tent shown in Table I. On the other hand, the fit will appear deceptively good.

(iii) Since the curvature in Fig. 2 is so hard to see, the prospects for separating different pole contributions in the region $t < 0$ are poor, at least with presently accessible energies.

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[13]This curvature is another expression of the shrinking of the $F(s,t)$ pattern. For the $\bar{p}-p$ case the curvature has the opposite sign, and $F(s,t)$ expands.

THEORY OF NONLEPTONIC HYPERON DECAY

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Several theories[1-11] of nonleptonic hyperon decay have been proposed, but only one, due to Pais,[1] is consistent with all of the observed relations[12-14]

$$\alpha_+ \approx \alpha_- \approx 0, \tag{1}$$

$$\alpha_\Lambda \approx -\alpha_0, \tag{2}$$

$$\alpha_\Xi \approx -\alpha_\Lambda, \tag{3}$$

between the asymmetry parameters in Σ, Λ, and Ξ decay. Pais uses the baryon doublet approximation[15] and Relation (1) to predict (2). Instead of (3), however, he obtains the weaker prediction:

$$|\alpha_\Xi| \approx |\alpha_\Lambda|. \tag{4}$$

Here we propose another theory in which (1), (2), and (3) are derived from time-reversal invariance, the $\Delta T = \frac{1}{2}$ rule,[12,16] and three doublet symmetries. Since the electromagnetic interaction satisfies two of these symmetries, we are able to make predictions about the weak electromagnetic decays $\Sigma^+ \to p + \gamma$ and $\Xi^- \to \Sigma^- + \gamma$.

In the doublet approximation,[1,15] the isotopic spin $\vec{T}$ is split into a doublet spin $\vec{I}$ and a K-spin $\vec{K}$,

$$\vec{T} = \vec{I} + \vec{K}, \tag{5}$$

and the baryons are grouped into four $I = \frac{1}{2}$ doublets:

$$N_1 = \begin{pmatrix} p \\ n \end{pmatrix}, \quad N_2 = \begin{pmatrix} \Sigma^+ \\ Y^0 \end{pmatrix}, \quad N_3 = \begin{pmatrix} Z^0 \\ \Sigma^- \end{pmatrix}, \quad N_4 = \begin{pmatrix} \Xi^0 \\ \Xi^- \end{pmatrix};$$

$$Y^0 = 2^{-1/2}(\Lambda^0 - \Sigma^0), \quad Z^0 = 2^{-1/2}(\Lambda^0 + \Sigma^0). \tag{6}$$

N_2 and N_3 form a $K = \frac{1}{2}$ doublet:

$$N_{23} = \begin{pmatrix} N_2 \\ N_3 \end{pmatrix}, \tag{7}$$

with $K_z = +\frac{1}{2}, -\frac{1}{2}$, respectively, and N_1 and N_4 have zero K spin. The assignments for π and K mesons are $(I = 1, K = 0)$ and $(I = 0, K = \frac{1}{2})$, respectively. To supplement this scheme, we introduce a hypercharge spin $\vec{U}$, with z component

$$U_z = \tfrac{1}{2}(B + S), \tag{8}$$

and observe that N_1 and N_4 form a $U = \frac{1}{2}$ doublet

$$N_{14}=\begin{pmatrix}N_1\\N_4\end{pmatrix}, \tag{9}$$

while N_2, N_3, and π are singlets. Two independent sets of spin-type matrices,[1] $\vec{\tau}$ and $\vec{\rho}$, are introduced to operate on these baryon spins; $\vec{\tau}$ acts only on the I spin and $\vec{\rho}$ on both K and U spins.[17]

From (5) and (8), the conservation of charge in nonleptonic hyperon decay, $\Delta Q=0$, becomes

$$\Delta I_z+\Delta K_z+\Delta U_z=0, \tag{10}$$

and the change in strangeness is given by

$$\tfrac{1}{2}\Delta S=\Delta U_z=\pm\tfrac{1}{2}. \tag{11}$$

These relations enable us to divide the decays into two groups, one obeying the selection rule $\Delta I_z=0$ and the other $|\Delta I_z|=1$:

$$\left.\begin{array}{l}n_2\to n_1+\pi\\ n_3\to n_4+\pi\end{array}\right\}\quad \Delta I_z=0, \tag{12a}$$

$$\left.\begin{array}{l}n_2\to n_4+\pi\\ n_3\to n_1+\pi\end{array}\right\}\quad |\Delta I_z|=1, \tag{12b}$$

where n_i denotes a member of the doublet N_i in (6).

Since the processes in (12) represent transitions from the doublets N_2, N_3 to N_1, N_4, we write the effective interaction Hamiltonian as

$$H_{\text{weak}}=\bar{N}_{23}\Gamma\pi N_{14}+\text{H.c.}, \tag{13}$$

where

$$\bar{N}_{ij}=(\bar{N}_i,\bar{N}_j),$$

and the operator Γ remains to be specified. If we assume the $\Delta T=\frac{1}{2}$ rule, then the most general parity-nonconserving expression for H_{weak} is

$$H_{\text{weak}}=H^{(0)}+H^{(1)}+H^{(2)},$$

$$H^{(0)}=\sum_{\lambda=V,A}\{\bar{N}_{23}\Gamma_\lambda\vec{\tau}\cdot\vec{\pi}\times[g_\lambda^{(0)}+h_\lambda^{(0)}\rho_3]N_{14}+\text{H.c.}\},$$

$$H^{(1)}=\sum_{\lambda=V,A}\{\bar{N}_{23}\Gamma_\lambda\vec{\rho}\cdot\vec{\pi}\times[g_\lambda^{(1)}+h_\lambda^{(1)}\rho_3]N_{14}+\text{H.c.}\},$$

$$H^{(2)}=\sum_{\lambda=V,A}\{\bar{N}_{23}\Gamma_\lambda\vec{\rho}\cdot(\vec{\tau}\times\vec{\pi}/i)\times[g_\lambda^{(2)}+h_\lambda^{(2)}\rho_3]N_{14}+\text{H.c.}\}. \tag{14}$$

Γ_λ indicates that the space-time interaction is either vector ($\lambda=V$), or axial vector ($\lambda=A$). We assume time-reversal invariance, and so the coupling constants $g_\lambda^{(i)}$ and $h_\lambda^{(i)}$ are real. Under charge conjugation, the behavior of H_{weak} is represented by

$$g_\lambda\to\xi_\lambda g_\lambda,\qquad h_\lambda\to\xi_\lambda h_\lambda, \tag{15}$$

where

$$\xi_\lambda=\begin{cases}+1, \lambda=V\\ -1, \lambda=A\end{cases}.$$

The selection rules for $H^{(k)}$ are as follows:

$$\text{for}\quad H^{(0)};\quad \Delta I=0,\quad \Delta K=\Delta U=\tfrac{1}{2};$$

$$\text{for } H^{(1)}, H^{(2)};\quad \Delta I=1,\quad \Delta K=\Delta U=\tfrac{1}{2}. \tag{16}$$

In $H^{(1)}$ and $H^{(2)}$, ΔK (determined by the factor $\bar{N}_{23}$) is combined with ΔI to a resultant $\Delta T=\frac{1}{2}$ by virtue of the scalar products of $\vec{\rho}$ with $\vec{\pi}$ and $\vec{\tau}\times\vec{\pi}$. This procedure is exactly analogous to combining the I spins of antinucleons and pions to a resultant $I=\frac{1}{2}$ by means of the operator $\vec{\tau}$ in $\bar{N}_1\vec{\tau}\cdot\vec{\pi}$. The factors $(g+h\rho_3)$ in (14) insure that, in general, N_1 and N_4 are coupled to N_{23} with different strengths.

To discover possible symmetry principles, we recall that, in the doublet approximation, the strong pion-baryon interaction is[1]

$$H_{\text{strong}}=\bar{N}_{14}\vec{\tau}\cdot\vec{\pi}N_{14}+\epsilon\bar{N}_{23}\vec{\tau}\cdot\vec{\pi}N_{23}, \tag{17}$$

where ϵ is a phase factor equal to ±1. It is invariant under

$$T'=\{N_{14}\to\xi_{14}N_{23},\ N_{23}\to\xi_{23}N_{14},\ \vec{\pi}\to\epsilon\vec{\pi},\ K^+\to\xi_{23}\xi_{14}K^+,\ K^0\to\xi_{23}\xi_{14}\bar{K}^0\}; \tag{18}$$

where, for the moment, the phase factors ξ_{14} and ξ_{23} are arbitrary, and also under

$$C'=C\times\{N_{23}\to\rho_2\tau_2N_{23},\ N_{14}\to\rho_2\tau_2N_{14},\ \pi^\pm\to-\pi^\mp,\ \pi^0\to-\pi^0,\ K^+\to K^-,\ K^0\to\bar{K}^0\}; \tag{19}$$

where C denotes charge conjugation. T' represents the interchange of K and U spins, and C' the product of charge conjugation with rotations of 180° about the y axes of $\vec{I}$-, $\vec{K}$-, and $\vec{U}$-spin spaces. C' is similar to G parity but, unlike G parity, it leaves (10) invariant for all values of ΔU_z. The K-meson transformations are such that, under certain conditions (see reference 1, section 5), the strong K-meson-baryon interactions will also be invariant under (18) and (19).

Let us now suppose that the weak pion-baryon interactions are invariant under T' and C'. Because T' implies

$$\langle \Sigma^+ | p\pi^0 \rangle = \xi_{23}\xi_{14}\epsilon \langle p | \Sigma^+\pi^0 \rangle = \xi_{23}\xi_{14}\epsilon \langle \Sigma^+ | p\pi^0 \rangle \quad (20)$$

(final-state interactions are neglected),[1] we choose the phase factors in (18) so that

$$\xi_{23}\xi_{14}\epsilon = +1 \quad (21)$$

and find that T' invariance eliminates $H^{(2)}$ and also the terms of $H^{(1)}$ with coupling constants $h_\lambda^{(1)}$ ($\lambda = V, A$) from (14). Then, using C' invariance, we can reduce the Hamiltonian to

$$H_{\text{weak}} = \overline{N}_{23}[\Gamma_V(g_V^{(0)}\vec{\tau}\cdot\vec{\pi} + g_V^{(1)}\vec{\rho}\cdot\vec{\pi}) + \Gamma_A h_A^{(0)}\vec{\tau}\cdot\vec{\pi}\rho_3]N_{14} + \text{H.c.} \quad (22)$$

The second term of (22), containing the factor $\vec{\rho}\cdot\vec{\pi}$, is the only one for which $\Delta I = 1$. Therefore, the decay $\Sigma^- \to n + \pi^-$, which is of type (12b), proceeds via a pure vector interaction, and its asymmetry parameter α_- is zero.

Expanding (22) in terms of particle fields, we see that $\Sigma^+ \to n + \pi^+$, which is of type (12a), can only arise from the first and third terms. If α_+ is to be zero, and H_{weak} parity nonconserving, then we must find an argument that implies

$$g_V^{(0)} = 0. \quad (23)$$

We therefore require the $\Delta I_z = 0$ part of H_{weak} to be invariant under

$$N_{23} \to i\rho_2 N_{23}, \quad N_{14} \to -i\rho_2 N_{14}. \quad (24)$$

This is equivalent to the quantum number transformation

$$K_z \to -K_z, \quad U_z \to -U_z, \quad I_z \to I_z, \quad (25)$$

and is only consistent with charge conservation (10) when ΔI_z vanishes. Our Hamiltonian then becomes

$$H_{\text{weak}} = \overline{N}_{23}[g_V^{(1)}\Gamma_V \vec{\rho}\cdot\vec{\pi} + h_A^{(0)}\Gamma_A \vec{\tau}\cdot\vec{\pi}\rho_3] + \text{H.c.} \quad (26)$$

We have now derived a Hamiltonian that implies (i) $\alpha_+ = \alpha_- = 0$ and (ii) that $\Sigma^- \to n + \pi^-$ is a pure vector interaction and $\Sigma^+ \to n + \pi^+$ a pure axial vector. The coupling constants $g_V^{(1)}$ and $h_A^{(0)}$ can be chosen so that the rates for these two decays are equal. Equation (26) also implies

$$\sqrt{2}\langle \Sigma^+ | p\pi^0 \rangle = -\sqrt{2}\langle \Xi^- | \Lambda\pi^- \rangle = \langle \Sigma^+ | n\pi^+ \rangle - \langle \Sigma^- | n\pi^- \rangle,$$
$$\sqrt{2}\langle \Lambda | p\pi^- \rangle = \langle \Sigma^+ | n\pi^+ \rangle + \langle \Sigma^- | n\pi^- \rangle, \quad (27)$$

from which the relations (2) and (3) follow.

The roles of vector and axial-vector interactions in (22) and (26) can easily be interchanged by introducing different phase factors (e.g., $N_{23} \to -\rho_2\tau_2 N_{23}$) in (19). Thus the specific assignments of V and A in (26) are not very important; what is important, however, is that the operators $\vec{\rho}\cdot\vec{\pi}$ and $\vec{\tau}\cdot\vec{\pi}\rho_3$ must, as a consequence of our symmetry principles, be associated with different space-time interactions. In this way we have derived the basic assumption of Pais[1] that behavior under space reflection is correlated with the change of doublet spin.

Having assumed that the strong and weak interactions are invariant under T' and C', we now observe that the interaction of mesons and baryons with the electromagnetic field A,

$$H_{\text{el}} = \{\tfrac{1}{2}\overline{N}_{23}(\rho_3 + \tau_3)N_{23} + \tfrac{1}{2}\overline{N}_{14}(\rho_3 + \tau_3)N_{14} + \overline{\pi}^+\pi^+ + \overline{K}^+K^+ - \text{c.c.}\}A \quad (28)$$

(c.c. denotes the charge conjugate of the preceding expression in the bracket), is automatically invariant under these transformations, provided C' is completed with $A \to -A$. Thus the influence of T' and C' will be manifest in the weak electromagnetic decays of hyperons.

To make predictions about these processes, we use the TCP theorem and time-reversal invariance to replace charge conjugation in (19) by space reflection. It then follows that

$$\langle \Sigma^+ | p\gamma \rangle_\pm = \mp\langle \Xi^- | \Sigma^-\gamma \rangle_\pm, \quad (29)$$

where the suffices + and - denote the parity-conserving and parity-nonconserving parts of $\langle A | B\gamma \rangle$, respectively. Equation (29) implies that the rates for $\Xi^- \to \Sigma^- + \gamma$ and $\Sigma^+ \to p + \gamma$ are equal, and that their asymmetry parameters[18] are opposite in sign. There are similar relations among the photon decays of neutral hyperons, but we cannot relate $\Lambda \to n + \gamma$ to $\Sigma^+ \to p + \gamma$ without making further assumptions.

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