Parity-Nonconserving $\Delta S=0$ Nonleptonic Weak Processes in the Presence of Second-Class Currents*

P. Herczeg
Department of Physics, Carnegie-Mellon University, Pittsburgh, Pennsylvania 15213
(Received 9 June 1972)

The possibility that currents with irregular charge-conjugation properties are present in the standard CP-invariant current×current nonleptonic weak Hamiltonian is considered. It is pointed out that the parity-violating $n\to p\pi^-$ amplitude is not determined then by the strangeness-changing amplitudes alone and may be considerably larger than in the conventional theory.

There is good evidence at present from nuclear-physics experiments for the existence of a parity-nonconserving nucleon-nucleon interaction with strength at the level of first-order weak interactions.[1] Detailed nuclear-physics calculations[2] show good agreement with experiment for the parity-forbidden α decay of the 8.88-MeV state of ^{16}O (Ref. 3); the predicted value of photon circular polarization in the most extensively studied parity-nonconserving electromagnetic transitions is too small, however, by two orders of magnitude.[4]

The input of these calculations is a parity-nonconserving nucleon-nucleon potential consisting of a pion-exchange term V_π^{weak} and a ρ-exchange potential V_ρ^{weak}, estimated on the basis of the Cabibbo theory.[5] The parity-violating $n\to p\pi^-$ amplitude $A(n_-^0)$ was determined from the sum rule

$$A(n_-^0)=\tfrac{1}{3}\sqrt{6}\,\tan\theta[A(\Xi_-^-)-2A(\Lambda_-^0)] \qquad (1)$$

relating $A(n_-^0)$ to the amplitudes of observed nonleptonic hyperon decays.[6] Relation (1) can be derived either using SU(3) and octet dominance[7] or SU(3), current algebra, and partial conservation of axial-vector current (PCAC).[8] The weak $NN\rho$ amplitude used was the one obtained using the factorization approximation.[9] Although there is disagreement among the various existing estimates of the ρ-exchange potential,[9,10] the good agreement of the calculations with experiment for the parity-forbidden α decay of ^{16}O, to which V_π^{weak}, being a pure isovector,[11] does not contribute,[12] indicates that the strength of V_ρ^{weak} cannot be much larger than the factorization value.

Without considering the possibility that the existing discrepancy may be resolved within nuclear physics, the present status of the subject points to theories in which $A(n_-^0)$ is considerably larger than in the Cabibbo theory. Certain theories involving neutral currents[13] belong to this category.[4,14]

In this note we shall consider the possibility that currents with irregular charge conjugation properties[15] are present in the standard CP-invariant current×current nonleptonic weak Hamiltonian.[16] The strangeness-conserving part of the Hamiltonian contains then an isovector term proportional to $\cos^2\theta$.[17] The resulting additional terms in the parity-violating ρ-exchange potential have been studied in the factorization approximation.[18] In nuclear matter, with a reasonable value of the induced pseudotensor form factor, their contribution turns out to be of the same order of magnitude as the ρ-exchange term associated with regular currents.[19]

Here we shall consider the one-pion-exchange term. We would like to point out that in the presence of irregular currents the parity-violating $n\to p\pi^-$ amplitude is considerably modified and may be much larger than in the conventional theory.[20]

We shall investigate the $\Delta S=1$ and $\Delta S=0$ $B\to B'\pi$ amplitudes in the framework of a model, discussed in Ref. 21, which assumes a particular form for the commutators of the regular axial charges with the irregular currents. The conclusions will be seen, however, to be of more general validity.

The nonleptonic weak Hamiltonian is assumed to have the standard form

$$H_{\text{NL}}=-(G/\sqrt{2})\tfrac{1}{2}(J_\mu J_\mu^* + J_\mu^* J_\mu)\,, \qquad (2)$$

where

$$J_\mu=J_{(110)\mu}\cos\theta+J_{(1/2\,1/2\,1)\mu}\sin\theta\,,$$
$$J_{(\nu)\mu}^*=J_{(\nu)\mu}^\dagger \quad \text{for} \quad \mu=1,2,3\,,$$
$$J_{\nu 4}^*=-J_{(\nu)4}^\dagger \quad (\dagger\equiv\text{Hermitian conjugation})\,,$$

and $J_{(110)}$ and $J_{(1/2\,1/2\,1)}$ are members of an octet $J_{(\nu)}$ which we shall assume now to be a mixture of a regular octet $J_{(\nu)}^R=V_{(\nu)}^R+A_{(\nu)}^R$ and an irregular octet with respect to charge conjugation.[15]

CP invariance of H_{NL} requires $J_{(\nu)}^R$ and $J_{(\nu)}^I$ to be Hermitian and anti-Hermitian, respectively:

$$J_{(\nu)}^{R*}=(-1)^{Q_\nu}J_{(-\nu)}^R\,,$$
$$J_{(\nu)}^{I*}=(-1)^{1+Q_\nu}J_{(-\nu)}^I\,. \qquad (3)$$

As a consequence of Eqs. (3), the SU(3) structure of H_{NL} is a mixture of $\{1\}$, $\{8s\}$, $\{8a\}$, $\{10\}$, $\{\overline{10}\}$, and $\{27\}$, that is, of *all* representations contained in the direct product $\{8\}\times\{8\}$.

In the soft-pion limit, the s-wave amplitude for the process $B\to B'+\pi(p)$ is given by[22]

$$\lim_{p\to 0}(2p_0)^{1/2}\langle B'\pi_{(-\nu)}|H_{NL}^{p.v.}|B\rangle = (f_\pi)^{-1}(-1)^{Q_\nu}\left\langle B'\left|\left[\int d^3x\, A^R_{(\nu)0}, H_{NL}^{p.v.}\right]\right|B\right\rangle . \tag{4}$$

Following Ref. 21, we shall assume that the transformation properties of $J_+^I=\frac{1}{2}(V^I+A^I)$ with respect to SU(3)×SU(3) generated by the regular octets V^R and A^R are the same as those of $J_+^R=\frac{1}{2}(V^R+A^R)$, i.e., that J_+^I belongs to the (8,1) representation. All the commutators in (4) are then specified, and the decay amplitudes of the various processes can be expressed in terms of seven reduced matrix elements. We obtain the following relations among the $\Delta S=1$ and $\Delta S=0$ amplitudes[23]:

$$A(\Lambda^0_-)+\sqrt{2}\,A(\Lambda^0_0)=0\,, \tag{5}$$

$$A(\Xi^-_-)-\sqrt{2}\,A(\Xi^0_0)=0\,, \tag{6}$$

$$A(\Sigma^-_-)-A(\Sigma^+_+)+\sqrt{2}\,A(\Sigma^+_0)=-2A(\Sigma^+_+)\,, \tag{7}$$

$$A(\hat{\sigma}^-_-)+A(\lambda^0_-)=\tfrac{1}{3}\sqrt{3}\,(\cot\theta-\tan\theta)[\delta_A(LS)+\tfrac{1}{2}\sqrt{6}\,A(\Sigma^+_+)]\,, \tag{8}$$

$$A(n^0_-)-A(\xi^-_-)-\sqrt{6}\,A(\hat{\sigma}^-_-)=\tfrac{1}{2}\sqrt{2}\,\tan^2\theta[A(\sigma^0_-)-A(\sigma^-_-)]+\tfrac{1}{2}\sqrt{2}(\tan\theta-\cot\theta)[\delta_A(LS)+\tfrac{1}{2}\sqrt{6}\,A(\Sigma^+_+)]\,, \tag{9}$$

$$A(n^0_-)+A(\xi^-_-)-\sqrt{2}\,A(\sigma^0_-)=2(\tan\theta-\cot\theta)A(\Sigma^+_+)+\tfrac{1}{2}\sqrt{2}\,(1-2\tan^2\theta)[A(\sigma^0_-)-A(\sigma^-_-)] + \tfrac{1}{2}\sqrt{2}\,(\cot\theta-\tan\theta)[\delta_A(LS)+\tfrac{1}{2}\sqrt{6}\,A(\Sigma^+_+)]\,, \tag{10}$$

$$A(n^0_-)-\tfrac{1}{3}\sqrt{6}\,\tan\theta[A(\Xi^-_-)-2A(\Lambda^0_-)]=(-\tfrac{1}{30}\sqrt{30}\,\tilde{M}^s_{\{8a\}}-\tfrac{1}{18}\sqrt{6}\,\tilde{M}^a_{\{8a\}})k_A+(\tfrac{1}{3}\tan\theta-\tfrac{2}{3}\cot\theta)A(\Sigma^+_+) + \tfrac{1}{6}\sqrt{2}\,\cot\theta[\delta_A(LS)+\tfrac{1}{2}\sqrt{6}\,A(\Sigma^+_+)]+\tfrac{1}{6}\sqrt{2}\,(1+\tan^2\theta)[A(\sigma^0_-)-A(\sigma^-_-)]\,, \tag{11}$$

$$A(\sigma^0_-)+\tfrac{1}{3}\tan\theta[A(\Lambda^0_-)+A(\Xi^-_-)]=-\tfrac{1}{9}\sqrt{3}\,k_A\tilde{M}^a_{\{8a\}}+\tfrac{1}{3}\sqrt{2}\,(\cot\theta-2\tan\theta)A(\Sigma^+_+) + \tfrac{1}{3}(1+\tan^2\theta)[A(\sigma^0_-)-A(\sigma^-_-)]+\tfrac{1}{6}(3\tan\theta-\cot\theta)[\delta_A(LS)+\tfrac{1}{2}\sqrt{6}\,A(\Sigma^+_+)]\,, \tag{12}$$

$$A(\sigma^0_-)-A(\sigma^-_-)=2\sqrt{2}\,\cot\theta\,A(\Sigma^+_+)+\tfrac{2}{15}\sqrt{15}\,k_A\cos^2\theta\,\tilde{M}_{\{10\}}\,, \tag{13}$$

$$\delta_A(LS)\equiv 2A(\Xi^-_-)-A(\Lambda^0_-)+\sqrt{3}\,A(\Sigma^+_0)=-\tfrac{1}{2}\sqrt{6}\,A(\Sigma^+_+)-\tfrac{1}{15}\sqrt{15}\,k_A\sin\theta\cos\theta(\tilde{M}_{\{10\}}+\tilde{M}_{\{\overline{10}\}})\,, \tag{14}$$

$$A(\Sigma^+_+)=k_A\sin\theta\cos\theta(\tfrac{1}{9}M_{\{27\}}-\tfrac{1}{30}\sqrt{30}\,\tilde{M}_{\{10\}})\,. \tag{15}$$

In the above equations $k_A=-(G/\sqrt{2})\sqrt{6}\,(f_\pi)^{-1}$, and $M_{\{k\}}$ and $\tilde{M}_{\{k\}}$ are the reduced matrix elements defined by

$$\langle B'_{(\nu_f)}|T^{\{k\}}_{c(\nu)}|B_{(\nu_i)}\rangle=\bar{u}(\vec{p}\,')\left[\sum_\gamma\begin{pmatrix}8 & k & 8_\gamma\\ \nu_i & \nu & \nu_f\end{pmatrix}M^\gamma_{\{k\}}\right]u(\vec{p})\,,$$
$$\langle B'_{(\nu_f)}|\tilde{T}^{\{k\}}_{c(\nu)}|B_{(\nu_i)}\rangle=\bar{u}(\vec{p}\,')\left[\sum_\gamma\begin{pmatrix}8 & k & 8_\gamma\\ \nu_i & \nu & \nu_f\end{pmatrix}\tilde{M}^\gamma_{\{k\}}\right]u(\vec{p})\,, \tag{16}$$

where $T^{\{k\}}_c$ and $\tilde{T}^{\{k\}}_c$ are the parity-conserving parts of

$$T^{\{k\}}=\sum_{\alpha\beta}\begin{pmatrix}8 & 8 & k\\ \alpha & \beta & \end{pmatrix}(J^R_{(\alpha)}J^R_{(\beta)}-J^I_{(\alpha)}J^I_{(\beta)})\,,$$
$$\tilde{T}^{\{k\}}=\sum_{\alpha\beta}\begin{pmatrix}8 & 8 & k\\ \alpha & \beta & \end{pmatrix}(J^R_{(\alpha)}J^I_{(\beta)}-J^I_{(\alpha)}J^R_{(\beta)})\,. \tag{17}$$

According to experiment[24] $A(\Sigma^+_+)\approx 0$ and $\delta_A(LS)\approx 0$, implying $M_{\{27\}}\approx\frac{3}{10}\sqrt{30}\,\tilde{M}_{\{10\}}$ and $\tilde{M}_{\{10\}}\approx\tilde{M}_{\{\overline{10}\}}$. Furthermore, in view of the good evidence[24] for the validity of the $\Delta T=\frac{1}{2}$ rule in p-wave Ξ decays and of the p-wave Lee-Sugawara relation, $\tilde{M}_{\{10\}}$ is likely to be negligibly small.[25] The right-hand sides of Eqs. (7)–(10) and (13)–(15) then vanish. We recover the $\Delta T=\frac{1}{2}$ sum rules for Λ, Ξ, and Σ decay [Eqs. (5)–(7)] and the $\Delta S=0$ relations[26] [Eqs. (8)–(10) and (13)]. The mixed sum rules relating the $\Delta S=0$ and $\Delta S=1$ amplitudes,[26] however, no longer hold. Instead we have

$$A(n^0_-)-\tfrac{1}{3}\sqrt{6}\,\tan\theta[A(\Xi^-_-)-2A(\Lambda^0_-)] = -k_A(\tfrac{1}{30}\sqrt{30}\,\tilde{M}^s_{\{8a\}}+\tfrac{1}{18}\sqrt{6}\,\tilde{M}^a_{\{8a\}})\,, \tag{18}$$

$$A(\sigma^0_-)+\tfrac{1}{3}\tan\theta[A(\Lambda^0_-)+A(\Xi^-_-)]=-\tfrac{1}{9}\sqrt{3}\,k_A\tilde{M}^a_{\{8a\}}\,. \tag{19}$$

As seen from Eq. (18), the knowledge of $A(\Lambda^0_-)$ and $A(\Xi^-_-)$ is not sufficient to determine the parity-violating $n \to p\pi^-$ amplitude. The reason for this is quite general and can be traced back to the SU(3) structure of the Hamiltonian (2). It turns out that the octet components of the $\Delta S=0$, $\Delta T=1$ and of the $\Delta S=1$, $\Delta T=\frac{1}{2}$ parts of H_{NL} are members of *different* octets:

$$H^{\Delta S=0}_{(100)} = -(G/\sqrt{2})(R^{\{8\}}_{(100)} + \tfrac{1}{3}\sqrt{3}\,\tilde{T}^{\{8a\}}_{(100)}),$$
$$H^{\Delta S=1}_{(1/2\;-1/2\;1)} = -(G/\sqrt{2})(-\sqrt{2})\cot\theta\, R^{\{8\}}_{(1/2\;-1/2\;1)}, \tag{20}$$

where

$$R^{\{8\}} = \sin^2\theta(\tfrac{1}{10}\sqrt{15}\,T^{\{8s\}} - \tfrac{1}{6}\sqrt{3}\,\tilde{T}^{\{8a\}}). \tag{21}$$

The $\Delta S=0$ and the $\Delta S=1$ amplitudes are therefore, in general, independent.

With the help of the sum rules (10) and (11), the right-hand side of Eq. (18) can be written as follows:

$$-k_A(\tfrac{1}{30}\sqrt{30}\,\tilde{M}^s_{\{8a\}} + \tfrac{1}{18}\sqrt{6}\,\tilde{M}^a_{\{8a\}}) = \tfrac{1}{2}[\sqrt{6}\,A(\hat{\sigma}^-_-) + \sqrt{2}\,A(\sigma^0_-)] - \tfrac{1}{3}\sqrt{6}\,\tan\theta[A(\Xi^-_-) - 2A(\Lambda^0_-)]. \tag{22}$$

Thus, the evaluation of this expression requires, in addition to $A(\Lambda^0_-)$ and $A(\Xi^-_-)$, the values of the parity-violating $\Sigma^- \to \Lambda\pi^-$ and $\Sigma^- \to \Sigma^0\pi^-$, or of any other two independent $\Delta S=0$ amplitudes.[27] In principle, these could be deduced, for example, from detailed studies of hypernuclear transitions. Such investigations are, of course, inconceivable at the present time.

Expressing $A(\Lambda^0_-)$ and $A(\Xi^-_-)$ in Eq. (18) in terms of the reduced matrix elements, we obtain

$$A(n^0_-) = \tfrac{1}{3}\sqrt{6}\,\tan\theta[A(\Xi^-_-) - 2A(\Lambda^0_-)][1 + \gamma(1 + \cot^2\theta)], \tag{23}$$

where

$$\gamma = (30\tilde{M}^s_{\{8a\}} + 10\sqrt{5}\,\tilde{M}^a_{\{8a\}})/(9\sqrt{5}\,M^s_{\{8s\}} + 15M^a_{\{8s\}} - 15\tilde{M}^s_{\{8a\}} - 5\sqrt{5}\,\tilde{M}^a_{\{8a\}}). \tag{24}$$

Barring accidental cancellations, γ can be of order unity.[16] For $|\gamma| \geq 1$ we have

$$|A(n^0_-)| \geq \cot^2\theta\,|\tfrac{1}{3}\sqrt{6}\,\tan\theta[A(\Xi^-_-) - 2A(\Lambda^0_-)]|$$
$$\approx 20\,|\tfrac{1}{3}\sqrt{6}\,\tan\theta[A(\Xi^-_-) - 2A(\Lambda^0_-)]|. \tag{25}$$

To summarize, if irregular currents participate in the nonleptonic weak interactions,[28] the weak one-pion-exchange potential is no longer determined by the $\Delta S=1$ hyperon decay amplitudes alone,[29] and may be considerably stronger than in the conventional theory. As a consequence, a limit on the magnitude of the induced pseudotensor form factor[19] cannot be obtained without information on the parity-violating $n \to p\pi^-$ amplitude. A large value of $A(n^0_-)$ may indeed be needed to remove the existing discrepancy between theory and experiment in the case of the parity-nonconserving electromagnetic transitions.

I would like to thank Professor L. Wolfenstein for helpful conversations and remarks. My thanks are due also to Professor P. Winternitz and Dr. S. K. Singh for their interest and stimulating discussions.

*Work supported in part by the U. S. Atomic Energy Commission.

[1]The present experimental situation is reviewed by F. Boehm, in Proceedings of the International Conference on Angular Correlation in Nuclear Disintegration, Delft, 1970 (unpublished). See also E. D. Lipson, F. Boehm, and J. C. Vanderleeden, Phys. Rev. C 5, 932 (1972). Reviews of the theory can be found in the following: R. J. Blin-Stoyle, in *Proceedings of the Topical Conference on Weak Interactions, CERN, 1969* (CERN, Geneva, 1969); E. M. Henley, Ann. Rev. Nucl. Sci. 19, 367 (1969); E. Fischbach and D. Tadić (unpublished).

[2]M. Gari and H. Kümmel, Phys. Rev. Letters 23, 26 (1969); E. M. Henley, T. E. Keliher, and D. U. L. Yu, Phys. Rev. Letters 23, 941 (1969).

[3]H. Hättig, K. Hünchen, and W. Wäffler, Phys. Rev. Letters 25, 941 (1970).

[4]M. Gari, O. Dumitrescu, J. G. Zabolitzky, and H. Kümmel, Phys. Letters 35B, 19 (1971); B. Desplanques and N. Vinh Mau, Phys. Letters 35B, 28 (1971).

[5]Of the members of the pseudoscalar and vector-meson nonets only $\pi^\pm$, $\rho^{\pm0}$, ω, and ϕ can contribute [G. Barton, Nuovo Cimento 19, 512 (1961)]. In the factorization approximation or in the field-algebra approach, ρ^0, ω, and ϕ give vanishing contribution (cf. Ref. 10). Higher-mass one-boson exchanges are neglected because of their short range and the nucleon hard core. For the same reason multiparticle exchanges need not be considered either. A possible exception is the two-pion-exchange contribution [R. J. Blin-Stoyle, Phys. Rev. 118, 1605 (1960); R. Lacaze, Nucl. Phys. B4, 657 (1968); W. K. Cheng, E. Fischbach, H. Primakoff, D. Tadić, and K. Trabert, Phys. Rev. D 3, 2289 (1971); D. Pignon, Phys. Letters 35B, 163 (1971)]. See also E. Fischbach and D. Tadić (unpublished).

[6]In Eq. (1) θ is the Cabibbo angle; $A(\Xi^-_-)$, for example,

is the s-wave amplitude for the process $\Xi^- \to \Lambda + \pi^-$.

[7]W. Kummer, Nuovo Cimento 36, 885 (1965); D. Tadić, Phys. Rev. 174, 1694 (1968).

[8]B. H. J. McKellar, Phys. Letters 26B, 107 (1967); E. Fischbach, Phys. Rev. 170, 1398 (1968); M. Feuer, Ph.D. thesis, Harvard University, 1969 (unpublished).

[9]F. C. Michel, Phys. Rev. 133, B329 (1964).

[10]E. Fischbach, D. Tadić, and K. Trabert, Phys. Rev. 186, 1688 (1969); P. Olesen and J. S. Rao, Phys. Letters 29B, 233 (1969); M. Feuer, Ph.D. thesis, Harvard University, 1969 (unpublished); E. Fischbach and D. Tadić (unpublished); B. H. J. McKellar and P. Pick, Phys. Rev. D (to be published); B. H. J. McKellar and P. Pick, Phys. Rev. D (to be published).

[11]G. Barton, Nuovo Cimento 19, 512 (1961).

[12]W. K. Cheng, E. Fischbach, H. Primakoff, D. Tadić, and K. Trabert, Phys. Rev. D 3, 2289 (1971).

[13]B. d'Espagnat, Phys. Letters 1, 209 (1963); G. Segrè, Phys. Rev. 173, 1730 (1968); T. D. Lee, Phys. Rev. 171, 1731 (1968); R. J. Oakes, Phys. Rev. Letters 20, 1539 (1968).

[14]R. F. Dashen, S. C. Frautschi, M. Gell-Mann, and Y. Hara, in *The Eightfold Way*, edited by M. Gell-Mann and Y. Ne'eman (Benjamin, New York, 1964), p. 254; B. H. J. McKellar, Phys. Rev. 178, 2160 (1968); Phys. Rev. Letters 21, 1822 (1968); E. Fischbach and K. Trabert, Phys. Rev. 174, 1843 (1968); D. Tadić, Phys. Rev. 174, 1694 (1968); E. Fischbach, D. Tadić, and K. Trabert, in *Proceedings of the Third International Conference on High-Energy Physics and Nuclear Structure, New York, N. Y., September, 1969,* edited by S. Devons (Plenum, New York, 1970).

[15]An octet of vector or axial-vector currents $J_{(\nu)\mu} \equiv J_{(TT_3Y)\mu}$ ($\mu=1, 2, 3, 4$) which goes into itself under charge conjugation satisfies (omitting in the following the Dirac indices)

$$CPJ_{(\nu)}(CP)^{-1}=\eta(J)(-1)^{Q_\nu}J_{(-\nu)},$$

where $(-\nu)\equiv(T-T_3-Y)$, $Q_\nu=T_3+\frac{1}{2}Y$, and $\eta(J)=\pm 1$, and is the same for all components of a given octet [Y. Dothan, Nuovo Cimento 30, 399 (1963); M. Gell-Mann, Phys. Rev. Letters 12, 155 (1964)]. $J_{(\nu)}=V_{(\nu)}+A_{(\nu)}$ will be called *regular* (J^R) if $\eta=+1$ and *irregular* (J^I) when $\eta=-1$. Under G parity

$$GPJ^R_{(1T_30)}(GP)^{-1}=-J^R_{(1T_30)}$$

and

$$GPJ^I_{(1T_30)}(GP)^{-1}=J^I_{(1T_30)},$$

so that $J^R_{(1T_30)}$ is a first-class and $J^I_{(1T_30)}$ a second-class current [S. Weinberg, Phys. Rev. 112, 1375 (1959)]. Regarding SU(3), we use the notation and conventions of J. J. de Swart [Rev. Mod. Phys. 35, 916 (1963)].

[16]At present there is no evidence against the existence of second-class currents. Recent studies of differences in the ft values of mirror β decays are suggestive of the possibility of their existence, but the alternative, that these effects are of electromagnetic origin, cannot be ruled out at this time [cf. M. A. B. Bég and J. Bernstein, Phys. Rev. D 5, 714 (1972) and references quoted therein].

[17]R. J. Blin-Stoyle and P. Herczeg, Phys. Letters 23, 376 (1966).

[18]R. J. Blin-Stoyle and P. Herczeg, Nucl. Phys. B5, 291 (1968).

[19]B. H. J. McKellar, in *Proceedings of the Third International Conference on High-Energy Physics and Nuclear Structure, New York, N. Y., September, 1969,* edited by S. Devons (Plenum, New York, 1970), p. 682; Phys. Rev. D 1, 2183 (1970); B. Eman, D. Tadić, F. Krmpotić, and L. Szybisz, Phys. Rev. C 6, 1 (1972).

[20]In the factorization approximation the effects of irregular currents on the one-pion-exchange term manifest themselves only through the term proportional to $g_S g_P$, the product of the induced scalar and the induced pseudoscalar form factors [cf. Ref. 18], which vanishes in the absence of an irregular vector current. Our conclusions are independent of the existence of irregular vector currents.

[21]P. Herczeg, Nucl. Phys. B4, 153 (1967).

[22]H. Sugawara, Phys. Rev. Letters 15, 870 (1965); M. Suzuki, Phys. Rev. Letters 15, 986 (1965).

[23]The $\Delta S=0$ s-wave amplitudes $n\to p+\pi^-$, $\Sigma^0\to\Sigma^++\pi^-$, $\Sigma^-\to\Sigma^0+\pi^-$, $\Sigma^-\to\Lambda+\pi^-$, $\Lambda^0\to\Sigma^++\pi^-$, and $\Xi^-\to\Xi^0+\pi^-$, respectively, by $A(n^0_-)$, $A(\sigma^0_-)$, $A(\sigma^-_-)$, $A(\hat{\sigma}^-_-)$, $A(\Lambda^0_-)$, and $A(\xi^-_-)$.

[24]Particle Data Group, Phys. Letters 39B, 1 (1972).

[25]In the baryon pole model for the p-wave amplitudes,

$$B(\Xi^-_-)-\sqrt{2}B(\Xi^0_0)\sim\tfrac{1}{9}M_{\{27\}}+\tfrac{1}{30}\sqrt{30}\,\tilde{M}_{\{10\}}\sim\tilde{M}_{\{10\}},$$

in view of $A(\Sigma^+_+)\approx 0$. Similarly, $\delta_B(LS)\sim\tilde{M}_{\{10\}}$ (cf. Ref. 21).

[26]E. Fischbach, Phys. Rev. 170, 1398 (1968).

[27]There are four reduced matrix elements: $M^s_{\{8s\}}$, $M^a_{\{8s\}}$, $\tilde{M}^s_{\{8a\}}$, and $\tilde{M}^a_{\{8a\}}$. In view of the sum rules (5)–(7) and (14) (with vanishing right-hand sides) only two of the four independent amplitudes are strangeness-changing.

[28]It is assumed that the Hamiltonian is CP-conserving so that the irregular currents satisfy relations (3). If both the regular and irregular currents are Hermitian and thus the Hamiltonian violates CP invariance [cf. N. Cabibbo, Phys. Letters 12, 137 (1964)], the octet components of the $\Delta S=0$ and the $\Delta S=1$ parts of the Hamiltonian are members of the same octet.

[29]This is to be contrasted with the situation in theories involving regular neutral currents (cf. Ref. 14), or in certain theories with CP nonconservation [cf. L. R. Ram Mohan, Phys. Rev. D 4, 825 (1971)]. Here the $\Delta S=0$, $\Delta T=1$ and the $\Delta S=1$, $\Delta T=\frac{1}{2}$ amplitudes are still related, although the sum rules may differ. In the d'Espagnat model, for example,

$$A(n^0_-)=-\tfrac{2}{3}\sqrt{3}\,a[A(\Xi^-_-)-2A(\Lambda^0_-)],$$

with $a=\frac{1}{2}\sqrt{2}\cot\theta$ [cf. E. Fischbach and K. Trabert, Phys. Rev. 174, 1843 (1968)]. The statement about the independence of the $\Delta S=0$, $\Delta T=1$ and $\Delta S=1$, $\Delta T=\frac{1}{2}$ amplitudes is true generally only if neutral irregular currents are not considered. If the Hamiltonian contains neutral irregular currents in addition to the charged ones [cf. P. Herczeg, Phys. Rev. D 4, 1239 (1971)], the $\Delta S=0$, $\Delta T=1$ and $\Delta S=1$, $\Delta T=\frac{1}{2}$ components can be arranged to belong to the same octet.