Core polarization effects in $(\vec{p},\pi^-)$ reactions on $^{12,14}C$

N. Nose and K. Kume
Department of Physics, Nara Women's University, Nara 630, Japan
(Received 26 December 1991)

Based on a two-nucleon model, we have studied the core polarization effects in near-threshold pion-production reactions $^{12,14}C(\vec{p},\pi^-)\ ^{13,15}O_{g.s.}$. The core polarization reduces the absolute value of the pion-production cross section over all angular directions, while the analyzing power distributions are quite insensitive to it.

PACS number(s): 25.40.Qa, 21.60.Cs, 24.70.+s, 27.20.+n

I. INTRODUCTION

Much experimental effort has been devoted to the study of the proton-induced pion-production reaction (p,π) on complex nuclei. In the near-threshold (p,π^-) reaction, highly selective excitations of the stretched two-particle–one-hole (2p-1h) high-spin states were observed in a number of nuclei [1–7]. For these stretched-state transitions, a strikingly similar pattern of the analyzing power distributions has been found experimentally. Recently, we have carried out the theoretical calculation based on a two-body pion-production model and have shown that the observed selective excitation of the 2p-1h high-spin states as well as their systematic pattern of the analyzing power can be understood as evidence of the enhancement of the two-body process $p+n\rightarrow\pi^- +pp(^1S_0)$ [8–10]. In the near-threshold (p,π^-) reaction, the final two-proton channels with small relative angular momentum are preferentially excited because of the large momentum transfer involved in the (p,π^-) reaction, and also the large angular momentum transfer to the nucleus is well accommodated by the stretched 2p-1h configuration. Although the two-body pion-production model succeeded in describing the stretched-state transitions, the reaction mechanism for non-stretched-state transitions is not thoroughly understood.

Previously, a clear isotope dependence of the cross section and the analyzing power was observed experimentally in the ground-state transitions $^{12,13,14}C(\vec{p},\pi^-)^{13,14,15}O_{g.s.}$ [11,12]. For the (p,π^-) reactions on $^{13,14}C$, the cross section and the analyzing power exhibit similar angular distributions, while they are quite different from those of ^{12}C. In particular, a sign difference between the analyzing powers of the reactions $^{12}C(p,\pi^-)^{13}O_{g.s.}$ and $^{13,14}C(p,\pi^-)^{14,15}O_{g.s.}$ is remarkable. Although these experimental results seem to support the dominance of the two-body pion-production processes $p+n\rightarrow\pi^-+pp$, semiclassical consideration leads to completely opposite signs of the analyzing powers [13]. In order to clarify this point, we have carried out the theoretical calculations for the reactions $^{12,14}C(\vec{p},\pi^-)^{13,15}O_{g.s.}$ based on a two-body pion-production model [14]. Our two-body calculation predicted the correct sign of the analyzing powers at forward direction, which is mainly due to the pion distortion effects. But the overall agreement with the experimental data was not so good as in the case of the stretched-state transitions, especially at backward direction where momentum transfer q is quite large. Unlike the case of stretched-state transitions, we expect the large momentum and the angular momentum mismatch in the ground-state transitions on carbon. If we assume that the reaction takes place around the nuclear surface due to strong absorption of the pion, the semiclassical value of the transferred angular momentum amounts to about $6\hbar$ which cannot be well matched to the small angular momenta of the initial and the final nuclear states. In fact, the strength of the ground-state transition is about an order of magnitude smaller than that of the transitions leading to high-spin states [12]. In the non-stretched-state transitions, therefore, we expect that the higher-order effects play an important role. Among several candidates of the higher-order corrections, we have studied the core polarization effects on the reactions $^{12,14}C(\vec{p},\pi^-)^{13,15}O_{g.s.}$. By core polarization, we mean the effects of the higher configurations outside the $0p$ shell-model space which were neglected in our previous work [14]. The core polarization effects for $0p$-shell nuclei have been studied for various nuclear reactions; the $M1$ form factors in $^{12,13}C$ [15,16], the beta decay in $A=12$ system [17], the polarized muon capture reaction $^{12}C(\mu^-,\nu_\mu)^{12}B$ [18], and the pion photoproduction reactions [19,20]. It has been shown that the effects of the core polarization are appreciable for these reactions especially at high-q region. For the (p,π^+) reaction on ^{12}C, Miller [21] calculated the contribution from the configurations outside the $0p$ shell in the final nuclear states. He showed that these effects are not negligible. A possible importance of the core polarization effects in the (p,π) reaction was also suggested [22,23], but no theoretical calculation has been done so far for the (p,π^-) reactions. Considering the large momentum transfer involved in the (p,π^-) reactions, it is necessary and worthwhile to evaluate the correction coming from core polarization.

In the present paper, we have calculated the first-order core polarization effects in the ground-state transitions $^{12,14}C(\vec{p},\pi^-)^{13,15}O_{g.s.}$. We used the effective interaction by Bertsch *et al.* (M3Y interaction) [24]. We also examined the interaction which was used in the analysis of $M1$ form factors in the electron scattering on ^{12}C [15]. Core polarization diagrams are evaluated perturbatively by taking into account the intermediate states up to $6\hbar\omega$ excitations.

In Sec. II, we briefly describe the pion-production model adopted in our calculation. In Sec. III, numerical results are given and are discussed.

II. CORE POLARIZATION

Let us consider the proton-induced pion-production reaction on complex nuclei

$$p+(A,Z)\rightarrow(A+1,Z+2)+\pi^- \ . \tag{1}$$

We restrict ourselves to the case of spinless target nucleus ($I=0$). The transition amplitude including the first-order core polarization effects can be expressed as the matrix elements of the two-body pion-production operator $\hat{O}(\mathbf{k})$:

$$\begin{aligned} T(I'I'_z\mathbf{k};\mathrm{p}\lambda) &= \langle \Phi_f|\hat{O}(\mathbf{k})|\Phi_i\psi_{\mathrm{p}\lambda}\rangle \\ &+ \langle \Phi_f|\hat{O}(\mathbf{k})\frac{Q}{E_i-H_0}(V-U)|\Phi_i\psi_{\mathrm{p}\lambda}\rangle \\ &+ \langle \Phi_f|(V-U)\frac{Q}{E_f-H_0}\hat{O}(\mathbf{k})|\Phi_i\psi_{\mathrm{p}\lambda}\rangle \ . \end{aligned} \tag{2}$$

Here, $\psi_{\mathrm{p}\lambda}$ is the wave function of the incident proton with momentum $\mathbf{p}$ and spin projection λ, and $\mathbf{k}$ the momentum of the emitted pion. Φ_i and Φ_f are the wave functions of the target and the residual nuclei. The spin and its z component of the final nucleus are denoted as I' and I'_z, respectively. All of the above state vectors are antisymmetrized. The operator Q projects on the states outside the $0p$ shell-model space. The first term in Eq. (2) can be written as [8]

$$\langle \Phi_f|\sum \hat{O}_{ij}(\mathbf{k})|\Phi_i\psi_{\mathrm{p}\lambda}\rangle = \sum C_{I_f}(j_a,j_b,j_c)(j_a I_f m_a M_f|I'I'_z)\frac{(-)^{j_a-m_a}}{\sqrt{[j_a]}}\langle (\psi_{j_b}\otimes\psi_{j_c})^{I_f M_f}|\hat{O}(\mathbf{k})|\psi_{j_a,-m_a}\psi_{\mathrm{p}\lambda}\rangle \ , \tag{3}$$

with

$$[j_a]=2j_a+1 \ , \tag{4}$$

where $C_{I_f}(j_a,j_b,j_c)$ are the shell-model amplitudes leading to the 2p-1h state $[(j_b\otimes j_c)^{I_f}j_a^{-1}]^{I'}$ with respect to the target nucleus. For the two-body pion-production process, we considered the π- and ρ-meson exchange diagrams as shown in Fig. 1. For the π-exchange diagrams, the s- and p-wave rescattering contributions are taken into account, both of which are important in the near-threshold region. The s-wave term is calculated by adopting the effective Hamiltonian by Koltun and Reitan [25] with off-shell extrapolation of the coupling strengths by Maxwell *et al.* [26]. The static Δ-isobar model is used to calculate the p-wave term. The detailed expression of the operator $\hat{O}(\mathbf{k})$ and the input parameters are found in Ref. [8].

The second and the third terms in Eq. (2) correspond to the first-order core polarization. The corrections to the diagram 1(a) in Fig. 1 are shown in Fig. 2. The dia-

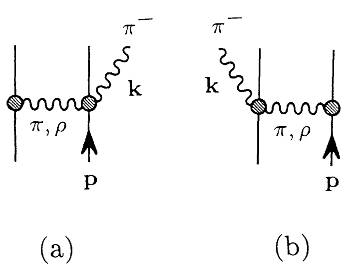

FIG. 1. Two-body pion-production processes with π and ρ exchange.

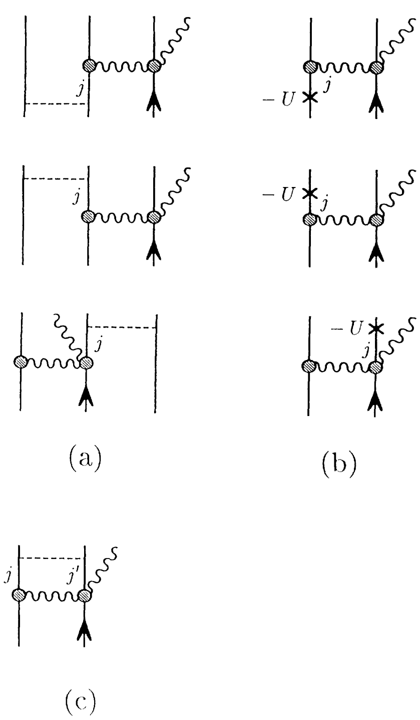

FIG. 2. First-order core polarization diagrams. Here, we only show the core polarization correction to the diagram 1(a) in Fig. 1. In practical calculations, corrections to the diagram 1(b) in Fig. 1 are also included. The diagrams 2(a) and 2(c) represent the first-order core polarization with 1p-1h and 2p-2h intermediate states, respectively. The single-particle orbits j and j' are those outside the $0p$ shell-model space. The diagrams 2(b) represent the one-body terms.

grams 2(a) come from the 1p-1h intermediate states and the diagram 2(c) from the 2p-2h intermediate states. We have subtracted the one-body terms 2(b) assuming the harmonic oscillator one-body potential U with oscillator parameter $b=1.64$ fm. In the case of the nucleus with a single nucleon outside the closed core, these one-body terms and the Hartree-Fock bubble terms cancel with each other. In practical calculations, the intermediate states are constructed as $[\psi_j \otimes \Phi_{A-1}]$ (1p-1h states) and $[(\psi_j \otimes \psi_{j'})^J \otimes \Phi_{A-2}]$ (2p-2h states), where Φ_{A-1} and Φ_{A-2} are the $0p$ shell-model wave function with $A-1$ and $A-2$ nucleons, respectively. Hereafter, j and j' represent the single-particle orbits outside the $0p$ shell-model space.

As an example, we show the nuclear matrix elements for the transition from the configuration

$$|\{(j_1^{n_1})^{J_1}(j_2^{n_2})^{J_2}\}^{I=I_z=0}\psi_{p\lambda}\rangle$$

to

$$|\{(j_1^{n_1})^{J_1''}(j_2^{n_2+1})^{J_2''}\}^{I'I_z'}\rangle$$

through the intermediate state

$$|[\{(j_1^{n_1-1})^{J_1'}(j_2^{n_2})^{J_2'}\}^{J_a}(jj')^{J_p}]^{I'I_z'}\rangle ,$$

corresponding to the third term in Eq. (2). Here, j_1 and j_2 represent the $0p_{1/2}$ and $0p_{3/2}$ orbits, respectively. For simplicity, we have suppressed the isospin indices. The relevant matrix elements can be calculated easily according to the standard shell-model technique as

$$\begin{aligned}
&\langle[\{(j_1^{n_1-1})^{J_1'}(j_2^{n_2})^{J_2'}\}^{J_a}(jj')^{J_p}]^{I'I_z'}|\sum\hat{O}(\mathbf{k})|[(j_1^{n_1})^{J_1}(j_2^{n_2})^{J_2}]^{I=I_z=0}\psi_{p\lambda}\rangle \\
&\quad=(-)^{n_2}(-)^{J_1-J_1'+j_1}\sqrt{2n_1}[j_1^{n_1-1}(J_1')j_1J_1|\}j_1^{n_1}J_1]\delta_{J_1J_2}\delta_{J_2J_2'}\delta_{J_aj_1} \\
&\quad\times\sum_{M_p,m_1}(j_1J_pm_1M_p|I'I_z')\frac{(-)^{j_1-m_1}}{\sqrt{[j_1]}}\langle(\psi_j\otimes\psi_{j'})^{J_pM_p}|\hat{O}(\mathbf{k})|\psi_{j_1,-m_1}\psi_{p\lambda}\rangle ,
\end{aligned} \tag{5}$$

and

$$\begin{aligned}
&\langle\{(j_1^{n_1})^{J_1''}(j_2^{n_2+1})^{J_2''}\}^{I'}|\sum V_{ij}|[\{(j_1^{n_1-1})^{J_1'}(j_2^{n_2})^{J_2'}\}^{J_a}(jj')^{J_p}]^{I'}\rangle \\
&\quad=(-)^{n_2}\sqrt{n_1(n_2+1)}[j_1^{n_1-1}(J_1')j_1J_1''|\}j_1^{n_1}J_1''][j_2^{n_2}(J_2')j_2J_2''|\}j_2^{n_2+1}J_2'']\sqrt{[J_1''][J_2''][J_a][J_p]} \\
&\quad\times\begin{Bmatrix} J_1' & J_2' & J_a \\ j_1 & j_2 & J_p \\ J_1'' & J_2'' & I' \end{Bmatrix}\langle(j_1j_2)^{J_p}|V|(jj')^{J_p}\rangle .
\end{aligned} \tag{6}$$

The transition amplitude $T(I'I_z'\mathbf{k};p\lambda)$ can be written as the sum of the two-body matrix elements of the 1p-1h type

$$\langle(\psi_j\otimes\psi_{j_2})^{I_f}|\hat{O}(\mathbf{k})|\psi_{j_1}\psi_{p\lambda}\rangle ,$$

$$\langle(\psi_{j_1}\otimes\psi_{j_2})^{I_f}|\hat{O}(\mathbf{k})|\psi_j\psi_{p\lambda}\rangle$$

and the 2p-2h type

$$\langle(\psi_j\otimes\psi_{j'})^{I_f}|\hat{O}(\mathbf{k})|\psi_{j_1}\psi_{p\lambda}\rangle .$$

Although the higher orbits j and j' are involved in these matrix elements, it should be noted that the angular momenta of the final two protons and the neutron hole are restricted to couple to the spin I' of the final nucleus. Thus the angular momentum mismatch still remains even if we consider the core polarization.

III. RESULTS AND DISCUSSION

According to the formulas described in Sec. II, the first-order core polarization effects on the reactions $^{12,14}C(\vec{p},\pi^-)^{13,15}O_{g.s.}$ are calculated. We adopted the $0p$ shell-model wave function by Hauge and Maripuu [27] in which the two-body interaction is constructed perturbatively from the Sussex matrix elements [28]. We have used the effective interaction by Bertsch *et al.* [24] (M3Y interaction) which is constructed to fit to the Sussex interaction. To see the dependence on the choice of effective interaction, we also examined the phenomenological one which was successfully applied to calculate the $M1$ form factor in electron scattering for ^{12}C [15]. This interaction consists of the central and tensor parts. The central potential is the Gaussian type with the Rosenfeld-type exchange mixture. The force range is assumed to be $r_c=1.6$ fm and the potential depth in the triplet-even state is taken to be $V_c=-60$ MeV. The tensor force is taken from the nucleon-nucleon interaction by Hamada and Johnston [29] with a radial cutoff at 0.7 fm. We used the value $\hbar\omega=15.4$ MeV to calculate the energy denominator. The distorted waves of the pion are generated from the pion-nucleus optical potential by the Michigan group [30–32] (MSU potential) and we used the proton-nucleus optical potential determined by In-

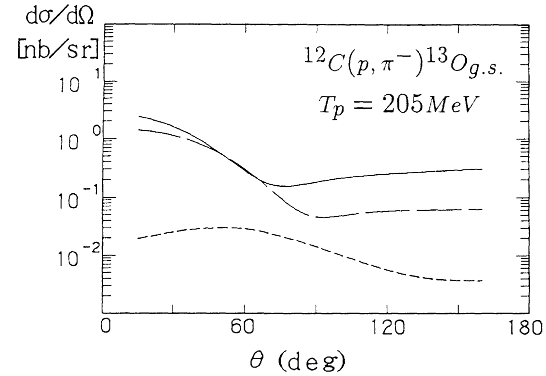

FIG. 3. Core polarization effects with intermediate states with $2\hbar\omega$ (solid line), $4\hbar\omega$ (long dashed line), and $6\hbar\omega$ (dashed line) excitations.

gemarsson *et al.* [33]. Since we are concerned with the two-body pion-production processes, a large number of two-body matrix elements contribute for the calculation of the core polarization. In the case of ^{12}C, the number of the relevant two-body matrix elements is about 350 if we take the intermediate states up to $6\hbar\omega$ excitations, and it is hard to include all of these matrix elements. We have, therefore, neglected the two-body matrix elements with the amplitude less than 5.0×10^{-3}. To ensure the validity of this approximation, we have calculated the contributions from the neglected matrix elements within $4\hbar\omega$ intermediate states. These two-body matrix elements scarcely affect the cross section and the analyzing power and can be safely neglected. In order to see the convergence of the calculation, we show the separate contribution up to $6\hbar\omega$ intermediate states for the reaction $^{12}\mathrm{C}(p,\pi^-)^{13}\mathrm{O}_{\mathrm{g.s.}}$ in Fig. 3. Here, we used the M3Y effective interaction. As is seen, the intermediate states with $2\hbar\omega$ and $4\hbar\omega$ components give important contributions, while the $6\hbar\omega$ component is at least about an order of magnitude smaller. We, therefore, take into account the intermediate states up to $6\hbar\omega$ excitations. Figure 4 shows the core polarization effects on the near-threshold reactions $^{12,14}\mathrm{C}(\vec{p},\pi^-)^{13,15}\mathrm{O}_{\mathrm{g.s.}}$ calculated by using the M3Y effective interaction. As is seen, the contributions from the core polarization are about an order of magnitude smaller than those from the $0p$ shell around the forward direction. For the ^{14}C, the cross section calculated by the $0p$ shell model decreases rapidly, and then the core polarization correction becomes comparable at backward direction. We can see that the core polarization reduces the absolute values of the reaction cross section over all angles but the angular distributions are not largely affected. The analyzing powers are also quite insensitive to the core polarization. These are the same for both the ^{12}C and the ^{14}C. For the analyzing power distributions, the $0p$-shell calculation predicts the sign change from positive (^{12}C) to negative (^{14}C) values at forward direction. This isotope dependence of the analyzing power comes from the pion distortion effects [14]. The core polarization scarcely affects these results. The two-body matrix elements coming from core polarization involve the single-particle orbits with large angular momenta j or large principal quantum numbers n. The contributions from these two-body matrix elements are, therefore, fairly large due to the momentum matching. But the amplitudes $C_{I_f}(j_a,j_b,j_c)$ corresponding to these matrix elements are quite small and, as a result, the core polariza-

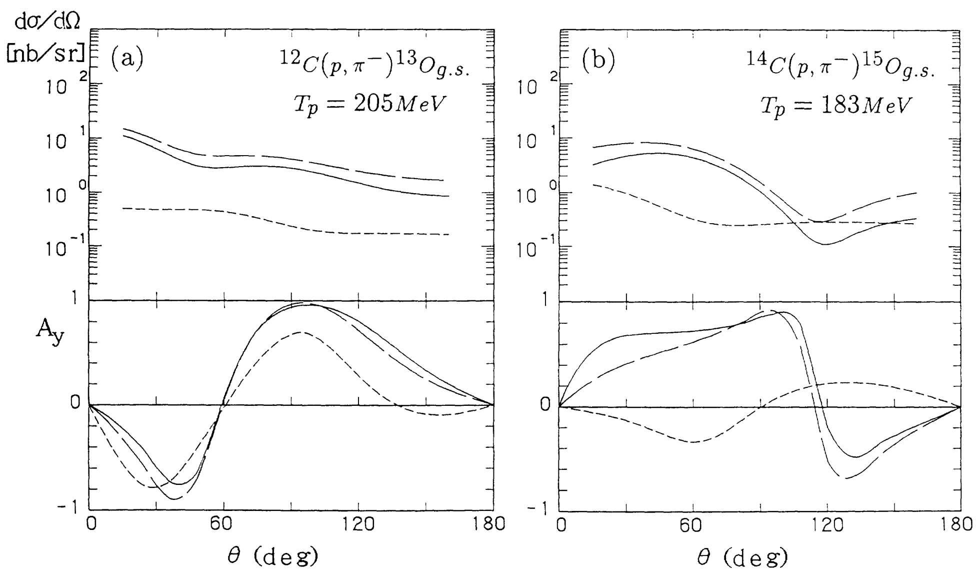

FIG. 4. Effects of core polarization on the reactions (a) $^{12}\mathrm{C}(\vec{p},\pi^-)^{13}\mathrm{O}_{\mathrm{g.s.}}$ at $T_p=205$ MeV and (b) $^{14}\mathrm{C}(\vec{p},\pi^-)^{15}\mathrm{O}_{\mathrm{g.s.}}$ at $T_p=183$ MeV. The long dashed lines correspond to the results with $0p$ shell-model wave function. The dashed lines are the results of the first-order core polarization only, and the solid lines are the sum of these contributions. Here, we used the M3Y effective interaction.

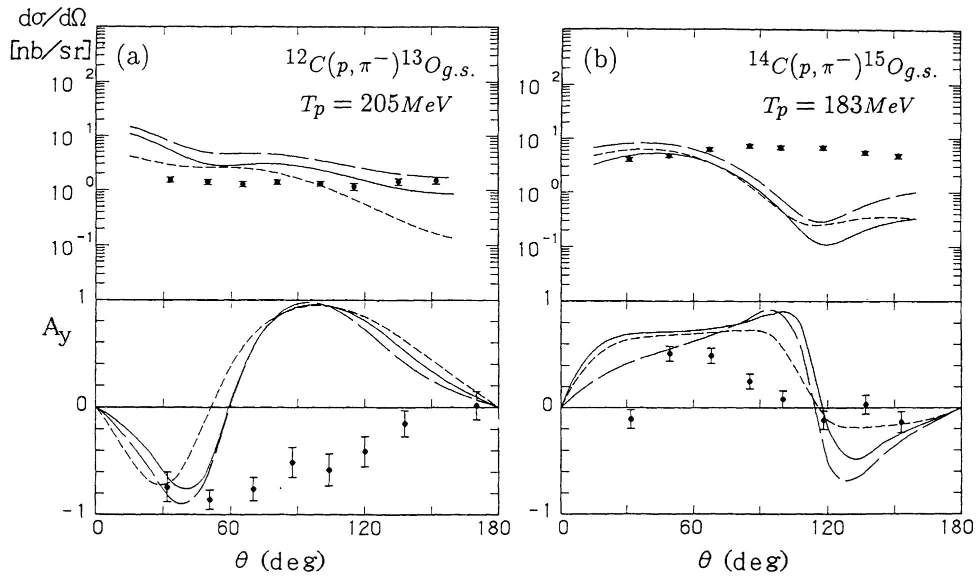

FIG. 5. Effects of core polarization on the reactions (a) $^{12}C(\vec{p},\pi^-)^{13}O_{g.s.}$ at $T_p=205$ MeV and (b) $^{14}C(\vec{p},\pi^-)^{15}O_{g.s.}$ at $T_p=183$ MeV. The long dashed lines correspond to the results with $0p$ shell-model wave function. The solid lines represent the results including the first-order core polarization with M3Y effective interaction and the dashed lines are those with the effective interaction used in Ref. [15] (see the text). The experimental data are taken from Ref. [11].

tion gives only moderate effects for the cross sections. Even if we consider the core polarization, the relevant 2p-1h matrix elements are restricted to couple to the spin I' of the final nucleus. Thus, the angular momentum mismatch still remains in the two-body distorted wave Born approximation amplitude. In order to see the dependence on the choice of the effective interaction, we examined the above-mentioned effective interaction used in Ref. [15]. The results are shown and are compared with the experimental data in Fig. 5. Because the two-body matrix elements calculated from the effective interaction in Ref. [15] are slightly larger than those calculated from M3Y interaction, the reduction of the cross section is somewhat larger. The analyzing power distributions are still quite insensitive to the core polarization.

We can see that the core polarization effects are non-negligible and the reaction cross sections are reduced over all angular directions. On the other hand, the analyzing power distributions are quite insensitive to the core polarization. Consequently, the discrepancies between theoretical values and the experimental data remain even if we take into account the core polarization effects.

IV. CONCLUSIONS

We have calculated the first-order core polarization effects on the near-threshold pion-production reactions $^{12,14}C(\vec{p},\pi^-)^{13,15}O_{g.s.}$. It is shown that the core polarization reduces the absolute values of the reaction cross section over all angular directions for both the ^{12}C and the ^{14}C. On the other hand, the analyzing power distributions are quite insensitive to the core polarization of the relevant nuclei. Thus the discrepancies between theory and experiments still remain for the ground-state transitions $^{12,14}C(\vec{p},\pi^-)^{13,15}O_{g.s.}$. Several refinements of the calculation should be necessary. The two-body operators $\hat{O}(\mathbf{k})$ used here might be oversimplified; we have neglected the distortion effects of the exchanged pions and we also treated the intermediate Δ isobar in a static approximation. There are several higher-order effects to be studied. The multistep processes which excite the high-spin intermediate states might play an important role. Since the reaction strengths leading to high-spin states are about an order of magnitude larger than that of the ground-state transition, it is probable that the two-step process with the intermediate high-spin states gives a non-negligible contribution. It is necessary to estimate these higher-order corrections before drawing definite conclusions.

The observed clear isotope dependence of the reaction cross section and the analyzing power was considered to be experimental evidence of the dominant two-body processes $p+n\rightarrow\pi^-+pp$ in the reactions $^{12,14}C(\vec{p},\pi^-)^{13,15}O_{g.s.}$. But it is not easy to understand them microscopically. More theoretical work should be necessary for a better understanding of the reaction mechanism.

We would like to thank T. Sato for allowing us to use the computer code for calculating the two-body matrix elements of the effective interaction.

[1] S. E. Vigdor, T. G. Throwe, M. C. Green, W. W. Jacobs, R. D. Bent, J. J. Kehayias, W. K. Pitts, and T. E. Ward, Phys. Rev. Lett. **49**, 1314 (1982).
[2] S. E. Vigdor, T. G. Throwe, M. C. Green, W. W. Jacobs, R. D. Bent, J. J. Kehayias, W. K. Pitts, and T. E. Ward, Nucl. Phys. **A396**, 61c (1983).
[3] M. C. Green, J. Brown, W. W. Jacobs, E. Korkmas, T. G. Throwe, S. E. Vigdor, T. E. Ward, P. L. Jolivette, and B. A. Brown, Phys. Rev. Lett. **53**, 1893 (1984).
[4] Z-J. Cao, R. D. Bent, H. Nann, and T. E. Ward, Phys. Rev. C **35**, 625 (1987).
[5] J. Kehayias, R. D. Bent, M. C. Green, M. Hugi, H. Nann, and T. E. Ward, Phys. Rev. C **33**, 1388 (1986).
[6] Z.-J. Cao, T. E. Ward, R. D. Bent, J. D. Brown, E. Korkmaz, H. Nann, and T. G. Throwe, Phys. Rev. C **35**, 825 (1987).
[7] T. G. Throwe, S. E. Vigdor, W. W. Jacobs, M. C. Green, C. W. Glover, T. E. Ward, and B. P. Hichwa, Phys. Rev. C **35**, 1083 (1987).
[8] K. Kume, Nucl. Phys. **A504**, 712 (1989).
[9] K. Kume, Nucl. Phys. **A511**, 701 (1990).
[10] K. Kume, Phys. Lett. B **271**, 17 (1991).
[11] W. W. Jacobs, T. G. Throwe, S. E. Vigdor, M. C. Green, J. R. Hall, H. O. Meyer, W. K. Pitts, and M. Dillig, Phys. Rev. Lett. **49**, 855 (1982).
[12] E. Korkmaz, S. E. Vigdor, W. W. Jacobs, T. G. Throwe, L. C. Bland, M. C. Green, P. L. Jolivette, and J. D. Brown, Phys. Rev. C **40**, 813 (1989).
[13] S. E. Vigdor, W. W. Jacobs, T. G. Throwe, and M. C. Green, Phys. Rev. Lett. **54**, 1204 (1985).
[14] K. Kume and N. Nose, Nucl. Phys. **A528**, 723 (1991).
[15] H. Sagawa, T. Suzuki, H. Hyuga, and A. Arima, Nucl. Phys. **A322**, 361 (1979).
[16] T. Suzuki, H. Hyuga, A. Arima, and K. Yazaki, Phys. Lett. **106B**, 19 (1981).
[17] K. Koshigiri, H. Ohtsubo, and M. Morita, Prog. Theor. Phys. **66**, 358 (1981).
[18] M. Fukui, K. Koshigiri, T. Sato, H. Ohtsubo, and M. Morita, Prog. Theor. Phys. **78**, 343 (1987).
[19] T. Sato, K. Koshigiri, and H. Ohtsubo, Z. Phys. A **320**, 507 (1985).
[20] T. Kobayashi, K. Takeshita, A. Kagaya, H. Tsubota, and K. Shoda, J. Phys. Soc. Jpn. **58**, 1570 (1989).
[21] G. A. Miller, Nucl. Phys. **A224**, 269 (1974).
[22] B. D. Keister and L. S. Kisslinger, Nucl. Phys. **A412**, 301 (1984).
[23] H. W. Fearing, in *Progress in Particle and Nuclear Physics,* edited by D. Wilkinson (Pergamon, Oxford, 1981), Vol. 7, p. 113.
[24] G. Bertsch, J. Borysowicz, H. McManus, and W. G. Love, Nucl. Phys. **A284**, 399 (1977).
[25] D. S. Koltun and A. Reitan, Phys. Rev. **141**, 1413 (1966).
[26] O. V. Maxwell, W. Weise, and M. Brack, Nucl. Phys. **A348**, 388 (1980); O. V. Maxwell and W. Weise, *ibid.* **A348**, 429 (1980).
[27] S. P. Hauge and S. Maripuu, Phys. Rev. C **8**, 1609 (1973).
[28] J. P. Elliott, A. D. Jackson, H. A. Mavromatis, E. A. Sanderson, and B. Singh, Nucl. Phys. **A121**, 241 (1968).
[29] T. Hamada and I. D. Johnston, Nucl. Phys. **34**, 382 (1962).
[30] K. Stricker, H. McManus, and J. A. Carr, Phys. Rev. C **19**, 929 (1979).
[31] K. Stricker, J. A. Carr, and H. McManus, Phys. Rev. C **22**, 2043 (1980).
[32] J. A. Carr, H. McManus, and K. Stricker-Bauer, Phys. Rev. C **25**, 952 (1982).
[33] A. Ingemarsson, O. Jonsson, and A. Hallgren, Nucl. Phys. **A319**, 377 (1979).