Core polarization effect on the discrete proton hole states of ^{205}Tl

Ramen Majumdar
Nuclear Physics Study Centre, 30/41 Attapara Lane, Calcutta 700050, India
(Received 31 March 1992)

The distribution of the discrete proton hole states of ^{205}Tl has been obtained within the hole-core vibrational scheme. The theoretical single-hole distribution pattern of these low-lying spin states has been discussed in the light of the recent experimental results based on the ^{208}Pb$(\vec{p},\alpha)$ reaction. Side by side, the shell-model hole states of the ^{205}Tl have been compared with the hole states of the ^{207}Tl to prove the rigid collectivity of the ^{206}Pb nucleus.

PACS number(s): 27.80.+w, 21.10.Pc, 21.10.Jx, 21.60.Cs

I. INTRODUCTION

In recent years various experimental data have been obtained to understand the shell-model aspect of the discrete low-lying neutron and proton hole states of ^{208}Pb [1–3] and ^{206}Pb [4]. The distributions of the $3s_{1/2}$ proton states of ^{207}Tl [1] and the $2f_{7/2}$ neutron state of ^{207}Pb [3] clearly indicate the nature of the deviation of the shell-model identity of these two discrete low-lying spin states. This invariably exhibits explicitly the strong interaction of the collective vibrational states of ^{208}Pb nuclei with single-particle proton or neutron states. The outcome of this interaction bears relevance to contemporary theoretical calculations on the shell-model aspect of discrete spin states in the vicinity of ^{208}Pb. We will discuss this theoretical aspect in Sec. IV. The present research deals with the effect of the core polarization on low-lying spin states of ^{205}Tl, and side by side, we will compare the results on ^{205}Tl with those of ^{207}Tl to prove the rigid collectivity of the ^{206}Pb core nucleus. Our theoretical work is based on recent experimental results on the fragmentation of the shell-model proton hole states of ^{205}Tl by the ^{208}Pb$(\vec{p},\alpha)$ reaction [5]. The proton hole states of ^{205}Tl have been mixed with the vibrational configurations of ^{206}Pb [6] to set up the relevant Hamiltonian matrices for the $\frac{1}{2}^+$, $\frac{3}{2}^+$, $\frac{5}{2}^+$, $\frac{7}{2}^+$, $\frac{9}{2}^+$, and $\frac{11}{2}^-$ states, and the matrices have been diagonalized to obtain the eigenvalues and eigenvectors [7]. The zero-order energies of the proton hole states have been obtained through the numerical solution of the second-order Schrödinger wave equation with existing potential parameters [8].

II. SHELL-MODEL ENERGIES OF THE PROTON HOLE STATES OF ^{205}Tl

To salvage the energies of the proton hole states of ^{206}Pb, we have solved the Schrödinger equation with a potential well consisting of the Woods-Saxon, spin-orbit, and Coulomb terms. The form of the potential well is

$$V=-\left[V_0\left\{1+\exp\left[\frac{r-r_0A^{1/3}}{a_0}\right]\right\}^{-1}+V_S\frac{1}{r}\frac{d}{dr}\left\{1+\exp\left[\frac{r-r_SA^{1/3}}{a_S}\right]\right\}^{-1}+\mathbf{L}\cdot\mathbf{S}\right]+V_C(r)\ , \tag{1}$$

where

$$V_C(r)=\begin{cases}\dfrac{1.44Z}{2R}\left[3-\dfrac{r^2}{R_C}\right] & \text{for } r<R_C\ ,\\ \dfrac{1.44Z}{4} & \text{for } r>R_C\ ,\end{cases}$$

and

$$R_C=r_CA^{1/3}\ .$$

III. HOLE-CORE COUPLING MODEL

The total Hamiltonian of the hole-core interaction scheme is [7]

$$H=H_p+H_{\rm vib}+H_{\rm int}\ . \tag{2}$$

Here $\langle H_p\rangle$ and $\langle H_{\rm vib}\rangle$ signify the energies (in MeV) of the shell-model proton states and the collective vibrational states of the ^{206}Pb, respectively. $H_{\rm int}$ is the interaction Hamiltonian. For the distortion of the collective core due to the core polarization effect, the main coupling term for $H_{\rm int}$ is linear in the vibrational amplitude $\alpha_{\lambda\mu}$. In the spherical nucleus, the coupling is the scalar product of the tensors $\alpha_{\lambda\mu}$ and $Y_{\lambda\mu}$ for the shape vibration and so $H_{\rm int}$ is

$$\begin{aligned}H&=(2\lambda+1)^{1/2}K(r)\left[\sum_\mu\alpha_{\lambda\mu}Y_\lambda(\theta,\phi)\right]_0\\&=(2\lambda+1)^{1/2}K(r)\left[\alpha_\lambda Y_\lambda(\theta,\phi)\right]_0\ .\end{aligned} \tag{3}$$

The zero-point amplitude for the individual $\lambda\mu$ mode of the collective vibrational state $(\alpha_\lambda)_0$ is

TABLE I. Potential parameters and shell-model energies of the proton hole states of ^{205}Tl from the core-polarization effect. $V_0=64.366$ MeV, $r_0=1.18$ fm, $a_0=0.64$ fm, $V_S=32.38$ MeV, $r_S=1.134$ fm, $a_S=0.785$ fm, and $r_C=1.184$ fm.

Hole states	$3s_{1/2}$	$2d_{3/2}$	$2d_{5/2}$	$1g_{7/2}$	$1g_{9/2}$
Energies	0.800	1.020	2.590	4.300	8.470
(MeV)	$2p_{1/2}$	$2p_{3/2}$	$1f_{5/2}$	$1f_{7/2}$	$1h_{11/2}$
	8.930	9.920	11.980	14.840	1.610

$$(\alpha_\lambda)_0\equiv[\langle n_\lambda=0|\alpha_\lambda^2|n_\lambda=0\rangle]^{1/2}$$
$$=\langle n_\lambda=1|\alpha_\lambda|n_\lambda=0\rangle=\frac{(4\pi G_\lambda)^{1/2}}{Z(\lambda+3)}\ ,$$
$$G_\lambda=\frac{B(E\lambda)}{B(E\lambda)_{\text{W.u.}}}\ . \tag{4}$$

The radial function $K(r)$ in H_{int} [Eq. (3)] is

$$K(r)=-r\frac{dV}{dr}\ . \tag{5}$$

The matrix elements of $K(r)$ are determined by numerical integration utilizing harmonic-oscillator wave functions for the proton shell-model orbitals. The Hamiltonian H has been diagonalized by using the wave function for the $J=j_1$ spin state as

$$\varphi_J=\sum_{\lambda j_2}a_{\lambda j_2}|n_\lambda;\lambda j_2,j_1\rangle\ . \tag{6}$$

Here j_2 and λ are the angular momenta for the proton hole state and vibrational state, respectively. Vectorially,

$$\boldsymbol{\lambda}+\mathbf{j}_2=\mathbf{j}_1\ . \tag{7}$$

The matrix elements of H_{int} in the $|\lambda j_2\rangle$ basis [Eq. (6)] is

$$\langle n_\lambda=1;\lambda j_2,j_1|H_{\text{int}}|n_\lambda=0,0j_1,j_1\rangle$$
$$=\langle j_2|K(r)|j_1\rangle\langle j_2\|Y_\lambda\|j_l\rangle(2j_1+1)^{-1/2}\langle\alpha_\lambda\rangle\ . \tag{8}$$

IV. RESULTS AND DISCUSSION

The proton hole states as exhibited in Table I have been combined with collective vibrational states, shown in Table II, to set up the Hamiltonian matrices of the $\frac{1}{2}^+$, $\frac{3}{2}^+$, $\frac{5}{2}^+$, $\frac{7}{2}^+$, $\frac{9}{2}^+$, and $\frac{11}{2}^-$ spin states of ^{205}Tl (Table III). The results of the diagonalization of the matrices are depicted in Tables IV–IX. First of all, we have calculated the energies of the shell-model hole states (Sec. III) with the potential parameters listed in Table I. The calculated energies have been optimized using a χ^2 minimization procedure for the diagonalization to reproduce the shell-model states of the $\frac{1}{2}^+$, $\frac{3}{2}^+$, $\frac{5}{2}^+$, and $\frac{11}{2}^-$ states of ^{205}Tl. The main contribution of the collective state for the fragmented spin state is also shown side by side in the same tables. The spreading of the discrete state indicates the strong interaction of the shell-model hole states with the collective excited states of the ^{206}Pb core nucleus. Comparison of the experimental results (Table III) with the theoretical results (Tables IV–IX) indicates that our calculated results support the experimental ones for the $3s_{1/2}$, $2d_{3/2}$, $2d_{5/2}$, and $1h_{11/2}$ hole states of ^{205}Tl. The experimental results listed in Table III show the nature and distribution of the hole states with regard to their shell-model identities. A quantitative comparison of the experimental results with the present theoretical ones is difficult to reproduce because of the uncertainties in the evaluation of the spectroscopic factors from the distorted-wave Born approximation differential cross sections $(d\sigma/dw)_{\text{DW}}$. Even the experimental spectroscopic factors differ appreciably if the calculations are carried out by finite-range distorted-wave Born approximations based on the double-folded α-particle potential instead of using the usual wave functions derived from general Woods-Saxon potentials [5]. The theoretical results further indicate that the hole strength might extend beyond the excitation energies quoted in the corresponding experimental results [5]. The physical reason for the spreading is due to the coincidence of the zero-order shell-model energies (Table I) for the above hole states with hole-coupled vibrational configurations. The fragments of the $1g_{9/2}$ shell-model hole state of ^{205}Tl must lie in a higher excitation energy region as its unperturbed energy centers around 8.47 MeV (Table I). In any event, the shell-model hole strengths of these two high-spin $1g_{7/2}$ and $1g_{9/2}$ hole states are convincingly damped out, as should have occurred, as both the vibrational states and shell-model hole states of ^{206}Pb span a broad excitation energy which forces the mixing of the $1g_{7/2}$ and

TABLE II. Vibrational states of ^{206}Pb. λ^π indicates the vibrational spin states. E is the energy in MeV, and $\langle\alpha_\lambda\rangle$ is the vibrational amplitude of the λ^π vibrational state.

λ^π	2_1^+	2_2^+	2_3^+	4_1^+	4_2^+	4_3^+	4_4^+
E	0.803	1.469	4.107	1.686	1.998	2.928	4.333
$\langle\alpha_\lambda\rangle$	0.030	0.007	0.021	0.014	0.008	0.010	0.018
λ^π	6^+	8^+	3^-	5_1^-	5_2^-	5_3^-	
E	4.357	4.580	2.648	2.782	3.558	3.772	
$\langle\alpha_\lambda\rangle$	0.015	0.008	0.041	0.008	0.007	0.013	

TABLE III. Experimental energies and spectroscopic factors of the fragmented states of ^{205}Tl. E is the energy in MeV, and S is the spectroscopic factors.

	E	S
$3s_{1/2}$ state	0.000	0.540
	1.218	0.070
	1.829	0.090
$2d_{3/2}$ state	0.207	1.000
	1.141	0.030
	1.342	0.090
	1.694	0.040
	2.222	0.040
	2.749	0.100
$2d_{5/2}$ state	0.619	1.450
	1.575	0.090
	1.860	0.170
	1.946	0.250
	2.838	0.200
	2.885	
$1g_{7/2}$ state	0.923	1.500
	2.302	0.170
	2.339	0.100
	2.885	
$1g_{9/2}$ state	1.430	1.400
	2.486	0.090
	2.630	0.300
	2.711	0.520
$1h_{11/2}$ state	1.483	0.540
	2.584	0.170

TABLE IV. $3s_{1/2}$ state of ^{205}Tl. E is the energy of the fragmented states in MeV, $a_{\lambda j_2}(\lambda, j_2)$ is the collective configuration of the states, and $a^2_{0j_2}$ is the shell-model hole strength of the fragments.

E	$a_{\lambda j_2}(\lambda, j_2)$	$a^2_{0j_2}$
0.008		0.826
1.979	$0.949\ (2_1^+, 2d_{3/2})$	0.068
2.489	$-0.996\ (2_2^+, 2d_{3/2})$	0.003
3.541	$-0.965\ (2_1^+, 2d_{5/2})$	0.042
4.403	$-0.994\ (5_1^-, 1h_{11/2})$	0.005
5.180	$-0.852\ (5_2^-, 1h_{11/2})$	0.005
5.420	$-0.975\ (5_3^-, 1h_{11/2})$	0.010
6.007	$-0.991\ (4_1^+, 1g_{7/2})$	0.005
6.741	$0.975\ (2_3^+, 2d_{5/2})$	0.010
8.658	$-0.997\ (4_4^+, 1g_{7/2})$	0.004
14.792	$-0.997\ (3^-, 1f_{5/2})$	0.005
17.563	$0.997(3^-, 1f_{7/2})$	0.005

TABLE V. $2d_{3/2}$ state of ^{205}Tl.

E	$a_{\lambda j_2}(\lambda, j_2)$	$a^2_{0j_2}$
0.523		0.803
1.709	$-0.841\ (2_1^+, 3s_{1/2})$	0.051
1.889	$-0.861\ (2_2^+, 2d_{3/2})$	0.078
2.268	$-0.995\ (2_2^+, 3s_{1/2})$	0.005
3.460	$-0.982\ (2_1^+, 2d_{5/2})$	0.026
6.919	$0.991\ (2_3^+, 3s_{1/2})$	0.004

TABLE VI. $2d_{5/2}$ state of ^{205}Tl.

E	$a_{\lambda j_2}(\lambda, j_2)$	$a^2_{0j_2}$
1.208	$0.683\ (2_1^+, 3s_{1/2})$	0.391
1.760	$0.823\ (2_1^+, 2d_{3/2})$	0.040
2.057	$0.508\ (2_1^+, 2d_{3/2})$	0.296
2.280	$0.962\ (2_2^+, 3s_{1/2})$	0.034
2.484	$-0.996\ (2_2^+, 2d_{3/2})$	0.004
3.460	$0.960\ (2_1^+, 2d_{5/2})$	0.026
4.558	$-0.611\ (4_2^+, 2d_{5/2})$	0.017
4.652	$-0.668\ (3^-, 1, h_{11/2})$	0.075
4.950	$0.909\ (2_3^+, 3s_{1/2})$	0.025
5.136	$-0.785\ (2_3^+, 2d_{3/2})$	0.007
5.177	$0.873\ (5_2^-, 1h_{11/2})$	0.013
5.402	$0.961\ (5_3^-, 1h_{11/2})$	0.012
5.528	$0.689\ (4_3^+, 2d_{5/2})$	0.009
5.986	$-0.996\ (4_1^+, 1g_{7/2})$	0.003
6.712	$0.989\ (2_3^+, 2d_{5/2})$	0.005
6.292	$0.699\ (4_4^+, 2d_{5/2})$	0.007

$1g_{9/2}$ states with several hole-vibrational mixed configurations of ^{205}Tl. Comparisons of theoretical results on the low-lying discrete states of ^{207}Tl with the experimental ones [2,9] as exhibited in Table X show explicit less quenching of the hole strengths from their main shell-model orbitals for the $3s_{1/2}$, $2d_{3/2}$, $2d_{5/2}$, and $1h_{11/2}$ states. For the $1g_{7/2}$ orbital, the proton knockout strength is explicitly far less than its sum-rule limit. This is true from experimental [2] and theoretical findings (Table X). We are not yet aware of the quantitative nature of the spreading of the $1g_{9/2}$ state of ^{207}Tl as regards its hole-strength distribution. In spite of this, strong damping of this $1g_{9/2}$ state for both ^{205}Tl and ^{207}Tl is ex-

TABLE VII. $1g_{7/2}$ state of ^{205}Tl.

E	$a_{\lambda j_2}(\lambda, j_2)$	$a^2_{0j_2}$
1.705	$-0.969\ (2_1^+, 2d_{3/2})$	0.050
2.446	$-0.867\ (4_1^+, 3s_{1/2})$	0.020
2.679	$0.975\ (4_1^+, 2d_{3/2})$	0.012
2.782	$-0.986\ (4_2^+, 3s_{1/2})$	0.004
3.002	$-0.991\ (4_2^+, 2d_{3/2})$	0.007
3.341	$0.936\ (2_1^+, 2d_{5/2})$	0.068
3.634	$0.737\ (4_3^+, 3s_{1/2})$	0.203
3.816	$-0.646\ (4_3^+, 3s_{1/2})$	0.194
3.976	$0.878\ (4_3^+, 2d_{3/2})$	0.070
4.054	$-0.987\ (2_2^+, 2d_{5/2})$	0.005
4.346	$0.613\ (3^-, 1h_{11/2})$	0.059
4.405	$-0.845\ (5_1^-, 1h_{11/2})$	0.047
4.488	$0.986\ (4_2^+, 2d_{5/2})$	0.007
5.287	$0.669\ (4_4^+, 2d_{3/2})$	0.033
5.464	$-0.597\ (4_3^+, 2d_{5/2})$	0.064
5.545	$-0.793\ (4_3^+, 2d_{5/2})$	0.067
6.007	$0.977\ (4_1^+, 1g_{7/2})$	0.019
5.771	$0.983\ (2_2^+, 1g_{7/2})$	0.009
6.397	$0.995\ (4_2^+, 1g_{7/2})$	0.004
6.987	$0.910\ (6^+, 2d_{5/2})$	0.019
7.229	$-0.995\ (4_3^+, 1g_{7/2})$	0.003
8.422	$0.991\ (2_3^+, 1g_{7/2})$	0.005
8.664	$0.871\ (6^+, 1g_{7/2})$	0.005
11.615	$-0.995\ (3^-, 2p_{1/2})$	0.006
14.654	$0.996\ (3^-, 1f_{5/2})$	0.003

TABLE VIII. $1g_{9/2}$ state of ^{205}Tl.

E	$a_{\lambda j_2}$	$a^2_{0j_2}$
3.325	0.986 $(2_1^+, 2d_{5/2})$	0.014
4.138	0.923 $(3^-, 1h_{11/2})$	0.029
5.104	−0.973 $(4_4^+, 3s_{1/2})$	0.006
5.325	0.706 $(2_1^+, 2d_{5/2})$	0.008
6.618	0.954 $(2_3^+ 2d_{5/2})$	0.037
6.982	0.896 $(4_4^+, 2d_{5/2})$	0.014
8.206	0.307 $(4_4^+, 2d_{5/2})$	0.487
8.417	−0.945 $(2_3^+, 1g_{7/2})$	0.037
8.716	0.765 $(6^+, 1g_{7/2})$	0.070
8.914	0.950 $(8^+, 1g_{7/2})$	0.042
9.479	−0.870 $(4_1^+, 1g_{7/2})$	0.159
9.939	0.992 $(2_2^+, 1g_{9/2})$	0.006
10.176	0.983 $(4_1^+, 1g_{9/2})$	0.017
10.468	0.996 $(4_2^+, 1g_{9/2})$	0.004
12.655	0.706 $(6^+, 2d_{5/2})$	0.010
12.716	−0.736 $(5_3^-, 2p_{1/2})$	0.008
12.859	−0.706 $(6^+, 1g_{9/2})$	0.018
17.521	0.994 $(3^-, 2f_{7/2})$	0.005

pected from our hole-core coupling model calculations. A comparison of the theoretical results on the discrete states of ^{205}Tl with that of ^{207}Tl leads to the conclusion of a strong collective vibrational aspect of the ^{206}Pb core nucleus.

There are various theoretical approaches to realize the distribution of the main hole states of both ^{207}Tl and ^{205}Tl. Within the framework of finite nuclear matter and the repulsive short-range potential among nucleons, the depletion of the shell-model hole strengths can be explained because of the scattering of the fermions from the low-lying states to the comparatively highly energetic shell-model orbitals. The theoretical calculations [10] on the Hartree-Fock occupational probabilities of the single-particle orbitals based on the random-phase approximation method and finite nuclear-matter calculations with the short-range and tensor components of the nucleon-nucleon interaction as well as the random-phase approximation calculation incorporated within the nuclear-matter distribution [11] explain the quenching of the shell-model strengths of only the dominantly excited low-lying spin states of ^{205}Tl. The energies and shell-model hole-strength distribution of the several low-lying

TABLE IX. $1h_{11/2}$ state of ^{205}Tl.

E	$a_{\lambda j_2}$ (λ, j_2)	$a^2_{0j_2}$
1.076		0.834
2.550	0.944 $(2_1^+, 1h_{11/2})$	0.085
3.077	−0.996 $(2_2^+, 1h_{11/2})$	0.003
3.310	0.989 $(4_1^+, 1h_{11/2})$	0.010
3.608	0.975 $(4_2^+, 1h_{11/2})$	0.004
4.576	−0.811 $(5_3^-, 3s_{1/2})$	0.003
5.338	−0.831 $(3^-, 2d_{5/2})$	0.016
5.377	0.853 $(5_1^-, 2d_{5/2})$	0.010
5.737	−0.978 $(2_3^+, 1h_{11/2})$	0.007
5.983	0.831 $(6^+, 1h_{11/2})$	0.008
11.149	−0.994 $(3^-, 1g_{9/2})$	0.004

TABLE X. Distribution of the proton hole states of ^{207}Tl. E is the energy in MeV and $a^2_{0j_2}$ is the spectroscopic factors. The value in parentheses indicates experimental estimates. For E (expt.), see Refs. [2,9].

State	E	$a^2_{0j_2}$
$3s_{1/2}$	0.000 (0.000)	0.940 (0.900)
$2d_{3/2}$	0.344 (0.350)	0.956 (0.950)
$2d_{5/2}$	1.502 (1.670)	0.860 (0.580)
$1g_{7/2}$	3.135 (3.470)	0.719 (0.450)
$1h_{11/2}$	0.912 (1.330)	0.734 (0.640)

discrete states of ^{207}Tl have been calculated [12] by the coupling of the vibrational states of ^{208}Pb with the hole orbitals. But the calculation is restricted to a very few low-lying states, and a significant distribution pattern has not been obtained. The quasiparticle phonon coupling model [13] and self-consistent mixing of the Hartree-Fock states with the phonon vibrational states arising from the random-phase approximation using the Skyrme III effective force [14] have been developed to explain the spreading of the spectroscopic factors of the hole orbitals of ^{207}Tl. But the broad based distribution of the fragmented states of both the two deep hole $1g_{7/2}$ and $1g_{9/2}$ states cannot be explained by these two theoretical approaches. This has been observed when a comparison of the theoretical results with the recent extensive experimental works on the proton hole orbitals of ^{207}Tl has been performed [2]. Very recently, the energies and single-hole strengths of a few discrete levels of ^{205}Tl have been obtained by the multistep shell-model method [15]. There, the matrices of the hole states have been constructed by the coupling of the proton hole states of ^{207}Tl with the two neutron hole states of ^{206}Pb connected by an effective two-body interaction of the Kuo-Brown type. The results obtained by this approach explain excellently the energies and single-hole components of the low-lying spin states. Though the calculation has been extended up to very-high-spin yrast states, still the spreading of the several hole orbitals, particularly the $1g_{7/2}$ and $1g_{9/2}$ ones, cannot be perceived by this theoretical approach. The other theoretical calculations based on the shell-model aspect [16,17] and semimicroscopic models [18,19] based on the coupling of the hole states with the phonon vibrational states (within a limited basis) of the core nucleus are mainly concerned with the main low-lying hole orbitals of ^{205}Tl.

Our theoretical approach, where a shell-model hole state can mix with the legion number of the collective vibrational states, is able to explain the broad attenuation of the hole strengths of all the discrete shell-model orbitals of ^{205}Tl. Also, none of the hole strengths of the low-lying states center around a particular excitation energy. In our earlier works on the deep hole states of ^{207}Tl [20], we have explained successfully the nature of the fragmentation by incorporating the high-lying vibrational states from the giant resonances of ^{208}Pb within the framework of the hole-core coupling scheme. But there we did not salvage the shell-model energies more realistically as we have done in the present calculations on both ^{205}Tl and

^{207}Tl. In the case of ^{205}Tl, the vibrational states are strongly collective, as has been seen in the present work, and we have allowed an interaction of the hole states with several vibrational configurations of ^{206}Pb. We are interested in the extraction of the spectroscopic factors and fragmented doorway collective states that strip off a sizable percentage of the shell-model hole strength from the main hole state. The Pauli blocking effect might interfere with the microscopic shell-model calculation of ^{205}Tl where the collective states split into the two shell-model hole states. But the effect does not even dominate much in that case inasmuch as the discrete hole states are well separated as has been realized from the present Table I.

In conclusion, we have attempted to explain the fragmentation of the discrete hole orbitals of ^{205}Tl by the hole-core coupling scheme [7]. The difficulty in explaining this physical aspect of the shell-model states of ^{208}Pb was faced by several authors a few years ago [2].

[1] E. N. M. Quint, J. F. J. van den Brand, J. W. A. den Herder, E. Jans, P. H. M. Keizer, L. Lapikas, G. van der Steenhoven, P. K. A. de Witt Huberts, S. Klein, P. Grabmayer, G. J. Wagner, H. Nann, B. Brois, and D. Gontte, Phys. Rev. Lett. **57**, 186 (1986).
[2] H. Langevin-Joliot, E. Gerlic, J. Guillot, and J. Van de Wiele, J. Phys. G **10**, 1435 (1984).
[3] S. Y. Vandirwerf, M. V. Harakeh, and E. N. M. Quint, Phys. Lett. B **216**, 15 (1989).
[4] M. C. Radhakrishna, N. G. Puttaswamy, H. Nann, J. D. Brown, W. W. Jacobs, W. P. Jones, D. W. Miller, P. P. Singh, and E. J. Stephenson, Phys. Rev. C **37**, 66 (1988).
[5] E. Gadioli, P. Guazzoni, S. Mattioli, L. Zetta, G. Graw, R. Hertenberger, D. Hofer, H. Kader, P. Schiemenz, R. Nen, H. Abele, and G. Staudt, Phys. Rev. C **43**, 2572 (1991).
[6] J. E. Fink, G. M. Crawley, J. A. Nolen, and R. T. Kowzes, Nucl. Phys. **A407**, 163 (1983).
[7] R. Majumdar, Phys. Rev. C **42**, 631 (1990).
[8] E. Rost, Phys. Lett. **26B**, 184 (1968).
[9] P. Grabmayr, G. Mairle, V. Schmidt-Rohr, G. P. A. Berg, J. Meissburger, P. Von Rossen, and J. L. Tain, Nucl. Phys. **A469**, 285 (1987).
[10] V. R. Pandharipande, C. N. Papanocolas, and J. Wambach, Phys. Rev. Lett. **53**, 1133 (1984).
[11] M. Jaminon, C. Mahaux, and H. Ngo, Nucl. Phys. **A440**, 228 (1985).
[12] I. Hamamoto, Phys. Rep. **10**, 63 (1974).
[13] V. G. Soloviev, Ch Stoyanov, and A. I. Vdovin, Nucl. Phys. **A399**, 141 (1983).
[14] Pham Van Thieu and N. Van Giai (unpublished).
[15] L. Rydstrom, J. Blomqvist, R. J. Liotta, and C. Pomar, Nucl. Phys. **A512**, 217 (1990).
[16] N. A. F. M. Poppelier and P. W. M. Glaudemans, Z. Phys. A **329**, 275 (1988).
[17] B. Silvestre-Brac and J. P. Brouser, Phys. Rev. C **24**, 717 (1981).
[18] N. Azziz and A. Covello, Nucl. Phys. **A123**, 681 (1969).
[19] J. W. Smits and R. H. Siemssen, Nucl. Phys. **A261**, 385 (1976).
[20] R. Majumdar, J. Phys. G **12**, 641 (1986).