Tests for Second-Class Currents in Nuclear Beta Decay*

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Semileptonic interactions with anomalous, i.e., "second-class," G-parity properities are uniquely correlated with certain characteristic terms in the amplitude for β decay between members of a common isotopic multiplet. The observational effects associated with second-class currents are worked out for several features of the β-decay spectrum as well as for β-γ correlation phenomena.

I. INTRODUCTION

A recent survey by Wilkinson and Alburger on rates of mirror pairs of β decays has stimulated renewed interest in the G-parity properties of strangeness-conserving semileptonic interactions.[1] A systematic trend in the data on ft values suggests, according to one possible interpretation, that contributions from hitherto undetected interactions with anomalous G-parity properties are being seen in allowed β decays. The G-parity classification scheme was first proposed by Weinberg, who noted that the dominant effects visible in all known $\Delta S=0$ semileptonic processes arise from "first-class" interactions.[2] Second-class interactions, even if they exist in nature, are expected to be kinematically suppressed in their contributions to nuclear β decay; and in the usual vector–axial-vector picture of current-current interactions, second-class contributions are wholly forbidden in such processes as $\pi^\pm \to \mu^\pm + \nu$ and $\pi^\pm \to \pi^0 + e^\pm + \nu$ decay – at any rate, at the level where electromagnetic violations of isospin conservation can be ignored. The best prospect for unambiguous detection of second-class interactions probably lies in the experimental study of high-energy quasielastic neutrino reactions, such as $\nu_\mu + n \to p + \mu^-$. However, in this note we want to elaborate on some tests, whose outlines already appear in the paper by Weinberg, which bear on possible second-class contributions to a particular sort of nuclear β decay process.

We restrict ourselves to the picture in which the basic interactions are supposed to couple the lepton current to weak hadronic currents of the vector and axial-vector sort, V_μ and A_μ. It is, of course, only the strangeness-conserving parts of these weak hadron currents that are of interest here and for which it is appropriate to make classifications with respect to G-parity properties. In Weinberg's terminology it is the parts of V_μ and A_μ which are, respectively, even and odd under the G-parity transformation which are "first class." The parts with reversed G-parity properties are relegated to the "second class." The simplest kind of test which bears on the existence of second-class interactions rests on the following argument: Consider a mirror pair of allowed β-decay processes, $\alpha \to \beta + e^- + \bar{\nu}$, $\tilde{\alpha} \to \tilde{\beta} + e^+ + \nu$. Here $\tilde{\alpha}$ is the mirror of α ($\tilde{\beta}$ the mirror of β), where $\tilde{\alpha}$ is obtained from α by a rotation through angle π about the 2 axis in isospin space (e.g., $\tilde{p}=n$, $\tilde{\Sigma}^+=\Sigma^-$, $\tilde{\Lambda}=\Lambda$). For the moment suppose that α and β belong to different isotopic multiplets. Insofar as electromagnetic violations of isospin conservation can be ignored, $\tilde{\alpha}$ and α would have the same masses (as would $\tilde{\beta}$ and β); and in the absence of second-class interactions the net rates for the two mirror processes would be identical. Of course, electromagnetic effects cannot be safely ignored, especially with respect to the mass shift and, hence, phase-volume differences which they induce; and with respect also to final-state Coulomb interactions. These dominant electromagnetic effects are, however, allowed for in the characterization of rate in terms of the ft parameter, which incorporates mass-shift and final-state Coulomb-interaction differences between the mirror decays, in leading approximation. It is therefore more nearly the ft values for the mirror decays that should be equal in the absence of second-class interactions. Nevertheless, since second-class interactions are expected, for either mirror decay process, to make contributions of order E_0/m (E_0 is the maximum electron energy, m the nucleon mass), one has to allow for possible first-class contributions in the same order – something which is not accomplished with the ft characterization.

A less ambiguous kind of test for second-class currents, one which we take up here, deals with allowed β decays in which the parent and daughter nuclei belong to a *common* isotopic multiplet (the most primitive example is neutron β decay itself). Here it is not necessary to compare two distinct mirror processes. Rather, for the given process one observes that certain characteristic terms in the amplitude can arise only from second-class currents and are forbidden, conversely, in the absence of second-class currents. It is then only a matter of seeing whether and how these terms can be discerned in various correlation and spectrum properties of the decay process.

We shall make certain approximations and assumptions in the present analysis and it is best to note these at the outset. (1) With two qualifications to be mentioned shortly, electromagnetic effects are ignored; so we have to assume that parent and daughter nuclei are sufficiently pure isotopically. Their electromagnetically induced mass difference, which is what makes the decay energetically possible, is, of course, allowed for in the kinematics through use of the physical masses. Dominant final-state Coulomb effects are also allowed for in the standard Fermi approximation. (2) We adopt the conserved-vector-current (CVC) hypothesis, operationally in the sense of ignoring the possibility of second-class *vector* currents.[3] In more detail, as we will note in the appropriate place, this hypothesis also fixes the size of the Fermi matrix element and relates a certain weak-magnetism term to magnetic moments. But these are separate issues. The main point is that only the question of possible second-class *axial* currents is being raised here. (3) We systematically ignore all effects beyond first order in the "recoil" parameter of smallness, E/M, where E is the electron energy, M the nuclear mass.

For definiteness we will discuss the case of (negative) electron decay; modifications appropriate to positron decay appear in the final formulas.[4] The reaction to be considered is

$$\alpha \to \beta + e^- + \nu ,$$

where the parent and daughter nuclei belong to a *common* isotopic multiplet. Let p_1, p_2, p, k denote, respectively, the four momenta of parent nucleus, daughter nucleus, electron, and neutrino. The parent and daughter masses are M_1 and M_2. Define

$$P = p_1 + p_2; \qquad q = p_1 - p_2 = p + k ,$$

$$M = \tfrac{1}{2}(M_1 + M_2); \qquad \Delta = M_1 - M_2 .$$

The β-decay amplitude is given by

$$T = (G/\sqrt{2})\cos\theta_C \langle \beta | V_\mu + A_\mu | d \rangle\, l_\mu , \tag{1}$$

where G is the usual weak-coupling constant ($G m_p^2 \approx 10^{-5}$) θ_C the Cabibbo angle, and l_μ the matrix element of the lepton current:

$$l_\mu = \bar{u}(p)\gamma_\mu (1 + \gamma_5) v(k) .$$

Let E and $\vec{p}$ denote the energy and three momentum of the electron and let $\hat{k}$ be a unit vector in the direction of the neutrino momentum, all of these quantities referring to the rest frame of the parent nucleus. The maximum possible electron energy is E_0,

$$E_0 = \Delta(1 + m_e^2/2M\Delta)/(1 + \Delta/2M) ,$$

where m_e is the electron mass. To first order in E/M the decay spectrum is given by

$$dw = \frac{|T|^2}{(2\pi)^5}\left(1 + \frac{3E - E_0 - 3\vec{p}\cdot\hat{k}}{M}\right)(E_0 - E)^2 pE\,dE\,d\Omega_e\,d\Omega_\nu . \tag{2}$$

To begin with a familiar situation, let us first specialize to the case of neutron β decay. Here, to first order in recoil effects and with second-class interactions allowed for only in the axial-vector current, we have

$$\langle \beta | V_\mu + A_\mu | \alpha \rangle l_\mu = \bar{u}(p_2)\left\{ -i\frac{g_V}{2M} P\cdot l + \frac{\mu}{2M}\sigma_{\mu\nu} l_\mu q_\nu + g_A \gamma_\mu \gamma_5 l_\mu + \frac{g_{II}}{2M}\sigma_{\mu\nu} l_\mu q_\nu \gamma_5 \right\} u(p_1). \tag{3}$$

The various form factors, which in general depend on the momentum transfer variable q^2, are to be evaluated at $q^2 = 0$. Here g_V is the vector coupling constant ($g_V = 1$ in the Cabibbo model); g_A is the axial-vector constant (in our convention $g_A \approx 1.2$ is positive); μ is a weak-magnetism coefficient given, in the CVC hypothesis, by $\mu = (\mu_p - \mu_n)$, where μ_p and μ_n are the *total* magnetic moments of proton and neutron. Finally, g_{II} is coefficient of a term which can arise only from a second-class axial-vector current. It is the question of testing for the presence of such a term that is our main concern here. The expression appearing in Eq. (3) can be reduced to a form which involves the matrix elements of Pauli spin operators acting in the space of nucleon spins. It is this version of the amplitude that we now want to generalize to arbitrary β decay.

We are dealing with the case where parent and daughter nuclei belong to a common isotopic multiplet, so that they have the same spin quantum number j. Let m and m' be the initial and final components of nuclear spin along some axis of quantization, and let $J_i (i = 1, 2, 3)$ be the components of the angular momentum operator acting in the space of nuclear spins. To first order in recoil quantities, and with allowance made for second-class contributions only from the axial-vector current, we can write the amplitudes in the following way[5]:

$$\langle\beta|V_\mu+A_\mu|\alpha\rangle l_\mu=-\frac{i}{2M}aP\cdot l\langle jm'|1|jm\rangle$$
$$-\frac{i}{2M}[j(j+1)]^{-1/2}\langle jm'|\tfrac{1}{2}[J_i,J_j]|jm\rangle$$
$$\times(2bl_iq_j+c\epsilon_{ij\lambda\nu}l_\lambda P_\nu-d\epsilon_{ij\lambda\nu}l_\lambda q_\nu)\ ,\qquad(4)$$

where repeated Latin indices are summed from 1 to 3, Greek indices from 1 to 4. Here, using standard conventions,

$$a=g_VM_F\ ,$$
$$c=g_AM_{GT}\ ,$$

where M_F and M_{GT} are the Fermi and Gamow-Teller matrix elements. In the CVC hypothesis, if I and I_3 are the quantum numbers describing total and the third component of isospin for the parent nucleus, we have

$$M_F=[(I-I_3)(I+I_3+1)]^{1/2}\ .$$

The term with the coefficient b is what has come to be described as a weak-magnetism effect. According to the CVC hypothesis we would have

$$b=A[(j+1)/j]^{1/2}M_F\mu\ .\qquad(5)$$

Here A is the mass number and μ is the isovector contribution to the total magnetic moment, measured in units of the *proton* magneton. Namely, let I_3 and I_3' be the quantum numbers for the third component of the isospin of any two members of the nuclear multiplet, and let $\mu(I_3)$ and $\mu(I_3')$ be the corresponding total magnetic moments. Then (in the convention where I_3 for a proton is $+\frac{1}{2}$)

$$\mu=\frac{\mu(I_3)-\mu(I_3')}{I_3-I_3'}\ .$$

Finally, the last term in Eq. (4) represents a contribution which can arise only from an "anomalous" axial-vector current. Namely, the coefficient d must vanish if the axial-vector current has unit isospin and odd G parity; or, more generally, if it has either odd isospin and odd G parity, or even isospin and even G parity. Of course, for β decay between members of an isotopic doublet only isovector currents can contribute; so the discovery of a nonvanishing d coefficient would clearly signal the existence of a second-class axial-vector current. For β decay between members of an isotopic triplet or quartet the d coefficient could alternatively arise from a first-class current with isospin two. However, since the primitive nucleon β-decay couplings involve isovector currents only, we may expect that effects coming from currents with higher isospin are suppressed in nuclear decays even where such currents can formally contribute. Moreover, some of the mirror decays which display ft discrepancies can involve only isovector currents. For these reasons we will refer to the d coefficient as bearing on the existence of second-class currents although the qualification mentioned above must be kept in mind. In any case it bears on the existence of possible currents which would certainly be anomalous in the present-day picture of weak interactions. In contrast to the situation regarding the coefficient b, we have no physical ideas which serve to relate the coefficient d to other physical parameters. Nevertheless, to provide a notational parallel, let us define a parameter μ_{II} according to

$$d\equiv A[(j+1)/j]^{1/2}M_F\mu_{II}\ .\qquad(6)$$

If the second- and first-class axial currents have, in some ill-defined sense, comparable strengths, one might expect that μ_{II} would be roughly comparable with μ.

We now want to turn to several spectral and correlation phenomena in β decay, parametrizing these effects in terms of the coefficients a, b, c, d. Quantitative determination of a and b can be used, via Eqs. (3) and (5), to test the CVC hypothesis; and experimental evidence for a nonvanishing coefficient d would signal the existence of second-class axial-vector currents. In the next section we report the spectrum, in electron energy and electron and neutrino directions, for β decay of unpolarized nuclei.[6] In Sec. III we report the spectrum, integrated over neutrino directions, for decay of polarized nuclei.[6] Finally, in Sec. IV we consider β-γ correlation effects for a situation in which the β-decay daughter nucleus itself undergoes a radiative decay.[7] Attention is restricted here to the case of dipole (electric or magnetic) radiation. In this connection, a suitable candidate might well be positron decay of $K^{36}(I=1, j=2^+)$ to the analog state A^{36}, followed by electric dipole radiation to the state $j=3^-$.

II. UNPOLARIZED NUCLEI

The spectrum in electron and neutrino variables for β decay from unpolarized nuclei has the structure

$$dw=F_\mp(Z,E)\frac{G^2\cos^2\theta_C}{(2\pi)^5}(E_0-E)^2pE\,dE\,d\Omega_e\,d\Omega_\nu$$
$$\times\left\{f_1(E)+f_2(E)\left(\frac{\vec{p}\cdot\hat{k}}{E}\right)+f_3(E)\left[\left(\frac{\vec{p}\cdot\hat{k}}{E}\right)^2-\frac{p^2}{3E^2}\right]\right\},\qquad(7)$$

where E is the total electron energy, E_0 the maximum allowed electron energy, $\vec{p}$ the electron momentum, and $\hat{k}$ a unit vector in the direction of the neutrino momentum. Dominant Coulomb effects are contained in the energy-dependent Fermi functions $F_{\mp}(Z, E)$, the upper sign applying to negative electron decay, the lower sign to positron decay. Similarly, in the following expression for the spectral functions f_i, upper signs refer to electron, lower signs to positron decay; and M denotes the mean mass of parent and daughter nuclei. We find

$$f_1(E) = a^2 + c^2 - \frac{2}{3}\frac{E_0}{M}(c^2 \pm cb \pm cd) + \frac{2}{3}\frac{E}{M}(3a^2 + 5c^2 \pm 2cb) - \frac{1}{3}\frac{m_e^2}{ME}(2c^2 \pm 2cb \pm cd);$$

$$f_2(E) = a^2 - \frac{1}{3}c^2 + \frac{2}{3}\frac{E_0}{M}(c^2 \pm cb \pm cd) - \frac{4}{3}\frac{E}{M}(3c^2 \pm cb) ;$$

$$f_3(E) = \frac{E}{M}(-3a^2 + c^2) . \tag{8}$$

In order to simplify the writing we have assumed that the coefficients a, b, c, d are all real, as they would be if time-reversal invariance holds true for β decay. To allow for the possibility that the coefficients are complex, one need only replace a^2 by $|a|^2$, c^2 by $|c|^2$, cb by $\text{Re}\, c^*b$, etc. The same strictures apply to the formulas that appear in the following sections, where again, for ease of writing, we suppose the coefficients to be real.

III. POLARIZED NUCLEI

In this section we report the spectrum, integrated over neutrino directions, for β decay of polarized (and oriented) nuclei. We suppose the parent nuclei to form an incoherent ensemble with respect to the spin projection m along an axis of quantization described by a unit vector $\hat{n}$. The mean polarization vector is, therefore,

$$\left\langle \frac{\vec{j}}{j} \right\rangle = \frac{\langle m \rangle}{j}\hat{n} .$$

The spectrum will also contain effects depending on nuclear orientation, as characterized by the departure of $\langle m^2 \rangle$ from the value $j(j+1)/3$ that obtains for an ensemble of randomly oriented spins. A convenient parameter is, therefore,

$$\Lambda_j = 1 - \frac{3\langle m^2 \rangle}{j(j+1)} .$$

The spectrum is

$$dw = 2F_{\mp}(Z, E)\frac{G^2\cos^2\theta_C}{(2\pi)^4}(E_0 - E)^2 pE\, dE\, d\Omega_e \times \left\{ f_1(E) + f_4(E)\frac{\vec{j}}{j}\cdot\frac{\vec{p}}{E} + f_5(E)\Lambda_j\left[\left(\frac{\hat{n}\cdot\vec{p}}{E}\right)^2 - \frac{p^2}{3E^2}\right]\right\}, \tag{9}$$

where, recall, $\hat{n}$ is a unit vector along the axis of polarization and orientation. The spectral function $f_1(E)$ has already been given in Eq. (8). For the other spectral functions we have

$$f_4(E) = \left(\frac{j}{j+1}\right)^{1/2}\left[2ac - \frac{2E_0}{3M}(ac \pm ab \pm ad) + \frac{2E}{3M}(7ac \pm ab \pm ad)\right] - \left(\frac{1}{j+1}\right)\left[\pm c^2 - \frac{2E_0}{3M}(\pm c^2 + cb + cd) + \frac{E}{3M}(\pm 11c^2 + 5cb - cd)\right];$$

$$f_5(E) = \frac{E}{2M}(c^2 \pm cb \mp cd) . \tag{10}$$

IV. β-γ CORRELATIONS

Here we consider a situation where the daughter nucleus produced in the β-decay process undergoes a radiative transition to a final state of spin j', with emission of dipole radiation (electric or magnetic). We are reverting again to the case where the β-decay parent nucleus is unpolarized[8]; and we again integrate over neutrino directions, considering the spectrum in its dependence on the electron variables and on the direction of the γ ray. The latter is characterized by a unit vector $\hat{K}$ along the direction of motion of the γ ray in the *laboratory frame* (rest frame of the β-decay parent nucleus). The spectrum reflects certain kinematic shift effects associated with the transformation to the lab frame from the rest frame of the β-decay daughter nucleus (where the radiative transition is most simply characterized). In the following spectrum formula it is the spectral function $g(E)$ that expresses these kinematic effects. The spectrum is given by

$$dw = F_{\mp}(Z, E)\frac{G^2\cos^2\theta_C}{(2\pi)^5}(E_0 - E)^2 pE\, dE\, d\Omega_e d\Omega_\gamma \times \left\{ f_1(E) + g(E)\frac{\hat{K}\cdot\vec{p}}{E} + \lambda_{j,j'} f_6(E)\left[\left(\frac{\hat{K}\cdot\vec{p}}{E}\right)^2 - \frac{p^2}{3E^2}\right]\right\}, \tag{11}$$

where the coefficient $\lambda_{j,j'}$ depends on the spins j and j' of the parent and daughter nuclei of the radiative process:

$$\lambda_{j,j'} = -(2j-1)/(j+1) \qquad j' = j+1 ,$$
$$= (2j+3)(2j-1)/j(j+1) \qquad j' = j,$$
$$= -(2j+3)/j \qquad j' = j-1 .$$

The spectral function $f_1(E)$ has been encountered previously. For the others we find

$$f_6(E) = \frac{E}{20M}(c^2 \pm cb \mp cd);$$

$$g(E) = \frac{2E_0}{3M}\left[-a^2 + \frac{c^2}{3}\left(1 - \frac{\lambda_{j,j'}}{10}\right]\right. \qquad (12)$$

$$\left. - \frac{4E}{3M}\left[a^2 + \frac{5c^2}{3}\left(1 - \frac{\lambda_{j,j'}}{100}\right)\right]\right. .$$

V. SUMMARY

In the approximation where recoil effects are ignored, the spectra described in the preceding sections are controlled solely by the coefficients a and c (Fermi and Gamow-Teller matrix elements). The structure in this allowed approximation is a familiar one: The spectral functions f_1, f_2, f_4 are energy independent and the remaining structure functions vanish. The effects of weak magnetism and of second-class currents arise in next approximation, as characterized by the small recoil parameter E/M. At this level the spectral functions f_1, f_2, f_4 acquire energy-dependent corrections, and the remaining functions no longer vanish. These qualitatively new recoil effects depend on all the coefficients a, b, c, d; but it may be supposed that a and c are well enough known from the effects produced in the allowed approximation. For weak magnetization the characterization of recoil phenomena in terms of the parameter E/M is somewhat misleading. As suggested in Eq. (5), in the CVC hypothesis the coefficient b is expected to be of order $A = M/m$, where m is the proton, M the nuclear mass. So the parameter of smallness for weak-magnetism effects is better regarded as E/m. Similarly, if the substantial ft discrepancies summarized by Wilkinson and Alburger are attributed to second-class currents, then the coefficient d would also be large (of order A, or even bigger). For not-too-light nuclei we therefore expect the coefficient b to be fairly large compared with a and c; and we may hope that d is similarly large. Nevertheless, for completeness we have carried along the genuinely small recoil terms, such as a^2E/M, in addition to the hopefully larger terms such as cdE/M, cbE/M, etc.

As is seen in the spectrum formulas, the weak-magnetism and second-class-current effects can in principle be separated experimentally. For example, the term linear in energy in the spectral function f_1 does not depend on the coefficient d. With a and c known, the energy dependence determines the coefficient b and in turn this provides a basis for testing the CVC hypothesis.[9] Similarly, the term linear in energy f_2 depends only on c^2 and cb; but the detection of a small energy dependence in an e-ν correlation experiment is no doubt excessively demanding. For experiments with polarized nuclei, the term linear in energy in the spectral function f_4 depends on all the coefficients a, b, c, d; with the first three of these known, this provides a basis for detecting the second-class-current coefficient d. For oriented nuclear spins the spectral function f_5 comes into play. This depends only on c^2, cb, and cd, a simple situation; the β-γ correlation function f_6, up to an unfortunately small numerical factor, measures the same combination of quantities. To summarize, with a and c regarded as known, from effects produced in leading approximation, the term linear in energy in f_1 determines cb, and d can then be determined from f_4, or $f_5 \propto f_6$.

It must be emphasized again that we have all along been restricting ourselves to β decay between members of a common isotopic multiplet. Within the CVC hypothesis, for allowed decays ($\Delta j = 0$, ± 1, "no") between members of different multiplets, the spectrum structure is again as described here, with the simplification $a = 0$ and with minor alterations in j-dependent factors if $\Delta j \neq 0$. But, for these more general β-decay processes the coefficient d does not any longer necessarily signal the presence of second-class currents.

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[1]D. H. Wilkinson, Phys. Letters 31B, 447 (1970); D. H. Wilkinson and D. E. Alburger, Phys. Rev. Letters 24, 1134 (1970); Phys. Letters 32B, 190 (1970).

[2]S. Weinberg, Phys. Rev. 112, 1375 (1958).

[3]R. P. Feyman and M. Gell-Mann, Phys. Rev. 109, 193 (1958).

[4]Although analog transitions typically involve positron emission, we take up the discussion for electron decay in order to parallel the familiar case of neutron β decay.

[5]For positron decay the amplitude $\langle\beta|V_\mu^\dagger + A_\mu^\dagger|\alpha\rangle$ is defined in an identical way, except that we choose to reverse the algebraic sign in front of the term with coefficient d.

[6]Such calculations have already been done in the case of neutron decay by S. M. Bilen'kii *et al.*, Zh. Experim. i Teor. Fiz. 37, 1758 (1959) [transl.: Soviet Phys.—

JETP 10, 1241 (1960)]. Our results merely generalize this work to arbitrary spin and include in addition the possiblilty of a second-class axial-vector-current contribution.

[7]Photon-electron correlations have previously been considered by J. Bernstein and R. R. Lewis, Phys. Rev. 112, 232 (1958). We here append the possible contribution of a second-class axial-vector current.

[8]Inclusion of orientation for the parent nucleus gives nothing new in the sense that all "nonkinematic" orientation-dependent correlations involving the photon direction [e.g., $J\cdot\hat{K}\hat{K}\cdot p$, $(J\cdot\hat{K}J\cdot p)^2$] depend upon the same combination, $c^2\mp cd\pm cb$, as does $[(\hat{K}\cdot p/E)^2-\frac{1}{3}(p/E)^2]$.

[9]Such a test has of course already been provided within the N^{12}, C^{12}, B^{12} system: M. Gell-Mann, Phys. Rev. 111, 362 (1958); Y. K. Lee, L. W. Mo, and C. S. Wu, Phys. Rev. Letters 10, 253 (1963).

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Cluster-Model Calculation of ^{6}Li(π^-, nn)^{4}He†

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The two-nucleon pion-absorption operator of Koltun and Reitan is applied to the pionic disintegration process of ^{6}Li. The target nucleus is represented by the cluster-model wave function of Tang, Wildermuth, and Pearlstein. The pion-absorption operator is of nonzero-range character and takes into account both the direct-absorption and the rescattering processes. The transition probability of the decay turns out to depend on the two parameters associated, respectively, with the intercluster distance and the size of the deuteron in ^{6}Li. The values of these two parameters are determined so as to fit the experimental momentum distribution of the two final neutrons to within experimental error. The possibility of the existence of a second peak in the momentum distribution is also discussed.

I. INTRODUCTION

In this paper we report the study of the pionic disintegration ^{6}Li(π^-, nn)^{4}He, of ^{6}Li,[1] where the negative pion is assumed to be absorbed at rest by the two nucleons in the (n, p, α) system constituting ^{6}Li. For the initial state we employ the cluster-model wave function (henceforth abbreviated as CMW) of Tang, Wildermuth, and Pearlstein.[2] This CMW was determined by the variational calculation with respect to the ground-state energy of ^{6}Li and contains three variational parameters, α, $\bar{\alpha}$, and β, respectively, related to the sizes of the α particle and deuteron and the distance between these two. The CMW was employed by Hansteen and his co-workers[3] in the study of the electrodisintegration of ^{6}Li in the Coulomb field of ^{197}Au. For the present pionic decay process of ^{6}Li, the CMW has already been used by Sakamoto,[4] but the pion-absorption operator he used contains features that are not in harmony with experimental results. Sakamoto's operator was the one used earlier by Eckstein[5] and others, and it is our belief that this operator has now been superceded by the more elaborate two-nucleon-absorption operator, first formulated by Woodruff[6] and recently improved by Koltun and Reitan.[7] Sakamoto's operator contains the zero-range factor representing the absorption of the pion only when the two nucleons overlap with each other. Consequently, the final answer for the transition probabilities turns out to be *independent* of the parameter $\bar{\alpha}$ representing the size of the deuteron, and hence the test of the CMW is rendered incomplete. Sakamoto's operator has the further defect that it takes into account only the single direct process of the pion absorption. Since the neutron in the (n, p, α) system cannot absorb the π^-, but can only scatter the π^-, it is clear that one must also take into account the two-step process, where the π^- is first scattered by the neutron and then is absorbed by the proton. Recently, Koltun and Reitan[7] formulated an elaborate pion-absorption operator (henceforth abbreviated as the K-R operator), which is free from these two defects of the operator used by Sakamoto and others. The K-R operator is more intricate than the well-known pseudovector interaction. It is a two-nucleon operator and is consistent with the threshold cross section of the reaction

$$p+p\rightarrow d+\pi^-, \quad (1)$$

which is equivalent to the absorption of the π^- at rest by a free deuteron. The K-R operator is constructed in such a way that it is in harmony with both S-wave and P-wave pions of the reaction (1). It follows, therefore, that the K-R operator