0^+-0^+ transition in charged photopion reactions

C. Bennhold and L. Tiator
Institut für Kernphysik, Universität Mainz, 6500 Mainz, Germany

A. M. Bernstein
Physics Department and Laboratory for Nuclear Science, Massachusetts Institute of Technology, Cambridge, Massachusetts 02139

S. S. Kamalov
Laboratory of Theoretical Physics, Joint Institute for Nuclear Research, Dubna, U.S.S.R.
(Received 13 April 1990)

In the near-threshold region for the $(\gamma,\pi^{\pm})$ reaction the spin-flip (Kroll-Ruderman) term dominates most nuclear reactions. The smaller non-spin-flip term is of considerable interest because it is very sensitive to the reaction dynamics and in particular to the Δ channel, even in the near-threshold region. We propose to study the non-spin-flip term via the $0^+ \rightarrow 0^+$ transition in the $^{14}C(\gamma,\pi^-)^{14}N$ reaction; the cross section is shown to be insensitive to the details of nuclear structure at low momentum transfer $Q \leq 0.5\ fm^{-1}$.

In the photoproduction of charged pions from threshold through the Δ region there are two main multipoles: the magnetic dipole which comes from a quark spin flip in the $N \rightarrow \Delta$ transition and the electric dipole which comes mainly from the Born terms.[1] The electric dipole operator near threshold is $\sigma \cdot \varepsilon \alpha \tau_\alpha$, where σ and τ_α are the nucleon spin and isospin and ε is the photon polarization. This is the Kroll-Ruderman term which involves a nucleon spin flip. By contrast, the photoproduction of neutral pions proceeds primarily via the resonant magnetic dipole term which is non-spin-flip and which vanishes near threshold.

Experimental studies of the $(\gamma,\pi^{\pm})$ reaction[2,3] have focused on the dominant Kroll-Ruderman term. The study of the non-spin-flip term in charged pion photoproduction has been hampered because of its small magnitude near threshold. This is shown in Fig. 1 comparing the contributions from the spin-flip $(S=1)$ and non-spin-flip $(S=0)$ amplitudes separately. Up to $T_\pi = 50$ MeV the $S=1$ term is larger by two orders of magnitude, whereas at higher energies the non-spin-flip term increases due to the delta. Experiments to observe it have been performed in the "mirror" transitions $^{13}C(\gamma,\pi^-)^{13}N_{g.s.}$ and $^{15}N(\gamma,\pi^-)^{15}O_{g.s.}$. In these cases there are two contributions, the $M1$ spin-flip term and the $E0$ non-spin-flip term. These have different angular distributions due to their different nuclear form factors. In particular, one can enhance the relative $E0$ contribution by performing experiments near the minimum of the $M1$ form factor. The data for the $^{13}C(\gamma,\pi^-)^{13}N$ reaction have indicated a small $E0$ contribution[4] while the experiment on the $^{15}N(\gamma,\pi^-)^{15}O$ reaction at $E_\gamma = 170$ MeV has shown the $E0$ term at the expected strength.[5] This apparent discrepancy has been interpreted as a nuclear-structure effect.[3]

In the present paper we propose that a direct measurement of the 0^+-0^+ transition in the $^{14}C(\gamma,\pi^-)^{14}N$ reaction would be both useful and feasible. The use of 0^+-0^+ transitions as in the Fermi (non-spin flip) transitions in β decay gives one the opportunity to make a measurement at all angles without the competition of a strong second transition. Unfortunately, it is hampered by the fact that few of these transitions are experimentally accessible. The case of ^{14}C is unusual because the target is unstable but the extremely retarded Gamow-Teller (spin-flip) β-decay matrix element between ^{14}C and the ^{14}N ground state[6] makes it possible to fabricate ^{14}C targets. This gives us a unique opportunity to observe the non-spin-flip term in the $^{14}C(\gamma,\pi^-)^{14}N$ reaction.

Due to its nuclear structure which is unique among p-shell nuclei, the $A=14$ system offers a number of interesting transitions to study (Fig. 2). The $M1$ transition from the ^{14}C to the ^{14}N ground state (and vice versa) is related to the retarded Gamow-Teller matrix element but is at higher momentum transfer. This transition has been studied extensively via the $^{14}N(\gamma,\pi^+)^{14}C$ reaction at Bates and Mainz. Note that the $^{14}C(\gamma,\pi^-)^{14}N$ and $^{14}N(\gamma,\pi^+)^{14}C$ reactions are not related by time-reversal invariance. They do have the same first-order matrix element and are related in the impulse approximation (IA) by the relative $\gamma p \rightarrow n\pi^+$ and $\gamma n \rightarrow p\pi^-$ reactions in the nuclear medium.

In addition to the retarded $M1$ transitions, there is a strong $M1$ transition that can be studied in the $^{14}C(\gamma,\pi^-)^{14}N$ reaction to the 1^+, $T=0$ level at 3.95 MeV in ^{14}N. This level is predicted to carry most of the Gamow-Teller sum rule in the $A=14$ system and has been seen in the $^{14}C(p,n)^{14}N$ reaction[7] with a strength that is consistent with the predicted strength.[8] This reaction is dynamically similar to the strong $M1$ transition in the $^{12}C(\gamma,\pi^{\pm})$ reaction. The dynamics of these processes should be under good control since, in the near-threshold region, they are dominated by the strong Kroll-Ruderman term and indeed the observed cross section is

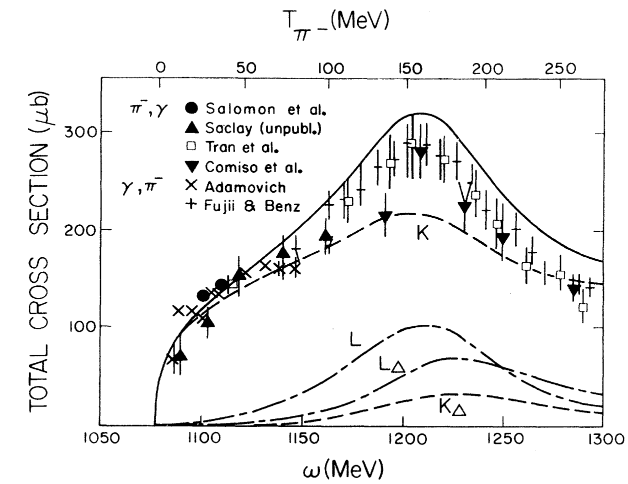

FIG. 1. Total cross section for $\gamma n \rightarrow \pi^- p$. The solid line shows the Born $+\Delta(1236)$ model of Blomqvist-Laget, the dash-dotted lines (L, L_Δ) give the non-spin-flip transitions ($S=0$), and the dashed lines (K, K_Δ) give the spin-flip transitions ($S=1$), for both Born $+\Delta$ and Δ separately. The figure is copied from Ref. 10.

in agreement with calculations for a pion energy of 32 MeV.[9]

The most complete and consistent way to take into account the strong pion-nucleus interaction in the final state is to use the distorted-wave impulse approximation (DWIA) in momentum space.[10,11] This allows a straightforward treatment of nonlocalities and propagators in elementary amplitudes such as the one developed by Blomqvist and Laget.[12] In this paper we will primarily use the methods of Ref. 10, which gives further details of our calculation. One feature of the calculation is that the nuclear structure and reaction dynamics appear in separate integrals. For our calculations we employ the pion optical potential by Stricker, McManus, and Carr,

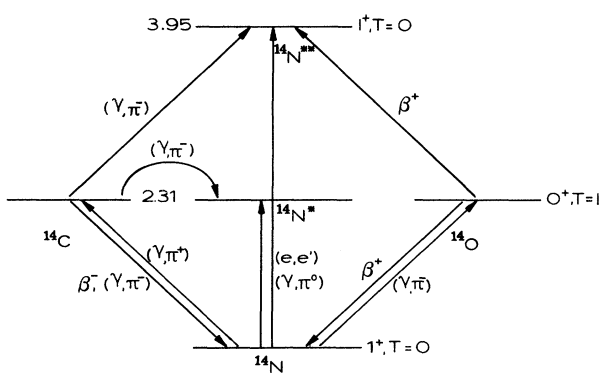

FIG. 2. Level diagram for the $A=14$ system with the transitions of interest.

which was obtained from the analysis of pion elastic scattering data.[13]

The elementary production operator can be divided into a spin-flip (spin-1) and non-spin-flip (spin-0) transition amplitude

$$t=(L+i\sigma\cdot\mathbf{K})\frac{\tau_{-\beta}}{\sqrt{2}}=\sum_{s,m_s} i^s(-1)^{m_s}\sigma^s_{-m_s}K^s_{m_s}\frac{\tau_{-\beta}}{\sqrt{2}} , \tag{1}$$

where $\sigma^0=1$ and $K^0=L$. In terms of the largest multipoles E_{0+} (electric dipole) and M_{1+} [magnetic dipole with the emerging pion in the resonant (3,3) channel], the non-spin-flip term L can be written as

$$L=2\hat{\mathbf{q}}\cdot(\hat{\mathbf{k}}\times\boldsymbol{\varepsilon})M_{1+} \tag{2}$$

and the spin-flip term $\mathbf{K}$ can be written as

$$\mathbf{K}=\boldsymbol{\varepsilon}(E_{0+}+M_{1+}\cos\theta_\pi)-\hat{\mathbf{k}}(\hat{\mathbf{q}}\cdot\boldsymbol{\varepsilon})M_{1+} , \tag{3}$$

where $\hat{\mathbf{k}}$ and $\hat{\mathbf{q}}$ are unit vectors in the direction of the incident photon and outgoing pion and $\boldsymbol{\varepsilon}$ is the photon polarization vector. For charged pion photoproduction in the near-threshold region where the E_{0+} term is dominated by the large Kroll-Ruderman amplitude the non-spin-flip term is considerably smaller; we note, however, that it is proportional to the resonant Δ magnetic multipole. This is why this term involves Δ dynamics even in the threshold region. Note also that for π^0 photoproduction where the E_{0+} multipole is weak, the spin-flip and non-spin-flip terms are comparable to each other and to the non-spin-flip term in charged pion photoproduction.

The $A=14$ system, depicted in Fig. 2, has been a well-known testing ground for nuclear-structure models within the $1p$ shell for years. The strongly suppressed β-decay rate of ^{14}C as well as the highly unusual β-decay spectra of ^{14}O and ^{14}C have aroused considerable interest. Most investigations have assumed the retarded β decay, which is a Gamow-Teller transition, to be due to a nearly complete cancellation among terms in the relevant $L=0$, $S=1$, $T=1$ one-body nuclear-structure matrix element. Further constraints are provided by other electromagnetic and weak observables such as the $M1$ form factors, the quadrupole moment, isovector and isoscalar magnetic moments, as well as some hadronic low-momentum-transfer reactions involving the $A=14$ system.

In the phenomenological analysis of elastic and inelastic electron scattering off ^{14}N by Huffman *et al.*[8] a $1p$-shell space is employed with no restriction on the form of the force to be used in this basis. The wave functions are constructed by coupling two $1p$-shell holes to an ^{16}O core. In a j-j coupling scheme, the most general wave functions are for the ground and 2.313 state in ^{14}N:

$$|0;J^\pi=1^+,T=0\rangle=a|1p^{-2}_{1/2}\rangle+b|1p^{-1}_{1/2}1p^{-1}_{3/2}\rangle +c|1p^{-2}_{3/2}\rangle , \tag{4}$$

$$|1;J^\pi=0^+,T=1\rangle=m|1p^{-2}_{1/2}\rangle+n|1p^{-2}_{3/2}\rangle , \tag{5}$$

normalized to $a^2+b^2+c^2=m^2+n^2=1$.

These parameters have been adjusted to yield the best fit to the elastic and inelastic $M1$ form factors (see Table I). In order to satisfactorily describe the high-Q enhance-

TABLE I. Comparison of configuration amplitudes. The parameters d and o are the $|2p_{1/2}1p_{1/2}^{-3}\rangle$ admixtures for the ground and excited states.

	a	b	c	d	m	n	o	a'	b'	c'
$H1$	0.977	−0.080	−0.195		0.553	−0.833		0.207	0.192	0.959
3.2%	0.918	−0.177	0.338	0.179	0.559	0.809	0.179			

ment of the inelastic form factor, Ref. 14 included an admixture of 3.2% of $|2p_{1/2}1p_{1/2}^{-3}\rangle$ configurations in both wave functions, representative of contributions from outside the $1p$ shell (see coefficients d and o of Table I). We will see below that the $E0$ part of pion photoproduction is sensitive to such admixtures for momentum transfer $Q \gtrsim 0.5\ \mathrm{fm}^{-1}$.

The most recent nuclear wave functions determined phenomenologically by a fit to the $M1$ form factor data of ^{14}N fall into two categories: Either they have a small but finite Gamow-Teller (GT) strength and underestimate the ^{14}C lifetime by several orders of magnitude, or the $L=0$, $S=1$, $T=1$ nuclear matrix element is almost zero yielding the proper logft value. Additional reactions are therefore needed to constrain the wave functions.

The photoproduction of low-energy pions proceeds via a reasonably well understood mechanism and, unlike electron scattering, is dominated by spin- and isospin-flip operators. Furthermore, the $^{14}\mathrm{N}(\gamma,\pi^+)^{14}\mathrm{C}_{\mathrm{g.s.}}$ cross section should be small at low Q if the $^{14}\mathrm{C}(\beta^-)$ decay rate suppression arises from a direct cancellation in the one-body matrix element. However, the experimental data measured in Mainz[15] and Bates[16] at $E_\gamma=173$ and 200 MeV are well described by wave functions with a finite GT matrix element.[15] This suggests that the GT matrix element is suppressed by a meson exchange contribution that is not present in (γ,π) reactions.

The wave function which determines the ^{14}N ground state and its excited states along with their isospin analogue states can be extended to the second excited state in ^{14}N at 3.95 MeV with the quantum number $1^+,0$ (see Fig. 2):

$$|2;J^\pi=1^+,T=0\rangle=a'|1p_{1/2}^{-2}\rangle+b'|1p_{1/2}^{-1}1p_{3/2}^{-1}\rangle+c'|1p_{3/2}^{-2}\rangle\ . \quad (6)$$

From normalization and orthogonality to the ground state, two of the three parameters are fixed, $a'^2+b'^2+c'^2=1$ and $aa'+bb'+cc'=0$. The remaining third parameter is then determined by constraining it to the strong Gamow-Teller β decay of $^{14}\mathrm{O}\to{}^{14}\mathrm{N}(1^+,0;3.95\ \mathrm{MeV})$ with a logft value of 3.11 ± 0.03 (Ref. 6).

We begin our discussion by comparing in Fig. 3 the magnitudes of the transitions to the ground state and two lowest excited states accessible in the reaction $^{14}\mathrm{C}(\gamma,\pi^-)^{14}\mathrm{N}$ as indicated in the level diagram of Fig. 2. Unless mentioned otherwise we will use the $H1$ wave function (Table I) for the calculations shown below. The reaction $^{14}\mathrm{C}(\gamma,\pi^-)^{14}\mathrm{N}_{\mathrm{g.s.}}$ is a pure $M1$ transition which for small angle cross section is suppressed by the very small Gamow-Teller (GT) matrix element discussed above. As has been previously discussed the small GT matrix element leads to a great sensitivity to the reaction dynamics. The angular distribution is similar to the process $^{14}\mathrm{N}(\gamma,\pi^+)^{14}\mathrm{C}$ with a deep minimum around 50°. Secondly, we show the process $^{14}\mathrm{C}(\gamma,\pi^-)^{14}\mathrm{N}^*$ which is the only pure $E0$ transition available in the p shell. The angular distribution at forward angles exceeds 1 μb while the second peak at backward angles may be harder to observe experimentally. The transition to the second excited state in ^{14}N at 3.95 MeV again has only an $M1$ contribution. Due to a much larger GT matrix element the cross section at forward angles exceeds the $M1$ transition to the ground state by nearly two orders of magnitude and almost reaches 10 μb. At backward angles both transitions have comparable cross sections.

As discussed in the introduction there is a special interest in the nuclear $E0$ excitation in photopion reactions because of its sensitivity to Δ dynamics, even near threshold. Figure 4 presents an angular distribution of the reaction $^{14}\mathrm{C}(\gamma,\pi^-)^{14}\mathrm{N}^*$ (2.31 MeV) at $E_\gamma=200$ MeV. To indicate the importance of the delta resonance even at this energy we have shown the Born terms separately. Their small contribution is due to the fact that in the ele-

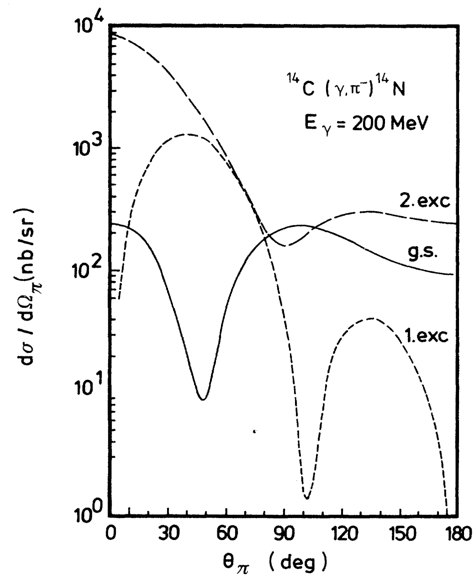

FIG. 3. Angular distribution for the reaction $^{14}\mathrm{C}(\gamma,\pi^-)^{14}\mathrm{N}$ at $E_\gamma=200$ MeV. The solid line shows the transition to the ground state ($0^+\to1^+$), the short-dashed line the transition to the first ($0^+\to0^+$), and the long-dashed line the transition to the second excited state of ^{14}N ($0^+\to1^+$).

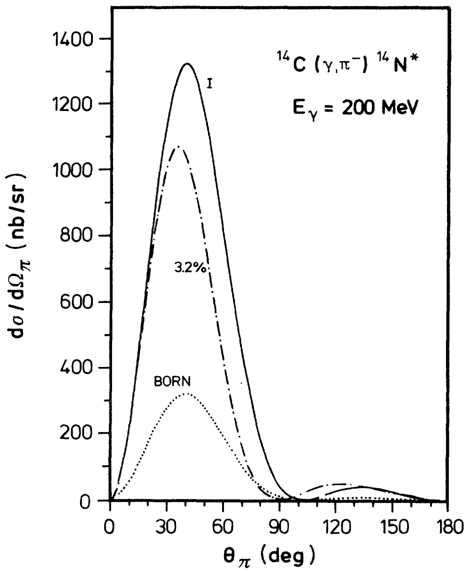

FIG. 4. Angular distribution for the 0^+-0^+ transition $^{14}C(\gamma,\pi^-)^{14}N^*(2.31\text{ MeV})$ at $E_\gamma = 200$ MeV. The solid (dotted) curve shows the full (Born terms only) calculation using BL-I (Ref. 12), while the dash-dotted line includes additional 3.2% of $2p$-shell admixture in the nuclear wave function.

mentary operator the non-spin-flip amplitude L is dominated by the delta. We compare a calculation performed within the $1p$-shell only with one that includes a $2p$-shell admixture of 3.2% (Table I). This is the empirical wave function[14] which improves the fit to the $^{14}N(e,e')^{14}N$ (2.31 MeV) for momentum transfers $Q > 1.5\text{ fm}^{-1}$. The latter calculation is reduced at the peak by almost 30% and has its minimum shifted to lower pion angles. Note that for pion angles $\leq 30°$ (momentum transfer $Q \leq 0.5\text{ fm}^{-1}$) the predicted cross sections are independent of the wave functions. The reason for this is that at low Q we are primarily sensitive to the wave-function normalization which is independent of configuration mixing. Therefore one has two separate regions. For low Q ($\leq 0.5\text{ fm}^{-1}$) the cross section is primarily sensitive to the reaction dynamics, particularly the Δ contribution. For larger Q the cross section is also dependent on the wave function. If measurements are performed in both regions, this could also provide nuclear-structure information.

It is of interest to discuss the previous attempts to determine the $E0$ contribution to mirror transitions in the (γ,π) reaction. Since the $M1$ and $E0$ transitions contribute incoherently, for small Q values the $M1$ part will dominate. In practice, it is only possible to determine the $E0$ part for $Q \gtrsim 1\text{ fm}^{-1}$ since for lower momentum transfers the $M1$ term is larger than the $E0$ contribution (note that this statement is model dependent since the relative shapes of the two contributions are determined by the wave functions but the relative magnitudes are determined by the reaction dynamics). This means that the wave-function-independent region $Q \leq 0.5\text{ fm}^{-1}$ cannot be reached by measurements using mirror transitions. For the $^{15}N(\gamma,\pi^-)^{15}O_{\text{g.s.}}$ reaction, the measured data are in good agreement with theoretical computations.[5] For the $^{13}C(\gamma,\pi^-)^{13}N_{\text{g.s.}}$ reaction the calculations tend to overpredict the $E0$ contribution[11,17] up to a factor of 2–3. Within the $1p$ shell the dominant $E0$ reduced density matrix element ($J=0, L=0, S=0$) is large and fixed by charge conservation and the nuclear Fermi transition, leaving no room for flexibility in nuclear structure. However, as shown in Ref. 3, the inclusion of selected multi-$\hbar\omega$ configurations, which are essential to describe electron-scattering form factors at higher momentum transfer,[14,18] can have a drastic effect on the $E0$ excitation. Only 6% of $2p$-shell admixture in the ground-state wave function of ^{13}C leads to the desired suppression of the $E0$ and agreement with experimental data.[3] The corresponding effect on the $M1$ contribution is far less pronounced. A very recent calculation[19] that employs a full shell-model wave function of ^{13}C with configurations up to the $2s$, $1d$, $2p$, and $1f$ shell confirms this more intuitive finding and yields a good overall agreement with all experimental data on $^{13}C(\gamma,\pi^-)^{13}N_{\text{g.s.}}$ and $^{13}C(\pi^+,\gamma)^{13}N_{\text{g.s.}}$. From this discussion it is clear that the measurement of the wave-function-independent Q region of $\leq 0.5\text{ fm}^{-1}$ for the $^{14}C(\gamma,\pi^-)^{14}N$ $0^+ \to 0^+$ transition will be vital to determine the dynamics of the $E0$ transition in nuclei.

Since ^{14}C is not a stable isotope, it is worthwhile discussing the experimental feasibility of making the proposed measurements. Targets of ^{14}C have been used for pion[20]- and electron[21]-scattering experiments. It is feasible to make targets up to approximately 1 g of ^{14}C (Ref. 22). For the $^{14}C(\gamma,\pi^-)^{14}N$ reaction about 1 MeV energy resolution is required to distinguish the final nuclear states. For pion kinetic energies of around 30 MeV where $dE/dx \sim 4\text{ MeV/(g/cm}^2)$ this means that the target thickness should not be larger than 0.2 g/cm^2. These experiments are best performed with tagged photon beams, which is feasible with the new cw electron accelerators. Using a tagged photon beam of $\sim 10^8$ photons/sec and a magnetic spectrometer with a solid angle of 30 msr and with a flight path of ~ 5 m (e.g., like the MEPS spectrometer at Bates) one would obtain a few counts/h for a cross section of 0.1 μb/sr. This would represent the lower limit of cross section that could be measured with this technique and would enable a measurement of the 0^+-0^+ transition at 200 MeV from 10° to 90°. If larger count rates are desired then large-area detectors such as plastic scintillators could be employed as long as the required resolution can be obtained. Furthermore, it is possible to make measurements with the bremsstrahlung end-point method as long as the excited-state cross sections are larger than the ground-state cross section. Again, this would enable one to measure the 0^+-0^+ transition at 200 MeV for angles up to 90°. In addition, the count rates would be much larger since untagged photon beams are more intense. We believe that either experimental solution is feasible.

In conclusion, we have presented calculations for the 0^+-0^+ non-spin-flip transition in $^{14}C(\gamma,\pi^-)^{14}N$ as well as $M1$ transitions to the ground and second excited state of

^{14}N in a momentum space DWIA framework. Using the reaction $^{14}C(\gamma,\pi^-)^{14}N$ $(0^+,1;2.31$ MeV) allows the measurement of the pure $E0$ transition in an experimentally clean way. This has been difficult to observe previously since in "mirror" transitions like $^{15}N(\gamma,\pi^-)^{15}O$ and $^{13}C(\gamma,\pi^-)^{13}N$ both spin-flip and non-spin-flip transitions contribute. There is no model-independent way to separate these two contributions unless polarization techniques are utilized. The effect of the delta channel is very large in the $E0$ multipole in the near-threshold region. The cross section for the 0^+-0^+ transition is insensitive to configuration mixing in the nuclear wave function for momentum transfer $Q \leq 0.5$ fm^{-1}. In particular, for this Q region, measurements of the $E0$ transition in muon transitions are not possible since the $M1$ contribution dominates in this region. For the other transition discussed in this paper the angular distributions for the process $^{14}C(\gamma,\pi^-)^{14}N(0^+,1;3.95$ MeV) are predicted in an almost model-independent way and should therefore agree well with the experiment. Finally, the reaction $^{14}C(\gamma,\pi^-)^{14}N_{g.s.}$ at low energies will provide an additional constraint on the magnitude of the disputed Gamow-Teller matrix element. It is therefore of great interest that this experiment be carried out as soon as possible.

ACKNOWLEDGMENTS

One of the authors (S.S.K.) would like to thank L. Tiator and D. Drechsel for the hospitality extended to him during his stay at the University of Mainz. This work was supported by the Deutsche Forschungsgemeinschaft (SFB 201).

[1]For a multipole analysis of the $N(\gamma,\pi)$ reaction see, e.g., F. A. Berends, A. Donnachie, and D. L. Weaver, Nucl. Phys. **B4**, 1 (1987); **B4**, 54 (1967); **B4**, 103 (1967).

[2]For recent reviews of the (γ,π) data, see A. M. Bernstein, in *Proceedings of the 7th Seminar on Electromagnetic Interactions of Nuclei at Low and Medium Energies, Moscow, 1988,* edited by G. M. Gurevitch (Institute for Nuclear Research, Academy of Science, Moscow, 1988); in *Hadronic Probes and Nuclear Interactions, Arizona State University,* Proceedings of the International Conference on Hadronic Probes and Nuclear Interactions, edited by J. R. Comfort, W. R. Gibbs, and B. G. Ritchie, AIP Conf. Proc. No. 133 (AIP, New York, 1985); P. Stoler, in *Proceedings of the 4th Workshop on Perspectives in Nuclear Physics at Intermediate Energies, Trieste, 1989,* edited by S. Boffi, C. Ciofi degli Atti, and M. Giannini (World Scientific, Singapore, 1990), p. 592.

[3]L. Tiator and C. Bennhold, in *Proceedings of the 4th Workshop on Perspectives in Nuclear Physics at Intermediate Energies* (Ref. 2), p. 581; C. Bennhold and L. Tiator, Phys. Lett. B **238**, 31 (1990).

[4]K. Shoda, A. Kagaya, O. Sasaki, S. Toyama, T. Kobayashi, and H. Tsubota, Phys. Lett. **169B**, 17 (1986); P. C. Dunn, A. Kaarsgaarn, J. H. Koch, P. K. A. de Witt Huberts, P. Stoler, P. F. Yergin, M. Seneviratne, and B. Schoch, Phys. Lett. B **196**, 434 (1987).

[5]A. Liesenfeld, B. Alberti, D. Eyl, G. Köbschall, A. W. Richter, K. Röhrich, Ch. Schmitt, L. Tiator, and V. H. Walther, Nucl. Phys. **A485**, 580 (1988).

[6]F. Aizenberg-Selove, Nucl. Phys. **A449**, 1 (1986).

[7]J. Rapaport *et al.*, Phys. Rev. C **36**, 500 (1987).

[8]R. L. Huffman, J. Dubach, R. S. Hicks, and M. A. Plum, Phys. Rev. C **35**, 1 (1987).

[9]Ch. Schmitt, K. Röhrich, K. Maurer, C. Ottermann, and V. H. Walther, Nucl. Phys. **A395**, 435 (1983).

[10]L. Tiator and L. E. Wright, Phys. Rev. C **30**, 989 (1984).

[11]R. A. Eramzhyan, M. Gmitro, S. S. Kamalov, and R. Mach, J. Phys. G **9**, 605 (1983).

[12]K. I. Blomqvist and J. M. Laget, Nucl. Phys. **A280**, 405 (1977).

[13]K. Stricker, H. McManus, and J. A. Carr, Phys. Rev. C **19**, 929 (1979); J. A. Carr, H. McManus, and K. Stricker-Bauer, *ibid.* **25**, 952 (1982).

[14]R. L. Huffman, R. S. Hicks, J. Dubach, B. Parker, M. A. Plum, G. Lahm, R. Neuhausen, and J. C. Bergstrom, Phys. Lett. **139B**, 249 (1984).

[15]K. Röhrich, L. Tiator, G. Köbschall, Ch. Reifferscheid, Ch. Schmitt, V. H. Walther, K. Weinand, and L. E. Wright, Phys. Lett. **153B**, 203 (1985); K. Röhrich, G. Köbschall, Ch. Reifferscheid, Ch. Schmitt, V. H. Walther, and K. Weinand, Nucl. Phys. **A475**, 762 (1987).

[16]B. Cottman *et al.*, Phys. Rev. Lett. **55**, 684 (1985); L. Ghedira *et al.*, Phys. Rev. C **41**, 653 (1990).

[17]T. Suzuki, T. Takaki, and J. H. Koch, Nucl. Phys. **A460**, 607 (1986).

[18]R. S. Hicks, J. Dubach, R. A. Lindgren, B. Parker, and G. A. Peterson, Phys. Rev. C **26**, 339 (1982).

[19]C. Bennhold and L. Tiator, Nucl. Phys. **A519**, 805 (1990).

[20]C. J. Harvey, H. W. Baer, J. A. Johnstone, C. L. Morris, S. J. Seestrom-Morris, D. Dehnhard, D. B. Holtkamp, and S. J. Greene, Phys. Rev. C **33**, 1454 (1986).

[21]M. Plum, R. A. Lindgren, J. Dubach, R. S. Hicks, R. L. Huffman, J. Alster, J. Lichtenstadt, M. Moinester, and H. Baer, Phys. Rev. C **40**, 1861 (1989).

[22]H. Baer, private communication.