A high-precision study of polarized proton scattering to low-lying states in 11B
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6. Model calculations

The following sections describe the results of model calculations which have been performed to reproduce the measured differential cross sections, analyzing powers and spin-flip probabilities whose extraction from the experiment has been described in chapters 4 and 5. The main ingredients of the calculations, namely the effective nucleon-nucleon interaction and the shell-model wave functions, are chosen by comparing calculated results to measured elastic cross sections and electromagnetic transition strengths (see sections 6.1 and 6.2, respectively). In section 6.3 the results obtained for elastic scattering and depolarization effects, which occur in elastic scattering of targets with a non-zero ground-state spin, will be presented. Section 6.4 discusses the results obtained for transitions to negative- and positive-parity states in $^{11}$B.

6.1 Choice of the effective nucleon-nucleon interaction

Several effective nucleon-nucleon interactions have been tested by comparing elastic scattering cross sections and analyzing powers calculated in a microscopic DWBA framework to experimental results. The $^{12}$C$(\tilde{p},p)$ reaction at a proton energy of 150 MeV has been chosen as a test case because, contrary to the $^{11}$B$(\tilde{p},p)$ reaction, no higher multipoles contribute to the elastic scattering cross section. In this case the necessary input from nuclear shell-model calculations is limited to the specification of the ground-state density of the target nucleus. The optical potential is generated by folding the effective interaction with the ground-state density (for details see, e.g., reference [38]). Occupation numbers of the single-particle states can be calculated in a $0\hbar\omega$ model space using the OXBASH shell-model code [19]. The resulting average number of particles in each orbit is: $1s_{1/2} = 2$, $1p_{3/2} = 3.267$ and $1p_{1/2} = 0.733$ and it is equal for protons and neutrons. A calculation of the $^{12}$C ground-state density in a complete $(0 + 2)\hbar\omega$ model space has also been performed, but did not produce a significant difference in the DWBA results. The single-particle wave functions were calculated using a harmonic oscillator potential. The range parameter $b$ of this potential can be estimated
from the root-mean-square charge radius \((\langle r^2 \rangle)^{1/2}\) of the nucleus using [108]

\[
(\langle r^2 \rangle)^{1/2} = b \sqrt{\frac{5Z - 4}{2Z}}.
\]

For \(^{12}\text{C}\) a radius of 2.47 fm has been assumed [109] leading to \(b = 1.68\) fm. Four different effective interactions were used in the calculations. As an example of a free \(t\)-matrix interaction, the parameterization of Franey and Love [30], which has been extensively used in the analysis of proton-scattering experiments, has been chosen. \(G\)-matrix type interactions, which include medium effects in the description of the nucleon-nucleon scattering, have been provided by Von Geramb [32], by Nakayama and Love [34] and by the Melbourne group; see, e.g., Dortmans and Amos [18]. The most recent of these density-dependent interactions is the one developed by the Melbourne group. It has been shown to perform well on a large number of target nuclei [110] and over a wide range of incoming proton energies [100]. Figure 6.1 shows the results of microscopic calculations of the \(^{12}\text{C}(\vec{p}, p)\) elastic scattering cross section and analyzing power performed with the code DWBA98 [28] using the effective interactions mentioned above (for an example and some explanations regarding the input files required for DWBA98, see appendix B). The Melbourne interaction yields a very close description of the shape of the differential cross section while overestimating its magnitude by about 20% for scattering angles \(\theta_{cm} \leq 30^\circ\). In contrast, the other interactions overestimate the measured cross section by up to a factor of two, especially towards smaller angles. In case of the analyzing power the best agreement is obtained using Von Geramb’s interaction. Generally speaking, the density-dependent interactions result in a closer description of the data than the free \(t\)-matrix parameterization of Franey and Love.

Because the use of the Melbourne interaction results in the closest reproduction of the shape and magnitude of the differential cross section, which is important when trying to extract transition strengths from the data, it will be used for all DWBA calculations presented in the following sections.

### 6.2 Choice of the model space

Microscopic distorted-wave calculations of inelastic transitions or of higher multipole contributions to elastic scattering require one-body transition densities (see section 2.1.5) which can be obtained from the nuclear shell model. The shell-model calculations presented in this thesis were performed using the OXBASH code [19] which comes with a wide variety of parameterizations.
6.2. Choice of the model space

Figure 6.1: Microscopic calculations of the differential cross section and analyzing power for elastic scattering of 150 MeV protons from $^{12}$C compared to the experimental results. The ‘Orsay-74’ data-set is taken from reference [95].

of the residual NN-interaction for different configuration spaces. To find out which of these parameterizations result in a suitable description of the transitions in question, electromagnetic transition strengths calculated from the model wave functions are compared to results from $(e,e')$ or $\gamma$-decay measurements available in the literature [11, 111]. For examples of the input sessions needed to perform the calculations, see appendix C.

Because of the different configuration spaces required for the description of transitions to negative- or positive-parity states, these two cases will be investigated separately.
Table 6.1: Experimental $B(O\lambda)$ values for $^{11}\text{B}$ ($J^\pi(gs) = 3/2^-$) from the literature [11] compared to oxbash results. Units are $e^2 fm^{2\lambda}$ for $B(E\lambda)$ values and $\mu_N^2 fm^{2\lambda-2}$ for $B(M\lambda)$ values.

<table>
<thead>
<tr>
<th>$E_x$ (MeV)</th>
<th>$J^\pi$</th>
<th>$O\lambda$</th>
<th>$B(O\lambda) \uparrow_{exp}$</th>
<th>$B(O\lambda) \uparrow_{SM}$</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td></td>
<td></td>
<td>CKII</td>
<td>MK3W</td>
</tr>
<tr>
<td>2.125</td>
<td>1/2$^-$</td>
<td>M1</td>
<td>0.54 ± 0.04</td>
<td>0.94</td>
</tr>
<tr>
<td></td>
<td></td>
<td>E2</td>
<td>2.6 ± 0.4</td>
<td>0.88</td>
</tr>
<tr>
<td>4.445</td>
<td>5/2$^-$</td>
<td>M1</td>
<td>0.80 ± 0.03</td>
<td>0.78</td>
</tr>
<tr>
<td></td>
<td></td>
<td>E2</td>
<td>21.3 ± 5.30</td>
<td>8.91</td>
</tr>
<tr>
<td>5.020</td>
<td>3/2$^-$</td>
<td>M1</td>
<td>1.15 ± 0.04</td>
<td>1.38</td>
</tr>
<tr>
<td>6.743</td>
<td>7/2$^-$</td>
<td>E2</td>
<td>3.7 ± 0.9</td>
<td>1.08</td>
</tr>
<tr>
<td>8.560</td>
<td>3/2$^-$</td>
<td>M1</td>
<td>0.073 ± 0.007</td>
<td>0.020</td>
</tr>
<tr>
<td>8.920</td>
<td>5/2$^-$</td>
<td>M1</td>
<td>0.749 ± 0.004</td>
<td>0.49</td>
</tr>
<tr>
<td></td>
<td></td>
<td>E2</td>
<td>1.6 ± 1.2</td>
<td>0.0003</td>
</tr>
</tbody>
</table>

6.2.1 Normal-parity states

Transitions to normal-parity states, i.e. states with the same parity as the ground state, can be described in a simple $0\hbar\omega$ model space using the parameterizations of the residual interaction provided by Cohen and Kurath [41]. A more realistic description, especially of electric-quadrupole transitions, can, however, be obtained by extending the model space to a $(0 + 2)\hbar\omega$ configuration.

The different parameterizations are referred to by their labels used in the oxbash input which are CKII for the $0\hbar\omega$ and MK3W for the $(0 + 2)\hbar\omega$ model space. The CKII interaction is identical to the $(8 - 16)2\text{BME}$ interaction from reference [41] while the MK3W interaction is constructed from a number of different parameterizations of the residual interaction for each of the participating sub-shells (for references see the file ‘label.dat’ from the oxbash code package). The latter interaction has successfully been applied to the analysis of inelastic proton scattering from $^{12}\text{C}$ [38].

Table 6.1 gives an overview of measured $B(O\lambda)$ values [11] versus results of the two different shell-model calculations. The definitions of the applied transition operators can be found in section 2.1.6. Bare $g$-factors and free charges were used in the calculations. The radial matrix elements were computed using harmonic-oscillator wave functions with $b = 1.69$ fm (obtained from reference [109]). Generally, there is a poor agreement between measured and calculated transition strengths. For the $B(M1)$ values both
calculations yield similar results, while the $B(E2)$ strengths are, as expected, better reproduced in the larger model space. The one-body transition densities used for the DWBA calculations of transitions to normal-parity states will therefore be calculated using the MK3W interaction.

### 6.2.2 Non-normal-parity states

Transitions to non-normal-parity states require calculations in at least a $1\hbar\omega$ model space, allowing for the excitation of nucleons from the 1s to the 1p and from the 1p to the 2s and 1d shells. A suitable parameterization of the residual interaction in this model space is the SPSDMWK interaction, which, like the MK3W interaction, is constructed from a number of other parameterizations (for references see the file `label.dat` from the OXBASH code package). The MK3W interaction itself may also be used and allows, in a restricted way, also for $3\hbar\omega$-type excitations. The results of the two types of shell-model calculations are compared to experimental $B(O\lambda)$ values from reference [111] in table 6.2. As for the normal-parity states, the matrix elements were calculated with bare $g$-factors and free charges. Except for the two states with the highest excitation energies, the SPSDMWK interaction produces better results and will be used in the DWBA calculations of non-normal-parity transitions in section 6.4.2.

### 6.3 Elastic scattering and depolarization effects

The description of elastic scattering of protons from some target nucleus in a DWBA framework requires an optical potential to convert the incoming
and outgoing plane waves into distorted waves. In the conventional approach the distorted waves are generated from a phenomenological optical potential (see section 2.1.4) for which the parameters are adjusted such that the calculation reproduces the measured cross sections and analyzing powers. The ambiguities inherent to this approach can be avoided in a fully microscopic model where the optical potential is constructed by folding the effective nucleon-nucleon interaction with the ground-state density of the target nucleus.

Target nuclei with a non-zero ground-state spin, such as $^{11}\text{B}$ with $J^\pi = 3/2^-$, allow for higher-order multipole contributions to the elastic scattering process. These higher-order contributions contain spin-flip parts which cause the depolarization parameter $D_{nn'}$ to deviate from unity. Measurements of depolarization effects in elastic scattering have, among others, been performed by Von Przewoski et al. [112] and Nakano et al. [113]. The theoretical description of the observed effects has up to now been attempted either in a completely phenomenological way, by adding spin-spin terms to the optical-model potential given in equation 2.45 which are then adjusted to reproduce the observed depolarization effects or in a semi-microscopic approach, where the $\Delta J = 0$ contribution to the transition is described by a phenomenological optical potential and the higher-multipole contributions follow from a microscopic DWBA analysis using one-body transition densities from a shell-model calculation.

In the following two subsections results from a conventional semi-microscopic analysis will be compared to calculations performed in a fully microscopic DWBA framework.

### 6.3.1 The $^{11}\text{B}$ case

Figure 6.2 shows the results of semi-microscopic DWBA calculations of the cross section, analyzing power and depolarization parameter of the elastic $^{11}\text{B}(\bar{p}, \bar{p})$ reaction at 150 MeV. The optical-model parameters needed for the calculation have been obtained from reference [95]. Although this parameter set is originally intended for elastic proton scattering from $^{12}\text{C}$, it gives a much better fit to the measurements than the parameters which have been derived for the $^{11}\text{B}$ case by Geoffrion et al. [101]. The calculation gives a good fit to the cross section and, in the angular region $\theta_{cm} \leq 30^\circ$, also to the analyzing power. This is, of course, to be expected, as the optical-model parameters are fitted to give a good reproduction of these observables. A real test of the model is provided by the depolarization parameter whose angular distribution is, except for the last data point, reproduced by the
semi-microscopic calculation (the experimental data on the depolarization parameter $D_{nn'}$ have been corrected for the offset given in table 5.3).

The results of the fully microscopic analysis are presented in figure 6.3. There is a very good agreement between the calculated and measured differential cross section. In case of the analyzing power the features of the distribution are reproduced while the magnitudes of the maxima are not described correctly. The overestimation of the second maximum is also observed in the semi-microscopic calculation and might be due to deficiencies in
the predicted analyzing powers of the higher-order multipole contributions. The angular distribution of the depolarization parameter $D_{nn'}$ is correctly reproduced over the complete measured region.

Both types of calculations predict that the largest contribution to the observed depolarization effect comes from the octupole part ($\Delta J = 3$) of the transition. It is unfortunate that no data points are available at larger scattering angles to fix the position of the maximum of the distribution.

For the $^{11}$B case, it can be concluded that the fully microscopic approach gives a valid description of the observables of the elastic scattering process.
6.3.2 Comparison to elastic scattering from $^{13}\text{C}$ and $^{15}\text{N}$

To verify the conclusions of the last section two more cases known from literature have been studied, namely $^{13}\text{C}$ and $^{15}\text{N}$ both having a ground-state spin $J^\pi = 1/2^-$. In these cases only the dipole ($\Delta J = 1$) part of the transition can change the spin orientation of the projectile and the observable depolarization effects are therefore expected to be smaller. The elastic $^{13}\text{C}(p,p)$ reaction has been measured by Von Przewoski et al. at an incoming proton energy of 72 MeV [112]. Data obtained in this experiment are shown together with the results of a fully microscopic analysis in

![Graph](image)

**Figure 6.4:** Elastic scattering of 72 MeV protons from $^{13}\text{C}$. The experimental data and the semi-microscopic calculation of the depolarization parameter are taken from reference [112].
Figure 6.5: Elastic scattering of 65 MeV protons from $^{15}$N. The experimental data have been taken from references [113, 114] and the semi-microscopic calculation of the depolarization parameter is taken from reference [112].

The results of a semi-microscopic analysis of the depolarization effect, which has been published by Von Przewoski et al. together with the experimental data, is indicated by the dotted line in the lower right plot of the figure. Except for a small shift in the calculated angular distributions of about $5^\circ$, which might be due to the use of a different harmonic oscillator parameter, the two approaches give similar results.

The $^{15}$N case has been studied by Nakano et al. [113, 114] at an incoming proton energy of 65 MeV and angular distributions of the analyzing power and the depolarization parameter have been published. Semi-microscopic
calculations of the depolarization have been performed by Von Przewoski et al. [112] and are compared to the results of a fully microscopic analysis in figure 6.5. The observed effect is much stronger than for $^{13}$C and, in contrast to the $^{11}$B and $^{13}$C cases, there are large discrepancies between the experimental data, the result of the fully microscopic calculation and the semi-microscopic approach. The size of the observed effect is explained by Von Przewoski et al. with the special structure of the $^{15}$N nucleus, which has a transverse form factor that is about ten times larger than for $^{13}$C [115]. The large difference between the fully microscopic and the semi-microscopic calculations at angles $\theta_{\text{cm}} \geq 40^\circ$ may have several origins: 1) Inappropriate optical-model parameters in the semi-microscopic calculations. Von Przewoski et al. note that the optical-model parameters applied in the calculation have been extrapolated from nearby nuclei and have not been obtained from a fit to $^{15}$N data directly. 2) Differences in the nuclear wave functions applied in the two calculations. 3) Shortcomings in the effective nucleon-nucleon interaction at large $q$-transfers.

A thorough understanding of the observed differences between the theoretical descriptions and the measurements of the depolarization effect in $^{15}$N requires further theoretical and experimental investigations.

### 6.4 Inelastic scattering from $^{11}$B

Because of its non-zero ground state spin $J^e = 3/2^-$ and isospin $T = 1/2$, inelastic transitions to states within $^{11}$B always consist of at least two different multipole contributions $\Delta J$ which, in turn, are made up of a number of possible combinations of orbital momentum transfer $\Delta L$, spin transfer $\Delta S$ and isospin transfer $\Delta T$. The allowed $JLST$ combinations follow the usual selection rules for angular momentum coupling with the additional restriction that the orbital momentum transfer $\Delta L$ has to be even for normal-parity transitions and uneven for non-normal-parity transitions. Table 6.3 gives an overview of the allowed $JLST$ combinations for transitions to the first excited negative- and positive-parity states in $^{11}$B.

The calculated contribution of a particular $J$-transfer to a certain transition depends on the shell-model wave functions of the initial and final states and on the relevant parts of the effective nucleon-nucleon interaction. The total cross section is given by an incoherent sum over the cross sections of the different multipole contributions

$$
\frac{d\sigma}{d\Omega} = \sum_{\Delta J} \left( \frac{d\sigma}{d\Omega} \right)_{\Delta J}
$$

(6.2)

Table 6.3: Examples of allowed $JLST$ combinations for normal-and non-normal-parity transitions in $^{11}\text{B}$.

<table>
<thead>
<tr>
<th>type</th>
<th>$\Delta J$</th>
<th>$\Delta L$</th>
<th>$\Delta S$</th>
<th>$\Delta T$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$M1$</td>
<td>1</td>
<td>0</td>
<td>1</td>
<td>0</td>
</tr>
<tr>
<td></td>
<td>1</td>
<td>0</td>
<td>1</td>
<td>1</td>
</tr>
<tr>
<td></td>
<td>2</td>
<td>1</td>
<td>0</td>
<td>1</td>
</tr>
<tr>
<td></td>
<td>1</td>
<td>2</td>
<td>1</td>
<td>1</td>
</tr>
<tr>
<td>$E2$</td>
<td>2</td>
<td>2</td>
<td>0</td>
<td>0</td>
</tr>
<tr>
<td></td>
<td>2</td>
<td>2</td>
<td>0</td>
<td>1</td>
</tr>
<tr>
<td></td>
<td>2</td>
<td>2</td>
<td>1</td>
<td>0</td>
</tr>
<tr>
<td></td>
<td>2</td>
<td>2</td>
<td>1</td>
<td>1</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>type</th>
<th>$\Delta J$</th>
<th>$\Delta L$</th>
<th>$\Delta S$</th>
<th>$\Delta T$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$E1$</td>
<td>1</td>
<td>1</td>
<td>0</td>
<td>0</td>
</tr>
<tr>
<td></td>
<td>1</td>
<td>1</td>
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<tr>
<td></td>
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<td>1</td>
<td>1</td>
<td>0</td>
</tr>
<tr>
<td>$M2$</td>
<td>2</td>
<td>1</td>
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<td>0</td>
</tr>
<tr>
<td></td>
<td>2</td>
<td>1</td>
<td>1</td>
<td>1</td>
</tr>
</tbody>
</table>

and the analyzing power and spin-flip probability are given by

$$A_n = \left(\frac{d\sigma}{d\Omega}\right)^{-1} \sum_{\Delta J} (A_n)_{\Delta J} \left(\frac{d\sigma}{d\Omega}\right)_{\Delta J}, \quad (6.3)$$

$$S_{m'm} = \left(\frac{d\sigma}{d\Omega}\right)^{-1} \sum_{\Delta J} (S_{m'm})_{\Delta J} \left(\frac{d\sigma}{d\Omega}\right)_{\Delta J}. \quad (6.4)$$

Because of the poor reproduction of experimentally observed transition strengths by the shell-model calculations (see tables 6.1 and 6.2) it is not to be expected that the magnitudes of the different multipole contributions to the measured cross sections are predicted correctly by the DWBA calculations. Instead the multipole parts of a certain transition are renormalized to fit the observed cross section and the resulting scaling factors are used to correct the calculated transition strengths (see chapter 7).

Table 6.4 gives an overview of allowed multipole contributions to the observed transitions and the scaling factors extracted in the analysis. The experimental data and the fits obtained from microscopic DWBA calculations of negative- and positive-parity states will be discussed in the following subsections.

6.4.1 Negative-parity states

Figures 6.6 and 6.7 display the measured angular distributions of cross sections, analyzing powers and spin-flip probabilities for the transitions to
Figure 6.6: Transitions to the first four negative-parity states in $^{11}$B. Only $J$-transfers which give a significant contribution to the observed angular distributions are plotted.
Table 6.4: Scaling factors obtained by fitting calculated multipole contributions to measured differential cross sections. The systematic error of the quoted numbers is about 20% unless a range of values is given for a certain scaling factor. The abbreviation ‘n.c.’ means that the particular \( J \)-transfer gives no significant contribution to the observed cross sections.

<table>
<thead>
<tr>
<th>( E_x ) [MeV]</th>
<th>( J^\pi )</th>
<th>( \Delta J = 0 )</th>
<th>( \Delta J = 1 )</th>
<th>( \Delta J = 2 )</th>
<th>( \Delta J = 3 )</th>
<th>( \Delta J = 4 )</th>
<th>( \Delta J = 5 )</th>
</tr>
</thead>
<tbody>
<tr>
<td>2.125</td>
<td>( 1/2^- )</td>
<td>—</td>
<td>0.3</td>
<td>0.7</td>
<td>—</td>
<td>—</td>
<td>—</td>
</tr>
<tr>
<td>4.445</td>
<td>( 5/2^- )</td>
<td>—</td>
<td>0.5</td>
<td>0.75</td>
<td>n.c.</td>
<td>n.c.</td>
<td>—</td>
</tr>
<tr>
<td>5.020</td>
<td>( 3/2^- )</td>
<td>n.c.</td>
<td>0.47</td>
<td>0.3</td>
<td>0.6</td>
<td>—</td>
<td>—</td>
</tr>
<tr>
<td>6.743</td>
<td>( 7/2^- )</td>
<td>—</td>
<td>—</td>
<td>0.7</td>
<td>n.c.</td>
<td>n.c.</td>
<td>n.c.</td>
</tr>
<tr>
<td>8.920</td>
<td>( 5/2^- )</td>
<td>—</td>
<td>1.2</td>
<td>0.5</td>
<td>n.c.</td>
<td>n.c.</td>
<td>—</td>
</tr>
<tr>
<td>7.286</td>
<td>( 5/2^+ )</td>
<td>—</td>
<td>0.1 — 1</td>
<td>0.2 — 0.3</td>
<td>0.7</td>
<td>( \leq 0.3 )</td>
<td>—</td>
</tr>
<tr>
<td>7.978</td>
<td>( 3/2^+ )</td>
<td>( \leq 1 )</td>
<td>( \leq 0.5 )</td>
<td>n.c.</td>
<td>( \leq 0.8 )</td>
<td>—</td>
<td>—</td>
</tr>
<tr>
<td>9.185</td>
<td>( 7/2^+ )</td>
<td>—</td>
<td>—</td>
<td>0.5</td>
<td>0.5</td>
<td>n.c.</td>
<td>n.c.</td>
</tr>
<tr>
<td>9.274</td>
<td>( 5/2^+ )</td>
<td>—</td>
<td>( \leq 0.4 )</td>
<td>n.c.</td>
<td>3 — 4</td>
<td>n.c.</td>
<td>—</td>
</tr>
</tbody>
</table>

Figure 6.7: Same as figure 6.6 except for transitions to the negative-parity states at \( E_x = 8.56 \) and \( E_x = 8.92 \) MeV in \(^{11}\)B.
negative-parity states together with fits obtained from the calculated multipole contributions to each state. As for $D_{m'}$ in elastic scattering from $^{11}$B, the experimental values for the spin-flip probability $S_{m'}$ have been corrected for the systematic offset given in table 5.3.

The magnitudes of the multipole contributions to a certain state were fixed by a 'fit-by-eye' procedure to the observed cross section. There were two reasons not to use an automatic least $\chi^2$-fit routine for the fit. Firstly, the focus of the analysis is on the extraction of $M1$-strengths from the data and it is therefore necessary to get a good description of the data points at the most forward scattering angles, while an automatic fit weighs the data points by their statistical errors only. Secondly, the manual procedure allows to take additional information provided by analyzing powers and spin-flip probabilities into account in cases where there is an ambiguity in the fit of the cross section data. The uncertainty of the fitting procedure is about 10%. Together with the systematic uncertainty of the measured cross sections (see section 4.4) the error on the resulting scaling factors displayed in table 6.4 is about 20%.

Over the observed angular region the cross sections of most of the negative-parity states are governed exclusively by the $M1$ and $E2$ parts of the transitions. In case of the $J^\pi = 3/2^-$ state at $E_x = 5.02$ MeV it was possible to fix also the $\Delta J = 3$ part which gives a significant contribution to the spin-flip probability at larger angles.

A special case is the state at $E_x = 8.56$ MeV for which Ajzenberg-Selove gives a tentative assignment of $J^\pi = 3/2^-$ [11]. The steep rise of the cross section towards small angles indicates that the transition is predominantly of $E0$ type. Because the shell-model calculations presented in this thesis do not produce a theoretical candidate for this state, one-body-transition densities which have been calculated for a $0^+ \rightarrow 0^+$ transition in $^{12}$C have been used to produce angular distributions of a pure $E0$ transition. As can be seen in the upper part of figure 6.7 the shapes of the resulting curves are in good agreement both with the experimentally observed cross section and the analyzing power. The assumption that the transition is governed by the $\Delta J = 0$ contribution and the assignment of a $J^\pi = 3/2^-$ spin to the state is therefore confirmed.

### 6.4.2 Positive-parity states

The good energy resolution and high statistics of the data presented in this thesis made it for the first time possible to observe positive-parity states in $^{11}$B using inelastic proton scattering. Because the non-normal-parity
transitions are much weaker than excitations to normal-parity states, the
extraction of spin-flip probabilities was only feasible for the strongest of
the four observed positive-parity states, i.e. the state $J^e = 5/2^+$ at $E_x = 7.286$ MeV.

Figure 6.8 displays results of the data analysis together with fits from DWBA
calculations. Compared to the normal-parity transitions more of the possible
$J$-transfers to a certain state play a role in the observed angular region.
This makes it impossible to obtain an unambiguous determination of the
magnitudes of the multipole contributions for a transition. Most of the
scaling factors listed in table 6.4 are therefore given as a range or a limit on
the contribution of a certain multipolarity to a state and the fits displayed
in figure 6.8 have to be regarded as examples for possible descriptions of the
states.

For a clear determination of the multipole contributions to non-normal-
parity transitions in $^{11}$B experimental data over a wider angular range than
covered in this experiment are required.
Figure 6.8: Same as figure 6.6 except for transitions to positive-parity states in $^{11}$B.