Theoretical Aspects of Nonleptonic Hyperon Decays*

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Recent experimental results on nonleptonic hyperon decays are studied on the basis of a doublet approximation for strong and weak interactions, with the implied suggestion that this higher symmetry may be more easily discernable in such reactions in which K-particles do not occur explicitly. The doublet approximation is characterized by a doublet spin I which is equal to $\frac{1}{2}$, 1, 0 for baryons, π, K, respectively and by a K spin. It is not necessary to assume that the strong K interactions are weak compared to the strong π interactions. For the mentioned reactions it is necessary to assume that the strong interactions which do not conserve I play a minor role compared to those which conserve I.

The following refinement of the nonleptonic $\Delta T=\frac{1}{2}$ rule is proposed. (T=isotopic spin.) The weak nonleptonic interactions consist of two parts $H^{(0)}$, $H^{(1)}$ with $\Delta I=0$, 1, respectively. In the doublet approximation $H^{(0)}$ and $H^{(1)}$ separately conserve parity in the presence of all strong π and K interactions. $H^{(0)}$ and $H^{(1)}$ together do not conserve parity, however. In addition to $\Delta I=1$, $H^{(1)}$ should in general satisfy a further constraint, but there are classes of graphs for which $\Delta I=1$ is sufficient.

Current $\times$ current structures for $H^{(0)}$ and $H^{(1)}$ are examined. Results of an earlier paper can be viewed as a special case of the $\Delta I=0$, 1 rule. The same is true for results obtained by Feldman, Matthews, and Salam and by Wolfenstein. The considerations of these authors can be extended to wider classes of graphs.

Odd relative helicity and the relation between rates for $\Lambda\to p+\pi^-$, $\Sigma^+\to p+\pi^0$ are consequences of the $\Delta I=0$, 1 rule only. So is the prediction that Ξ decay is strongly P nonconserving.

The parity properties of $H^{(0)}$, $H^{(1)}$ are sufficient conditions. It is a delicate question whether they are necessary. For a subset of graphs they are not necessary, but this set seems arbitrary. If it is assumed that the parity conditions are necessary, the schizon scheme is ruled out.

It is noted that the nonleptonic weak interactions may be generated by the strong interactions in terms of the following prescription. $H^{(1)}$ is generated by assuming that the $\pi(K)$ fields have small $K(\pi)$ components. An $H^{(0)}$ is generated by assuming that the doublets $N_1(N_2)$ have small $N_2(N_1)$ components; likewise for N_3 and N_4. Further, it is observed that one can construct a non-electromagnetic $\Delta T=\frac{3}{2}$ interaction which is small in the sense that it only contributes to $K_{\pi2}{}^+$ to the extent that the doublet approximation is *not* valid.

I. INTRODUCTION

BEYOND the demonstration of the existence of isotopic spin (T) and strangeness (S) rules, the study of strong reactions has so far taught us little about more intimate connections between the varieties of strongly interacting particles. Attempts to consider some of the new particles as composites in terms of others have till now not produced any insight which cannot as well be reached by assuming that any one baryon, say, is neither less nor more elementary than any other. In this paper we continue[1] to adopt this last view. From this standpoint one may try to further interconnect particles and interactions by asking for stronger symmetries than those which yield T and S conservation. It is known that such symmetries cannot exist rigorously.[2] Nor is there thus far any indication that some of the strong interactions are relatively weak compared to others so that expansions in the former might be a useful procedure. Conjectures that one part of the strong interactions possesses symmetries stronger than another therefore have had as yet to remain in a speculative stage. Unless we find some qualitative clues, the strong interaction problems appear to be in somewhat of a deadlock. It is the purpose of this paper to discuss certain weak decay reactions which may possibly provide us with such a clue about the strong interactions.

The reactions we have in mind are

$$\Sigma^+\to n+\pi^+,\quad (A^+,\alpha^+),\tag{1.1}$$

$$\to p+\pi^0,\quad (A^0,\alpha^0),\tag{1.2}$$

$$\Sigma^-\to n+\pi^-,\quad (A^-,\alpha^-).\tag{1.3}$$

We shall often refer to these reactions as $\Sigma_+{}^+$, $\Sigma_0{}^+$, $\Sigma_-{}^-$, respectively. Their amplitudes and asymmetry parameters will be denoted by A, α, labeled as indicated. Experimental results[3] are compatible with the requirement of the $\Delta T=\frac{1}{2}$ rule that A^+, A^-, $A^0\sqrt{2}$ shall form a triangle. If we neglect final-state interactions, this triangle can conveniently be drawn in the so-called (s,p) plane.[4]

Concerning the $\Delta T=\frac{1}{2}$ rule (which to this author seems neither less nor more mysterious than the $\Delta T=0$ rule of the strong interactions), we shall adopt the same view as in a previous paper.[5] For the purpose of this study, the rule will be supposed to be rigorous for nonleptonic decays. At the same time we do not wish to prejudge the question whether deviations from $\Delta T=\frac{1}{2}$ are electromagnetic only.

Experiment further indicates[3] that the Σ triangle is oriented in a rather special way which can be expressed by

$$\alpha^+\simeq 0,\quad \alpha^-\simeq 0.\tag{1.4}$$

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[1] A. Pais, Phys. Rev. **86**, 663 (1952).
[2] A. Pais, Phys. Rev. **110**, 574, 1480 (1958).
[3] The most recent results are given by Cork, Kerth, Wenzel, Cronin, and Cool, Phys. Rev. **120**, 1000 (1960).
[4] M. Gell-Mann and A. Rosenfeld, in *Annual Review of Nuclear Science* (Annual Reviews, Inc., Palo Alto, California, 1957), Vol. 7, p. 454.
[5] A. Pais, Nuovo cimento, **18**, 1003 (1960). This paper will be referred to as I.

$\Delta T = \frac{1}{2}$ implies that if Σ_+^+ is nearly pure $s(p)$ wave, then Σ_-^- is nearly pure $p(s)$ wave. It is presently not known which is which. This can be decided experimentally.[6]

Equation (1.4) is rather remarkable. It shows that *insofar as* Σ_+^+ *and* Σ_-^- *are concerned* it may be a good approximation to say that Σ^+ and Σ^- each do have a well-defined parity relative to the π-nucleon system; and that the parity of Σ^+, whatever it is, is opposite to that of Σ^-.

Now either a system has a well-defined parity relative to another or it hasn't, so what does "a good approximation" mean? To see this, note the following. It is easy to give examples[7] of a weak interaction H which leads to a Σ triangle so oriented that $\alpha^+ = \alpha^- = 0$. But this by no means solves the problem.

Consider Σ^- decay as a first example. Let H be such that Σ_-^- is pure p wave, say. However, the strong interactions generally allow Σ^- to be part of the time a Σ^+, for example, $\Sigma^+ + 2\pi^-$. During this time H can induce s-wave decay Σ_+^+. The final $n\pi^-$ state can then be reached by strong reabsorption of a $\pi^+\pi^-$ pair. The net result is an s-wave contribution to Σ_-^-. Similarly, virtual Σ_0^+ decay would give a mixed (s,p) contribution to Σ_-^-. In other words, even if the weak interactions properly orient the triangle, the strong interactions in general do not respect this orientation. That is, to stick with the example, unless we could provide a reason which would inhibit virtual $\Sigma^- \to \Sigma^+$ transitions. This is possible, though not rigorously.

As a next example consider the sequence

$$\Sigma^- \to \Lambda + \pi^- \to (p + \pi^-) + \pi^- \to n + \pi^-,$$

where the strongly P-nonconserving Λ decay is involved and for the rest only P-conserving strong interactions. Why is the Σ^- nearly impervious to this nonconservation?

We would like to point out that the so-called doublet approximation[8] (DA; also known as restricted symmetry) provides a natural though not rigorous answer to these questions. Here one assumes even $\Sigma\Lambda$ parity, neglects the $\Sigma\Lambda$ mass difference, and puts

$$\Lambda = (Y^0 + Z^0)/\sqrt{2}, \quad \Sigma^0 = (-Y^0 + Z^0)/\sqrt{2}. \tag{1.5}$$

The baryons then regroup in terms of four doublets; see Eq. (2.2) below. The reason that $\Sigma^- \to \Sigma^+$ is inhibited is [Sec. 2(a)] that they belong to different doublets which in the DA do not intercombine. The reason that Σ^- is not affected by P nonconservation in Λ decay is that Σ^- cannot combine with Y^0 for the same reason. Thus, if the Y^0, Z^0 "parts" of Λ separately are P-conserving in decay, then Σ^- will stay P-conserving. While if Y^0, Z^0 have opposite parity in decay, the physical particle Λ will be strongly parity nonconserving.

These remarks may serve to indicate the general approach planned in this paper. [On purpose we have not included Σ^+ decay in these few examples, as there a more delicate problem arises, see Sec. 2(b)]. We shall endeavor to arrange things so that we get exact P conservation for Σ_+^+ and Σ_-^- in the approximation where a stronger symmetry than charge independence is supposed to hold. As has been shown,[2] the DA is the weakest symmetry stronger than charge independence. Thus the DA is the natural starting point. If we succeed we shall be able to assign to Σ^-, for example, a parity relative to the π-nucleon system, but only to the extent that the DA holds. When we talk of P conservation in certain decays, we shall always refer to a situation where a symmetry higher than charge independence is assumed. It is therefore not implied that P conservation in Σ_+^+ and Σ_-^- would be as good as it is in atomic physics for example.

The following must be strongly emphasized. The quest for stronger symmetries has so far most often seemed a self-inflicted agony. A symmetry first set up has subsequently to be broken. The difference in the present case is that the assumed symmetry leads to interesting physical conclusions which, at least so far, are not in qualitative disagreement with experiment. Even so, the question always remains: What is the influence on the decays in question of those strong interactions which do not respect the DA? We shall not face this question in this paper. As long as it is not answered, the work does not represent a theory but a program.

Remark. For reasons given in I we lay the emphasis on the near parity conservation of Σ_+^+ and Σ_-^- and consider the near equality of their rates as more of an accident. No doubt this equality will eventually be a vital clue as well. However, it does not seem to raise such a qualitative puzzle as the parity aspect does.

In this paper we shall apply the DA both to π and K couplings. We do this mainly to emphasize that on the whole it is not the essential point which *kinds* of fields and interactions follow the DA. It is not relevant therefore whether K and π couplings are of the same order of strength or not. What is relevant, on the other hand, is of course the role of the interaction which breaks down the doublet symmetry. In Sec. 2(a) we treat the π couplings in the DA and define the relevant quantum numbers.

In Sec. 2(b) a proposed refinement of the $\Delta T = \frac{1}{2}$ rule is stated. It is suggested that the nonleptonic interactions consist of two parts $H^{(0)}$ and $H^{(1)}$ which separately conserve parity but which clash when taken together. To separate any parity-nonconserving interaction in two separately P-conserving parts is of course totally trivial. One could do the same for β decay.

[6] See, for example, I, Eq. (15).

[7] $H = g_1(2^{\frac{1}{2}}\bar{n}\gamma_\lambda\Sigma^+\partial_\lambda\pi^+ + \bar{p}\gamma_\lambda\Sigma^+\partial_\lambda\pi^0) + g_2(2^{\frac{1}{2}}\bar{n}\gamma_\lambda\gamma_5\Sigma^-\partial_\lambda\pi^- + \bar{p}\gamma_\lambda\gamma_5\Sigma^+\partial_\lambda\pi^0)$ satisfies all requirements. See S. Bludman, Phys. Rev. **115**, 468 (1959). To avoid misunderstanding, it should be stressed that Bludman discusses such interactions as effective rather than as primitive couplings.

[8] See reference 2. Also A. Pais, Phys. Rev. **112**, 624 (1958), Sec. II. Section V of this last paper contains a formal basis for the **(I,K)** quantum numbers in terms of the orthogonal 4-group.

What is not trivial in the present case is that these separate parts are simultaneously subject to a condition in terms of a quantum number other than parity, namely, the doublet spin. We therefore suggest that doublet spin and parity properties are correlated in a definite way.

The main theorems on Λ decay are given in Sec. 2(d), on Ξ decay in Sec. 4.

In certain instances the DA will be insufficient for the purpose of obtaining the desired parity properties. We shall then consider two (not mutually exclusive) approaches: (a) the use of further invariance arguments, see Sec. 2(c); and (b) the investigation of special types of virtual transitions, see Sec. 6. In this second approach we follow ideas due to Feldman, Matthews, and Salam[9] and to Wolfenstein[10] and try to generalize their results.

In the work of FMS[9] some emphasis is laid on the differences of the dispersion approach as compared to the Lagrangian methods used by others. Rather than to underline such differences, the present work aims to emphasize above all what such varied techniques actually have in common. As has been stated in I, the essence of the problem seems to be the establishment of shared symmetries of weak and strong interactions. In Sec. 6(c) the connection between the FMS results and the present argument will, in fact, be established through the analysis in terms of symmetry arguments of the weak vertices used by these authors.

In Sec. 5 we discuss K-particle effects in the DA. In particular we show in Sec. 5(d) that such reactions as $K\rightarrow\pi$, $K^0\rightarrow 2\pi$, $K^+\rightarrow 3\pi$ can be described in the DA terms of the weak interactions $H^{(0)}$, $H^{(1)}$ introduced in Sec. 2(b) even though these interactions separately conserve parity. Here we meet with a very essential point that has been brought out by the work of Wolfenstein.[10] There are, in fact, specific graphs, see Sec. 6(a), which give pure opposite parity contributions to Σ_+^+ and Σ_-^-. These graphs follow the doublet rule $\Delta I=0$ or 1. But for these Wolfenstein graphs the condition that $H^{(0)}$, $H^{(1)}$ are P conserving [see Sec. 2(b)], though sufficient, *is not necessary*. On the other hand, we shall also see in Sec. 6(a) that we can retain Wolfenstein's results but extend them to a larger class of graphs if indeed, as proposed in Sec. 2(b), $H^{(0)}$ and $H^{(1)}$ separately conserve parity.

The great importance of this question lies in the following. By an argument given in I, if $H^{(0)}$ and $H^{(1)}$ separately do conserve parity, then an incompatibility exists between the rule proposed in Sec. 2(b) and the schizon scheme[11]; see also Sec. 3(f).

While most of the arguments summarized before do not have reference to specific structures of the weak interactions beyond their doublet spin properties, a certain interest attaches to the question, how can these interactions be brought in current $\times$ current form. This problem is dealt with in some detail in Sec. 3 where it is shown that the present work goes beyond I in two respects: (a) In I we used global symmetry from the start, in this paper the weaker DA is sufficient in many instances; (b) the (js,jt) coupling used in I represents a special choice for $H^{(0)}$, $H^{(1)}$. Other possibilities are noted.

The concluding Sec. 7 is mainly devoted to a few general remarks on (a) the possible generic connection between weak and strong interactions, (b) possible nonelectromagnetic deviations from $\Delta T=\frac{1}{2}$, (c) the question of the leptonic decays. We shall state which remarks made in I apply to the particular (js,jt) coupling scheme used there, and which remarks have a wider validity.

If the present approach is correct, a new question arises. Why should the DA manifest itself as a useful symmetry in nonleptonic hyperon decay but not in the reactions studied previously?[2] These last reactions all involve real K's, the decays virtual K's only. We are, therefore, led to surmise that where K particles appear only as a virtual cloud, the DA is more easily discernible. Results on π-hyperon scattering[12] may perhaps shed further light on this point.

2. THE DOUBLET APPROXIMATION (DA)

(a) Strong π Interactions

The DA for this coupling has been discussed elsewhere.[2] We briefly state the main points. It is necessary for the existence of this approximation that the (Σ,Λ) parity be even. The π couplings are considered under the neglect of the (Σ,Λ) mass difference and are of the typical form,

$$H_\pi=[G_1\bar{N}_1\boldsymbol{\tau}\gamma_5N_1+G(\bar{N}_2\boldsymbol{\tau}\gamma_5N_2+\bar{N}_3\boldsymbol{\tau}\gamma_5N_3)+G_4\bar{N}_4\boldsymbol{\tau}\gamma_5N_4]\boldsymbol{\pi}, \quad (2.1)$$

where

$$N_1=\begin{pmatrix}p\\n\end{pmatrix},\ N_2=\begin{pmatrix}\Sigma^+\\Y^0\end{pmatrix},\ N_3=\begin{pmatrix}Z^0\\\Sigma^-\end{pmatrix},\ N_4=\begin{pmatrix}\Xi^0\\\Xi^-\end{pmatrix}, \quad (2.2)$$

and where Y^0, Z^0 are given in Eq. (1.5). The space-time structure of the couplings is immaterial for the argument, as long as the pseudoscalar nature of $\boldsymbol{\pi}$ is guaranteed. In fact, we are not even committed to the form (2.1). What then is the essence of the DA?

First of all, Eq. (2.1) implies that we are free to rotate the τ spin together with π. $\boldsymbol{\tau}$ is the isotopic spin for N_1, N_4 but not for Σ, Λ. We call it the doublet spin, to which we refer in general as I. For $\boldsymbol{\pi}$ we have $T=I=1$. Secondly, we are free to rotate in the "(N_2,N_3)

[9] G. Feldman, P. Matthews, and A. Salam, Phys. Rev. **121**, 302 (1961). We refer to this paper as FMS.
[10] L. Wolfenstein, Phys. Rev. **121**, 1245 (1961).
[11] T. D. Lee and C. N. Yang, Phys. Rev. **119**, 1410 (1960).
[12] Alston *et al.*, Phys. Rev. Letters **5**, 520 (1960).

plane." We can unite N_2 and N_3 to

$$N=\begin{pmatrix}N_2\\N_3\end{pmatrix}, \tag{2.3}$$

and assign to this "doublet" (each component of which is multicomponent itself) a spin $K=\frac{1}{2}$, with $K_3=+\frac{1}{2}$ $(-\frac{1}{2})$ for N_2 (N_3). The relation between $\mathbf{T}$, $\mathbf{I}$, and $\mathbf{K}$ is

$$\mathbf{T}=\mathbf{I}+\mathbf{K}. \tag{2.4}$$

We shall see later (Sec. 5) that if K particles participate in the DA we have for them $I=0$, $T=K=\frac{1}{2}$. We call K the K spin. The DA for strong π coupling is now defined generally by the statement that (I,K,I_3,K_3) are good quantum numbers. One may simultaneously and independently apply to I and K the usual rules of the vector addition model.

Observe that H_π not only conserves baryons, but also conserves individual doublets. Thus the DA guarantees that virtual transitions $\Sigma^\pm\to\Sigma^\mp+\pi$ mesons are forbidden. This is just the inhibition we are after.

(b) $\Delta T=\frac{1}{2}$ and $\Delta I=0$, 1

We shall now suppose that the nonleptonic decay interactions do not only satisfy $\Delta T=\frac{1}{2}$ but more specifically that they also have ΔI properties. To see what such a statement means, it is instructive to reason by analogy with the $\Delta T=\frac{1}{2}$ rule itself.

If we say that nonleptonic decay reactions satisfy $\Delta T=\frac{1}{2}$, we may also say that the decay interaction satisfies $\Delta T=\frac{1}{2}$. This trivial statement is, of course, due to the fact that the strong interactions satisfy $\Delta T=0$, so that they cannot modify the ΔT properties of the weak interaction. Electromagnetic corrections are of the $\Delta T=1$ type and make the $\Delta T=\frac{1}{2}$ rule impure. Likewise, if we say that a weak interaction has a certain ΔI, this is only meaningful to the extent that the strong interactions have $\Delta I=0$—that is the DA. "Corrections" to the DA will make the ΔI rules for weak interactions impure. (As we have stated in the Introduction, we shall not discuss here the influence of such distortions on the decay processes.)

If we assign a ΔI to a weak interaction, the latter should be expressable in terms of doublets, inasfar as baryons are concerned. We are therefore working in an approximation in which the weak and strong interactions share the doublet symmetry.

All nonleptonic, Λ, Σ, Ξ decays have $|\Delta K_3|=\frac{1}{2}$. It is easily seen that actually $\Delta K=\frac{1}{2}$. As $\Delta T=\frac{1}{2}$, it follows that $\Delta I=0$ or 1, according to the vector relation

$$\Delta\mathbf{T}=\Delta\mathbf{I}+\Delta\mathbf{K}. \tag{2.5}$$

The most general decay interaction H, therefore, is of the form,

$$H=H^{(0)}+H^{(1)}, \tag{2.6}$$

where $H^{(0)}$, $H^{(1)}$ are characterized by $\Delta I=0$, 1, respectively.

For Σ decays the ΔI_3 assignments are as follows:

$$\begin{aligned}&\Sigma^-\to n+\pi^-: && |\Delta I_3|=1,\\ &\left.\begin{aligned}\Sigma^+&\to n+\pi^+\\ &\to p+\pi^0\end{aligned}\right\}: && \Delta I_3=0.\end{aligned} \tag{2.7}$$

We note that according to Eq. (2.7), $H^{(0)}$ allows Σ^+ decays but forbids Σ_-^-. $H^{(1)}$ allows, in general, all three decays, because $\Delta I=1$ implies $\Delta I_3=\pm1$, 0. Thus we can generally write $H^{(1)}$ as

$$H^{(1)}=H_1^{(1)}+H_{-1}^{(1)}+H_0^{(1)}, \tag{2.8}$$

where the subscript refers to the ΔI_3 value.

As $H^{(1)}$ allows all Σ modes, the first question is whether we could (a) restrict ourselves to the $H^{(1)}$ term in Eq. (2.6), and (b) choose the parity structure of $H_0^{(1)}$ to differ from $H_{\pm1}^{(1)}$ in such a way that Σ_+^+ (which proceeds via $H_0^{(1)}$) has opposite parity compared to Σ_-^- (which proceeds via $H_{-1}^{(1)}$). This in itself is indeed feasible but it would be in violent contradiction with $\Delta T=\frac{1}{2}$. For clearly Σ_0^+ would now also be parity conserving. Instead of a Σ triangle we would therefore have two amplitudes aligned along the s(or p) axis, the third aligned along the p(or s) axis in the (s,p) plane.

Let us digress for a moment from the main program which is to understand the parity properties of Σ decays if $\Delta T=\frac{1}{2}$ is assumed to be strictly valid. It may be worth while to note that the discussion of Eq. (2.8) shows that one can conceive of (nonelectromagnetic) violations of the $\Delta T=\frac{1}{2}$ rule of such a nature that the Σ triangle does no longer exist rigorously, while yet the parity conservation in the $n\pi^\pm$ channels remains intact. Such violations should be relatively small, however, as deviations from $\Delta T=\frac{1}{2}$ do not seem to be large.

We now return to the main program and try a different tack. Suppose that we could find an argument *additional* to the $\Delta I=1$ specification of $H^{(1)}$, in such a way that $H^{(1)}$ would contribute to Σ_0^+ but not to Σ_+^+. Then Σ_+^+ would go via $H^{(0)}$ only; Σ_-^- goes anyway via $H^{(1)}$ only; while Σ_0^+ would go via both $H^{(0)}$ and $H^{(1)}$. We shall come back at length to the construction of this additional argument. Accepting for the moment that this can be found, we can clearly achieve the parity properties of Σ decays by the following hypothesis.

$\Delta I=0$, 1 *rule*. The weak nonleptonic interactions consist of two parts $H^{(0)}$, $H^{(1)}$ with $\Delta I=0$, 1, respectively. $H^{(0)}$ and $H^{(1)}$ separately conserve parity, but clash when taken together. That is to say, $H^{(0)}+H^{(1)}$ does not conserve parity. For $H^{(1)}$, Σ_+^+ is to be inhibited by an argument *additional* to $\Delta I=1$.

This rule interlocks doublet spin and parity. As we shall see in Sec. 3, the *js*, *jt* couplings of I are special examples of $H^{(0)}$, $H^{(1)}$. The parity condition on $H^{(0)}$, $H^{(1)}$ is a sufficient condition. There may perhaps be

accidents where this condition would not be necessary, for examples see Sec. 6.

The next task is to find the additional argument concerning $H^{(1)}$. This is a more subtle problem and there are at least two avenues of approach which are by no means mutually exclusive. (1) While $\Delta I=1$ specifies $H^{(1)}$ insufficiently, nevertheless $\Delta I=1$ is adequate by itself if in the calculation of Σ^+-decay probabilities there are specific virtual transitions which are strongly predominant. This is conceivable, see Sec. 6. (2) For $H^{(1)}$ we need a stronger symmetry than the DA to reach our goal. On the one hand, it is distinctly unsatisfactory to employ strong symmetries. On the other hand, the particular symmetry we shall invoke in the next subsection will allow us at once to tie parity conservation in Σ_+^+ and Σ_-^- to parity nonconservation in Λ decay.

(c) Further Discussion of $H^{(1)}$

It does not affect $H^{(0)}$ if we subject $H^{(1)}$ to an additional symmetry argument. Indeed it is typical for the weak processes which concern us here that without loss of rigor we may subject $H^{(1)}$ and $H^{(0)}$ to different invariance requirements as long as $H^{(1)}$ (or $H^{(0)}$) shares that invariance with the strong interactions. This is true as $H^{(1)}$ and $H^{(0)}$ can never interfere because weak interactions are considered to first order only.

We can even go further. By the same token $H_{\pm1}^{(1)}$ and $H_0^{(1)}$ do not interfere. We seek for an argument which inhibits Σ_+^+, a reaction which can proceed via $H_0^{(1)}$. We shall state the argument in terms of a symmetry shared by $H_0^{(1)}$ and the strong interactions. By our reasoning it is entirely immaterial whether $H_{\pm1}^{(1)}$ shares this additional symmetry or not. (As it happens, it does not.) Now Σ_-^- proceeds via $H_{-1}^{(1)}$ only. Therefore, it remains true that this reaction is P-conserving in the DA without any additional argument.

In general we may say that the shared invariance of strong interactions and a partial weak interaction is a legitimate tool because we deal with problems linear only in the weak interactions.

The additional argument on $H_0^{(1)}$ is now that it shares with the strong interactions invariance for

$$N_1\rightarrow i\epsilon_2\tau_2N_2,\quad \pi^\pm\rightarrow -\epsilon\pi^\mp,$$
$$N_2\rightarrow i\epsilon_1\tau_2N_1,\quad \pi^0\rightarrow -\epsilon\pi^0, \tag{2.9}$$

where the ϵ's are phase factors equal to ±1. Equation (2.9) leaves Eq. (2.1) invariant provided we neglect the N_1-N_2 mass difference while

$$\epsilon=\pm1,\quad \text{for}\quad G_1=\pm G. \tag{2.10}$$

This alternative for ϵ corresponds in essence to $G^\pm$ symmetry defined[13] and discussed in I. In addition we must rotate N_3 and N_4 appropriately and, for $\epsilon=-1$ but not for $\epsilon=+1$ we must apply $N_3\leftrightarrow N_4$. For what follows in this section we need not say more about the doublets 3 and 4.

If $H_0^{(1)}$ shares the invariance under Eq. (2.9), then

$$(\Sigma^+|n\pi^+)^{(1)}=\epsilon_1\epsilon_2\epsilon(n|\Sigma^+\pi^-)^{(1)}, \tag{2.11}$$

where the superscript (1) indicates that we refer to the transition brought about by $H_0^{(1)}$ (in the presence of H_π). We require

$$\epsilon_1\epsilon_2\epsilon=-1, \tag{2.12}$$

and now apply an argument (already used in I) which was first employed by Treiman[14] in a similar context. Namely, under the neglect of final state interactions

$$(n|\Sigma^+\pi^-)^{(1)}=(\Sigma^+|n\pi^+)^{(1)}. \tag{2.13}$$

Hence it follows from Eqs. (2.11–13) that $(\Sigma^+|n\pi^+)^{(1)}$ vanishes under the stated conditions. The argument thus amounts to the following. The amplitude in question is a function of m_p, m_Σ and the momentum transfer $\Delta=(q_\Sigma-q_p)^2$. Under the conditions stated $(m_p=m_\Sigma)$ this function is equal to minus itself for all values of Δ.

It is by no means obvious that $H^{(1)}$ can be constructed so as to satisfy Eqs. (2.9) and (2.12). In fact we shall see in Sec. 3 that several expressions for $H^{(1)}$, specified by $\Delta I=1$ only, will have to be discarded if Eqs. (2.9) and (2.12) hold true. Thus the argument restricts the dynamical form of the weak interactions. There are, however, several possibilities for $H^{(1)}$ which do satisfy the requirements. From now on I call these the allowed forms of $H^{(1)}$.

(d) Two Theorems on P-Nonconservation in Λ Decay

(A) It follows from the $\Delta I=0$, 1 rule that parity is not conserved in Λ decay in the same approximation that parity conserved in Σ_+^+ and Σ_-^-, provided $H^{(1)}$ is of the allowed form. Proof: Equation (2.9) also implies

$$(Y^0|p\pi^-)^{(1)}=\epsilon_1\epsilon_2\epsilon(p|Y^0\pi^+)^{(1)}. \tag{2.14}$$

Hence the argument which led to $(\Sigma^+|n\pi^+)^{(1)}=0$ also gives $(Y^0|p\pi^-)^{(1)}=0$. Thus $Y^0\rightarrow p+\pi^-$ proceeds only via $H^{(0)}$. On the other hand, $Z^0\rightarrow p+\pi^-$ is a $\Delta I_3=-1$ transition and can therefore proceed only via $H^{(1)}$. But $H^{(0)}$ and $H^{(1)}$ clash in parity. Therefore from Eq. (1.5) it follows that P is not conserved in Λ decay.

Actually this result can be sharpened considerably.

(B) It follows from the $\Delta I=0$, 1 rule and for allowed $H^{(1)}$ that

$$\alpha^0=-\alpha_\Lambda, \tag{2.15}$$

where the left (right) side of this equation refers to the helicity of the proton in $\Sigma^+\rightarrow p\pi^0$, $(\Lambda\rightarrow p\pi^-)$.

[13] See reference 5, Eq. (12). $G^\pm$ symmetry implies $G_1=\pm G=G_4$. Here we do not always need to state the relative magnitude of G_4. It is possible but unnecessary to consider more general phase factors in Eqs. (2.9–10).

[14] S. B. Treiman, Nuovo cimento **15**, 916 (1960). I am indebted to Professor Treiman for discussions on this point.

To prove this statement, we note that by (A) the transition $Y^0 \rightarrow p\pi^-$ proceeds via $H^{(0)}$ only. But $H^{(0)}$ has $\Delta I=0$ and therefore satisfies doublet charge symmetry, as a result of which

$$(Y^0|p\pi^-)=(\Sigma^+|n\pi^+). \tag{2.16}$$

This result is a consequence of the DA only. Equation (2.16) was also obtained in I but under much more restrictive conditions.[15]

The second part of the proof consists in showing that

$$(Z^0|p\pi^-)=(\Sigma^-|n\pi^-), \tag{2.17}$$

provided $H^{(1)}$ is of the allowed form. Equation (2.15) follows from Eqs. (2.16) and (2.17) by an argument given in I.

The verification of Eq. (2.17) has to wait till Sec. 3(d). We have now, in fact, pushed the argument as far as is feasible independently of the structure of $H^{(0)}$ and $H^{(1)}$. The next task is to consider these dynamical structures more closely.[16] Concerning the symmetries used, we shall arrive at the following conclusions. (a) For Σ_-^- to be P conserving, the DA is sufficient. (b) The same is true for Eq. (2.16) which makes up half of the relation (2.15). (c) For Σ_+^+ to be P conserving, the additional argument [see Eqs. (2.11–13)] goes beyond the DA but the Ξ-nucleon mass difference may retain its actual value. (d) The same is true for Eq. (2.17). Not until Sec. 6 shall we see that the conditions mentioned under (c), (d) may also be weakened to the DA symmetry if certain virtual transitions predominate.

3. DOUBLET SPIN STRUCTURES FOR WEAK INTERACTIONS

(a) Baryon Currents. Examples

To begin with, we consider $\Delta T=\frac{1}{2}$, $|\Delta S|=1$ interactions of the form (baryon current, $\Delta S=0$)$\times$(baryon current, $|\Delta S|=1$). The results so obtained are immediately applicable to more general situations. In this section we shall have no need to specify the space-time structure of currents.

Consider first $\Delta S=0$, $T=1$ currents without any DA assumption. These are bilinear in (N_1,N_1); (Σ,Σ); (Σ,Λ); or (N_4,N_4) and are easy to write down. We shall now use the following device. *Even in the presence of the* Σ, Λ *mass difference* we shall express the currents in terms of $N_1, \cdots, N_4$ of Eq. (2.2). We consider Y^0, Z^0 as mathematical constructs defined by Eq. (1.5) in terms of the real particles Λ, Σ^0.

Of course, we shall use expressions in terms of doublets with the ulterior motive to go to the DA. It is, however, quite essential to realize the following. If we assume the $\Delta T=\frac{1}{2}$ rule to be rigorous (barring electromagnetism), then any weak interaction should satisfy $\Delta T=\frac{1}{2}$ not only in the DA but also in the actual split (Σ,Λ) situation. The device just mentioned guarantees from the outset that this requirement is met. The need to bear this point in mind was first emphasized by Treiman.[14]

With the exclusion of a remark to be made in Sec. 7, we shall ignore in this paper the question whether the procedure just mentioned is strictly necessary. This question is tied to whether or not deviations from $\Delta T=\frac{1}{2}$ are indeed purely electromagnetic.

The most general form for $\Delta S=0$, $T=1$ baryon currents is

$T=1$:

$$\alpha_1\bar{N}_1\boldsymbol{\tau}N_1+\alpha(\bar{N}_2\boldsymbol{\tau}N_2+\bar{N}_3\boldsymbol{\tau}N_3)+\alpha_4\bar{N}_4\tau N_4+\alpha'\mathbf{j}', \tag{3.1}$$

where $\boldsymbol{\tau}$ is again the doublet spin. We shall use the notation

$$\tau^{\pm}=\frac{1}{\sqrt{2}}(\tau_1\pm i\tau_2);\quad j^{\pm}=\frac{1}{\sqrt{2}}(j_1\pm ij_2). \tag{3.2}$$

$\mathbf{j}'$ is given by

$$\begin{aligned} &j^{+\prime}=\bar{N}_2N_3\sqrt{2};\quad j^{-\prime}=\bar{N}_3N_2\sqrt{2};\\ &j_3{}'=-(\bar{N}_2N_2-\bar{N}_3N_3). \end{aligned} \tag{3.3}$$

Consider next the $|\Delta S|=1$, $T=\frac{1}{2}$ current. Write it first in terms of the physical baryons, then transcribe to the doublet language. The most general result is

$$|\Delta S|=1,\quad T=\tfrac{1}{2}:\quad \beta_1s_1+\beta_2s_2+\beta_3s_3+\beta_4s_4, \tag{3.4}$$

with

$$\begin{aligned} &s_1=\begin{pmatrix}\bar{N}_3N_1\\ \bar{N}_2N_1\end{pmatrix};\quad s_2=\begin{pmatrix}\bar{N}_4N_2\\ -\bar{N}_4N_3\end{pmatrix},\\ &s_3=\begin{pmatrix}\bar{N}_3\tau_3N_1+\bar{N}_2\tau^-N_1\sqrt{2}\\ -\bar{N}_2\tau_3N_1+\bar{N}_3\tau^+N_1\sqrt{2}\end{pmatrix};\\ &s_4=\begin{pmatrix}-\bar{N}_4\tau_3N_2+\bar{N}_4\tau^-N_3\sqrt{2}\\ -\bar{N}_4\tau_3N_3-\bar{N}_4\tau^+N_2\sqrt{2}\end{pmatrix}. \end{aligned} \tag{3.5}$$

For any s we shall often use the notation

$$s=\begin{pmatrix}s^-\\ s^0\end{pmatrix};\quad \bar{s}=\begin{pmatrix}s^+\\ \bar{s}^0\end{pmatrix}. \tag{3.6}$$

The method of first writing down currents in terms of the usual baryons and then transcribing to doublets is too roundabout. A simpler procedure is the following.

(b) Baryon Currents. General Method

Just as $\boldsymbol{\tau}$ acts on the doublet spin components we introduce $\boldsymbol{\rho}$ which acts on the K-spin components of

[15] See reference 5, Eqs. (32) and (33).

[16] In I we consider an instance (called the first possibility) where Eq. (2.15) holds true, even though $\alpha^+=\alpha^-=0$ is not explicitly realized. This is the case for G^- symmetry and pure js coupling. This shows that if Eq. (2.15) were to be in agreement with experiment, it would not be possible to conclude with definiteness that the weak interactions are of the form $H^{(0)}+H^{(1)}$. Nevertheless it is gratifying that this form of H leaves no ambiguity in the value of this relative helicity.

N, see Eq. (2.3);

$$\rho_1=\begin{pmatrix}0 & 1\\ 1 & 0\end{pmatrix},\quad \rho_2=\begin{pmatrix}0 & -i\\ i & 0\end{pmatrix},\quad \rho_3=-\begin{pmatrix}1 & 0\\ 0 & -1\end{pmatrix},$$
$$\rho^{\pm}=\frac{1}{\sqrt{2}}(\rho_1\pm i\rho_2). \tag{3.7}$$

Note the minus sign in the definition of ρ_3. We can now write Eq. (3.3) in the compact form, $\mathbf{j}'=\bar{N}\boldsymbol{\rho}N$.

Thus we can look upon the $T=1$ current of Eq. (3.1) as follows. The first three terms correspond to $I=1$, $K=0$; we can in fact write the α term of Eq. (3.1) as $\alpha\bar{N}\boldsymbol{\tau}N$. The j' term has $I=0$, $K=1$. This is an example of the general rule that we will get all currents by a vector addition procedure of all possible $\mathbf{I}$ and $\mathbf{K}$ to the desired $\mathbf{T}$. This formal procedure is independent of the (Σ,Λ) mass difference. (Of course, only then can I and K serve as good quantum numbers if this difference is neglected.) The following lemmas will be obvious:

(1) Any bilinear baryon current; whether $\Delta S=0$ or 1 can only have $I=0$ or 1.

(2) S-conserving currents have $K=0$ or 1.

(3) $|\Delta S|=1$ currents have $K=\frac{1}{2}$.

Thus one can write down the following complete list of currents.

$\Delta S=0$:

$$T=0;\quad I=K=0:\qquad \rho=\eta_1\bar{N}_1N_1+\eta\bar{N}N+\eta_4\bar{N}_4N_4, \tag{3.8}$$

$$I=K=1:\qquad \rho'=\bar{N}\boldsymbol{\rho}\boldsymbol{\tau}N, \tag{3.9}$$

$$T=1;\quad I=1,\, K=0:\qquad \mathbf{j}=\alpha_1\bar{N}_1\boldsymbol{\tau}N_1+\alpha\bar{N}\boldsymbol{\tau}N+\alpha_4\bar{N}_4\boldsymbol{\tau}N_4, \tag{3.10}$$

$$I=0,\, K=1:\qquad \mathbf{j}'=\bar{N}\boldsymbol{\rho}N, \tag{3.11}$$

$T=2$; $I=K=1$. Current is v^{T_3}, where

$$\begin{aligned}
v^2&=\bar{N}\rho^-\tau^-N,\\
v^1&=-(1/\sqrt{2})\bar{N}(\rho^-\tau_3+\rho_3\tau^-)N,\\
v^0&=(\sqrt{\tfrac{2}{3}})\bar{N}[\rho_3\tau_3-\tfrac{1}{2}(\rho^-\tau^++\rho^+\tau^-)]N,\\
v^{-1}&=(1/\sqrt{2})\bar{N}(\rho^+\tau_3+\rho_3\tau^+)N,\\
v^{-2}&=\bar{N}\rho^+\tau^+N.
\end{aligned} \tag{3.12}$$

$|\Delta S|=1$:

$$T=\tfrac{1}{2};\quad I=0,\, K=\tfrac{1}{2}:\qquad s=\beta_1s_1+\beta_2s_2, \tag{3.13}$$

$$I=1,\, K=\tfrac{1}{2}:\qquad s'=\beta_3s_3+\beta_4s_4, \tag{3.14}$$

[See Eq. (3.5).]

$T=\frac{3}{2}$; $I=1$, $K=\frac{1}{2}$. Current is $u^{T_3}=\zeta_1u_1^{T_3}+\zeta_4u_4^{T_3}$;

$$\begin{aligned}
&u_1^{\frac{3}{2}}=-\bar{N}_3\tau^-N_1,\qquad u_4^{\frac{3}{2}}=-\bar{N}_4\tau^-N_2,\\
&u_1^{\frac{1}{2}}=\sqrt{\tfrac{2}{3}}[\bar{N}_3\tau_3N_1-(1/\sqrt{2})\bar{N}_2\tau^-N_1],\\
&u_4^{\frac{1}{2}}=\sqrt{\tfrac{2}{3}}[\bar{N}_4\tau_3N_2+(1/\sqrt{2})\bar{N}_4\tau^-N_3],\\
&u_1^{\frac{1}{2}}=\sqrt{\tfrac{2}{3}}[\bar{N}_2\tau_3N_1+(1/\sqrt{2})\bar{N}_3\tau^+N_1],\\
&u_4^{-\frac{1}{2}}=\sqrt{\tfrac{2}{3}}[-\bar{N}_4\tau_3N_3+(1/\sqrt{2})\bar{N}_4\tau^+N_2],\\
&u_1^{-\frac{3}{2}}=\bar{N}_2\tau^+N_1,\qquad u_4^{-\frac{3}{2}}=-\bar{N}_4\tau^+N_3.
\end{aligned} \tag{3.15}$$

All currents have the indicated T properties for the actual values of the Σ, Λ masses. We next discuss the possible structures of $H^{(0)}$ and $H^{(1)}$ in terms of these currents. As $H^{(0)}$, $H^{(1)}$ have definite I properties, we shall now have to use the currents in their true doublet form.

(c) Structure of $H^{(0)}$

Let (I_0,K_0) denote the (I,K) values of the $\Delta S=0$ current; likewise (I_1,K_1) refers to $\Delta S=1$. To construct an $H^{(0)}$ we need $I_0=I_1=0$ or 1. In either case K_0 can be equal to zero or one. Thus there are four possibilities.

$$I_0=I_1=K_0=0:\qquad \rho(s^0+\bar{s}^0). \tag{3.16}$$

[See Eqs. (3.6), (3.8), (3.13).]

$$I_0=I_1=0;\quad K_0=1:\qquad \bar{N}\rho^+Ns^--(1/\sqrt{2})\bar{N}\rho_3Ns^0+\text{h.c.} \tag{3.17}$$

[See Eqs. (3.7), (3.11), (3.13).]

$$I_0=I_1=1;\quad K_0=0:\qquad \mathbf{j}\mathbf{t}, \tag{3.18}$$

$$\mathbf{t}=\zeta_1(\mathbf{t}_{12}+\mathbf{t}_{21})+\zeta_4(\mathbf{t}_{34}+\mathbf{t}_{43}),\qquad \mathbf{t}_{ij}=\bar{N}_i\boldsymbol{\tau}N_j. \tag{3.19}$$

Here we meet a typical recoupling problem. Couple the currents of Eqs. (3.10) and (3.15) together to $\Delta T=\frac{1}{2}$. The answer is

$$\mathbf{j}\mathbf{t}-(\sqrt{2}/3)(j^+s'^--(1/\sqrt{2})j^0s'^0+\text{h.c.}),$$

where the term in brackets itself has $\Delta T=\frac{1}{2}$; see Eq. (3.14). Equation (3.18) is the jt coupling discussed in I.

$I_0=I_1=1$, $K_0=1$:

$$\bar{N}\rho^+\boldsymbol{\tau}N(\beta_1\bar{N}_3\boldsymbol{\tau}N_1+\beta_2\bar{N}_4\boldsymbol{\tau}N_2)-(1/\sqrt{2})\bar{N}\rho_3\boldsymbol{\tau}N(\beta_1\bar{N}_2\boldsymbol{\tau}N_1-\beta_2\bar{N}_4\boldsymbol{\tau}N_3)+\text{h.c.} \tag{3.20}$$

There is an obvious structural connection between Eqs. (3.17) and (3.20).

(d) Allowed Structure of $H^{(1)}$

As was stated in Sec. 2(c), we mean by this an interaction with $\Delta I=1$ and which shares with the strong interactions the invariance under Eq. (2.9) with the condition (2.12).

The transformation (2.9) is the product of an I-spin rotation and a $1\leftrightarrow 2$ substitution. The currents of

Eqs. (3.9), (3.11), (3.12) all contain $\bar{N}_2N_3$ and $\bar{N}_3N_2$. Such terms cannot possibly respect $1\leftrightarrow 2$. Hence, ρ', $\mathbf{j}'$, and v cannot appear in the allowed structures of $H^{(1)}$. There remain three possibilities, all of which with $K_0=0$.

$$I_0=1,\ I_1=0:\quad j^+s^- -(1/\sqrt{2})j^0s^0+\text{h.c.} \tag{3.21}$$

This is the interaction called (js) in I.

$$I_0=1,\ I_1=1:\quad j^+s'^- -(1/\sqrt{2})j^0s'^0+\text{h.c.} \tag{3.22}$$

We shall call this interaction (js').

$$I_0=0,\ I_1=1:\qquad \rho(s'^0+\bar{s}'^0). \tag{3.23}$$

We denote this coupling by $(\rho s')$. It is easily shown that the interactions (3.21–23) do share the invariance for the transformation (2.9) and do obey the condition (2.12) provided that

$$\left.\begin{array}{ll}(js): & \alpha_1=\pm\alpha\\ (js'): & \alpha_1=\mp\alpha\\ (\rho s'): & \eta_1=\pm\eta\end{array}\right\}\quad\text{for}\quad \epsilon=\pm 1. \tag{3.24}$$

Thus, see Eq. (2.10), allowed structures for $H^{(1)}$ exist for both G^+ and G^- invariance.

We are now ready to derive Eq. (2.17) and thus complete the proof of Theorem (B), Sec. 2(d). One shows, in fact, that under the conditions (3.24), $H_{-1}^{(1)}$ shares with H_π the invariance under

$$\begin{array}{ll} N_3\to i\epsilon_1'\tau_2N_1, & \pi^\pm\to -\epsilon\pi^\mp,\\ N_1\to i\epsilon_3'\tau_2N_3, & \pi^0\to -\epsilon\pi^0,\end{array} \tag{3.25}$$

provided the phases satisfy

$$\epsilon_1'\epsilon_3'\epsilon=-1. \tag{3.26}$$

Equation (2.17) follows from Eqs. (3.25–26) by the same argument used in connection with Eqs. (2.11–13). Hence, Eq. (2.17) has been derived both for G^+ and G^- invariance.[16]

It may be useful to state the conditions under which the helicity relation (2.15) holds. For Eq. (2.16) the DA is sufficient. For Eq. (2.17) it is insufficient. Following the various transformations, one concludes that Eq. (2.17) nevertheless holds in the presence of the true Ξ-nucleon mass difference for G^+ symmetry. G^- necessitates full global symmetry.

(e) π Currents. Space Parity and $G^\pm$ Symmetry

The baryon currents of the foregoing sections may be completed with meson currents. Here we consider the π field only. (Currents involving K's occur in Sec. 5.) For π we have $I=1$, $K=0$. The only baryon current to which π terms may be added is j of Eq. (3.10). Representatives are (∂ denotes a spatial derivative)

$$\mathbf{j}_{(\pi)A}=\partial\boldsymbol{\pi};\quad \mathbf{j}_{(\pi)V}=\boldsymbol{\pi}\times\partial\boldsymbol{\pi}. \tag{3.27}$$

The $\Delta I=0$, 1 rule and all that follows hold true for $\mathbf{j}_A\to\mathbf{j}_A+\mathbf{j}_{(\pi)A}$; $\mathbf{j}_V\to\mathbf{j}_V+\mathbf{j}_{(\pi)V}$ as long as the DA *only* is invoked. But for the stronger $G^\pm$ symmetries used in Eq. (3.24) something new happens. According to Eq. (3.24) the $\mathbf{j}$ current which enters in the $G^\pm$ case is[17]

$$\mathbf{j}^\pm=\bar{N}_1\boldsymbol{\tau}N_1\pm\bar{N}\boldsymbol{\tau}N+\bar{N}_4\boldsymbol{\tau}N_4. \tag{3.28}$$

Under the transformations $(N_1\leftrightarrow N_2;\ N_3\leftrightarrow N_4)$; $(N_1\leftrightarrow N_3;\ N_2\leftrightarrow N_4)$ we have $\mathbf{j}^\pm\to\pm j^\pm$ while for $G^\pm$: $j_{(\pi)A}\to\pm j_{(\pi)A}$; $\mathbf{j}_{(\pi)V}\to\mathbf{j}_{\pi(V)}$. Hence for $G^\pm$ we may add $j_{(\pi)A}$ to $j_A^\pm$ and still all arguments of Sec. 3(d) hold true. But while we can add $j_{(\pi)V}$ to j_V (for G^+), we can*not* add $j_{(\pi)V}$ to j_V (for G^-). This indicates that we can further restrict the allowed structure of $H^{(1)}$ by arguments concerning conserved currents.

The phenomenon just described is due to the fact that with respect to the group G^- (not G^+) a nontrivial parity is introduced in isotopic space. For π currents this isotopic parity is linked to the space parity.

(f) Structure of the Nonleptonic Decay Interaction

This completes the survey of the doublet spin structure of $H^{(0)}$ and $H^{(1)}$. The next question is how one guarantees that these two interactions separately are P-conserving but clash when taken together. In I this question has been discussed for the particular choice $H^{(0)}=jt$, $H^{(1)}=js$ [see I, Eq. (48)] but all arguments apply equally well to any allowed $H^{(0)}$, $H^{(1)}$. The same is true for the possibility mentioned in I that all S-nonconserving currents are either all pure V or all pure A. Furthermore, a general argument was given in I that showed the parity clash idea and the universal Fermi interaction (total current)×(total current) to be incompatible. In particular the $\Delta I=0$, 1 rule and the schizon scheme are mutually exclusive. See, however, the remarks on this question in Sec. 6(a).

The particular coupling scheme discussed in I is clearly not unique. It is not the purpose of the present paper to express preferences for one or another form of $H^{(0)}$, $H^{(1)}$. It is remarkable, however, to note the following. The $\Delta I=0$, 1 rule and a coupling of ($\Delta S=0$, $T=1$)×($|\Delta S|=1$, $T=\frac{1}{2}$) currents are actually compatible provided the S-conserving current does not only contain $I=1$, $K=0$ but also $I=0$, $K=1$. In fact, from Eqs. (3.17) and (3.21) we derive the following. If we couple the $T=1$ current,

$$\bar{N}_1\gamma_\lambda\boldsymbol{\tau}N_1\pm\bar{N}\gamma_\lambda(\boldsymbol{\tau}+\text{const}\,\boldsymbol{\varrho}\gamma_5)N+\bar{N}_4\gamma_\lambda\boldsymbol{\tau}N_4, \tag{3.29}$$

to the $|\Delta S|=1$, $T=\frac{1}{2}$ current s given by Eq. (3.13), then all requirements of the $\Delta I=0$, 1 rule can be fulfilled. In Eq. (3.29) we have exemplified a space-time structure to which corresponds a pure (V or A) S-violating current. It is interesting to note that in such a coupling scheme (unlike the one discussed in I) the S-nonconserving current is purely of the $T=\frac{1}{2}$ kind.

[17] For simplicity we put $\alpha_4=\alpha_1$.

Finally, we note that the relation given in Eq. (4) of I between the rates of Σ^+ and Λ decay is a consequence of the $\Delta I=0$, 1 rule rather than of the more specific form of coupling used in that paper.

4. THEOREM ON P-NONCONSERVATION IN Ξ DECAY

In accordance with the DA we consider the amplitude for $\Xi^- \to \Lambda+\pi^-$ as the sum of the amplitudes for

$$\Xi^- \to Y^0+\pi^-, \quad \Delta I_3=-1, \quad \Delta I=1, \tag{4.1}$$

$$\Xi^- \to Z^0+\pi^-, \quad \Delta I_3=0. \tag{4.2}$$

The ΔI_3, ΔI as far as specified are also obvious. But the ΔI for reaction (4.2) cannot be fixed by an argument similar to the one given for $Y^0 \to p+\pi^-$ in Sec. 2(d).

However, let us ask if it is possible to relate Ξ decay to the Σ decays. For this we should relate $(\Xi^-|Y^0\pi^-)$, $(\Xi^-|Z^0\pi^-)$ to $(\Sigma^+|n\pi^+)$, $(\Sigma^-|n\pi^-)$. Such relations are possible if and only if we consider situations more degenerate than the DA. The general method for judging what the possibilities are is the following. The DA implies that the I group may be used. Additional degeneracy implies additional invariance for substitutions between such doublets as are taken degenerate. We ask for the possible and minimal degeneracies which relates Ξ to Σ decay. These are the following.

(a) Let the strong interactions share with $H^{(1)}$ the invariance for

$$N_4 \to \xi_3 N_3, \quad N_2=\xi_1 N_1, \quad \pi \to \epsilon\pi. \tag{4.3}$$

It follows that

$$(\Xi^-|Y^0\pi^-)=\epsilon\xi_1\xi_3(\Sigma^-|n\pi^-). \tag{4.4}$$

(b) Let the strong interactions share with $H^{(0)}$ the invariance for

$$\begin{aligned} &N_4 \to i\xi_2'\tau_2 N_2, \quad N_3 \to i\xi_1'\tau_2 N_1, \\ &\pi^\mp \to -\epsilon\pi^\pm, \qquad \pi^0 \to -\epsilon\pi^0; \end{aligned} \tag{4.5}$$

then

$$(\Xi^-|Z^0\pi^-)=\epsilon\xi_1'\xi_2'(\Sigma^+|n\pi^+). \tag{4.6}$$

As $(\Sigma^+|n\pi^+)$ obeys $\Delta I=0$, it follows from Eq. (4.6) that the reaction (4.2) proceeds via $\Delta I=0$ under the stated conditions. Thus, according to Eq. (1.6) the minimal conditions which relate Ξ to Σ decay imply that parity is violated in $\Xi^- \to \Lambda\pi^-$ in the same approximation that parity is conserved in Σ_+^+ and Σ_-^-.

This is the counterpart of theorem (A) for Λ decay, see Sec. 2(d). Theorem (B) has no analog. That is, unlike the phases in Eqs. (2.16), (2.17), those in Eqs. (4.4), (4.6) are not uniquely fixed.

We show this by one counter-example. Let $H^{(0)}$ be of the type jt. Then one easily shows: $\epsilon\xi_1'\xi_2'=\pm1$ for $\zeta_1=\pm\zeta_4$ [see the definitions in Eq. (3.19)]. Let $H^{(1)}$ be of the type js. One proves: $\epsilon\xi_1\xi_3=\pm1$ for $\beta_1=\pm\beta_2$. Apparently one needs not only an argument about the structure of the interaction, but an even more detailed argument about the structure of the S-nonconserving current. At any rate we have $|\alpha_\Xi|=|\alpha_\Lambda|$ for full global symmetry.

5. INFLUENCE OF K-PARTICLE EFFECTS

(a) Strong K Couplings

It is the purpose of this section to study the extent to which the $\Delta I=0$, 1 rule and all subsequent statements can be upheld in the presence of strong K interactions, but only inasfar as the latter respect the DA. Thus we need in particular those strong K couplings which satisfy $\Delta I=0$, just like H_π of Eq. (2.1). For the present we do not speculate on whether these specific K couplings do or do not form a major part of all K interactions. Rather do we ask, if they exist what is their influence.

To start with, we follow a procedure similar to the one of Sec. 3(a). Let H_K be the most general K coupling bilinear in baryons,[18] linear in K. Transcribe this general coupling in terms of doublets, without implying any mass degeneracy. Write K as a spinor,

$$K=\begin{pmatrix} K^+ \\ K^0 \end{pmatrix}.$$

Then

$$H_K=H_K^{(0)}+H_K^{(1)}, \tag{5.1}$$

$$H_K^{(0)}=\bar{s}K+s\bar{K}, \quad (\Delta I=0), \tag{5.2}$$

$$H_K^{(1)}=\bar{s}'K+s'\bar{K}, \quad (\Delta I=1). \tag{5.3}$$

Here s and s' are the same structures as already introduced in Eqs. (3.13–14). (Of course the constants β in those equations have a different magnitude here.) The products in Eqs. (5.2–3) are in the usual sense of spinor multiplication, $s\bar{K}=s^-K^-+s^0\bar{K}^0$, etc.

If we now go to the DA, then these two couplings have distinct ΔI properties, as indicated. To verify this, remember the assignments of Eqs. (3.13–14) for s, s' and use $I=0$; $T=K=\frac{1}{2}$ for the K particles.

Thus $H_K^{(0)}$ respects the DA and is an interaction discussed previously.[19] (The S_1, S_2 of previous work[2] are equal to $\frac{1}{2}S+K_3$; $\frac{1}{2}S-K_3$, respectively.) $H_K^{(1)}$ breaks the DA through interference with itself and with $H_K^{(0)}$. These are consequences of the ΔI assignment, not of the particular trilinear structure which has only been mentioned to exemplify the argument.

We now ask if the previous results hold true if we include $H_K^{(0)}$ in the strong interactions and to what extent the answers depend on the characteristic relative parities of K-particle physics. These are (1) the parity $P(K^+)$ of charged K's relative to Λ nucleon, (2) the parity $P(\Xi)$ of cascade relative to nucleon, and (3) the parity $p(K)$ of charged relative to neutral K particles.[8]

It is therefore necessary to write out the interaction

[18] See reference 2, Eqs. (6)–(9).
[19] See reference 2, Eq. (12).

(5.2) in some more detail. We have

$$H_K{}^{(0)}=[F_1\bar{N}_1O_{12}N_2+F_2\bar{N}_3O_{34}N_4]K_0 +[F_1\bar{N}_1O_{13}N_3-F_2\bar{N}_2O_{24}N_4]K^+ +\text{h.c.}, \quad (5.4)$$

where all O operators in essence represent either 1 or $i\gamma_5$. The parity possibilities are completely specified as follows:

$$\begin{aligned} &O_{13}=i\gamma_5(1) \quad \text{for} \quad P(K^+) \text{ odd (even)},\\ &P(\Xi)\text{ even: } O_{12}=O_{34},\quad O_{13}=O_{24},\\ &P(\Xi)\text{ odd: } O_{12}\neq O_{34},\quad O_{13}\neq O_{24},\\ &p(K)\text{ even: } O_{12}=O_{13},\quad O_{24}=O_{34},\\ &p(K)\text{ odd: } O_{12}\neq O_{13},\quad O_{24}\neq O_{34}. \end{aligned} \quad (5.5)$$

Note that in principle one can dispose independently over $P(\Xi)$ and $p(K)$.

We now observe that the assignment $\Delta I=0$ to $H_K{}^{(0)}$ is independent of $P(K^+)$, $P(\Xi)$, $p(K)$. This is true because each of the four terms of Eq. (5.4) satisfy $\Delta I=0$ individually. Hence, as explained in Sec. 2(b), it remains a meaningful procedure to assign a ΔI to a weak interaction, whatever the parities are which enter in $H_K{}^{(0)}$.

It follows, therefore, that the $\Delta I=0$, 1 rule of Sec. 2(b) can be maintained as long as the additional argument for $H^{(1)}$ can be upheld in the presence of $H_K{}^{(0)}$. We shall see presently that this argument can be fully maintained if $P(\Xi)$ is even, but that some P-nonconservation may occur if $P(\Xi)$ is odd. All arguments will turn out to be independent of $p(K)$, however.

A remark on this latter parity is in order here. If $p(K)$ is odd, then I is still a good quantum number, but K is not. As has been noted elsewhere,[8] odd $p(K)$ implies deviations from $\Delta T=0$ in strong interactions, hence from $\Delta T=\frac{1}{2}$ in weak interactions. Thus, odd $p(K)$ could only then be a possibility if the virtual K interactions would play only a minor role in nonleptonic hyperon decays.

(b) Nonleptonic Hyperon Decays

In this section we discuss the additional argument for $H^{(1)}$ in the presence of $H_K{}^{(0)}$ and also the extension of the argument for Λ and Ξ decays in the presence of this I-conserving strong K interaction. In all these instances the reasoning follows the same pattern. Wherever we have used symmetries stronger than the DA, we ask if these symmetries can be extended to include shared invariance for $H_K{}^{(0)}$.

(a) Equation (2.11). Complete Eq. (2.9) as follows: $N_3\to i\epsilon_4\tau_2N_4$, $N_4\to i\epsilon_3\tau_2N_3$, $K^0\to\epsilon_1\epsilon_2\bar{K}^0$, $K^+\to\epsilon_1\epsilon_3K^+$, $\epsilon_1\epsilon_3=\epsilon_2\epsilon_4$. Equation (2.11) remains true in the presence of $H_K{}^{(0)}$ if $F_1=F_2$ and $P(\Xi)$ is even. For odd $P(\Xi)$ charged K couplings cause a deviation. The same applies to Eq. (2.14).

(b) Equation (2.16) remains valid.

(c) Equation (2.17). Complete Eq. (3.25) with $N_2\to i\epsilon_4'\tau_2N_4$, $N_4\to i\epsilon_2'\tau_2N_2$, $K^0\to -\epsilon_1'\epsilon_2'K^0$, $K^+\to \epsilon_1'\epsilon_3'K^-$, $\epsilon_1'\epsilon_3'=\epsilon_2'\epsilon_4'$. With $F_1=F_2$, $P(\Xi)$ even, Eq. (2.17) remains valid. K^0 couplings cause a deviation for odd $P(\Xi)$.

(d) Equation (4.4). Complete Eq. (4.3) with $N_3\to \xi_4N_4$, $N_1\to\xi_2N_2$, $K^+\to\xi_1\xi_3K^+$, $K^0\to\xi_1\xi_2\bar{K}^0$, $\xi_1\xi_3=\xi_2\xi_4$. Equation (4.3) remains valid for $F_1=F_2$, even $P(\Xi)$. For odd $P(\Xi)$ the K^+ couplings cause a deviation.

(e) Equation (4.6). Complete Eq. (4.5) with $N_2\to i\xi_4'\tau_2N_4$, $N_1\to i\xi_3'\tau_2N_3$, $K^0\to -\xi_1'\xi_2'K^0$, $K^+\to \xi_1'\xi_3'K^-$, $\xi_1'\xi_3'=\xi_2'\xi_4'$. Equation (4.6) remains valid for $F_1=F_2$, even $P(\Xi)$. For odd $P(\Xi)$ the K^0 couplings cause a deviation.

All results stated in this section are independent of $p(K)$.

(c) $K\pi$ Currents

Direct $K\pi$ transitions can be brought about through $K\pi$ terms in the currents which enter the weak interactions. This mechanism is additional to the one already met in the foregoing. Their inclusion does not change the essence of the previous argument.

We have specified $I=1$, $K=0$ for π, $I=0$, $K=\frac{1}{2}$ for K. Hence the only baryon currents to which $K\pi$ terms be added are t and s'; see Eqs. (3.14) and (3.19):

$$\begin{aligned} \mathbf{t}&\to\mathbf{t}+\mathbf{t}_K,\\ s'&\to s'+s_K', \end{aligned} \quad (5.6)$$

where (up to a constant)

$$\begin{aligned} \mathbf{t}_K&=\boldsymbol{\pi}(K^0+\bar{K}^0),\\ s_K'&=\begin{pmatrix} K^+\pi^0+K^0\pi^+\sqrt{2}\\ K^+\pi^-\sqrt{2}-K^0\pi^0\end{pmatrix}. \end{aligned} \quad (5.7)$$

Note that $\mathbf{t}_K$ involves neutral K particles only. The discussion of K^2 currents follows similar lines.

(d) Weak Nonleptonic K Transitions

Now that strong K and π interactions as well as decay couplings have been defined with respect to the DA, nonleptonic K decays can be discussed in this approximation. Once again, we do not insist that such a description of the reactions represents the complete picture. Rather do we pose the following conditional problem. What can be said about $H^{(0)}$ and $H^{(1)}$ if it were true that deviations from the DA modify only slightly the description of nonleptonic K decay?

Note that a transition $K^\pm\to$ system with $S=0$ is necessarily of the type $\Delta I=1$. For if it were a $\Delta I=0$ transition, then the final state would have $Q=0$. In particular

$$K^+\to 2\pi^++\pi^- \quad\text{or}\quad \pi^++2\pi^0 \quad \text{have} \quad \Delta I=1, \quad (5.8)$$

$$K^-\to\pi^-,\quad \bar{K}^0\to\pi^0 \qquad \text{have} \quad \Delta I=1. \quad (5.9)$$

On the other hand,

$$K_1^0 \rightarrow \pi^+ + \pi^- \quad \text{or} \quad 2\pi^0 \quad \text{have} \quad \Delta I = 0. \tag{5.10}$$

Thus the θ^0 and $\tau^\pm$ modes proceed via $H^{(0)}$ and $H^{(1)}$, respectively. ($K_2^{(0)}$ goes via $H^{(1)}$.) Representative graphs are shown in Fig. 1. For each of the two reactions one graph goes via nucleons, the other via cascades.

Let us now continue to assume that $H^{(0)}$ and $H^{(1)}$ conserve parity. Then, in general, the ΔI properties of these couplings are no sufficient guarantee for actually allowing the reactions (5.8) and (5.10) to occur. Using the O operators of Eq. (5.4), we find

$$\begin{aligned} &K_{\pi 3}^+ \text{ and } K \rightarrow \pi \text{ allowed via} \\ &\quad \text{nucleons if} \quad O_{13} = 1(i\gamma_5) \quad \text{for} \quad H^{(1)} = 1(i\gamma_5), \\ &\quad \text{cascades if} \quad O_{24} = 1(i\gamma_5) \quad \text{for} \quad H^{(1)} = 1(i\gamma_5). \end{aligned} \tag{5.11}$$

These conditions are evidently general and do not depend on the particular graphs shown in Fig. 1. Likewise for $K_{\pi 2}^0$ we have

$$\begin{aligned} &K_{\pi 2}^0 \text{ allowed via} \\ &\quad \text{nucleons if} \quad O_{12} = 1(i\gamma_5) \quad \text{for} \quad H^{(0)} = i\gamma_5(1), \\ &\quad \text{cascades if} \quad O_{34} = 1(i\gamma_5) \quad \text{for} \quad H^{(0)} = i\gamma_5(1). \end{aligned} \tag{5.12}$$

To see when $K_{\pi 2}^0$ and $K_{\pi 3}^+$ are both allowed in the DA, we must distinguish two cases.

(a) $P(\Xi)$ even. The conditions are:

$$\text{if } O_{12} = O_{13} = 1, \quad \text{then } H^{(0)} = i\gamma_5,\ H^{(1)} = 1; \tag{5.13}$$

$$\text{if } O_{12} = O_{13} = i\gamma_5, \text{ then } H^{(0)} = 1, \quad H^{(1)} = i\gamma_5; \tag{5.14}$$

$$\text{if } O_{12} \neq O_{13}, \qquad \text{then } H^{(0)} = H^{(1)}. \tag{5.15}$$

Equation (5.15) implies that if $p(K)$ were odd it would be impossible to assume that all four decays:

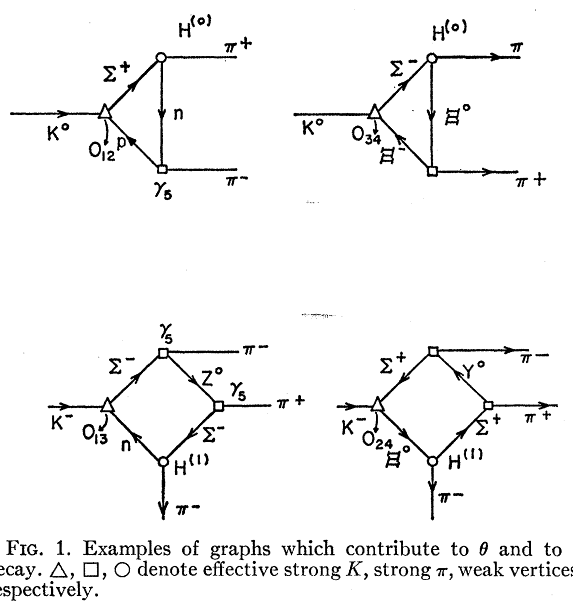

FIG. 1. Examples of graphs which contribute to θ and to τ decay. △, □, ○ denote effective strong K, strong π, weak vertices, respectively.

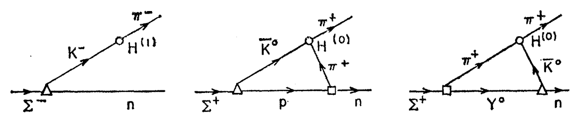

FIG. 2. K contributions which conserve parity if $H^{(0)}$, $H^{(1)}$ violate parity (Wolfenstein).

$K_{\pi 2}^0$, $K_{\pi 3}^+$, Σ_+^+, and Σ_-^- are well approximated by the DA. While Eqs. (5.13), (5.14) show that for even $p(K)$ the assumption of P-conserving but clashing $H^{(0)}$, $H^{(1)}$ is adequate to describe all these decays in the DA.

(b) $P(\Xi)$ odd. The transitions (5.8–10) are always allowed. Thus even $P(\Xi)$ excludes odd $p(K)$; odd $P(\Xi)$ does not exclude odd $p(K)$.

We conclude the following. It is *possible* to endow $H^{(0)}$, $H^{(1)}$ with all the properties stated in the $\Delta I = 0$, 1 rule of Sec. 2(b) and have both $K_{\pi 2}^0$ and $K_{\pi 3}^+$ allowed in the DA. For this last purpose it is of course totally *unnecessary* to assume that these two separate weak interactions do conserve parity. The parity properties of $H^{(0)}$, $H^{(1)}$ become only manifest in the general discussion of Σ decays given in the foregoing.

But even in this last respect we must make a proviso. It is conceivable that also for Σ decays we could allow $H^{(0)}$ and $H^{(1)}$ to be parity nonconserving provided that certain dynamical accidents happen.

It cannot be anybody's purpose to give a complete theory of accidents. Let us nevertheless consider some specific examples in the next section.

6. QUESTIONS OF DOMINANT VIRTUAL TRANSITIONS

(a) Dominance of K_π and $K_{\pi 2}$

Consider the contribution to Σ_-^- and Σ_+^+ from the graphs drawn in Fig. 2. The Σ^- graph was first considered by FMS. The possibility to have P conservation if K_π and $K_{\pi 2}$ dominate was first noted by Wolfenstein, who also stated the relevance of the DA with regard to the $\Sigma^+ \rightarrow Y^0$ graph. We add the following comments:

(1) From the present point of view we deal here with special $H^{(0)}$ and $H^{(1)}$ transitions as indicated in the figure, see Eqs. (5.9–10).

(2) These graphs provide examples of P-conserving contributions even if $H^{(0)}$ and $H^{(1)}$ do not separately conserve parity, see the discussion of Sec. 5(d). Thus if such graphs would entirely dominate Σ_+^+ and Σ_-^-, the argument given in I concerning the incompatibility of the schizon scheme with the parity structure of the nonleptonic weak interactions would *not* apply.

(3) We consider next examples of graphs which enter in the same order as those of Fig. 2 but which give P violation if $H^{(0)}$, $H^{(1)}$ are P-nonconserving. Consider the weak K_π and $K_{\pi 2}$ transitions of Eqs. (5.9–10) as brought about by

$$\begin{aligned} &K^- \rightarrow n + \bar{p}(\text{weak}); \quad n + \bar{p} \rightarrow \pi^-(\text{strong}), \\ &\bar{K}^0 \rightarrow n + \bar{n}(\text{weak}); \quad n + \bar{n} \rightarrow \pi^+ + \pi^-(\text{strong}). \end{aligned} \tag{6.1}$$

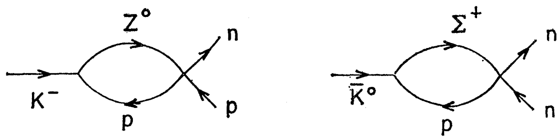

FIG. 3. Contributions to $K^- \to n\bar{p}(\Delta I=1)$ and to $\bar{K}^0 \to n\bar{n}(\Delta I=0)$.

This is only one of many ways in which K_π and $K_{\pi 2}$ vertices can be generated but Eq. (6.1) will suffice to illustrate the point. The weak links of Eq. (6.1) are drawn in Fig. 3. These links do not only suffice to generate the graphs of Fig. 2 but also those of Fig. 4. It is obvious that these last graphs give in general (that is, barring still further accidents) P-nonconserving contributions to the decays if $H^{(0)}$, $H^{(1)}$ do not conserve parity.

(4) However, it follows from the argument given in Sec. 5(d) that it means no restriction to the Wolfenstein argument to let $H^{(0)}$ and $H^{(1)}$ be P-conserving. But in turn, if this parity condition is satisfied, we are also guaranteed that the graphs of Fig. 4 conserve parity.

(5) Consider the graphs of Figs. 2 and 4 from the point of view of the $\Delta I=0$, 1 rule. As explained in Sec. 2(b), we need *in general* an additional argument concerning $H^{(1)}$. But for the graphs under consideration *no argument beyond the DA is necessary.*

(6) We conclude the following. To achieve P-conservation in Σ_+^+ and Σ_-^-, the Wolfenstein graphs are acceptable whether or not $H^{(0)}$ and $H^{(1)}$ separately conserve parity. If they do conserve parity, we can extend without further ado the Wolfenstein argument so as to include the graphs of Fig. 4. The DA is not only necessary but also sufficient for this subset of transitions.

The connection between the Wolfenstein model and the $\Delta I=0$, 1 rule has thus been established by focusing the attention on the properties of the weak vertices. A similar argument will make clear the connection between this rule and the results of FMS.

(b) Single Baryon Dominance

The considerations of this section were entirely stimulated by the results of FMS. Following these authors, we consider the particular chains of weak and strong interactions drawn in Fig. 5. The strengths of the strong vertices are expressed in terms of G_1 and G, see Eq. (2.1). In the spirit of FMS, one may consider G_1 and G as renormalized constants, however. As is evident from the real and virtual baryon states indicated in Fig. 5, we study the problem in the DA only. This is no restriction as compared to FMS.

The vertices have in general a momentum structure. With FMS we shall ignore this for the strong vertex but not for the weak vertices X_i $(i=1,2,3)$. The reason for insisting on this dependence is that most models for nonleptonic decays depend very sensitively on the hyperon nucleon mass difference. For example, an effective interaction $\bar{\Lambda}\gamma_\alpha(1+\gamma_5)p\partial_\alpha\pi$ gives an $\alpha_\Lambda \sim 0.9$ but would give $\alpha_\Lambda=0$ if we neglect $m_\Lambda - m_p$. Such mass difference effects are not fully exhibited for constant X_i, as we shall see.

The $\Delta T=\frac{1}{2}$ rule implies

$$X_3=X_2-X_1, \tag{6.2}$$

and relates the transition matrices $\mathfrak{M}$ as follows.

$$\mathfrak{M}^- = \mathfrak{M}^+ - \mathfrak{M}^0\sqrt{2}. \tag{6.3}$$

Put

$$\begin{aligned} X_i &= a_i + b_i\gamma_5, \\ a_i &= A_i + i\gamma q C_i, \\ b_i &= B_i + i\gamma q D_i, \end{aligned} \tag{6.4}$$

where A–D are functions of the 4-momentum transfer $-q^2$. Apart from common factors, we have then

$$\mathfrak{M}^+ = \gamma_5 \frac{\xi^+}{M-1} - \frac{\eta^+}{M+1}, \tag{6.5}$$

$$\sqrt{2}\mathfrak{M}^0 = \gamma_5 \frac{\xi^0}{M-1} - \frac{\eta^0}{M+1}, \tag{6.6}$$

$$\mathfrak{M}^- = \gamma_5 \frac{\xi^+ - \xi^0}{M-1} - \frac{\eta^+ - \eta^0}{M+1}, \tag{6.7}$$

$$\xi^+ = A_1(M^2) - \epsilon A_2(1) - MC_1(M^2) + \epsilon C_2(1), \tag{6.8}$$

$$\eta^+ = B_1(M^2) + \epsilon B_2(1) - MD_1(M^2) + \epsilon D_1(1), \tag{6.9}$$

$$\xi^0 = A_1(M^2) - \epsilon A_1(1) - MC_1(M^2) + \epsilon C_1(1), \tag{6.10}$$

$$\eta^0 = B_1(M^2) + \epsilon B_1(1) - MD_1(M^2) + \epsilon D_1(1), \tag{6.11}$$

$$\epsilon = G/G_1. \tag{6.12}$$

The nucleon mass has been put $=1$, M is the Σ mass.

So far we have only used the DA for the strong vertices. We ask what happens if we let the weak ones share this symmetry. *This establishes a relation between* X_1 *and* X_2, *but this relation depends on the* ΔI *structure of the weak interactions.*

Consider first a general $H^{(0)}$ interaction. This shares with the strong interactions the doublet charge symmetry property. Thus $(\Sigma^+|p)=(Y^0|n)$, that is

$$X_1 = X_2, \quad (\Delta I=0). \tag{6.13}$$

For $H^{(1)}$ the situation is more complex as can be seen with reference to the three kinds of $\Delta I=1$ couplings of Eqs. (3.21–23). One finds

$$\begin{aligned} &\text{for}\quad js,\ \rho s'\ (\Delta I=1){:}\quad X_1=-X_2, \\ &\text{for}\qquad js'\ (\Delta I=1){:}\quad \text{no simple relation.} \end{aligned} \tag{6.14}$$

A further investigation of js' did not reveal any chance reason why $X_1=\pm X_2$ might be valid for this case. We exclude js' from the following, not because there is any

argument against this interaction but because we have nothing to say about it.

Equations (6.8–11) now give

for $\Delta I=0$:

$$\xi^+=\xi^0=A_1(M^2)-\epsilon A_1(1)-MC_1(M^2)+\epsilon C_1(1), \quad (6.15)$$

$$\eta^+=\eta^0=B_1(M^2)+\epsilon B_1(1)-MD_1(M^2)+\epsilon D_1(1); \quad (6.16)$$

for $\Delta I=1$ (not js'):

$$\xi^+=A_1(M^2)+\epsilon A_1(1)-MC_1(M^2)-\epsilon C_1(1), \quad (6.17)$$

$$\eta^+=B_1(M^2)-\epsilon B_1(1)-MD_1(M^2)-\epsilon D_1(1), \quad (6.18)$$

while ξ^0, η^0 are still given by Eqs. (6.10–11).

(c) Further Comments. The FMS Model

(1) Consider the special case

$$\begin{aligned} &\text{All} \quad C, D=0,\\ &\text{All} \quad A, B \text{ independent of } q^2,\\ &\qquad\ \epsilon=\pm 1. \end{aligned} \quad (6.19)$$

This last restriction is just the one to $G^\pm$ symmetry. It follows now from Eqs. (6.5–12) that $\mathfrak{M}^0$ is pure $s(p)$ wave for $\epsilon=+1(-1)$, in either case in contradiction with the large asymmetry observed in $\Sigma_0{}^+$. Equation (6.19) is one of the assumptions of FMS. Hence these

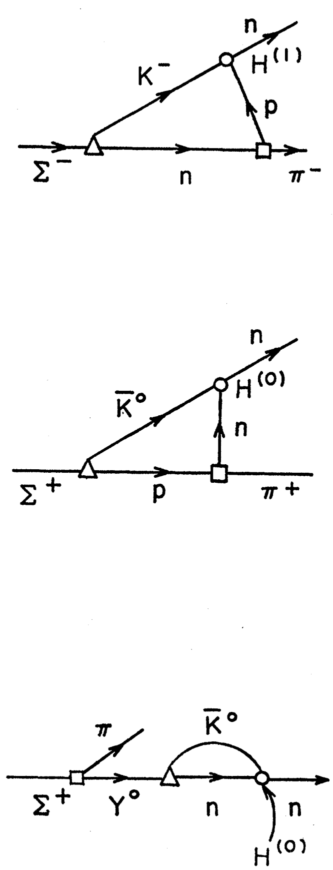

FIG. 4. K contributions which generally violate parity if $H^{(0)}$, $H^{(1)}$ violate parity.

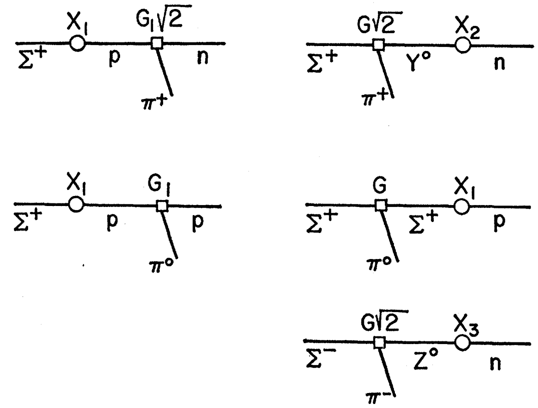

FIG. 5. FMS graphs for transitions via a single baryon.

authors need further decay mechanisms, for which they choose the K_π graphs discussed in Sec. 6(a).

(2) In the FMS treatment, the contributions to the Σ^- graph of Fig. 5 are zero. It can now easily be traced back how this comes about. FMS assume that the weak vertex has the form $a_Y+b_Y\gamma_5$ where a_Y, b_Y are constants and are taken to be the same for all graphs. In the present language, FMS assume that Eq. (6.13) is true, that is, at least for baryon contributions they make the assumption of a pure $\Delta I=0$ interaction. To this there is of course no objection. It is important to note, however, that it is therefore implicitly contained in the work of FMS that not all weak vertices can be iterates of a pure $(I=1, \Delta S=0)\times(I=\frac{1}{2},|\Delta S|=1)$ coupling, in accordance with the more general reasoning of previous sections.

(3) To obtain a parity-pure answer for $\mathfrak{M}^+$, FMS proceed by using the same $\Delta I=0$ vertex as was mentioned before. We know from the general discussion that, if $H^{(0)}$ conserves parity, then $\mathfrak{M}^+$ does the same. Here FMS follow a different course which is probably somewhat more restrictive. On the one hand, they do not assume that $H^{(0)}$ conserves parity; on the other, they assume Eq. (6.19) to be true. This indeed leads to a P-conserving $\mathfrak{M}^+$, namely, $s(p)$ wave for $\epsilon=+1(-1)$, see Eqs. (6.15–16) and (6.19).

(4) Thus the FMS model is equivalent to an effective $H^{(1)}$ of the type K_π and an effective $H^{(0)}$ of the type $\bar{N}_2N_1$.

(5) It is instructive to note that there are other solutions to Eqs. (6.5–12) which give the desired properties of $\mathfrak{M}^\pm$. These are obtained by dropping Eq. (6.19).

Example: Take an effective $H^{(0)}$ with $A_1=B_1=C_1=0$ and a nonvanishing constant D_1, denoted by $D_1{}^{(0)}$. Take further an effective $H^{(1)}$ with $B_1=C_1=D_1=0$ and a nonvanishing and constant A_1, denoted by $A_1{}^{(1)}$. Put $\epsilon=-1$. Then we obtain from Eqs. (6.15–18)

$$\mathfrak{M}^+=D_1{}^{(0)}, \quad \mathfrak{M}^-=-2A_1{}^{(1)}\gamma_5(M-1)^{-1}. \quad (6.20)$$

Hence, without using K_π, we have obtained a non-

vanishing P-conserving $\mathfrak{M}^{\pm}$ with the required opposite parities. Note moreover that the rates for $\Sigma_+{}^+$ and $\Sigma_-{}^-$ are of the desired same order of magnitude if $D_1{}^{(0)} \simeq A_1{}^{(1)}$. In fact these rates equal each other for $A_1{}^{(1)} = 1{,}3D_1{}^{(0)}$.

Equation (6.20) is given as an example, not as a proposal. It shows that even if single pole dominance is correct, it is not at all obvious what one should conclude from that.

Moreover, Eq. (6.20) provides a second example of a situation where the DA is enough for $H^{(1)}$. Indeed, under the conditions stated to obtain Eq. (6.20) the contributions of $H_1{}^{(1)}$ to $\Sigma_+{}^+$ vanish, see Eqs. (6.17–18). Hence there is once more no need for the additional argument referred to in Sec. 2(b).

7. CONCLUDING REMARKS

The previous sections have dealt with the structure of weak nonleptonic interactions in a purely descriptive way. The question arises, if the $\Delta I=0$, 1 rule is correct, could one see why these interactions consist of two parts, $H^{(0)}$ and $H^{(1)}$?

We would like to present a speculation on this question. The strong interactions (2.1) and (5.2) which make up the DA can be written as

$$\mathbf{j}\boldsymbol{\pi}+(\bar{s}K+s\bar{K}), \tag{7.1}$$

where j denotes the isotopic structure of the π-field source. Assume that the $\pi(K)$ fields have small $K(\pi)$ components,[20]

$$\boldsymbol{\pi}\rightarrow\boldsymbol{\pi}+\Omega_1\mathbf{K},\quad \mathbf{K}=(K^+,K^-,K_1{}^0),$$
$$\begin{pmatrix}K^+\\ K^0\end{pmatrix}\rightarrow\begin{pmatrix}K^+\\ K^0\end{pmatrix}+\Omega_2\begin{pmatrix}\pi^+\\ 2^{-\frac{1}{2}}\quad\pi^0\end{pmatrix}, \tag{7.2}$$

where Ω_1, Ω_2 may be operators proportional to a weak coupling constant, but are independent of T, I, or K. Then

$$\mathbf{j}\boldsymbol{\pi} \text{ generates } \mathbf{j}\Omega_1\mathbf{K},$$
$$\bar{s}K+s\bar{K} \text{ generates} \tag{7.3}$$
$$s^+\Omega_2\pi^+ + s^-\Omega_2\pi^- + (1/\sqrt{2})(s^0+\bar{s}^0)\Omega_2\pi^0.$$

These couplings are both of the type js, hence, $H^{(1)}$, see Eq. (3.21). We might say that the Eqs. (7.2) generate a "schizon scheme[11] without schizons."

We can generate $H^{(0)}$ by

$$N_1\rightarrow N_1+\Omega_{12}N_2,\quad N_3\rightarrow N_3+\Omega_{34}N_4,$$
$$N_2\rightarrow N_2+\Omega_{21}N_1,\quad N_4\rightarrow N_4+\Omega_{43}N_3. \tag{7.4}$$

In this way

$$\mathbf{j}\boldsymbol{\pi} \text{ generates } \mathbf{t}\boldsymbol{\pi}, \tag{7.5}$$

which is a jt coupling, hence, $H^{(0)}$, see Eq. (3.18).

[20] Similar to the way small nonstatic electric (magnetic) phenomena accompany almost pure static magnetic (electric) phenomena.

One can proceed likewise with the K-couplings and finds

$$\bar{s}K+s\bar{K} \text{ generates } \bar{N}_2(N_2K^0+N_3K^+)+\cdots, \tag{7.6}$$

where the term written out is a representative of several terms, all of the type $H^{(0)}$.

Thus far we have only considered $\Delta T=\frac{1}{2}$. The small electromagnetic deviations from this rule give $\Delta T=\frac{3}{2}$ and larger. The doublet picture provides another way of inhibiting $\Delta T=\frac{1}{2}$. This is perhaps of interest for reactions such as $K^+\rightarrow\pi^++\pi^0$ for which $\Delta T=\frac{3}{2}$ (or $\frac{5}{2}$). Consider, namely, a decay interaction with $\Delta K=\frac{3}{2}$, $\Delta I=1$, $\Delta T=\frac{3}{2}$ which is possible in accordance with Eq. (2.5). This still does not allow $K_{\pi2}{}^+$ decay in the DA, as there $\Delta K=\frac{1}{2}$, $\Delta I=2$. If the DA gets broken, however, only the ΔT-rules survive, however, so that $K_{\pi2}{}^+$ can now take place. In this sense we may perhaps be justified to call this nonelectromagnetic $\Delta T=\frac{3}{2}$ effect a "small" effect. A dimensionless parameter which characterizes the DA is $\delta=(M_\Sigma-M_\Lambda)/M_\Lambda$. In the DA, $\delta=0$, its actual value is $\delta\simeq0.067$. If we consider δ as a measure for the amplitude ratio $K_{\pi2}{}^+/K_{\pi2}{}^0$, then the ratio of rates would be $\sim\delta^2=0.005$, a suggestive order of magnitude.

The common orders of magnitude of all weak processes suggests a common generic mechanism of leptonic and nonleptonic decay couplings. If the latter are generated by the strong interactions, one may ask if the same should not be true of the leptonic decays. I do not think that such a question can be answered in a theory which does not account for the law of baryon conservation.

Finally, we summarize what is general, what is special about the conclusions obtained in I, where the particular choice $H^{(1)}=js$, $H^{(0)}=jt$, was studied.

In I we obtained parity clash for G^+ symmetry. In this paper we have shown that this is also possible for G^- invariance, while at the same time Eq. (2.15) holds true.[16]

In I we explored the consequence of two further assumptions, namely, (a) The S-nonconserving baryon currents are either all pure V (i.e., γ_λ) or pure A (i.e., $\gamma_\lambda\gamma_5$). (b) The same S-nonconserving currents intervene in both leptonic and nonleptonic decays with $|\Delta S|=1$.

It was shown in I that (a) implies odd $P(\Xi)$. This conclusion is not specific for the (js,jt) coupling scheme. It was also shown in I that (a) and (b) imply the occurrence of leptonic $\Delta T=\frac{3}{2}$ transition. This last conclusion is more specifically true only if $H^{(0)}$ contains jt.

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