Differential Analyzing Power in \( pp' \) Scattering on a \( ^{28} \text{Si} \) Nucleus in the Case of the Excitation of High-Spin Particle–Hole States

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Abstract—The experimental energy dependence of the differential analyzing power for \( 5^1_1, T = 0 \) and \( 6^1_1, T = 1 \) levels in the \( ^{28} \text{Si} \) nucleus is compared with the results of the calculations based on the DWBA-91 code. Information obtained for the nuclear structure from an analysis of inelastic scattering is discussed.

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1. INTRODUCTION

Investigation of the problem of the suppression of excitation strengths for anomalous-parity high-spin states is one of the objectives pursued in exploring the excitation of such states by protons. In this context, so-called stretched excited states are appealing. Within the concept of \( 1\hbar \omega \) transitions in the shell model, a single particle–hole configuration contributes to the excitation of such states. For the \( ^{28} \text{Si} \) nucleus, these are the \( 6^1_1, T = 0 \) excited state at 11.58 MeV and the \( 6^1_1, T = 1 \) excited state at 14.36 MeV [1, 2]. For these excitations, the particle–hole configuration has the form \( (f_{7/2} \cdot d_{5/2}^{-1}) \). However, the experiment reported in [3] revealed that the respective excitation strength is strongly fragmented. This fragmentation can be explained within the shell model under the assumption that different subshells are mixed in the ground state of the nucleus [4, 5]. In particular, configurations like \( (d_{5/2}^{10} \otimes s_{1/2}^{1})_{J=0} \) and \( (d_{5/2}^{10} \otimes d_{3/2}^{2})_{J=0} \) are admixed to the dominant shell-model configuration \( (d_{5/2}^{12})_{J=0} \) of the \( ^{28} \text{Si} \) nucleus.

In [5], Carr et al. showed that it is rather difficult to separate experimentally fragments of spin–parity \( 6^- \), because they can often be misidentified as \( 5^- \), \( T = 0 \) or \( 5^- \), \( T = 1 \) excitations. In the present study, we find, for the example of the excitation of the \( 6^1_1, T = 1 \) and \( 5^1_1, T = 0 \) (9.70 MeV) levels, that the energy dependence of the differential analyzing power \( \sigma A_y \) (the product of the differential cross section and the analyzing power \( A_y \)) [6] is strongly different in spin–flip and non-spin–flip processes in the \( ^{28} \text{Si} \) nucleus; therefore, this observable provides an additional tool for identifying individual fragments of high-spin states.

Previously, it was indicated in [7, 8] that, for the \( ^{28} \text{Si} \) nucleus, experimental data on \( \sigma A_y \) depend differently on energy in the case of the excitation of normal-parity high-spin states (\( 5^- \)) and in the case of the excitation of anomalous-parity (\( 6^- \)) high-spin states. This article presents the first-ever comparison of the energy dependence of experimental data on this quantity with the results of calculations performed within the distorted–wave method. Of the two highest spin magnetic states in this nucleus, the \( 6^1_1 \) states of isospin \( T = 0 \) and \( T = 1 \), only the second state is subjected to a comprehensive analysis because only in that case is available experimental information reliable in what is concerned with the admixture of neighboring levels [9]. The experimental data used in the present study were obtained in [10] at \( E_p = 100 \) and 180 MeV. The data at \( E_p = 65 \) and 500 MeV were obtained in [11] and [12], respectively.

2. METHOD USED AND RESULTS OF THE CALCULATIONS

The differential cross section and analyzing powers for inelastic proton scattering were calculated by using the DWBA–91 code [13], which is based on the distorted–wave Born approximation. The distorted waves in the elastic channel were calculated with the aid of the phenomenological optical potential presented in [10–12]. We used the parametrization from [14] for the effective nucleon–nucleon interaction potential. This potential depends on the

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Amplitudes $A_{i(1,j_1)j_2)}$ of particle–hole configurations in the wave function for the $5^-_1$, $T = 0$ level (the indices “1” and “2” label a particle and hole, respectively; given here are the doubled values of the total angular momentum of the particle and the hole)

<table>
<thead>
<tr>
<th>$(l_1j_1)(l_2j_2)$</th>
<th>$(j\bar{T})(d\bar{3})$</th>
<th>$(d\bar{3})(j\bar{T})$</th>
<th>$(j\bar{T})(d\bar{5})$</th>
<th>$(d\bar{5})(j\bar{T})$</th>
<th>$(j\bar{T})(d\bar{5})$</th>
<th>$(d\bar{5})(j\bar{T})$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$(f\bar{T})(d\bar{3})$</td>
<td>0.558</td>
<td>-0.050</td>
<td>0.175</td>
<td>0.041</td>
<td>-0.121</td>
<td>0.046</td>
</tr>
</tbody>
</table>

nuclear–matter density at the point of nucleon–nucleon interaction and involves the spin–dependent central, the spin–orbit, and the tensor component.

In the DWBA-91 code, the incident-proton and intranuclear-nucleon wave functions are fully symmetrized. In the present study, we also perform a comparison with the results of the calculations performed on the basis of the LEA code, which employs the distorted-wave impulse approximation. In that case, wave functions are symmetrized only for two nucleons involved in inelastic interaction, while exchange processes are considered in the zero-range approximation. Previously, a semimicroscopic analysis of the energy dependence of observables of inelastic proton–proton scattering for high-spin states was performed in [15, 16], but this was done only for the case of normal parity ($5^-_1$ and $4^-_1$) states. The respective calculations were performed on the basis of the LEA code. As to the excitation of anomalous-parity levels, tensor forces play an important role there [2], in which case the zero-range approximation leads to large errors in relation to approaches where exchange processes are taken accurately into account. For the case involving the excitation of stretched anomalous-parity states, it is therefore of interest to compare the two approaches in question to taking into account exchange processes.

Calculations within the distorted-wave method require knowing, in addition to the effective potential of nucleon–nucleon interaction, the wave functions for the excitations being considered. Since the spins of the $6^-_1$, $T = 1$ and $5^-_1$, $T = 0$ states are high, their wave functions admit the representation in the form of one or a few particle–hole configurations making a significant contribution to relevant transitions. As a criterion of the correctness of the excited-state wave functions used in our calculations, one can employ electromagnetic form factors obtained in inelastic electron–nucleus scattering.

In [15, 16], the form factors for inelastic proton scattering were calculated within the folding model [17]. In this approach, it is sufficient to know the transition charge density, which is extracted from an analysis of inelastic electron–nucleus scattering.

In the case of the excitation of normal–parity levels, the transition density is dominated by its mass component [2]. As for the excitation of anomalous–parity levels, the relevant inelastic-transition form factors depend on the spin and current components of the transition density. It does not seem possible to separate these components unambiguously on the basis of an analysis of electron–scattering data.

For stretched states, the transition current density is determined by the shell structure of the nucleus. If it is necessary to perform a full antisymmetrization within the approach based on the distorted-wave Born approximation, knowledge of the excited-state wave function rather than knowledge of the transition mass density alone is required even in the case of the excitation of normal–parity levels.

As was indicated above, the factor characterizing the suppression of the total excitation strength for the $6^-_1$, $T = 1$ level is determined by the shell structure of the ground state of the nucleus. Since a transition from the $1d_{5/2}$ to the $1f_{7/2}$ level is necessary for the excitation of the level being considered, the population of the $1d_{5/2}$ shell in the ground state of the nucleus determines directly the integrated suppression factor (without allowance for fragmentation). From calculations performed within the random-phase approximation generalized to nuclei featuring unfilled shells [19], it follows that the wave function calculated for the $5^-_1$, $T = 0$ level is sensitive to the wave function for the ground state of the $^{28}\text{Si}$ nucleus. All of the calculations in [19] differed only in the choice of ground-state wave function. The ground state of the nucleus was calculated either on the basis of the shell model with matrix elements chosen differently for the nucleon–nucleon interaction or by the Hartree–Fock method. For the $5^-_1$ level, a comparison with experimental data on inelastic electron–nucleus scattering revealed that the ground-state wave function calculated on the basis of the shell model with the Kuo matrix element [20], in which case the population of the $1d_{5/2}$ shell is minimal and is equal to 0.5593, is the best one. The amplitudes of the particle–hole configurations for the $5^-_1$ level constructed for this ground state are given in the table according to the calculations based on the random-phase approximation. They were determined by using the formula

$$A_{i(1,j_1)j_2)} = \langle J_f M_f | A(l_1j_1,l_2j_2; JM)|J_iM_i)\rangle,$$

$$A(l_1j_1,l_2j_2; JM) = (-1)^{j_2-m_2} \langle j_1m_1j_2-m_2|JM)a_{l_1j_1,m_1}a_{l_2j_2,m_2}}$$

where $(l_1j_1)$ and $(l_2j_2)$ are the quantum numbers (orbital angular and total angular momenta) of,