Can relativistic pionic stripping explain (p,π^+) reactions?*

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The relativistic pionic stripping formalism is used to study pion production data on ^{12}C and ^{40}Ca in order to determine the appropriate form of the pion-nucleon vertex and to determine whether pionic stripping is the dominant mechanism for pion production.

[NUCLEAR REACTIONS Proton induced pion production, relativistic pionic stripping, $^{12}C(p,\pi^+)^{13}C$, $^{40}Ca(p,\pi^+)^{41}Ca$; calculated angular distributions, energy dependence (near threshold).]

I. INTRODUCTION

At present the mechanism for the $A(p,\pi^+)A+1$ reactions is not well established. Some[1] favor the pionic stripping mechanism in which [in analogy with the (p,γ) process] the incident proton shakes off a positive pion before being captured. Others,[2] noting that two or more nucleons are likely to be ejected from nuclei after pion capture, have proposed more complicated reaction mechanisms for (p,π^+). Attempts to pin down the reaction mechanism by comparison with existing experimental data are complicated by the large momentum transfer and consequent lack of nuclear structure information as well as by ambiguities[3] inherent in the pionic stripping mechanism.

The pionic stripping controversy stems from different nonrelativistic reductions for the pion-nucleon vertex. The most frequently discussed choices are the static model

$$\Gamma_{\pi N}=\frac{f}{m_\pi}\vec{\sigma}\cdot\vec{\nabla}_\pi \tag{1}$$

and the Galilean-invariant model

$$\Gamma_{\pi N}=\frac{f}{m_\pi}\vec{\sigma}\cdot\left[\vec{\nabla}_\pi+\frac{\omega_\pi}{2M}(\vec{\nabla}_N-\overleftarrow{\nabla}_N)\right]. \tag{2}$$

To circumvent the ambiguity in the nonrelativistic pion-nucleon vertex, a relativistic pionic stripping formalism has been suggested.[4,5] The cross section in this theory results from the relativistic T matrix element

$$T_{fi}=g\sqrt{2}\int d^4x\,\delta(x_0)\psi^{\omega}_{Jm}(x)\Gamma\psi^{\rm in}_{\vec{p},\lambda}(x)\phi^{\rm out}_{\vec{k}}(x) \tag{3}$$

in which $\psi^{\omega}_{Jm}(\chi)$ is a Dirac shell-model wave function of the captured proton, Γ is the relativistic pion-nucleon vertex operator, and $\vec{p}$ is the momentum of the incident proton. The vector k is the final momentum of the pion so that $q=p-k$ is the four-momentum transfer.

Using plane waves (PWIA) for the pion $\phi^{\rm out}_{\vec{k}}$ and the proton and a Dirac spinor $u_\lambda(\vec{p})$ for the latter, this T matrix element can be reduced to

$$T_{fi}=C\left(T_{\rm S}\mathcal{Y}^{\dagger\omega}_{Jm}(\hat{q})\frac{\vec{\sigma}\cdot\vec{k}}{p_0+M}\chi_\lambda+T_{\rm NS}\mathcal{Y}^{\dagger-\omega}_{Jm}(\hat{q})\chi_\lambda\right) \tag{4}$$

in the rest frame of the final nucleus. M is the nucleon mass, $\mathcal{Y}^{\omega}_{Jm}(\hat{q})$ is the central field spinor associated with the Dirac shell-model wave function, and $\vec{\sigma}$ and χ_λ are the Pauli spin matrix and spinor, respectively. The factors $T_{\rm S}$ and $T_{\rm NS}$ are called the static and nonstatic form factors by virtue of the fact that the first term of Eq. (4) resembles the matrix element resulting from the nonrelativistic pion-nucleon vertex operator, Eq. (1). The cross section resulting from Eq. (4) in the laboratory system is

$$\frac{d\sigma}{d\Omega}=\frac{(2J+1)kMM_{A+1}g^2}{2\pi[(P\cdot P_A)^2-M^2M_A{}^2]^{1/2}}\times\left(\frac{k^2}{(P_0+M)^2}|T_{\rm S}|^2+|T_{\rm NS}|^2-\frac{2\hat{q}\cdot\vec{k}}{P_0+M}T_{\rm S}T_{\rm NS}\right) \tag{5}$$

with the πNN coupling $g^2/4\pi=14.6$.

The form factors $T_{\rm S}$ and $T_{\rm NS}$ were shown to be[4]

$$T_{\rm S}=F_l(q);\quad T_{\rm NS}=\omega G_{l'}(q)-\frac{q}{P_0+M}F_l(q) \tag{6}$$

for the pseudoscalar coupling $(\Gamma=\gamma^5)$, and

$$T_S = \frac{1}{2M}\left[(P_0+M-k_0)F_l(q) - \omega q G_{l'}(q)\right],$$

$$T_{NS} = \frac{k_0}{2M}\left(\omega G_{l'}(q) + \frac{q}{P_0+M}F_l(q)\right) - \omega\frac{(\vec{k})^2+2\vec{k}\cdot\vec{q}}{2M(P_0+M)}G_{l'}(q) \tag{7}$$

for pseudovector coupling ($\Gamma = \gamma^5\gamma^\mu\partial_{\mu\pi}/2M$). The functions $F_l(q)$ and $G_{l'}(q)$ are the momentum space radial large and small component wave functions and are related to the coordinate space wave functions through the Fourier-Bessel transformation

$$F_l(q) = \int_0^\infty rF(r)j_l(qr)dr;$$

$$G_{l'}(q) = \int_0^\infty rG(r)j_{l'}(qr)dr, \tag{8}$$

where $l = J+\frac{1}{2}\omega$ and $l' = J-\frac{1}{2}\omega$. The functions $F(r)$ and $G(r)$ are components of the Dirac wave function

$$\psi^\omega_{Jm}(t=0,r) = \frac{1}{r}\begin{pmatrix} F(r)\mathcal{Y}^\omega_{Jm}(\theta,\phi) \\ iG(r)\mathcal{Y}^{-\omega}_{Jm}(\theta,\phi)\end{pmatrix} \tag{9}$$

which has parity $P=(-1)^{J+\omega/2}$ ($\omega=\pm1$). The wave function ψ^ω_{Jm} is an eigenfunction of the shell-model Dirac Hamiltonian

$$H_{SM}\psi^\omega_{Jm} = E^\omega_J\psi^\omega_{Jm} \tag{10}$$

which is assumed to be local and to commute with the total angular momentum ($\vec{J}$), parity (P), and time reversal (T) operators:

$$H_{SM} = \vec{\alpha}\cdot\vec{p} + \gamma^0[M+U_S(r)+\gamma^0U^0_V(r) - \gamma^0\gamma^r U^r_t(r)]. \tag{11}$$

The matrices $\vec{\alpha}$, γ^0, and $\gamma^r = \vec{\gamma}\cdot\hat{r}$ are the usual Dirac matrices. The functions U_S, U^0_V, and U^r_t are called the scalar, fourth component of the four vector, and tensor potentials, respectively, and result from self-consistent calculations[6] or more phenomenological shell-model parametrizations.[7]

The momentum space radial wave functions $F_l(q)$ and $G_{l'}(q)$ satisfy a pair of coupled integral equations:

$$\omega qF_l(q) = -\int_0^\infty q'^2dq'F_l(q')V^{(t)}_{ll'}(q',q) + (M+E)G_{l'}(q) + \int_0^\infty q'^2dq'G_{l'}(q')V^{(-)}_{l'}(q',q), \tag{12}$$

and

$$-\omega qG_{l'}(q) = \int_0^\infty q'^2dq'G_{l'}(q')V^{(t)}_{ll'}(q,q') + (M-E)F_l(q) + \int_0^\infty q'^2dq'F_l(q')V^{(+)}_l(q',q),$$

in which the momentum space potentials are defined as follows:

$$V^{(t)}_{ll'}(q,q') = \frac{2}{\pi}\int_0^\infty r^2dr\, j_l(qr)iU^r_t(r)j_{l'}(q'r),$$

$$V^{(-)}_{l'}(q,q') = \frac{2}{\pi}\int_0^\infty r^2dr\, j_{l'}(qr)[U_S(r)-U^0_V(r)] \times j_{l'}(q'r), \tag{13}$$

$$V^{+}_l(q,q') = \frac{2}{\pi}\int_0^\infty r^2dr\, j_l(qr)[U_S(r)+U^0_V(r)]j_l(q'r).$$

In the present work some of the corrections to the pionic stripping formalism outlined above are studied. The available experimental data near the pion production threshold together with nuclear wave functions resulting from relativistic self-consistent field calculations are then used in an attempt to constrain the appropriate relativistic pion-nucleon vertex operator. The question of whether pionic stripping is the dominant reaction mechanism is also investigated.

II. NONRELATIVISTIC APPROXIMATIONS

It is instructive to consider nonrelativistic approximations to Eqs. (6) and (7) in order to compare them with our exact results. The most obvious method of generating a nonrelativistic approximation is to use Eqs. (12) to eliminate the small component radial wave function $G_{l'}(q)$ in favor of the large component $F_l(q)$ which can then be viewed as a nonrelativistic wave function. For example, one might assume that the two potential terms in the first of Eqs. (12) are weak enough to neglect. Further, the eigenvalues E of the final nucleon in (p,π^+) are always within a few MeV of the nucleon mass M so that $E\approx M$ is not a bad approximation. This sequence of approximations was denoted the weak potential approximation (WPA) in Ref. 4 where it was shown that the WPA applied to the pseudovector coupling model [Eq. (7)] leads to the Galilean-invariant model, while the same approximation applied to the pseudoscalar coupling model [Eq. (6)] leads to an average between the Galilean-invariant and static models. The authors of Ref. 5 considered what was essentially the WPA for pseudoscalar coupling, but also showed that if one goes further and ignores the effect of the potential on the eigenvalue [$E\approx(q^2+M^2)^{1/2}$ instead of $E\approx M$], then the static model results. Clearly the old ambiguities[3] involving the nonrelativistic reduction of the pion-nucleon vertex can be easily generated within the present formalism.

It is not the purpose of the present work to resolve this ambiguity; rather we suggest that the relativistic formalism be used. Indeed, the relativistic self-consistent theory[6] generates potentials which fail to satisfy the WPA so that nonrelativistic approximations of any kind are suspect. This may be understood by considering the potential terms in the first of Eqs. (12). Clearly either a small potential of the type $V_{ll'}^{(t)}$ or a potential of the type $V_{l'}^{(-)}$ which is not everywhere small in comparison to $2M$ will invalidate the WPA. The simplest self-consistent calculations[6] achieve nuclear saturation via an attractive scalar potential (U_S) and a repulsive vector potential (U_V^0) such that the sum $U_S + U_V^0$ is a weakly attractive well (~50 MeV) while the difference $U_S - U_V^0$ is of the order of the nucleon mass. This particular combination of single-particle potentials also leads to spin-orbit splittings of nuclear single-particle states which are in qualitative agreement with experiments. Should the nuclear spin-orbit splittings not be generated through such a vector-scalar cancellation, the only other mechanism which can generate such splittings in a local potential framework is the tensor potential $V_{ll'}^{(t)}$. Either mechanism will lead to a failure of the WPA.

III. CORRECTIONS TO *PWIA*

In distorted wave theory the proton and pion wave functions $\psi_{\vec{p},\lambda}^{\mathrm{in}}$ and $\phi_{\vec{k}}^{\mathrm{out}}$ in Eq. (3) are generated from optical model potentials of the A and $A+1$ nuclei, respectively. Near the pion production threshold ($T_\pi < 50$ MeV) the pion optical model potential is rather weak so that the pion wave function $\phi_{\vec{k}}^{\mathrm{out}}$ differs from a plane wave mainly through the distortion of the Coulomb potential. Thus, it is reasonable to use a plane wave for the pion and multiply the resulting cross section by the Coulomb barrier-penetration factor[8]:

$$N_\pi = \frac{2\pi n_\pi}{e^{2\pi n_\pi} - 1}\,; \quad n_\pi = \frac{Z}{137}\,(1+m_\pi{}^2/k^2)^{1/2}\,. \tag{14}$$

While nonrelativistic optical model potentials for protons are well established for a variety of nuclei, the relativistic treatment of the optical model problem has just begun.[9] We are thus forced to treat the absorption and distortion of the incident proton by crude techniques. The Coulomb distortion is included via the corresponding Coulomb-barrier-penetration factor of the proton. The absorption due to the imaginary part of the proton optical potential will undoubtedly affect the large and small components of the wave function differently. Our lack of a model to account for this leads us to multiply the overall cross section by a factor of $\frac{1}{2}$ to crudely account for the absorption. The factor is suggested by a nonrelativistic analysis of the closely related (γ, p) reaction in this energy region.[10]

It remains to calculate the distortion of the incident proton due to the non-Coulomb part of the real optical model potential. For the present work, it is assumed that this real optical model potential is equivalent to the self-consistent shell-model potential. This assumption is made plausible by the fact that successful relativistic optical model analyses based on a combination of vector and scalar potentials do not suffer from the striking energy dependence of the real central potential parameters which characterize the nonrelativistic analyses.[9] Rather than calculate the full scattering wave function for our self-consistent potential, we have incorporated the main effect of the distorting potential by orthogonalizing the incident plane wave with respect to the shell model wave functions ψ_{Jm}^{ω} into which the nucleon is captured:

$$\psi_{\vec{p},\lambda}^{\mathrm{in}}(x) = e^{-ip_0 t}\left[U_\lambda(\vec{p})e^{i\vec{p}\cdot\vec{r}} - \sum_m A_m(\vec{p})\psi_{Jm}^{\omega}(\vec{r})\right], \tag{15}$$

where

$$A_m(\vec{p}) = \int \psi_{Jm}^{\dagger\omega}(\vec{r})U_\lambda(\vec{p})e^{i\vec{p}\cdot\vec{r}}d^3r\,. \tag{16}$$

This correction leads to a T matrix element:

$$\begin{aligned} T_{fi} = C\Big\{ & T_S \mathcal{Y}_{Jm}^{\dagger\omega}(\hat{q})\,\frac{\vec{\sigma}\cdot\vec{k}}{P_0+M}\,\chi_\lambda + T_{\mathrm{NS}}\mathcal{Y}_{Jm}^{\dagger\omega}(\hat{q})\chi_\lambda \\ & + \sum_{m'} N_D C(k)\left[F_l(p) + \omega\,\frac{P_0-M}{P_0+M}\,G_{l'}(p)\right] \\ & \times \mathcal{Y}_{Jm}^{\dagger\omega}(\hat{p})\chi_\lambda \mathcal{Y}_{1\,m-m'}^{*}(-\hat{k})\Big\}, \end{aligned} \tag{17}$$

where

$$N_D = \frac{(-)^{J+m}}{\sqrt{\pi}\,(2J+1)[(2J+2)(2J+1)(2J)]^{1/2}} \times \begin{pmatrix} J & J & 1 \\ -m & m' & m-m' \end{pmatrix} \tag{18}$$

and

$$C(k) = \int_0^\infty F_l(r)G_{l'}(r)j_1(kr)dr\,. \tag{19}$$

The evaluation of the cross section resulting from Eq. (17), involving standard angular momentum recoupling, is straightforward but lengthy. Our calculations for the $^{12}C(p,\pi^+)^{13}C$ reaction at 185 MeV for the $\frac{1}{2}^-$ ground state and $\frac{1}{2}^+$ first excited state of ^{13}C indicate that this distortion results in a less than 1% correction to the differential cross section over all angles. This surprisingly small correction results from two independent features of the last term in Eq. (17). First, the term is proportional to a factor whose dominant contribution is from the large component shell-model wave function F_l evaluated at the momentum of the incident proton p which is $\approx 3.1\ \mathrm{fm}^{-1}$, i.e., over twice the Fermi momentum of a nucleus. Second, the term is also proportional to $C(k)$, a term that vanishes identically at threshold ($k=0$). For an incident proton energy of 185 MeV, the pion momentum is about $0.5\ \mathrm{fm}^{-1}$ and since the first peak of the spherical Bessel function j_1 occurs at 2.1, the wave functions in the integrand of Eq. (19) are most heavily weighted near 4.2 fm, almost twice the rms radius of ^{12}C. Thus $C(k)$ is also very small in our region of interest and these distortion effects from the self-consistent potentials appear to be negligible for the incident proton.

IV. CALCULATIONS

One of the major criticisms of pionic stripping calculations to date is that too much freedom has been taken with the parameters of the single-particle wave functions. The spin-orbit potential and the well diffuseness have been varied in attempts to influence the high momentum components of the wave functions without sufficient regard to what these parameter variations do to the low momentum components. Experiments such as elastic electron scattering from nuclei contain much more information about nuclear wave functions than just their rms radii and it is imperative that we constrain our theories to fit this low momentum data before testing them in the high momentum region probed by (p, π^+). For this reason the present calculations are performed with a self-consistent nuclear model[11] for which the parameters were chosen specifically to fit the electromagnetic form factor of ^{40}Ca. A similarly good fit was also obtained for the form factor of ^{208}Pb with the same parameters. Unfortunately, this model does not extrapolate well to the region of light nuclei such as ^{12}C where much of the (p, π^+) data exist. Part of the problem arises from the increased importance of the spurious center of mass motion for lighter nuclei, other inaccuracies probably arise from a failure of self-consistent field methods to accurately describe collections of just a few nucleons. To partially compensate for this problem, we have scaled our wave functions in the ^{12}C region by a factor which results in the correct rms radius for ^{12}C when applied to our self-consistent calculation.

In Figs. 1–3 are shown the results of calculations for the ground and first two excited states of the $A+1$ system for the ^{12}C$(p, \pi^+)^{13}$C reaction. The most striking feature of these results (also noted in Ref. 4) is the large difference (~two orders of magnitude) between our relativistic results with pseudoscalar (γ^5) and pseudovector ($\gamma^5\gamma^\mu$) coupling for the πNN vertex. In Ref. 4 it was explained that this difference comes mainly from the different forms for the nonstatic form factor T_{NS} in pseudoscalar [Eq. (6)] and pseudovector [Eq. (7)] coupling. The T_{NS} of Eq. (7) is guaranteed to be appropriately small due to the factor $k_0/2M$ which multiplies its dominant term, while the T_{NS} in Eq. (6) can be similarly small only through a significant cancellation between its two terms. In fact, such a cancellation occurs for those potentials which satisfy the WPA but breaks down for the self-consistent potentials which change the relation between the large and small components F and G and lead to a complete domination of the cross section by the nonstatic part of the T matrix element for pseudoscalar coupling.

In a recent paper[12] Friar has shed light on this difference between pseudoscalar and pseudovector coupling in our relativistic PWIA (RPWIA) theory; he proves an equivalence theorem (under certain conditions) between the T matrix elements of pseudoscalar and pseudovector coupling. In the presence of strong single-particle potentials of the type $U_S(r)$ and $U_t^r(r)$, the equivalence theorem

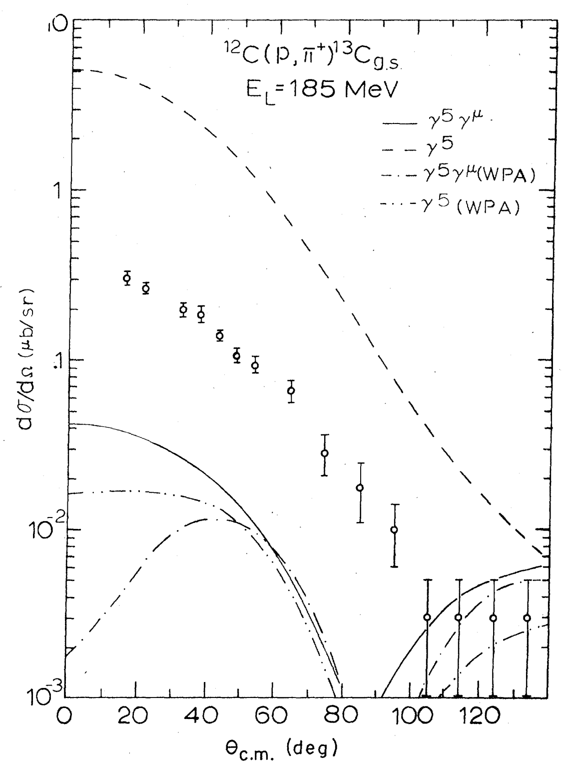

FIG. 1. Angular distributions for the reaction ^{12}C$(p, \pi^+)^{13}$C$_{g.s.}$. Theoretical curves were calculated with a $^1P_{1/2}$ wave function.

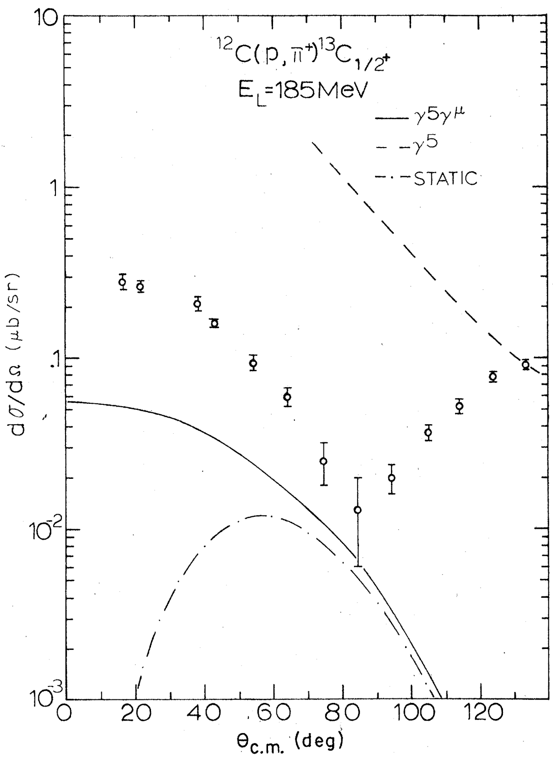

FIG. 2. Angular distributions for the reaction $^{12}C(p,\pi^+)^{13}C_{3.08\,\mathrm{MeV}}$. Theoretical curves were calculated with a $^2s_{1/2}$ wave function.

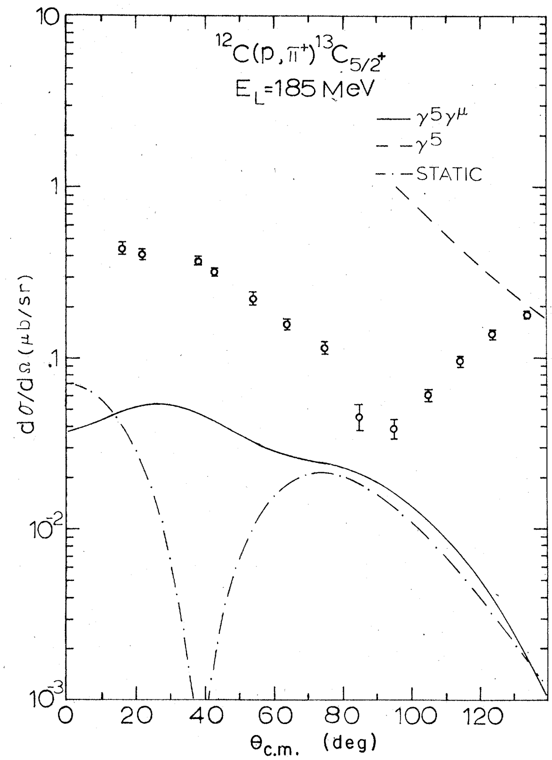

FIG. 3. Angular distributions for the reaction $^{12}C(p,\pi^+)^{13}C_{3.80\,\mathrm{MeV}}$. Theoretical curves were calculated with a $^1d_{5/2}$ wave function.

is badly broken. In fact,

$$T_{fi}(\gamma^5\gamma^\mu)=T_{fi}(\gamma^5)+ig\sqrt{2}\int\bar{\psi}^\omega_{Jm}(\vec{r})\,\frac{U_S(r)}{M}\,\gamma^5U_\lambda(\vec{p})e^{i\vec{q}\cdot\vec{r}}d^3r-g\sqrt{2}\int\bar{\psi}^\omega_{Jm}\,\frac{U_t^r(r)}{M}\,\gamma^0\gamma^rU_\lambda(\vec{p})e^{i\vec{q}\cdot\vec{r}}d^3r\,, \qquad (20)$$

where we have ignored the fact that the plane wave for the proton is not a solution of the Dirac Hamiltonian in the self-consistent field. Friar correctly surmised that the major difference between our results in Ref. 4 resulted from the last two (seagull) terms in Eq. (20) and not from our use of the plane wave.

This explanation of the difference between pseudoscalar and pseudovector coupling in terms of seagull diagrams in which the πNN and external potential vertices overlap (see Fig. 4) puts the theory of pion production or absorption on a single nucleon on an analogous footing with the low energy pion-nucleon scattering theory in which the partial conservation of axial vector current (PCAC) may be incorporated in terms of pure pseudovector coupling theory or via a pseudoscalar coupling theory which is augmented by meson nonlinearities like those of the σ model.[13] This analogy strongly suggests that we adopt the pure pseudovector coupling for our relativistic pionic stripping theory.

If we can put any faith in our self-consistent

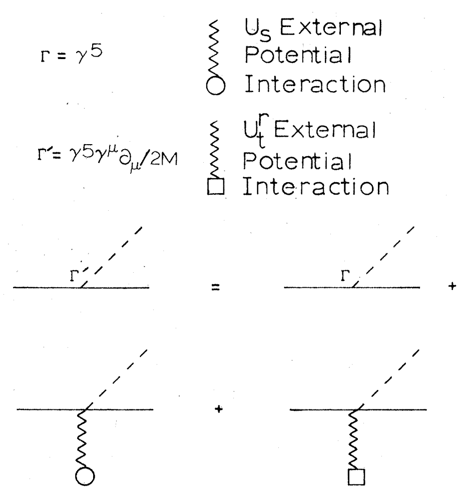

FIG. 4. Graphical expression for the equivalence theorem for pseudovector and pseudoscalar coupling of a pion to a nucleon in an external field.

wave functions for ^{12}C, then the data[14] in Figs. 1–3 clearly rule out a pure pseudoscalar vertex for the πNN coupling. The uniform underestimate of the data by the pseudovector vertex probably reflects a failure of the self-consistent model, but could also indicate a need for a mixture of pseudoscalar and pseudovector couplings or a need for a different reaction mechanism.

The results for the nonrelativistic WPA approximation are shown in Fig. 1. The pseudovector model in WPA may be viewed as the Galilean-invariant model where the large component of our relativistic wave function plays the role of a nonrelativistic wave function. The pseudoscalar model in WPA represents an average between the static and Galilean-invariant models under the same assumption. The most striking effect of the WPA is that most of the seagull terms which break the equivalence theorem are eliminated. There still remains a substantial disagreement between the two models at forward angles. The pseudovector result in WPA displays the prominent dip at forward angles which also characterizes the Galilean-invariant model.[1] The pseudoscalar result in WPA lacks this dip and represents a better approximation (at forward angles) to the relativistic pseudovector result than does the pseudovector WPA result. The pure static approximation (not shown), obtained by retaining only the first term of Eq. (4), gives an almost perfect fit to the relativistic pseudovector result below 60° in Fig. 1. This is a purely accidental fit as is shown in Figs. 2 and 3 where the pure static result fails badly in reproducing the relativistic pseudovector result at forward angles.

Since the self-consistent model is not particularly appropriate for light nuclei, it is more interesting to study (p, π^+) data on medium and heavy nuclei. The data[15] for the reaction $^{40}Ca(p,\pi^+)^{41}Ca_{g.s.}$ at 185 MeV are particularly appropriate since this is the region for which the model parameters were chosen. As can be seen from Fig. 5, our relativistic pseudovector result yields a reasonably good fit to the data below 110°. This fit alone is not sufficient evidence to justify the acceptance of the relativistic pionic stripping mechanism with pseudovector coupling as the dominant mechanism for proton induced pion production in this energy region. Angular distributions at other energies and for other medium to heavy (^{208}Pb for example) nuclei would be very useful for this purpose.

Some authors[5,8] have suggested that the energy dependence near threshold and at forward angles (where q changes very slowly) may be particularly useful for establishing the reaction mechanism for (p, π^+). Recently, data[16] have been taken for constant $q_{c.m.} \approx 2.43$ fm^{-1} at $E = 152$–164 MeV for the

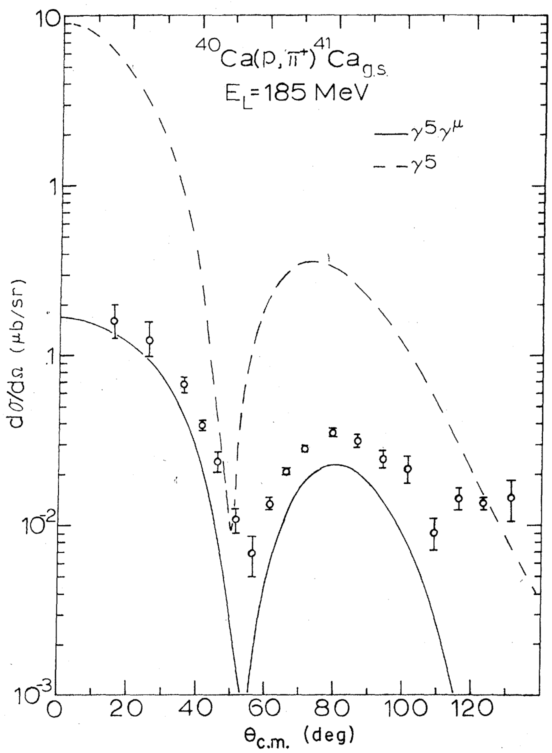

FIG. 5. Angular distributions for the reaction $^{40}Ca(p,\pi^+)^{41}Ca_{g.s.}$. Theoretical curves were calculated with a $^1f_{7/2}$ wave function.

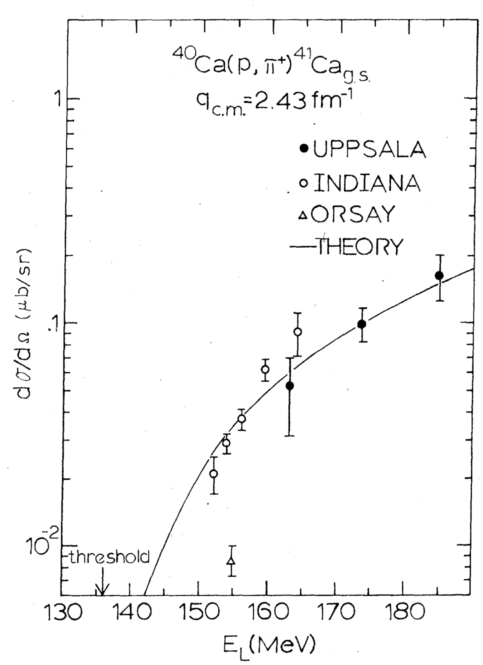

FIG. 6. Energy dependence of the $^{40}Ca(p,\pi^+)^{41}Ca_{g.s.}$ cross section at $q_{c.m.} = 2.43$ fm^{-1} near pion production threshold.

$^{40}Ca(p,\pi^+)^{41}Ca_{g.s.}$ reaction. Unfortunately, they are in disagreement with older data[17] at $E = 154$ MeV so that the experimental situation is somewhat unclear at the present time. Nevertheless we show in Fig. 6 our results for the energy dependence of $^{40}Ca(p,\pi^+)^{41}Ca$ at $q_{c.m.} = 2.43\ fm^{-1}$ with the pseudovector coupling model. Our results agree with the data of Ref. 16 at $E = 155$ MeV where the datum of Ref. 17 is a factor of 2 lower. Our slope does not fit the data of Ref. 17 inasmuch as our results are also in agreement with Ref. 15 at $E = 185$ MeV. It should be pointed out however, that most of the variation in our calculation for $d\sigma/d\Omega$ in Fig. 6 results from the Coulomb-barrier-penetration factor of the pion [Eq. (14)] which will be present no matter what the reaction mechanism is.

V. SUMMARY AND CONCLUSIONS

Some crude estimates of the effects of absorption and distortion have been made for the relativistic plane wave impulse approximation for pionic stripping. The most important of these corrections is the Coulomb distortion of the pion which significantly depresses the cross section near threshold. A surprisingly unimportant correction is that due to the distortion of the proton by its strong interaction with the nucleus. This effect was estimated by orthogonalizing the plane wave with respect to the bound state into which the nucleon is captured and was found to be negligible. This result may be a quirk of the orthogonalized plane wave method, yet it is consistent with nonrelativistic optical model analyses which suggest that the real potential becomes repulsive at about 200 MeV.[18]

The corrected plane wave theory in conjunction with a relativistic self-consistent calculation designed to fit the electromagnetic form factors of medium to heavy nuclei is used to generate angular distributions for the $^{12}C(p,\pi^+)^{13}C$ and $^{40}Ca(p,\pi^+)^{41}Ca_{g.s.}$ reactions at 185 MeV. The experimental data definitely rule out the pure pseudoscalar coupling for the πNN vertex which constantly gives a cross section at least one order of magnitude too large. The pure pseudovector coupling result does not violate the upper bound set by the cross section and even gives a good fit to the cross section for the ^{40}Ca nucleus. These results are analogous to the PCAC results for pion-nucleon scattering where the pseudovector coupling is also preferred over the pure pseudoscalar coupling.

The ^{40}Ca results are also consistent with the somewhat unclear experimental picture of the energy dependence of (p,π^+) cross sections at forward angles near threshold. This behavior is basically due to the Coulomb distortion of the pion wave function and cannot be taken as a success of pionic stripping over other mechanisms.

In conclusion we have removed the major ambiguities in calculations of the pionic stripping mechanism and shown that pionic stripping and accurate nuclear wave functions are capable of explaining (p,π^+) cross sections in a limited energy domain. As more experimental data become available on medium and heavy nuclei, the question can be finally resolved as to whether pionic stripping is the dominant mechanism for (p,π^+) reactions.

Note added: Since submitting this manuscript for publication, we have received two data points at energies of 163 and 173 MeV for the forward angle $^{40}Ca(p,\pi^+)^{41}Ca$ angular distribution from the Uppsala group. These data fall nicely along our theoretical curve and have been added to Fig. 6. We wish to thank Professor B. Höistad for providing us with these data prior to their publication.

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