-square values calculated for the DWBA description of the analysing power ($\theta<100^\circ$) for the $3^+$ (1.41 MeV) state in $^{56}$Fe. In the DWBA calculations the FG potentials were used.

dependence is reasonably described if we assume deformations of the spin-orbit potential, which increase with decreasing projectile energy. Chi-square values for the analysing powers at 15.3 MeV, 17.2 MeV, 20.4 MeV and 24.6 MeV were calculated for several ratios of the spin-orbit and central deformation, for scattering angles smaller than $100^\circ$. These values are given in figure 5.14. Beside the increasing ratio with decreasing energy also an increase in the uncertainty of this ratio is obvious. At 15.3 MeV the minimum $\chi^2$-value remains much larger than that obtained for the other energies, and clearly the fit shown in figure 5.13 for $\beta_c = \beta_{s.o.}$ is much worse than those for the higher bombarding energies. We also analysed the experimental results of the Saclay group at 18.6 MeV (GI67) and 19.6 MeV (He69) with a reasonable interpolation of our fixed geometry potentials. The spin-orbit deformation
obtained from their data together with those of the present analysis are presented in figure 5.15. For comparison we give in the same figure the spin-orbit deformations obtained from analyses on $^{54}$Fe at 30 MeV (Ka70) and on $^{62}$Ni between 16 MeV and 25 MeV (Gl67). This picture shows a clear energy dependence for the $^{54}$Fe values which is not likely to be produced by a direct reaction process only. Thus these results do not agree with the suggestion of Raynal (Ra71), that the larger spin-orbit deformations for the $^{2+}$ (1.41 MeV) state in $^{54}$Fe are due to the iso-spin dependence of the two-body spin-orbit interaction.

If we consider the differential cross sections at 15.3 MeV and 17.2 MeV it is obvious that they are not described so well. Although
the bombarding energies are amply above the (p,n) threshold, the rise of the cross sections at backward angles, especially at 15.3 MeV, suggests contributions of compound nucleus effects. However, we expect such contributions not to be large at the forward scattering angles, because there the direct reaction contribution shows a large maximum. Moreover, compound nucleus effects would have drastically reduced the analysing powers at backward angles, which is not observed in our experimental data. We notice that for all bombarding energies the slopes of the cross sections are found to be much more flat as compared to the results for $^{56}$Fe and the Ni isotopes.

The $2^+$ state (2.96 MeV) in $^{54}$Fe (fig. 5.16 and fig. 5.13).

Contrary to the observations for the first $2^+$ state in $^{54}$Fe both the analysing powers and the differential cross sections are nicely fitted at 17.2 MeV, 20.4 MeV and 24.6 MeV. For these energies the experimental data for the analysing powers resemble those obtained for the first $2^+$ state in $^{56}$Fe. This resemblance dramatically disappears at 15.3 MeV, where the analysing power for the second $2^+$ state shows a shape which does not agree with an $L=2$ transfer. Nevertheless the differential cross section is reproduced very well. However, this may be somewhat fortuitous, since the rise in the deformation parameter with decreasing energy (Cf. table 5.4) shows that a contribution additional to the calculated direct process should be present.

Calculations with the best fit optical model parameters from our experiment and those from the experiment of Lombardi et al. (Lo72) and the global potential of Becchetti and Greenlees have been performed. In all these calculations we found a pattern for the analysing power similar to that shown by the solid curve given in figure 5.14. Like demonstrated also a change in the deformation of either the imaginary potential or the spin-orbit potential will not be successful; all the calculations show a maximum around 120°, which is not observed in our experimental data. Moreover we mention that no changes in the theoretical curves were obtained if the optical potential in the exit channel was adjusted for the lower energy of the scattered proton.
Fig. 5.16 Differential cross section and analysing power for inelastic scattering of protons ($\epsilon_p = 17.2, 20.4, 24.6$ MeV) by $^{54}$Fe. The curves are obtained from macroscopic DWBA calculations using the FG potentials.
The $2^+$ (3.16 MeV) state, $4^+$ (2.56 MeV, 3.84 MeV and 4.28 MeV) states in $^{54}$Fe and the $2^+$ (2.65 MeV) state in $^{56}$Fe.

Due to much higher beam currents in our experiments at 15.3 MeV we have obtained also results for these weaker excited states with reasonable statistics. These results are shown in the lower part of figure 5.13, in figure 5.17 and in the lower part of figure 5.11 together with various macroscopic DWBA calculations.

Our results for the differential cross section of the $4^+$ (2.56 MeV) and $4^+$ (3.84 MeV) agree well with those measured by Gray et al. (Gr67) at 17.9 MeV. Their results on the shape of the differential cross section for the $2^+$ (3.16 MeV) state is also comparable to ours. Nevertheless we find a significantly larger deformation parameter (0.09) than the value 0.06 as reported by Gray. We remind that a similar effect is noticed from our deformation parameters for the $2^+$ (2.96 MeV) state.

It is seen from the figures that the angular distribution of the analysing power for the second and third $2^+$ states in $^{54}$Fe and the second $2^+$ state in $^{56}$Fe are much alike and that they certainly do not agree with the general pattern expected from an $L=2$ DWBA calculation.

The experimental data for the $4^+$ states in $^{54}$Fe at 3.84 and 4.28 MeV are described not so badly, especially if we regard the difficulties met for the excited $2^+$ states. We mention that it has been suggested (Gr67) that the large differential cross section for the state at 2.56 MeV beyond 90° may be due to the contribution of an additional $0^+$ level near this excitation energy.

In comparing the analysing powers measured at 15.3 MeV one gets the impression that those for the second and third $2^+$ state in $^{54}$Fe and the second $2^+$ state in $^{56}$Fe are more alike to that for the $4^+$ states in $^{54}$Fe than to that for the lowest excited $2^+$ state in $^{54}$Fe and $^{56}$Fe. Similar observations have been reported recently by Plavko et al. (Pl77). Mainly based on experimental data for $^{58}$Ni at 18.6 MeV they stated that for states with different spin or parity but with approximately the same excitation energy analysing powers with almost the same shape are obtained. Supporting the idea of Raynal they suggested an explanation based on the dominance of either proton or neutron configurations in the excitation process. In our case, we think such an explanation unrealistic in view of the large energy dependence, which is abundantly clear for the second $2^+$ state in $^{54}$Fe.
Fig. 5.17 Differential cross section and analysing power for inelastic scattering of protons ($E_p = 15.3$ MeV) by $^{54}$Fe. The curves are obtained from macroscopic DWBA calculations using the FG potentials.
5.3.2 Discussion

In this section we showed that in general a standard DWBA analysis with a macroscopic form factor can give a satisfactory description of the analysing powers and differential cross sections of inelastic scattering by nuclei around $A=60$ and bombarding energies at around 20 MeV. The previously found anomaly for the inelastic scattering from the lowest excited $2^+$ state in $^{54}$Fe is reproduced and moreover it is found to depend largely on the projectile energy. At an energy of 15.3 MeV it is seen that the usual description of this anomaly in terms of a larger deformation of the spin-orbit potential fails to reproduce the experimental data. The data on the second $2^+$ state in $^{54}$Fe show even more clearly that large changes occur in the analysing power with decreasing bombarding energy.

From our DWBA calculations large changes with bombarding energy are not expected from a direct reaction process. It is therefore most probable that an additional energy dependent contribution to the reaction amplitude is responsible for the observed deficiencies in the DWBA calculations. At the lower energies compound nucleus contributions cannot be ruled out, especially at the backward angles. However, at the forward scattering angles, where the anomaly for the $2^+$ (1.41 MeV) state is most pronounced, we expect these contributions to be small. As shown by Von Geramb (Ce73) semi-direct reaction amplitudes may be caused by the excitation of giant resonances which occur as doorway states in the target nucleus. Recently it has been shown that by the inclusion of such effects the analysing power even at forward scattering angles can change appreciably (Le75). A suggestion to explain the anomalous behaviour in this way has been made by Van Hall (Ha77). We mention that the existence of resonances at this high excitation energy in medium weight nuclei is supported by measurements of the elastic scattering of polarized protons from $^{56}$Fe, which reveal a broad resonance around a proton energy of 14.8 MeV (We76).

At 15.3 MeV we observed also difficulties in describing the data obtained for the excitation of the second and third $2^+$ state in $^{54}$Fe and the second $2^+$ state in $^{56}$Fe. Regarding these last two cases it would be interesting to see whether these difficulties also persist at higher bombarding energies.
REFERENCES

Ki75  H.J. Kim, Nucl. Data Sheets 16 (1975) 317.
Ra71 J. Raynal, Structure of nuclei, Trieste Lectures 1971, p. 75
Ve70 H. Verheul, Nucl. Data Sheets B3 no. 5-6 (1970) 161.
6.1 Introduction

During the past decade considerable effort has been made to describe the inelastic scattering of protons with microscopic models of nuclear structure. In this description both the effective nucleon-nucleon interaction, which produces the inelastic transition, and the spectroscopy of the nuclear states are involved. In order to derive suitable expressions for the effective interaction, the differential cross sections for the excitation of low lying states in $^{89}$Y (Aw67, St67, Be68, Sc69, Ge71, Wh72), $^{207}$Pb (Gl68, Wa73, Wa75, Sc77) and $^{209}$Bi (Sc74) have been studied extensively, since these states are expected to be well described by simple shell model configurations. However, these studies reveal that the active participation of the core nucleons has to be included in the transition density, both to improve the fit to the experimental data and to obtain a reasonable strength for the effective interaction. Nevertheless these reactions are generally considered to provide a convenient test for the character of the effective interaction.

The main object of our experiment is to furnish an additional test for the microscopic description of the inelastic scattering by $^{89}$Y. The inclusion of the analysing powers may impose more stringent requirements on the parameters for the effective interaction. The ground state of $^{89}$Y is believed to be well described by a single proton in a $2p_{1/2}$ orbit outside a closed $^{88}$Sr core. The lowest excited states are believed to be described predominantly by single proton excitations. Since protons scattered from valence protons are more strongly affected by the two-body spin-orbit (TBSO) interaction than protons scattered from valence neutrons (Lo72) the analysing powers measured for the low lying states can be used to study the influence of this TBSO interaction. This study eventually may be helpful to clarify to what extent this interaction can account for the observed large analysing powers reported for the first $2^+$ states in closed neutron shell nuclei such as $^{54}$Fe, $^{90}$Zr and $^{92}$Mo (Gl67, Gl69, He69, Ha77).
In our experiment we measured angular distributions of the differential cross sections and analysing powers for the ground state, the excited states at 0.91 MeV, 1.51 MeV, 1.74 MeV and 2.22 MeV and for the unresolved multiplets at 2.54 MeV, 2.87 MeV and 3.12 MeV. In section 6.2 we discuss some of the experimental details of our experiment. The elastic scattering has been analysed with the optical model. The results of this analysis is presented in section 6.3. The inelastic scattering data are compared to macroscopic DWBA descriptions in section 6.4.1. Our microscopic model analysis for the states at 0.91 MeV, 1.51 MeV and 1.74 MeV is described in section 6.4.2.

6.2 The experiment

Our experimental arrangement and method has been discussed in detail in chapters 2 and 3. The experimental data were measured in one continuous beam run of about 70 hours. Typical values for the beam current on the target were 30 nA and the polarization of the beam was about 78%. Differential cross sections and analysing powers have been measured at 43 different scattering angles ranging from 22.5° up till 165°. The target was about 3 mg/cm² thick. An example of the obtained energy spectra is shown in figure 6.1. The overall energy resolution was about 100 keV for the spectra measured in transmission mode (generally from 22.5° to 90°) and about 150 keV for the spectra measured in reflection. The analysis of the energy spectra was done with the automatic peak search programme POESPAS (BL75) and in some cases, mainly for the weaker excited states, by using the PDP-9 with its CRT display.

Since we did not use a particle identification system the α-particles from the \(^{89}\text{Y}(p,\alpha)^{86}\text{Sr}\) (Q = 1.69 MeV) reaction were not separated from the proton spectra. However, the cross sections for this reaction are known to be at least a factor 5 till 10 smaller than the cross sections for the weaker excited states in \(^{89}\text{Y}\) (Ve74, Co75). Due to the large energy spread caused by the relatively thick target the α-spectra produced a smooth background so that their contributions are largely corrected for by the rectilinear background subtraction we performed around each proton peak.
Fig. 6.1 Energy spectrum of protons scattered by $^{89}$Y. The lower and upper part of the spectrum correspond to the different spin orientations of the incoming beam.

The first excited state at 0.91 MeV is very weak and moreover located on the tail of the elastic peak. Therefore the background subtraction formed the major source of the uncertainties in the data derived for this level at the most forward scattering angles. The accuracy of the data for the state at 1.74 MeV at angles from 22.5° till about 30° might be affected by a peak near 1.80 MeV, which is known to arise from the inelastic scattering from $^{28}$Si ($E_x = 1.78$ MeV) within the detector volume (Th67). Both these peaks at 1.74 MeV and 1.78 MeV were fitted simultaneously by assuming the same peak shape.

Large contributions from the elastic scattering by $^{12}$C and $^{16}$O contaminants in the target appeared in the energy spectra. Where these peaks overlapped with the $^{89}$Y peaks a correction was made using the polyethylene and mylar spectra which were measured for the same scattering angle. Sometimes these corrections resulted in such large errors that the data were left out. Also much smaller contributions were found from other contaminants whose mass numbers were identified by the kinematic shift with scattering angle to be about 19, 28, 32 and 40. In overlap situations with the states at 0.91 MeV, 1.51 MeV
and 1.74 MeV the contributions from the last three contaminants were subtracted whereas the data corresponding to an overlap with \( A=19 \) were omitted. For \( A=28 \) and \( A=32 \) we used the angular distributions calculated from the optical model parameters derived by Blair et al. (8170) for \(^{28}\)Si at 20.1 MeV assuming the contaminants to be \(^{28}\)Si and \(^{32}\)S. For \( A=40 \) a similar procedure was followed assuming the contaminant to be \(^{40}\)Ca and using the parameters of Bray et al. (8171) for the scattering by \(^{40}\)Ca at 21.0 MeV. Both sets of parameters were obtained in analysis which included the analysing powers. Only the \( A=28 \) and \( A=32 \) contaminants were found to contribute significantly: in some extreme cases as much as 30%.

The errors shown with the experimental data in this chapter include those due to counting statistics, an estimate of the error made in the background subtraction as well as an estimate of the error made in the subtraction of the contributions of the contaminants.

6.3 Elastic scattering.

The elastic scattering has been analysed in terms of the phenomenological optical model using the search programme OPTIMO (Vo72). The parameters of Beochetti and Creemles (Cf. section 4.3.1.) were used as starting values in the parameter search. Relative errors of 3% were assigned to the differential cross section. An absolute error of 0.01 was taken

![Graph](image)

**Fig. 6.2** Differential cross section and analysing power for the elastic scattering of 21.1 MeV protons by \(^{88}\)Y. The curves through the experimental points are calculated from the best fit optical model parameters.
for the analysing power unless the experimental error exceeded that value. The absolute normalization of the experimental differential cross section was used as an additional parameter in the search process.

The best fit to the experimental data is shown in figure 6.2. The corresponding parameters are listed in the first column of table 6.1. We mention that inclusion of a volume absorption term resulted in a value very close to zero, whereas no appreciable improvement in the quality of the fit was obtained. We found the parameters and the normalization factor from our search to be significantly different from the values deduced by Hulstman from differential cross section data measured at 20.73 MeV (Hu73). Therefore we performed separate searches on our differential cross section and analysing power. The parameters from these searches are given in the second and third column of table 6.1. Clearly the best fit parameter set is largely determined by the analysing power. The parameters and normalization factor, resulting from the search on the differential cross section only, are found to be in close agreement to those reported by Hulstman. In our analysis of the inelastic scattering we used the best fit parameters, which follow from the more complete analysis.

### Table 6.1 Optical potential parameters for \( ^{86}Y(p,p') \) \( E_p = 21.1 \) MeV 

<table>
<thead>
<tr>
<th>( V_0 ) (MeV)</th>
<th>( r_0 ) (fm)</th>
<th>( a_0 ) (MeV)</th>
<th>( V_1 ) (MeV)</th>
<th>( r_1 ) (fm)</th>
<th>( a_1 ) (MeV)</th>
<th>( V_{20} ) (MeV)</th>
<th>( r_{20} ) (fm)</th>
<th>( a_{20} ) (MeV)</th>
<th>( \chi^2 )</th>
<th>( \chi^2_{\nu} )</th>
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<td>Best fit</td>
<td>56.0</td>
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<td>0.76</td>
<td>6.43</td>
<td>1.31</td>
<td>0.71</td>
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<td>0.90</td>
<td>0.72</td>
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<td>441</td>
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<td>( \sigma ) only</td>
<td>51.4</td>
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<td>0.89</td>
<td>0.71</td>
<td>9652</td>
<td>329</td>
</tr>
</tbody>
</table>

\( * \) Coulomb radius \( r_0 = 1.25 \) fm.

The best fit parameters are listed in the first column of table 6.1. We mention that inclusion of a volume absorption term resulted in a value very close to zero, whereas no appreciable improvement in the quality of the fit was obtained. We found the parameters and the normalization factor from our search to be significantly different from the values deduced by Hulstman from differential cross section data measured at 20.73 MeV (Hu73). Therefore we performed separate searches on our differential cross section and analysing power. The parameters from these searches are given in the second and third column of table 6.1. Clearly the best fit parameter set is largely determined by the analysing power. The parameters and normalization factor, resulting from the search on the differential cross section only, are found to be in close agreement to those reported by Hulstman. In our analysis of the inelastic scattering we used the best fit parameters, which follow from the more complete analysis.

### 6.4 Inelastic scattering

#### 6.4.1 Macroscopic DWBA analysis

In this section the experimental differential cross sections and analysing powers are compared with the results of macroscopic DWBA calculations. The best fit optical model parameters were used to calculate the elastic scattering wavefunctions both for the entrance
and exit channel; a non-locality correction was not included. From calculations with the Becchetti Greenlees parameters we found no noticeable differences if the optical model parameters were corrected for the lower energy in the exit channel. Unless otherwise stated the real, imaginary and spin-orbit part of the optical potential were deformed by taking the same deformation parameter for each term.

For the spin-orbit part we used the full Thomas coupling. The calculations were performed with a DWBA programme written by L.D. Tolsma and adapted by Verhaar et al. (Ve72) to include the full Thomas spin-orbit coupling.

The experimental angular distributions together with the macroscopic descriptions are given in figure 6.3 through 6.5. The deformation parameters, which correspond to the solid curves in our figures, are compared with values reported from several other experiments in table 6.2. We mention that all of these values include the statistical factor \([(2l+1)/(2l+1)(2l+1)]^1\), and thus they correspond to the \(l^l= values\) defined in chapter 4 by (4.27). The spin and parity assignments are those suggested from these other experiments and from the \((n,\alpha')\) experiment of Alster et al. (Al66), the \(^{90}\text{Zr}(d,^3\text{He})^{86}\gamma\) experiment of Freedon et al. and the \((n,n',\gamma)\) experiment of Patzer and Schrotroth (Pa64).

From the high resolution work of Hulstman it is known that the states that we found at 2.54 MeV and 2.87 MeV comprise both a triplet of states at 2.53 MeV (\(L=3\)), 2.565 MeV (\(L=5\)), 2.621 MeV (\(L=5\)) and 2.872 MeV (\(L=3\)), 2.882 MeV (\(L=2\)) and 2.893 MeV (probably \(L=4\)); the state

<table>
<thead>
<tr>
<th>(E^{\text{lab}}) (MeV)</th>
<th>(L)</th>
<th>(J^\pi)</th>
<th>(\alpha)</th>
<th>((\alpha\pi))</th>
<th>((\alpha\pi))</th>
<th>((\alpha\pi))</th>
</tr>
</thead>
<tbody>
<tr>
<td>2.54</td>
<td>2</td>
<td>0.032</td>
<td>0.008</td>
<td>0.008</td>
<td>0.008</td>
<td>0.008</td>
</tr>
<tr>
<td>2.565</td>
<td>3</td>
<td>0.015</td>
<td>0.005</td>
<td>0.015</td>
<td>0.015</td>
<td>0.015</td>
</tr>
<tr>
<td>2.621</td>
<td>(3/2)</td>
<td>0.032</td>
<td>0.008</td>
<td>0.008</td>
<td>0.008</td>
<td>0.008</td>
</tr>
<tr>
<td>2.872</td>
<td>(3/2)</td>
<td>0.032</td>
<td>0.008</td>
<td>0.008</td>
<td>0.008</td>
<td>0.008</td>
</tr>
<tr>
<td>2.893</td>
<td>(4)</td>
<td>0.032</td>
<td>0.008</td>
<td>0.008</td>
<td>0.008</td>
<td>0.008</td>
</tr>
</tbody>
</table>

\(^2\text{Peterson and Alster (Pa64)}; ^3\text{Hulstman (Hu72)}; ^4\text{Freeman et al. (Fr68)}; ^5\text{Sect et al. (Se69)}; ^6\text{Penrose analyzation.}\)
at 3.12 MeV is known to be a doublet: 3.105 MeV (L=2) and 3.137 MeV (L=2) (Hu73). These multiplets were analysed in our experiment as well as by Benenson (Be68) and Peterson (Pe68) under the assumption that we are dealing with one level only, instead of a multiplet.

The $\beta_3$-values derived from the inelastic proton scattering at the various bombarding energies are found to be in close agreement. Slight deviations are most likely introduced by the uncertainty in the absolute normalization in these experiments. All our deformation parameters are consistent with the values reported for 24.5 MeV. The value found for the 1.74 MeV state may be somewhat inaccurate, in Benenson's and our experiment, due to the difficulties encountered at forward scattering angles in the subtraction of the peak near 1.80 MeV. However, it is striking that the deformation parameters for the three lowest excited states are significantly larger than those at the 61.2 MeV bombarding energy.

The several fits obtained for each state are discussed successively below.

The 0.91 MeV state (fig. 6.3):

Although it is known from the experiments at 24.5 MeV and 61.2 MeV that a microscopic description for this state seems to be more appropriate, the macroscopic description corresponding to an L=5 transfer is given for comparison. Except for scattering angles smaller than 40° the differential cross section is reproduced reasonably, whereas the analysing power is described surprisingly well over the entire angular region. Increasing the spin-orbit deformation has only a slight effect on both the analysing power and the differential cross section and does not result in appreciable improvements of the fit, as can be seen from figure 6.3.

The 1.51 MeV and 1.74 MeV state (fig. 6.3):

These excited states have spin and parity 3/2$^-$ and 5/2$^-$ respectively. Since these are just the values expected from the coupling of a $2p_1/2$ proton to the $2^+$ state at 1.84 MeV of the $^{88}$Sr core, a description in terms of the weak coupling excited core model seems to be appropriate. Indeed the effective deformation parameters for these states are roughly proportional to $\sqrt{2I_z+1}$. However, the transition strength to this doublet exhausts only 30% of the strength found by Stautberg et al. (St67) for the $2^+$ state in $^{88}$Sr. This disagreement was also reported
Fig. 6.3  Differential cross section and analysing power for the inelastic scattering of 21.1 MeV protons by $^{89}$Y. The curves are obtained from macroscopic DWBA calculations using the best fit optical model parameters.
from the $^{88}\text{Sr}(\alpha,\alpha')$ and $^{89}\text{Y}(\alpha,\alpha')$ experiments of Alster et al. (A166).

Nevertheless, by the close resemblance of the differential cross sections the question remained open whether these data should be described within the weak coupling or the microscopic approach. In comparing the analysing powers for these two states large differences are seen in the angular region around $120^\circ$. This clearly indicates that a weak coupling description of these states is inadequate.

The collective model description, assuming $L=2$, reproduces the differential cross sections rather well whereas only a good description is obtained for the analysing power for the state at $1.74$ MeV except for the most backward angles. Further improvement results by taking a two times larger spin-orbit deformation. Such a value is consistent with those found at approximately the same bombarding energy for the first $2^+$ state in other closed neutron shell nuclei like $^{90}\text{Zr}$ and $^{92}\text{Mo}$ (G169).

The $2.22$ MeV, $2.54$ MeV and $2.87$ MeV states (fig. 6.6).

These strongly excited states of which the last two are unresolved triplets are quite good described assuming $L=3$ transfer. The fits to the state at $2.87$ MeV are somewhat poorer. However, a better agreement both for the cross section and the analysing power is obtained if we take into account the $L=2$ transition strength given by Hulstman. As shown for the $2.22$ MeV state the large positive values for the analysing powers at the forward scattering angles could be described if we took a two times smaller deformation for the imaginary part of the optical potential. Thereby also a slight shift in the angular distribution for the analysing powers results, which reproduces the data for scattering angles till $100^\circ$ definitely better at the cost of poorer agreement between $120^\circ$ and $150^\circ$. We found that a change in the deformation of the spin orbit potential did not improve the fit significantly.

The state at $3.12$ MeV (fig. 6.5).

Assuming $L=2$ transfer, both the angular distributions are fitted reasonably well. A larger deformation of the spin orbit potential improves the fit to the analysing power appreciably. At backward angles the same difficulties are encountered as for the description of the $1.74$ MeV state.
Fig. 6.4 Differential cross section and analysing power for the inelastic scattering of 21.1 MeV protons by $^{89}$Y. The curves are obtained from macroscopic DWBA calculations using the best-fit optical model parameters.
6.4.2 Microscopic DWBA analysis

In the microscopic description of the lowest excited states we will use the following simple proton shell model configurations.

\[
\begin{align*}
|1/2^-, g.s.\rangle &= |2p_{1/2} > |{}^{88}\text{Sr}\rangle \\
|9/2^+, 0.91 \text{ MeV} \rangle &= |1g_{9/2} > |{}^{88}\text{Sr}\rangle \\
|3/2^-, 1.51 \text{ MeV} \rangle &= |2p_{3/2}^{-1}; 2p_{1/2} > |{}^{88}\text{Sr}\rangle \\
|5/2^-, 1.74 \text{ MeV} \rangle &= |1f_{5/2}^{-1}; 2p_{1/2} > |{}^{88}\text{Sr}\rangle
\end{align*}
\]

Thus the ground state is interpreted as a single \(2p_{1/2}\) proton outside the \(^{88}\text{Sr}\) core. The first excited state is then described by exciting this proton to the \(1g_{9/2}\) shell, whereas the second and third state are obtained by promoting a \(2p_{3/2}\) and a \(1f_{5/2}\) proton respectively to the \(2p_{1/2}\) state. This simple model is confirmed by the spectroscopic factors derived by Freedon et al. from their \(^{90}\text{Zr} (d,\alpha)^{89}\text{Y}\) experiment (Pr69), and moreover, for the ground and the first excited state by the spectroscopic factors obtained from a \(^{88}\text{Sr} (^3\text{He},d)^{89}\text{Y}\) experiment by Stautberg et al. (St67).
It is known from the studies on the differential cross sections, at the various energies, that the above mentioned configurations are oversimplified and that one has to take into account appreciable core polarization contributions. In the present analysis this is done by using the collective model approach as proposed by Love and Satchler (Lo67). Thus the transition amplitude is composed of a valence part, which describes the single particle excitation, and a core polarization term, having a strength characterized by means of the core coupling parameter $y_L$. In principle this parameter can be derived from the effective charges needed to describe the measured electromagnetic transition rates to these states. However, for the first excited state this value has not been measured whereas the uncertainties in the values for the second and third excited state are large (Pe68). We therefore adjusted the core coupling parameters in order to obtain the absolute magnitude of the differential cross section data at the forward scattering angles. Thereby, of course, this study provides only a limited test of the microscopic DWIA description. But for a first estimate of the importance of the several terms in the transition amplitude we rather take this convenient parametrization of the core components then employing an extensive full-microscopic calculation.

The single proton wavefunctions used in the calculation of the transition density are taken as eigenstates of a Saxon-Woods potential with radius $r=1.25 \text{ fm}$, diffuseness $a=0.65 \text{ fm}$ and a spin-orbit potential defined by (4.12) with the same geometry as the central potential and a strength $V_{\text{s.o.}} = 6.25 \text{ MeV}$. The Coulomb potential associated with a uniformly charged sphere with the same radius as the Saxon Woods potential was included. For the binding energy of the proton in the $2p_{1/2}$ shell $7.0 \text{ MeV}$ is taken, which is approximately equal to the proton separation energy. The binding energies for the proton in the $2p_{3/2}$ and $1f_{5/2}$ shell were taken in agreement with the separation energy and the observed excitation energies. The well depth of the Saxon Woods potential was adjusted to obtain the mentioned binding energies. Non-locality corrections have not been included.

The binding energies and the potential well depths are given in table 6.3.

For the central part of the effective nucleon-nucleon interaction we used the long range part of the even state Hamada-Johnston potential with a separation distance of $1.05 \text{ fm}$. Since the excited and scattered
Table 6.3 Binding energies and corresponding well depths

<table>
<thead>
<tr>
<th></th>
<th>$2p_{1/2}$</th>
<th>$1g_{9/2}$</th>
<th>$2p_{3/2}$</th>
<th>$1f_{5/2}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>E.E. (MeV)</td>
<td>7.0</td>
<td>6.1</td>
<td>8.5</td>
<td>8.7</td>
</tr>
<tr>
<td>V (MeV)</td>
<td>56.01</td>
<td>57.95</td>
<td>55.96</td>
<td>55.49</td>
</tr>
</tbody>
</table>

particles both are protons this central interaction is restricted to the singlet even component of the interaction. In some of our calculations we included non central forces i.e. a two body spin-orbit interaction and a two body tensor interaction. As the two body spin-orbit force is known to be strongest between like nucleons it might be expected to contribute significantly to the excitation process. We mention that the inclusion of a spin-orbit interaction has been shown to have a distinct effect on the analysing powers for the first excited $2^+$ state in $^{54}$Fe and $^{90}$Zr (Ra71). For these non-central parts of the interaction we took those derived by Eikemeier and Hackenbroich (Ei71) from experimental data on free nucleon-nucleon scattering.

The programme MEPHISTO (Ge73) was used to perform the calculations. In this programme the knock-on exchange contributions are calculated exactly. In the collective description of the core contributions the real, imaginary and spin-orbit part of the optical potential have been deformed. For the latter the Oak-Ridge type of coupling was used (Fr67). The best fit parameters given in section 6.3 were used to calculate the elastic scattering wavefunctions, non-locality corrections have not been included.

The 0.91 MeV state (fig. 6.6 and fig. 6.7).

In the microscopic description a proton is excited from the $2p_{1/2}$ shell to the $1g_{9/2}$ shell. For the direct central part of the interaction this results in the following participating triads (LSJ): (505), (515), (514) and (314). In figure 6.6 the total valence contribution (direct+exchange) and the core contribution of the (LSJ)=(505) type is shown. We found that the (314) triad largely dominates the valence contribution, whereas as large core polarization term is needed to achieve a reasonable description of both the experimental distributions.
Fig. 6.6 Experimental data compared with DWBA calculations, for the excitation of the 0.91 MeV state, including the central nucleon-nucleon interaction (VALENCE) and the macroscopic core polarization (CORE).

Fig. 6.7 Experimental data compared with DWBA calculations, for the excitation of the 0.91 MeV state, including the central nucleon-nucleon interaction ($V$), the macroscopic core polarization ($C$), the two-body spin-orbit force (TBSO) and the two-body tensor force (TBTF).
The core-coupling strength needed in our experiment is $\gamma_5 = 0.00042 \text{ MeV}^{-1}$, which is close to the value derived from the $^{89}\gamma(p,p')$ experiment at 61.2 MeV. (Wh72). The angular distributions comprising both the valence and core polarization contributions tends to be dominated by the latter. Only at the forward scattering angles and for the more backward angles the valence part is seen to contribute significantly. This, as compared to the macroscopic analysis, results in a better description for the differential cross section around $30^\circ$ and a better description for the analysing power around $150^\circ$.

In figure 6.7 the results of a calculation including either a TBSO force or the TBTF are given. The inclusion of these non-central parts has only a small effect on the differential cross section, whereas the results for the analysing power suggest a smaller contribution of the TBTF. In view of the large (macroscopic) core contribution it seems to be premature to draw definite conclusions from these observations. We mention that a microscopic analysis of the core contributions by Petrovich (Pe70) shows that the contribution to the differential cross section from the (314) triad is reduced by about a factor 0.6 and that from the (505) triad is enhanced about a factor of 0.7.

**I, 51 MeV state (fig. 6.8, 6.9 and 6.10)**

This state is considered as an excitation of a proton from the $2p_{3/2}$ shell to the $2p_{1/2}$ shell. Thus the following triads may contribute to the direct valence term: (011), (211), (202) and (212). The calculations presented in figure 6.8 show that the $L=0$ transfer is highly favoured and that the valence contribution alone reproduces the absolute magnitude of the differential cross section until about $120^\circ$. Clearly the angular distribution of the experimental cross section shows much more structure than the calculated one. The curve calculated for the valence contribution to the analysing power gives a better description in the angular region where the macroscopic description fails (around $120^\circ$), however, in the region around $70^\circ$ the macroscopic description is superior. Inclusion of a slight core polarization contribution with a strength parameter $\gamma_2 = 0.00032 \text{ MeV}^{-1}$, which is well in between the lower and upper value as deduced from the effective charges (Pe68), fills up the minimum in the differential cross section at the backward angles considerably, whereas in the same angular region the fit to the analysing power becomes somewhat worse.
Fig. 6.9 Experimental data compared with DWBA calculations, for the excitation of the 1.51 MeV state, including the central nucleon-nucleon interaction (VALENCE) and the macroscopic core polarization (CORE).
As shown in figure 6.9 the analysing power is reproduced much better if the TBSO interaction is included. The differential cross section is only little affected by this interaction and the description remains poor. No appreciable improvements are achieved neither for the cross section or the analysing power if the contribution of the TBF is taken into account. In figure 6.9 the significant effect from the TBSO interaction is clearly demonstrated. However, a marginal note should be made with regard to this observation since we found that it was essential to include the spin-orbit interaction and the core polarization contribution together. So we obtained definitely more structure in the angular distribution of the analysing power in agreement with the experimental data.

In principle the effective interaction contains an imaginary part (Sh59). Moreover, it has been shown, among others by Terrien (Te73), that the addition of the entire macroscopic imaginary coupling to the real microscopic term greatly improves the agreement between the experimental and theoretical differential cross section. Similar observations are reported for the analysing power (Ho72). In order to indicate the influence of a complex interaction we included the macroscopic imaginary formfactor ($L=2$; $E_2=0.038$) in our calculations. The results of these calculations are given in figure 6.10. It is seen that a good fit to the analysing power is obtained if both the two-body spin-orbit force and the imaginary formfactor are included, whereas the shape of the differential cross section is hardly affected. These results suggest that the imaginary term in the core polarization contribution may be responsible for the better fit to the analysing power given by the dashed curve in figure 6.9.

The 1.74 MeV state (fig. 6.11)

This state is considered to be excited by lifting a proton from a $1f_{5/2}$ orbit to a $2p_{1/2}$ orbit. Thus the triads involved in this transition are (202), (212), (213) and (413). As shown in fig. 6.11 this single particle excitation alone hardly contributes to the total differential cross section. Consequently in this apparent simple microscopic description a large core polarization term has to be added. The core-coupling parameter needed here to reproduce the differential cross section at the forward scattering angles is about 0.0010 MeV$^{-1}$, which is slightly higher than the upper value reported by Peterson (Pe68). Because of this large core-polarization term
Fig. 6.10 Experimental data compared with DWBA calculations, for the excitation of the 1.61 MeV state, including the central nucleon-nucleon interaction (V), the macroscopic imaginary term (Im), the two-body spin-orbit force (TBST) and the two-body tensor force (TBTF).

Fig. 6.11 Experimental data compared with DWBA calculations, for the excitation of the 1.74 MeV state, including the central nucleon-nucleon interaction (VALENCE) and the macroscopic core polarization (CORE).
the information obtained about the effective interaction from this reaction is rather poor. Nevertheless it can be observed from figure 6.11 that the analysing power in the angular region between 50° and 120° is described better if the microscopic central term is included. We mention that no significantly different results are obtained if the non-central parts of the interaction are included.

6.5 Conclusions.

The macroscopic model of inelastic scattering is found to give a good description of the differential cross section and the analysing power for the excitation of the states at 1.74 MeV (L=2), 2.22 MeV (L=3) and the unresolved multiplets at 2.54 MeV (L=3), 2.87 MeV (L=2+L=3) and 3.12 MeV (L=2). The fit to the analysing power for the states at 1.74 MeV and 3.12 MeV is further improved if a two times larger deformation of the spin-orbit potential is taken. This value is consistent with that obtained by Glasshausser et al. (G169) in the analysis of the analysing powers for the first 2+ state in other closed neutron shell nuclei in this mass region, measured at about the same bombarding energy.

A simple microscopic description of inelastic scattering has been attempted for the three lowest excited states at 0.91 MeV, 1.51 MeV and 1.74 MeV. It is found that a (macroscopic) core polarization term needs to be included, in order to achieve a reasonable description of both the differential cross section and the analysing power. For the state at 1.74 MeV and also, to less extent, for the state at 0.91 MeV, these core polarization contributions are so large that they greatly obscure the influence of the two-body interaction. Contrary to this a noticeable effect is observed in the description for the analysing power for the 1.51 MeV state. Whereas the macroscopic description fails to reproduce this analysing power a reasonable fit is obtained in our microscopic calculation, provided a two-body spin-orbit interaction and a macroscopic L=2 component, either complex or imaginary, are included. However, the magnitude of the differential cross section is overestimated if the non-central parts or the macroscopic component are included. Thus a renormalization of the transition density, which is dominated by the S=1 part of the interaction, is needed. We remark that this does not necessarily imply a smaller strength for the S=1 part of
the interaction, since it has been shown in some particular cases that in a microscopic description of core polarization effects the $S=1$ contribution to the differential cross section is reduced by about of a factor of two (P.73). As we found in our calculation the $S=1$ contribution to be an order of magnitude larger than the $S=0$ contribution, it remains to be seen whether a more elaborate microscopic analysis will end in significantly different results.

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Pa64 D.M. van Pattser and S.M. Schafroth, Nucl. Phys. 50 (1964) 113
Ra71 J. Raynal, Structure of Nuclei, Trieste Lectures 1971, p. 75.
SAMENVATTING

Dit proefschrift omvat een experimentele studie van de verstrooiing van gepolariseerde protonen aan atoomkernen. De experimenten zijn uitgevoerd met behulp van de gepolariseerde ionenbron van het Eindhovense A.V.F. cyclotron. Een beschrijving van de experimentele opstelling wordt in hoofdstuk 2 gegeven. In hoofdstuk 3 wordt nader ingegaan op de gevolgde meetprocedure.

De inelastische verstrooiing van protonen wordt, voor niet te sterk aangeslagen toestanden, in het algemeen goed beschreven door macroscopische DWBA berekeningen. Het is echter bekend dat het analyserend vermogen voor de verstrooiing van protonen, met een energie rond de 20 MeV, naar de eerste aangeslagen $2^+$-toestand van kernen met een gesloten neutronenschil (N=28 en N=50) een afwijkend gedrag vertoont in vergelijking met naburige kernen. In macroscopische DWBA berekeningen kan dit afwijkende gedrag alleen beschreven worden wanneer voor de kernen met een gesloten neutronenschil de spin-baan potentiaal tweé tot drie maal sterker gedeformeerd wordt dan de centrale potentiaal. Aangezien de betreffende aangeslagen toestanden in grote mate beschreven zullen worden door de proton exitaties, zou de grotere deformatie voor de spin-baan potentiaal kunnen duiden op een sterke iso-spin afhankelijkheid van de microscopische spin-baan wisselwerking. Bij protonenergien van 30 MeV en 40 MeV blijkt het afwijkende gedrag niet meer zo duidelijk aanwezig te zijn. Ten einde een veronderstelde energie-afhankelijkheid nader te bestuderen zijn verstrooiingsexperimenten uitgevoerd aan $^{54}$Fe, $^{56}$Fe ($E_p=15.3, 17.2, 20.4, 24.6$ MeV) en $^{58}$Ni, $^{60}$Ni en $^{62}$Ni ($E_p=20.4, 24.6$ MeV).

De resultaten van deze experimenten worden tezamen met macroscopische DWBA berekeningen beschreven in hoofdstuk 5. Uit de gegevens voor de eerste aangeslagen toestand ($2^+, 1.41$ MeV) in $^{54}$Fe blijkt duidelijk dat er energie-afhankelijke effecten een rol spelen. Een goede beschrijving van het analyserend vermogen wordt verkregen met bij 24.6 MeV een 1.5 maal grotere deformatie voor de spin-baan potentiaal, terwijl bij 17.2 MeV een 3 maal grotere deformatie nodig is. Bij 15.3 MeV is een beschrijving van het analyserend vermogen met behulp van een grotere deformatie van de spin-baan potentiaal niet goed meer mogelijk. De geconstateerde energie-afhankelijkheid voor de deformatie van de spin-baan potentiaal duidt er op dat een verklaring gezocht dient.
te worden in een additionele energie-afhankelijke bijdrage.

Het is bekend dat microscopische berekeningen eveneens met weinig succes het analyserend vermogen voor de $2^+$ (1.41 MeV) toestand in $^{54}$Fe beschrijven. Dit kan eventueel te wijten zijn aan de onvolledige golffuncties die voor de beschrijving van dit niveau gebruikt zijn. Ook de informatie over de niet centrale termen in de effectieve nucleon-nucleon wisselwerking is nogal beperkt. Om een nadere bestudering van deze wisselwerking mogelijk te maken zijn verstrooiingsexperimenten uitgevoerd aan $^{89}$Y. De laagst aangeslagen toestanden in $^{89}$Y kunnen redelijk beschreven worden door eenvoudige shellenmodel golffuncties. Daarom juist lijkt de bestudering van de verstrooiing naar deze toestanden een goede test te bieden voor de effectieve nucleon-nucleon wisselwerking. De experimentele resultaten en de macroscopische en microscopische DWBA berekeningen worden beschreven in hoofdstuk 6.

In de microscopische berekeningen wordt in het algemeen slechts een redelijke aanpassing verkregen aan de differentiële doorsneden en de analyserende vermogens wanneer naast de bijdrage van een één-deeltjes excitaatie ook de macroscopische bijdrage ten gevolge van de core-polari-arisatie in rekening gebracht wordt. De invloed van de laatste genoemde bijdrage blijkt zo groot te zijn dat de informatie verkregen over de effectieve wisselwerking beperkt is. Een microscopische beschrijving van de core-polarisatie zou daarom gewenst zijn.
Het tot stand komen van dit proefschrift zou niet mogelijk geweest zijn zonder de inzet van valen.


De cyclotron-bedrijfsgroep ben ik zeer erkentelijk voor hun zorg en moeite waardoor steeds een goed werkend cyclotron verkregen werd.

In verband met de verzorging van dit proefschrift dank ik: Marjike Schilstra voor het typen van het manuscript, Ruth Cuyters, Harm Rozema en Wim Gudden voor het tekenen van de figuren en Peter Teunisse voor de vervaardiging van de foto's daarvan.
LEVENSLoop

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STEELLINGEN

behorend bij het proefschrift van

J.P.M.G. Melssen

Eindhoven, 16 mei 1978
Uit verstrooiingsexperimenten met protonen met energieën tussen 15 MeV en 25 MeV blijkt duidelijk dat het anomalie gedrag in het analyserende vermogen voor de $2^+$ (1.41 MeV) toestand in $^{54}$Fe energieafhankelijk is. Hierdoor is een verklaring van deze anomalie, uitsluitend gebaseerd op een iso-spin afhankelijke spin-baas wisselwerking, uitermate onwaarschijnlijk.

J. Raynal, Structure of Nuclei, Trieste Lectures 1971, p. 75.

Dit proefschrift, hoofdstuk 5.

Bij het ontwerpen van een systeem voor de radiale injectie van deeltjes in een cyclotron verdient het aanbeveling om in ruime mate aandacht te besteden aan de transmissie van de versnelde bundel in de eerste omlopen in het cyclotron.


De meetduur in experimenten met gepolariseerde protonen wordt in hoge mate bepaald door de beschikbare bundelintensiteit. Ter verkrijging van een betere energiedefinitie zal de bundelintensiteit in het algemeen drastisch verminderen. Het nadeel van de langduriger meettijden lijkt niet op te wegen tegen het voordeel van de minder tijdrovende en in het algemeen betrouwbaardere analyse van de verkregen energiespectra.

Het is zaak dat naast experimenten met nieuwe structuren in het onderwijs uitvoerig nagegaan wordt welke resultaten volgen uit veranderingen in de onderwijsmethode binnen de huidige structuur.

Gelet op het voortdurend hoge peil van de werkloosheid is het onaanvaardbaar dat bezuinigingsmaatregelen worden doorgevoerd door middel van personeelstop.
Het is onjuist te beweren dat het houden van internationale sport-ontmoetingen in niet democratisch bestuurde landen een continuering van de politieke situatie in die landen bevordert.

Ter verduidelijking van de verkeerssituatie op kruisingen of splitsingen van wegen van gelijke orde is het aan te bevelen de plicht om voorrang te verlenen aan alle bestuurders die van rechts naderen ook op te leggen aan bestuurders van motorvoertuigen.

Het verdient aanbeveling om de uiterste verkoopdatum vermeld op de verpakking van voedingsmiddelen te vervangen door een uiterste consumptiedatum. Deze datum dient op een niet voor misverstand vatbare wijze te worden aangegeven.

De attractiviteit van de professionele voetbalsport zou bevorderd kunnen worden door de premies voor de spelers te koppelen aan de recette.