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π^+ photoproduction near threshold on ^{3}He

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The positive-pion photoproduction yield on ^{3}He was measured near threshold. The transition matrix element of this process is extracted with a $\pm 1.5\%$ accuracy. We discuss the relation of our result, firstly, with magnetic electron scattering on ^{3}He and ^{3}H, secondly, with the properties of pionic ^{3}He atom.

[NUCLEAR REACTION ^{3}He(γ, π^+)^{3}H, measured σ, $E=1-5$ MeV.]

I. INTRODUCTION

The contribution of mesonic degrees of freedom to the various processes in nuclei can be investigated with the highest confidence in the case of the two- and three-nucleon systems, for which accurate wave functions, based on realistic nucleon-nucleon potentials[1] are available. Of particular interest are the spin-flip operators which govern the magnetic electron scattering and the low energy pion photoproduction.

In the case of the two-nucleon system, the magnetic electrodisintegration cross section at threshold has been computed in impulse approximation[2] and found to be lower (by about 20% at $q^2 \simeq m_\pi^2$) than the measured one.[3] This discrepancy is removed by taking into account the mesonic degrees of freedom in the nucleus. On the other hand, the experimental cross section[4] for the reaction

$$\gamma + d \to n + n + \pi^+ \tag{1}$$

is well reproduced in the plane wave impulse approximation[5] (PWIA), in contrast to the case of the threshold electrodisintegration. Thus, at least for the two-nucleon system, the spin-flip form factor is more reliably determined by the threshold photoproduction process.

We have extended this investigation to the ^{3}He nucleus by carefully measuring the cross section for the reaction

$$\gamma + {}^3\text{He} \to {}^3\text{H} + \pi^+ \tag{2}$$

in an energy range from 1 to 5 MeV above threshold, relative to the cross section on the corresponding reaction on a free proton

$$\gamma + p \to n + \pi^+ . \tag{3}$$

From our data we extract the ratio of the slopes of the threshold cross sections for reactions (2) and (3). This result is compared with data existing on the (π^-,^{3}He) pionic atom: Panofsky ratio, radiative branching ratio, and total width of the 1s level.

The available microscopic theories take into account in the effective photoproduction amplitude only the leading spin-flip term $E_{0+}\vec{\sigma}\cdot\vec{\epsilon}$, so that no definite conclusion can be drawn as yet on the eventual role of meson exchange currents in reaction (2).

Nevertheless, we confirm the inadequacy of the procedure which uses magnetic electron scattering data[6] to predict the pion photoproduction at threshold.

From a different point of view, our result can be expressed in terms of the (π^+,^{3}H) coupling constant, which is found to be in agreement with recent determinations coming from various experiments.

II. THE EXPERIMENT AND ITS RESULT

The setup and the experimental method are similar to the one described in detail elsewhere.[8] We shall comment only on the peculiarities of the present experiment. The two detector sets consisted of two Čerenkov Lucite counters 20 mm thick and a third one 40 mm thick. The targets were stainless steel cylinders 110 mm long, 40 mm in diameter, filled, respectively, with liquefied ^{3}He and hydrogen; the steel windows hit by the photon beam were 2/100 mm thick. The liquid content of the cells was known by measuring the pressure drop in the gas vessels. These were isolated from the targets at the end of the filling procedure; the vapor pressure above the liquid was recorded and used to determine the density. We checked the superficial mass of both targets by a transmission measurement using 511 keV photons from a collimated ^{22}Na source. The quantity relevant for our analysis is the ratio of the densities of the two targets: ^{3}He at 1.90 °K and hydrogen at 17.2 °K. Our measurement yielded for this ratio the value 1.08 ± 0.03, in good agreement with the value 1.10 obtained from densities quoted in the literature,[7] and which we adopt in the data analysis. The ^{4}He contamination in the ^{3}He was found to be 1.4% by mass spectrography; the ^{1}H purity was 99.8%.

The yield per target nucleus at electron energy E_e for a given integrated beam intensity is

$$A = \frac{\epsilon\Omega}{4\pi}\int_{E_{\rm th}}^{E_e} R(E)\sigma(E)B(E,E_e)dE,$$

where $E_{\rm th}$ is the reaction threshold, $\epsilon\Omega/4\pi$ the detection efficiency of the e^+ from the $\pi^+\rightarrow\mu^+\rightarrow e^+$ decay, $R(E)$ the fraction of the muons stopping in the target, $\sigma(E)$ the total π^+ photoproduction cross section for photon energy E in the laboratory system, and $B(E,E_e)$ the bremsstrahlung spectrum produced by the electron beam of energy E_e and $\Delta p/p = 0.003$ hitting a 0.01 radiation length tungsten foil.

The quantity R changes with the pion energy and is different for helium-3 and hydrogen. The escape probability of pions and muons was estimated by a Monte Carlo procedure including the complete setup geometry; the results of calculation were approximated by

$$R(p) = 1 - 0.0015(E - E_{\rm th})^2,$$

$$R(^3{\rm He}) = 1 - 0.0035(E - E_{\rm th})^2,$$

where the energy in the laboratory system is expressed in MeV.

Near threshold the cross sections are conveniently parametrized by

$$\sigma(p) = a(p)q/k,$$

$$\sigma(^3{\rm He}) = a(^3{\rm He})Sq/k,$$

where q and k are the pion and photon momenta in the center-of-mass system; $S = 2\pi\gamma/(e^{2\pi\gamma}-1)$, where $\gamma = e^2/\hbar v$, describes the effect of the Coulomb repulsion between a pointlike tritium residual nucleus and the pion of velocity v. The quantity $a(^3{\rm He})$ contains the nuclear transition matrix element, the distortion of the pion wave by the strong interaction with the nucleon, and the (negligible) influence of the finite extension of the nuclear charge; because of the small size of the nuclear states involved, $a(^3{\rm He})$ has been taken as a constant in our analysis (this point has been discussed in Ref. 8). We used the bremsstrahlung shape $B(E,E_e)$ of Jabbur and Pratt[9] to compute the integrals $I(E_e) = \int R(E)(q/k)SB(E,E_e)dE$ for both target nuclei. Inserting these quantities in the expression of the theoretical yield, we obtain

TABLE I. Photoproduction yield per nucleus for the hydrogen and ^{3}He targets at different values of the nominal bremsstrahlung end point energy E_e.

E_e (MeV)	$A(p)$ (a.u.)	E_e (MeV)	$A(^3{\rm He})$ (a.u.)
152.5	1.05 ± 0.19	145	1.73 ± 0.20
153	4.24 ± 0.39	145.5	3.74 ± 0.29
154	16.4 ± 0.79	146.5	12.3 ± 0.58
155	32.0 ± 1.1	147.5	23.3 ± 0.70
156	54.3 ± 1.4	148.5	37.0 ± 0.93

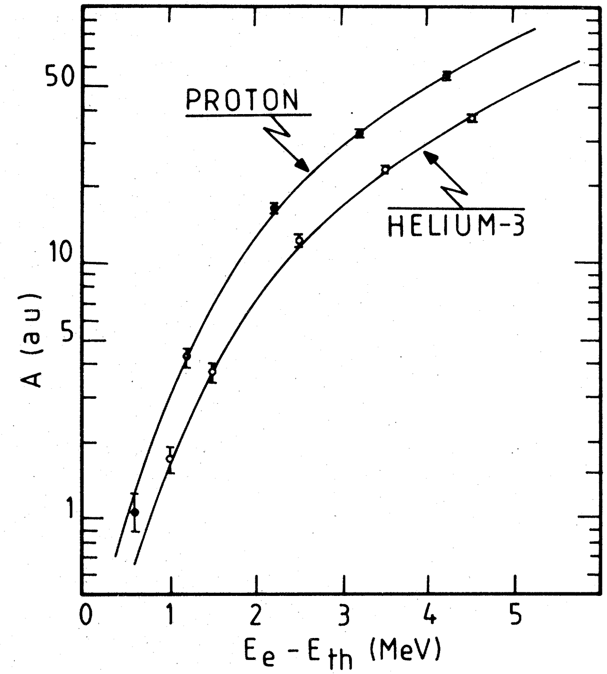

FIG. 1. Measured π^+ photoproduction yields per nucleus as a function of the excess energy above threshold in the laboratory system for hydrogen and ^{3}He. Solid lines are the calculated yields giving the best fit to the experimental data.

$$A(p) = \frac{\epsilon\Omega}{4\pi}a(p)I_p(E_e)$$

and

$$A(^3{\rm He}) = \frac{\epsilon\Omega}{4\pi}a(^3{\rm He})I_{\rm He}(E_e).$$

These functionals were adjusted to our experimental data (Table I and Fig. 1), allowing for an energy shift between the nominal and the true values of E_e as explained in Ref. 8. The fit yields,[10] with a normalized $\chi^2 = 1.0$,

$$a(^3{\rm He})/a(p) = 0.62 \pm 0.02. \qquad (4)$$

Using the value $a(p) = (201 \pm 7)\ \mu$b,[11] we can reconstruct the absolute total cross section as a function of energy, with a relative accuracy of 5% (Fig. 2).

III. EVALUATION OF THE TRANSITION MATRIX ELEMENT

Here we restrict our interpretation to a microscopic description. We neglect all many-body effects except those involving the pion rescattering and obtain

$$a(^3\mathrm{He}) = \lim_{q=0} \frac{1}{S}\frac{k_3}{q}\sigma(\gamma\,^3\mathrm{He} \to {}^3\mathrm{H}\pi^+)$$
$$= 4\pi\,|E_{0+}(n\pi^+)|^2 \frac{1+\mu/M_1}{1+\mu/M_3}$$
$$\times \frac{1+k_1/M_1}{1+k_3/M_3} C_+^2\,|M_+(Q_+^2)|^2 . \qquad (5)$$

C_+ is the modification of the π^+ photoproduction amplitude computed in the plane wave approximation, due to the multiple scattering. μ, M_1, and M_3 are the masses of the pion, nucleon, and $3N$ nucleus, respectively; k_1 and k_3 are the photon energies at threshold in the center-of-mass system for reactions (3) and (2).

The matrix element is defined by

$$|M_+(Q_+^2)|^2 = \frac{1}{|E_{0+}(n\pi^+)|^2}$$
$$\times \tfrac{1}{4} \sum_{m_i m_j \lambda} |\langle ^3\mathrm{H}|\sum_{j=1}^{3} 0_j e^{-i\vec{Q}_+\cdot\vec{r}_j}|^3\mathrm{He}\rangle|^2 . \qquad (6)$$

0_j is the full one-body photoproduction operator which contains, in addition to the leading $E_{0+}\vec{\sigma}_j\vec{\epsilon}_\lambda\tau_j^{(-)}$, momentum dependent terms. $Q_+^2 = 0.481\ \mathrm{fm}^{-2}$ is the momentum transfer at threshold ($q=0$) in the center-of-mass system; as it is a fixed quantity, the matrix element is independent of the π^+ emission angle.

Since by definition $a_p = 4\pi|E_{0+}|^2$, our result yields

$$|M_+(Q_+^2)|^2 = (0.52 \pm 0.02)/C_+^2 .$$

Let us turn now to the evaluation of C_+. Considering the magnitude of the average internucleon distance, one could expect a non-negligible correction to the PWIA. This problem is generally handled by using a distorted-wave impulse approximation (DWIA). The pion plane wave $e^{i\vec{q}\cdot\vec{r}_j}$ is replaced in the matrix element by its solution $\phi_\pi(r_j)$ in the combined strong and Coulomb potentials. The real part of the strong optical potential is generated in first order by the sum of the scattering amplitudes on all nucleons. In a production process, this prescription means that in the first order, the pion is allowed to scatter on the nucleon which is its own source. This effect, however, is already included when one uses the effective production amplitude on the nucleon. Thus it is inconsistent to describe a production process by an effective production amplitude together with a first order optical potential. The correct procedure is to perform a multiple scattering calculation. For such a light nucleus as the (^{3}He $-$ ^{3}H), the distinction is not academic. In terms of the pion nucleon scattering lengths a, the first order

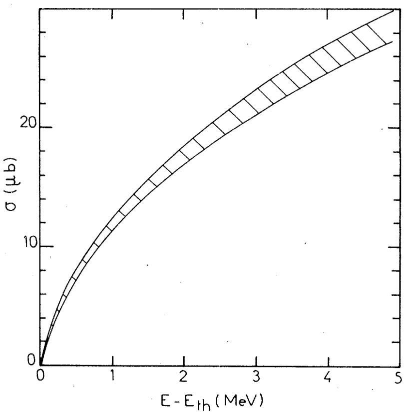

FIG. 2. Cross section of the $(\gamma\pi^+)$ reaction on ^{3}He deduced from this experiment as a function of the excess energy above threshold in the laboratory system. The shaded area represents the experimental error taking into account the uncertainty both of this experiment and that of $a_p = 4\pi|E_{0+}(n\pi^+)|^2$.

optical potential contains $2a(\pi^+n) + a(\pi^+p) \simeq a(\pi^+n)$ because of the near cancellation of the scattering length on an $(n$-$p)$ pair. The remaining $a(\pi^+n)$ is attractive and produces an increase over the PWIA cross section by 8%, as found in Ref. 12. If we compute, instead, the multiple scattering expansion, the correction to the matrix element induced by a first rescattering is

$$(1+\mu/M_1)[a(\pi^+p) + a(\pi^+n)]\left\langle \frac{1}{r_{12}} \right\rangle \simeq -0.0065 ,$$

a reduction from the PWIA cross section by about 1%. Here r_{12} is the internucleon distance in the $3N$ nucleus.

As the first order corrections to the processes discussed in this paper are generally small, we pursued our calculation[13] up to the second order, using the fixed scatterer approximation which predicts correctly the $(\pi^- d)$ scattering length. We also verified that the third order scattering amplitude is negligible.

For the case of reaction (2) this procedure yields

$$C_+^2 = 0.99 . \qquad (7)$$

Thus we obtain

$$|M_+|^2 = 0.52 \pm 0.02 \text{ at } Q_+^2 = 0.481\ \mathrm{fm}^{-2} . \qquad (8)$$

IV. COMPARISON WITH ELECTRON SCATTERING DATA

Magnetic electron scattering on nuclei is mainly a nucleon spin-flip process. If the orbital contribution and the mesonic corrections were negligible, the measured magnetic form factor would reduce to the spin-flip form factor. Similarly, if the momentum dependent terms in the photoproduction amplitude had no influence, this process could be described by the same transition spin-flip form factor, cf. Eq. (6).

From electron scattering experiments,[6] one extracts the magnetic transition form factor (^{3}He $\to$ ^{3}H) at $Q_+^2 = 0.481$ fm^{-2},

$$|F(Q_+^2)|^2/|F(0)|^2 = 0.61 \pm 0.02 ,$$

which yields, after unfolding the nucleon magnetic form factor, the body form factor

$$|F_m(Q_+^2)|^2 = 0.64 \pm 0.02 . \tag{9}$$

This is 20% higher than our value (8). Whatever the reason for this difference, which was noted also for the $2N$ system, it is worthwhile to emphasize that at the level of high accuracy of the existing data the two processes cannot be unified by a common form factor.

V. RELATION WITH THE π^- RADIATIVE CAPTURE IN ^{3}He

For reasons which will be discussed in Sec. VI, the squared matrix element $\frac{1}{2}|M_-|^2$ for the reaction

$$\pi^- + {}^3\text{He} \to {}^3\text{H} + \gamma \tag{10}$$

is expected to be practically equal to $|M_+|^2$. (The factor $\frac{1}{2}$ takes into account the different statistical factors in the two reactions.) Only many-body effects breaking the charge symmetry could influence differently M_+ and M_-.

M_- can be obtained from the 1s orbit Panofsky ratio P^{1s}. The measured value P is related to the charge exchange and radiative transition rates, λ_{nl} (exchange) and λ_{nl} (radiative), from all pion orbits (n,l), and is not equal to P^{1s}, the Panofsky ratio from the 1s orbit. In the case of such a light nucleus as ^{3}He, where pions are mostly captured from s orbits, further important simplifications occur:

(i) The estimated radiative and charge exchange branching ratios are much smaller for p orbits than for s orbits.

(ii) The charge exchange and radiative branching ratios from s states can be assumed to be independent of the principal quantum number n.

Then one can write for ^{3}He

$$P_3\text{measured} \simeq P_3^{1s} = \frac{\lambda_{1s}(\text{exchange})}{\lambda_{1s}(\text{radiative})}$$

and, more precisely, using calculated transition rates[14] from the 2p orbit,

$$P_3 < P_3^{1s} < 1.03 P_3 .$$

The M_- matrix element is related to P_3^{1s} and P_1, the Panofsky ratio in hydrogen[15]:

$$|M_-(Q_-^2)|^2 = 2\frac{P_1}{P_3^{1s}}\frac{k_1}{k_3}\frac{M_3+\mu-\omega_{03}}{M_1+\mu-\omega_{01}}\frac{M_1+\mu-k_1}{M_3+\mu-k_3} \times |M_{-0}(q_{03}^2)|^2\frac{C_{-0}^2}{C_-^2} , \tag{11}$$

where q_{0i} and ω_{0i} are the momenta and energies in the c.m. system of the outgoing π^0 in the case of nucleus $i=1$ (proton) or $i=3$ (^{3}He). M_{-0} is the charge exchange matrix element. At the small momentum transfer $q_{03}^2 = 0.027$ fm^{-2}, $|M_{-0}|^2$ is taken to be equal to $1-\frac{1}{3}q_{03}^2\langle R_{\text{ch}}^2\rangle = 0.97$, the squared charge form factor[6] of the $3N$ nucleus.

The multiple scattering correction factors for charge exchange and radiative capture, computed as explained above, are

$$C_{-0}^2 = 0.92 \text{ and } C_-^2 = 0.98 .$$

Using the most recent determinations of the Panofsky ratios,[16]

$$P_1 = 1.548 \pm 0.009 ,$$

$$P_3 = 2.82 \pm 0.07 ,$$

we obtain

$$\tfrac{1}{2}|M_-(Q_-^2)|^2 = 0.56 \text{ at } Q_-^2 = 0.474 \text{ fm}^{-2} .$$

Considering the variation of the electromagnetic form factors of the $3N$ nucleus[6] from $Q_-^2 = 0.474$ fm^{-2} to $Q_+^2 = 0.481$ fm^{-2}, one can estimate that $|M_-|^2$ would decrease by a negligible amount of 0.8%. The agreement between this value and our determination (8) is reasonably good.

It is interesting to exploit the link between P_3, the radiative branching ratio R_γ, and the total width Γ_{1s} of the 1s state of the pionic ^{3}He atom. The radiative branching ratio from the 2p orbit, weighted by the capture probability W_{2p}, is negligible compared to those from s orbits. Assuming, as above, that the branching ratio does not depend on the principal quantum number, one finds

$$R_\gamma\text{measured} \simeq R_\gamma^{1s} \simeq \frac{\lambda_{1s}(\text{radiative})}{\lambda_{1s}(\text{total})}\sum_n W_{ns} .$$

The total width of the 1s state is expressed as follows:

$$\Gamma_{1s} = |M_-(Q_-^2)|^2 C_-^2 \sum_n W_{ns}\frac{1}{R_\gamma^{1s}}\frac{\hbar c k_3}{\mu}\left(1-\frac{k_3}{M_3+\mu}\right) \times \left(1+\frac{\mu}{M_1}\right)^2 |\psi_\pi(0)|^2 4\pi |E_{0^+}(p\pi^-)|^2 .$$

$|\psi_\pi(0)|^2$ is the central density of the pion wave in a point-Coulomb field (the correction due to the finite size of ^{3}He is of the order of 2×10^{-4}).

Using

$$R_\gamma = (6.6 \pm 0.8) \times 10^{-2} \quad \text{(Ref. 15)},$$

$$4\pi |E_{0+}(p\pi^-)|^2 = (254 \pm 30) \times 10^{-30}\ \text{cm}^2 \quad \text{(Ref. 11)},$$

one obtains

$$\Gamma_{1s} = (32 \pm 5) \sum_n W_{ns}\ \text{eV}. \qquad (12)$$

32 ± 5 eV represents the upper limit of Γ_{1s} as determined by the radiative branching ratio. If one assumes that $\sum_n W_{ns} = 0.84$ as in the case of ^{4}He,[17] one gets $\Gamma_{1s} = (27 \pm 4)$ eV.

Γ_{1s} has been directly measured and found to be equal to (42 ± 14) eV (Ref. 18) and (68 ± 15) eV.[19] Only the first of these direct determinations is in agreement with (12), provided that $\sum_n W_{ns}$ is close to 1.

VI. COMPARISON WITH PREDICTED VALUES OF THE SPIN-FLIP FORM FACTOR

The available theoretical predictions fall in two categories:

(i) The calculations[20] using for $M(Q^2)$ the form factor extracted from magnetic electron scattering experiments. We have shown in Sec. IV that the present data disprove this approach.

(ii) The microscopic calculations[14,12] of the spin-flip form factor $F_{sf}(Q^2)$ computed from $3N$ wave functions obtained from various $2N$ potentials. At threshold, $M_+(Q^2)$ reduces to $F_{sf}(Q^2)$ only in the approximation of "frozen" nucleons. However, the Fermi motion of the nucleons introduces, in the interaction Hamiltonian, terms depending on the nucleon momentum. This effect has been evaluated[21] for the reaction $\gamma + {}^6\text{Li} \rightarrow \pi^+ + {}^6\text{He}$ and has resulted in a 10% decrease of the theoretical cross section. As the average momentum of a proton in ^{3}He and in the $1p$ shell in ^{6}Li are almost equal,[22] one expects a similar reduction in the present case.

This remark applies equally well to π^- radiative capture because the relative importance of the momentum dependent and main terms is approximately the same for π^- and π^+ photoproduction on the nucleon; furthermore, for the capture from s orbits, the terms linear in the pion momentum contribute negligibly, as shown in Ref. 23. In addition to the above mentioned deficiency of the theory, it should be noted (see Table II) that different N-N potentials lead to widely different values for the form factor.

Hence it is impossible to draw any firm conclusion at this time.

TABLE II. Theoretical values of the squared spin-flip transition form factor $F_{sf}(Q_+^2)$ at transfer $Q_+^2 = 0.481$ fm^{-2} for different $3N$ wave functions. P_D is the D state percentage in the $3N$ wave function.

Wave function	I[a]	II[a]	III[a]	MT13[b]	SSC[b]	RSC[b]
P_D(%)	0	4.9	9.2	0	7.9	9.3
$\lvert F_{sf}(Q_+^2)\rvert^2$	0.63	0.57	0.53	0.57	0.52	0.49[c]

[a] From Ref. 14.
[b] From Ref. 12.
[c] There is a misprint in Ref. 12 which was pointed out to us by the authors; we quote here the correct value.

VII. THE PION-^{3}He COUPLING CONSTANT

The present experiment also yields a direct determination of the pion-^{3}He coupling constant f_3,[24] following an approach already applied to the $A = 6$ system.[25] One obtains $f_3^2 = 0.043 \pm 0.002$, a sizable reduction compared to the pion-nucleon coupling $f_1^2 = 0.080$.[11]

Recent polological treatments of the $p + {}^3\text{H} \rightarrow n + {}^3\text{He}$ (Ref. 24) and ${}^3\text{He} + {}^3\text{H} \rightarrow {}^3\text{He} + {}^3\text{H}$ (Ref. 27) reactions provide additional, though less precise, indications for this reduction: $f_3^2 = 0.045 \pm 0.015$ (Ref. 26) and $f_3^2 = 0.055 \pm 0.020$ (Ref. 27), respectively.

The reduction of the pion-nucleus coupling compared to the pion-nucleon one, as deduced from our result, is in contradiction with the enhancement expected from the customary procedure to extrapolate the variation of the axial nuclear form factor from the spacelike region, where it is measured, to the timelike point $q^2 = -m_\pi^2$.[28]

As this reduction seems to increase with the size of the nuclear system ($f_3^2 = 0.04$, $f_6^2 \cong 0.02$, $f_{12}^2 \cong 0.006$), it may be attributed qualitatively to an increasing shadowing of the pion-nucleon coupling by the nucleus.[26]

Let us also mention that our result can be used to estimate the contribution of the breakup channels to the Goldberger-Treiman relation applied to the $3N$ nucleus.

VIII. CONCLUSION

The experiment described in this paper illustrates the precision and reliability obtained in the determination of the matrix element for threshold pion photoproduction.

By comparing our result with magnetic electron scattering data, we show that a unique spin-flip form factor cannot account for both processes. This fact suggests that many-body contributions affect the two reactions differently. An interesting extension of our work would be the determination of the matrix element M_+ at larger momentum transfer by measuring the cross section at threshold for the reaction $e + {}^3\text{He} \to e' + {}^3\text{H} + \pi^+$.

We have also shown that our result is consistent with data related to stopped pion radiative capture. The Panofsky ratio agrees well with our measurement when pion rescattering is correctly taken into account in the estimation of the charge exchange rate.

We call attention to the importance of achieving a complete and coherent calculation, treating, in parallel, the ^{3}He and the deuterium case which would yield quantitative information on the many-body contribution to π photoproduction.

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