Axial Currents in Nuclei

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It is shown that the time component of axial currents in nuclei can be given a simple description in terms of a one-pion exchange diagram with its structure constrained by a soft-pion theorem. We suggest the directional correlation measurements in nuclear β decay as a possible means to test the theory.

By far the best evidence for mesonic currents in nuclei is found in the radiative capture of thermal neutron by proton $n+p\to d+\gamma$,[1] and here, soft-pion theorems (or in general chiral-symmetry arguments) play a crucial role in providing a model-independent description of the dominant one-pion exchange electromagnetic current.[2] No such test has been found so far for the axial current. The main reasons for this glaring difference are twofold: Firstly, weak processes involving deuteron and low-momentum transfer are not easily measurable; and secondly, in all cases that have been considered, the effect of mesonic currents was looked for in the *space* component of the axial current where soft-pion theorems are rendered powerless by the important role of the isobar $\Delta(1232)$ and by short-range correlations.[3] In this Letter we argue that the situation is drastically different with the *time* component (e.g., axial density) for which chiral-symmetry arguments are as relevant and powerful as in the electromagnetic case and we then suggest an experimental means to test this assertion.

One important requirement that should be met by any successful theory of meson-exchange currents is *that one-pion exchange process be dominant over other shorter-ranged processes such as multipion or heavier-meson exchanges.* Were this not so, one could not hope to have a simple description of meson-exchange phenomena: The theory would be complicated and most likely unreliable because of our incomplete knowledge of nuclear interactions at short distances. It is instructive to see how nicely this requirement is met by the isovector $M1$ operator governing the radiative neutron capture. Consider the one-pion exchange graph Fig. 1. Because of short-range correlations, only the pion with small four-momentum contributes significantly to the matrix element and heavy-meson exchanges are suppressed, so that the important quantity to consider is the amplitude M(current + nucleon → pion + nucleon) in the situation where the pion is soft.[2,3] The soft-pion theorem (combined with current algebra) tells us that for the isovector vector current V_μ^3, the amplitude M^V is described, apart from the standard nucleon Born graphs (with gradient πNN coupling), by the two graphs of Fig. 2(a). These are the graphs that contribute to the exchange current. Now the first (seagull) graph—more important of the two—is the matrix element of the operator $(1/F_\pi)\epsilon_{3\alpha\beta}A_\mu^\beta$, where F_π is the pion decay constant ($\simeq 0.67m_\pi$) and A_μ^β is the axial current with charge index $\beta=1,2,3$. The important point here is that the time ($\mu=0$) and space ($\mu=1,2,3$) components have

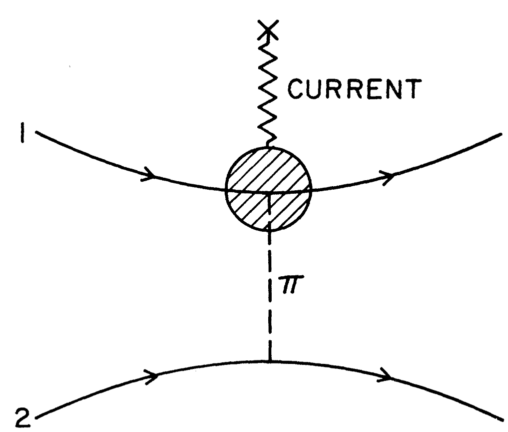

FIG. 1. Two-body pion-exchange current.

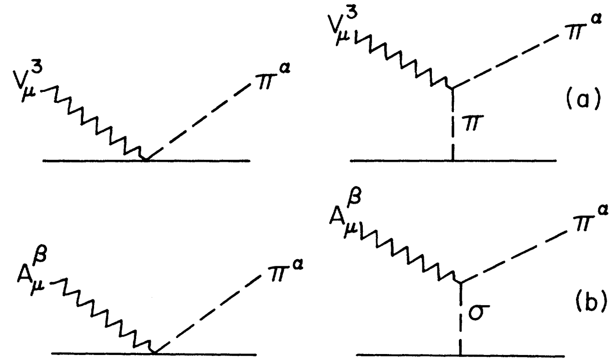

FIG. 2. Soft-pion amplitudes for the blob of Fig. 1 that are *not* included in the impulse approximation; (a) for the vector current; (b) for the axial current.

quite different magnitudes: The time component is $\sim O(p/M)$ and the space component is $\sim (g_A/F_\pi)\epsilon_{3\alpha\beta}(\tau^\beta/2)\vec{\sigma} \sim O(1)$, where $\vec{p}$ is the nucleon momentum, and M the nucleon mass. In contrast, the single-particle operator goes like $\frac{1}{2}\tau^3\gamma_\mu$, so that nonrelativistically it is of $O(1)$ for $\mu=0$ and $O(p/M)$ for space components. Thus for the space components, the vertex M^V is *intrinsically* enhanced relative to the single-particle operator.[4] Although this naive power counting does not tell us *a priori* how large two-body operators can be relative to single-particle operators, it nevertheless suggests the possibility that the space components of the meson-exchange current be dominated by the one-pion exchange mechanism with the excitation of Δ playing a relatively minor role. This has been confirmed by the remarkable success of the calculations[1]; the short-range operators due to heavy-meson exchanges, etc., indeed contribute negligibly. It also follows from such arguments that due to an intrinsic suppression of M^V for $\mu=0$, there is no good reason to believe that other graphs involving Δ's and heavy mesons or multipions should not be as important as the one-pion exchange contribution for the description of isovector charge form factors in nuclei.

We make similar arguments for the axial current. Chiral-symmetry considerations suggest that the amplitude M^A for the axial current can be reliably given by the graphs of Fig. 2(b) (apart from the nucleon Born graphs with gradient πNN coupling) *provided that a soft-pion situation holds.* The σ, chiral partner of π, is supposed to be very massive, so that for the current with small momentum the σ-exchange term can be ignored. Thus the sole contribution comes from the seagull term given by the matrix element of the vector current $(1/F_\pi)\epsilon_{\beta\alpha\gamma}V_\mu{}^\gamma$. Its nonrelativistic structure is $F_\pi{}^{-1}\epsilon_{\beta\alpha\gamma}\frac{1}{2}\tau_\gamma \sim O(1)$ for $\mu=0$ and $O(p/M)$ for $\mu=1,2,3$. In contrast, the single-particle axial-current operator $ig_A\frac{1}{2}\tau_\beta\gamma_\mu\gamma_5$ goes as $O(p/M)$ for $\mu=0$ and $O(1)$ for $\mu=1,2,3$. Therefore it is now the *time* component in which the mesonic amplitude M^A is intrinsically enhanced relative to single-particle processes and hence that offers, barring accidental cancellations, a *prima facie* possibility to be dominated by the one-pion exchange mechanism. A rough estimate shows that the Δ contribution to this component is suppressed by a factor of as much as 10^2. Similarly there is no reason to believe that one-pion-exchange currents would provide a reliable description of the space component of the axial current, i.e., meson-exchange effects in the Gamov-Teller matrix element. We suggest this as a reason why one-pion-exchange current has failed so far in accounting for the quenching of g_A in light nuclei.[5]

Given the structure of M, it is straightforward to derive the effective nuclear operators.[2] The results are

$$A_0{}^\pm(\vec{x}) = A_0{}^{(1)\pm}(\vec{x}) + A_0{}^{(2)\pm}(\vec{x}), \tag{1a}$$

$$A_0{}^{(1)\pm}(\vec{x}) = -g_A\sum_i\tau_i{}^\pm(\vec{\sigma}_i\cdot\vec{p}_i/M)\delta(\vec{x}-\vec{x}_i) + g_A\sum_i\tau_i{}^\pm(\vec{\sigma}_i\cdot\vec{k}/2M)\delta(\vec{x}-\vec{x}_i), \tag{1b}$$

$$A_0{}^{(2)\pm}(\vec{x}) = \frac{m_\pi{}^2 g}{8\pi M F_\pi}\sum_{i<j}(\tau_i\times\tau_j)^\pm[\vec{\sigma}_i\cdot\hat{r}\delta(\vec{x}-\vec{x}_j) + \vec{\sigma}_j\cdot\hat{r}\delta(\vec{x}-\vec{x}_i)]\frac{\exp(-m_\pi r)}{m_\pi r}\left(1+\frac{1}{m_\pi r}\right), \tag{1c}$$

where $\vec{r}=\vec{x}_i-\vec{x}_j$, $\vec{k}$ is the momentum carried by the axial current, and $\vec{p}$ the initial momentum of the nucleon making the transition. Now how does one "see" this mesonic contribution $A_0{}^{(2)}$? Total rates, whether β decay or muon capture, are not useful since they are usually dominated by the space component of the axial current with the important exception of 0^+-0^- transitions[6] which are a potential source of information on the time component. Another favorable case would be a weak process involving large energy transfer and small momentum transfer, so that single-particle processes are kinematically suppressed.[7] However, it is probably difficult to separate out the time component of the axial current from the rest. The s-wave pion absorption might be useful but it is not clear how accurately one can extract the relevant matrix elements.

One place where the effect of $A_0^{(2)}$ can be studied is the angular correlation measurements in β decay.[8] These are difficult experiments but offer a clearer interpretation. As an illustration,[9] consider the celebrated mass-12 β decays: $^{12}B(1^+;\text{ground}) \to {}^{12}C(0^+;\text{ground})$ and $^{12}N(1^+;\text{ground}) \to {}^{12}C(0^+;\text{ground})$. The transition matrix elements can be suitably parametrized in terms of nuclear form factors[10]:

$$\langle 0^+;p_1|V_\alpha^\pm(0)|1^+;p_2\rangle = \epsilon_{\alpha\beta\gamma\delta}k_\beta\xi_\gamma P_\delta(F_M/4AM^2), \quad (2)$$

$$\langle 0^+;p_1|A_\alpha^\pm(0)|1^+;p_2\rangle = \xi_\alpha(F_A^{(1)} \pm F_A^{(2)}) + k_\alpha(\xi\cdot k)(F_p^{(1)} \pm F_p^{(2)}) + (P_\alpha/2M)(\xi\cdot k)(2AM)^{-1}(F_T^{(1)} \pm F_T^{(2)}), \quad (3)$$

where $k=p_1-p_2$, $P=p_1+p_2$, ξ is a spin-1 polarization vector, and A the mass number. The form factors that change sign for $\beta^\mp$ transitions belong to the second-class current while all others to the first class.[11] The angular correlation between the initial spin direction and the β-particle momentum, that is measured in the form of polarization and alignment, determines the asymmetry parameters $\alpha_\mp$ for $\beta^\mp$ decays[12] which in terms of nuclear form factors have the form (to lowest order in k)

$$\alpha_\mp \simeq \tfrac{2}{3}[\pm(F_M - F_T^{(2)}) - F_T^{(1)}]/2MF_A^{(1)}; \quad (4)$$

while the difference between the $\alpha_\pm$ bears information on the second-class current form factor,[8] the sum singles out the quantity of our interest:

$$S \equiv \alpha_+ + \alpha_- \simeq -\tfrac{4}{3}F_T^{(1)}/2MF_A^{(1)}. \quad (5)$$

The form factor $F_T^{(1)}$ is dominated by A_0, the contribution from $\vec{A}$ being negligible. Similar arguments apply to other mirror decays.[9]

The test of our theory in the mass-12 system is, at the present, made difficult for two reasons. First, the experimental situation is not yet clear as the available data do not concur[12]: $S_{\text{exp}} = (0.94 \pm 0.85)/M$ from polarization and $S_{\text{exp}} = -(2.4 \pm 0.6)/M$ from alignment. Second, we are not sure how much confidence one can have in the calculated value of the single-particle matrix element $[A_0^{(1)}]$. This contrasts with the case of the process $n+p \to d+\gamma$ in which the single-particle $M1$ matrix element is very accurately known. Clearly a definitive confrontation of our theory with experiments would require resolving both of these problems. To have a rough idea of the magnitudes involved, we have estimated the matrix elements using the Cohen-Kurath intermediate-coupling wave functions[13] obtaining $S_{\text{exchange}}/S_{\text{impulse}} \approx 0.4$ with an uncertainty of roughly 30%. Even granting a larger uncertainty, the mesonic effect is seen to be quite substantial supporting the qualitative arguments given above.[14]

In conclusion, we should emphasize the need for accurate values of $\alpha_\pm$; in view of the important issue on the meson degrees of freedom in nuclei, experimental efforts to improve on the available data are urged independently of the second-class-current problem.

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[4]Note that the single-particle $M1$ operator remains dominant partly because of the "accidental" enhancement factor, the isovector magnetic moment of the nucleon, which multiplies the operator. There is no such enhancement factor in the case of the axial current.

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[6]Our preliminary estimate indicates that the mesonic contribution may easily modify the decay rate of $^{16}N(0^-) \to {}^{16}O(0^+) + e^+ + \nu$ (and its inverse muon absorption rate) by a factor of 2 or more. We hope to return to this potentially interesting process in a later paper.

[7]E.g., see J. Bernabéu, T. E. O. Ericson, and C. Jarlskog, Phys. Lett. 69B, 161 (1977).

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[14]The calculated value comes out to be $S_{theo} = (-3.4 \sim -3.0)M^{-1}$ consistent with the alignment result. In view of the uncertainties mentioned above, the quality of agreement or disagreement cannot be assessed at present.

Comparison of the Reactions $^{12}C(\pi^+,p)^{11}C$ and $^{12}C(p,d)^{11}C$ near the Same Momentum Transfer

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Observed transitions in the reaction $^{12}C(\pi^+,p)^{11}C$ at 49.3 MeV confirm the importance of multistep processes in the reaction mechanism. The relative intensity of the transitions is nearly the same as in the reaction $^{12}C(p,d)C^{11}$ at 700 MeV, except for an apparent isospin selection rule. Transitions to low-lying states in ^{10}C via the reaction $^{12}C(\pi^+,d)$ have also been observed.

Angular distributions have been measured for several transitions in the reaction $^{12}C(\pi^+,p)^{11}C$ at $T_\pi = 49.3$ MeV and have been compared to the $^{12}C(p,d)^{11}C$ data at $T_p = 700$ MeV.[1] Both reactions are classified as neutron-pickup reactions, but the similarity in their reaction mechanism is not adequately known. The comparison is made at the same momentum transfer, between 2.3 and 3.2 fm^{-1}, for both reactions. The probability is small that a bound neutron in the target has so large a momentum. Therefore, rather than a single-step pickup process it is more likely that the momentum transfer occurs in multiple steps involving off-shell particles. For example, the incoming projectile experiencing the strong nuclear field may acquire an off-shell momentum ($p^2 \neq E^2 - m^2$). This effect on the incoming and outgoing projectile is included in a standard distorted-wave Born-approximation (DWBA) calculation using optical potentials.[2] One goal of this experiment is to add the (π^+,p) data to the extensive study[3,4] of the pion optical potential near $T_\pi = 50$ MeV. Various optical potentials which explain the available elastic- and inelastic-pion-scattering data have very different off-shell behavior and predict widely different cross sections for the (π^+,p) reactions.[2]

The large momentum transfer is also likely to result in multiple changes in the structure of the nucleus. There are states of ^{11}C that are predominantly a neutron hole weakly coupled to the collective 2^+ (4.44 MeV) or 3^- (9.63 MeV) states of ^{12}C. A reasonable model of the transitions to these states is a second-order DWBA calculation coupling the inelastic (collective) channels to the ^{11}C neutron-pickup channels. Such a model has been used successfully[1,5] for the reaction $^{12}C(p,d)^{11}C$ at $T_p = 700$ MeV. An alternative approach[6] to a coupled-channels calculation for the (π^+,p) reaction is to use an effective pion-two-nucleon interaction, $H_{NN\pi}$, that includes single-rescattering contributions. The latter approach is also applicable to two-step transtions which do not have strong soupling to well-defined intermediate states, such as the two-step transition to the $T = \frac{3}{2}$, 12.5-MeV state of ^{11}C (see below). The study of these various transitions will enhance our understanding of the (π^+,p) reaction process.

The experiment was performed at the LEP channel at the Clinton P. Anderson Meson Physics Facility (LAMPF), using a stack of eight intrinsic germanium crystals.[7] Approximately 5×10^6 pions per second with 50 MeV kinetic energy bombarded natural carbon targets with thicknesses between 0.35 and 0.7 g/cm^2. The proton spectra, using proton identification from the range-energy relationship, are about 99.9% pure. Three multiwire proportional chambers were placed in front of the detector to trace the trajectory of each particle. The absolute cross section was normalized to the known[8] $d(\pi^+,p)p$ cross section

