Joint description of weak radiative and nonleptonic hyperon decays in broken SU(3)

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(Received 14 December 2005; published 28 April 2006)

We give a joint description of weak radiative (WR) and nonleptonic (NL) hyperon decays (HD) in broken SU(3). The two groups of decays are linked via $SU(2)_W$ spin symmetry and vector-meson dominance (VMD). We use experimental information on the parity-conserving NLHD amplitudes to fix the corresponding WRHD amplitudes. With the latter known, the data on the WRHD branching ratios and asymmetries permit us to determine the parity-violating WRHD amplitudes in terms of two parameters corresponding to the two-quark and single-quark transitions. We obtain a good description of the data, and, in particular, a large $\Sigma^+ \to p\gamma$ asymmetry. Then, using the $SU(2)_W$ + VMD route we determine the non-soft-meson correction term in the parity-violating NLHD amplitudes. The latter is shown to subtract a substantial amount from the current-algebra commutator thus leading towards the resolution of the *S*:*P* discrepancy in NLHD.

DOI: 10.1103/PhysRevD.73.076005 PACS numbers: 11.40.−q, 11.30.Hv, 13.30.−a, 14.20.Jn

I. INTRODUCTION

For a long time weak hyperon decays have been presenting us with a couple of puzzles (see [1,2]). These have been in particular: the problem of the *S*:*P* ratio in the nonleptonic hyperon decays (NLHD) and the issue of a large negative asymmetry in the $\Sigma^+ \to p\gamma$ weak radiative hyperon decay (WRHD).

Both problems emerged several decades ago. The first (*S*:*P*) problem is the inconsistency between the size of matrix elements of the parity-conserving (p.c.) Hamiltonian in between ground-state baryon states, as estimated from the p.c. (*P*-wave) NLHD amplitudes, and the size of the same matrix elements when estimated from the parity-violating (p.v., *S*-wave) NLHD amplitudes via PCAC in the standard soft-pion limit. The two estimates differ by a factor of around 2 or a little bit larger, depending on the details of the models used.

The second problem emerged when first experiments hinted that the $\Sigma^+ \to p\gamma$ asymmetry is large [3]. The large size of this asymmetry was unexpected since a theorem proved by Hara [4] stated that in the SU(3) limit the relevant parity-violating amplitude should vanish. For broken SU(3), having in mind the size of hadron-level SU(3)-breaking effects elsewhere, one would expect this asymmetry to be of the order ± 0.2, and not of the order of -1 (the present experimental number is -0.76 ± 0.08). The situation was further confounded by a number of theoretical calculations which violated Hara's theorem (even) in the SU(3) limit (see Ref. [2]).

Some time ago it was pointed out [2] that the status of Hara's theorem can be clarified through the measurement of the $\Xi^0 \to \Lambda\gamma$ decay asymmetry. By yielding a large and negative value of -0.78 ± 0.19 for this asymmetry, the recent NA48 experiment [5,6] has decided very clearly in favor of the theorem.

The experimental result of the NA48 collaboration permits us to conclude that theoretical results which violate Hara's theorem in the SU(3) limit constitute artefacts of the relevant approaches. This concerns both the quark-level calculations of Kamal and Riazuddin [7], and the hadron-level calculation of the present author [8]. However, the origins of the artefacts are different in these two approaches.

In the quark model of Ref. [7] Hara's theorem is violated because in the calculations the intermediate photon-emitting quark enters its mass shell. Thus, this quark is treated as a free ordinary particle. This leads to a nonvanishing *nonlocal* contribution and violates Hara's implicit assumption that the relevant transition be described in a language of a *local* hadron-level theory [9].

The result of Ref. [8] follows from Ref. [10] when the description of weak p.v. couplings of vector mesons *V* to baryons *B* provided by [10] is supplied with the idea of vector-meson dominance (VMD). In Ref. [10] the $B' \to VB$ weak p.v. amplitudes are obtained by the application of $SU(6)_W$ to the *full* $B' \to PB$ weak p.v. amplitudes (*P*-pseudoscalar mesons), with the latter determined from experimental data on nonleptonic hyperon decays. If the applicability of VMD is accepted, the *way* in which the $SU(6)_W$-related $B' \to VB$ counterparts of the $B' \to PB$ amplitudes are determined in [10] must be incorrect.

In fact, Ref. [10] considers contributions to the p.v. NLHD amplitudes coming from the current-algebra (CA) commutator term only. It is the application of $SU(2)_W$ spin symmetry to this contribution which ultimately leads to terms violating Hara's theorem. In general, however, the p.v. NLHD amplitudes contain two terms: the CA commutator term and the correction term which should vanish in the soft-pion limit. If the latter term is not small for physical pion momentum, then its $SU(2)_W$-related counterpart in WRHD is not small either and could be important in

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the description of WRHD. The observed sign and size of the $\Xi^0 \to \Lambda\gamma$ asymmetry permits us to make definite conclusions concerning not only Hara's theorem, but also—via the $SU(2)_W$ + VMD route—the size and sign of the non-soft-pion term in nonleptonic hyperon decays. As observed in the approach in [11], in which in the parity-violating sector the SU(3) symmetry was exact, WRHD permits us to establish that the correction term subtracts a substantial part from the CA commutator contribution, thus working towards the resolution of the old S:P problem in NLHD.

In the present paper we introduce explicit SU(3) breaking into the parity-violating sector of the scheme of [11], and show that despite the fact that the p.v. $\Sigma^+ \to p\gamma$ amplitude vanishes in exact SU(3), in broken SU(3) this amplitude is comparable in size to other SU(3)-unsuppressed p.v. WRHD amplitudes. As a result we obtain a large $\Sigma^+ \to p\gamma$ asymmetry. Our description of the branching ratios and asymmetries in weak radiative hyperon decays is in good agreement with the data. Although it deviates from the experimental data more than the corresponding description of NLHD, it reproduces both the large size of all observed asymmetries, and provides a fair description of the branching ratios. In addition, it predicts a substantial positive asymmetry in the $\Xi^- \to \Sigma^-\gamma$ decay. We also show that when SU(3) is broken in the parity-violating sector, then the non-soft-pion contribution to NLHD [obtained from WRHD via the $SU(2)_W$ + VMD route] is of proper sign and order of magnitude to resolve the S:P problem.

II. GENERAL

If we write the effective Lagrangian for nonleptonic hyperon decay $B_i \to B_f\pi$ as

$$\bar{u}_f(A + B\gamma_5)u_i\Phi_\pi, \tag{1}$$

where A (B) denotes the parity-violating (parity-conserving) amplitude, the decay rate is given by

$$\Gamma = \frac{1}{4\pi}\frac{k_\pi}{m_i}(E_f + m_f)[|A|^2 + |\bar{B}|^2], \tag{2}$$

where E_f, $m_f(m_i)$ are energy and mass for the final (initial) baryon, k_π is pion momentum, and

$$\bar{B} = \sqrt{\frac{E_f - m_f}{E_f + m_f}}B. \tag{3}$$

The asymmetry is

$$\alpha = \frac{2\,\mathrm{Re}(A^*\bar{B})}{|A|^2 + |\bar{B}|^2}. \tag{4}$$

Similarly, if the effective Lagrangian for weak radiative hyperon decay $B_i \to B_f\gamma$ is written as

$$\bar{u}_f i\sigma_{\mu\nu}(p_f - p_i)^\nu(C + D\gamma_5)u_iA^\mu, \tag{5}$$

with C (D) being the parity-conserving (violating) amplitude, then the decay rate is given by

$$\Gamma = \frac{1}{\pi}\left(\frac{m_i^2 - m_f^2}{2m_i}\right)^3[|C|^2 + |D|^2], \tag{6}$$

and the asymmetry is

$$\alpha = \frac{2\,\mathrm{Re}(C^*D)}{|C|^2 + |D|^2}. \tag{7}$$

Theoretical models of hyperon decays may relate some or all of the four amplitudes A, B, C, D. We start with the parity-conserving sector and the relation between amplitudes B and C.

III. PARITY-CONSERVING AMPLITUDES

The parity-conserving NLHD amplitudes are known to be well described by the pole model with the ground-state $(56, 1/2^+)$ baryons in the intermediate state. By $SU(2)_W$ spin symmetry one expects that the same model (supplied with the VMD assumption) is adequate for the description of the p.c. WRHD amplitudes. In this section we present our version of this approach.

A. Nonleptonic decays

In the ground-state baryon pole model the explicit dependence of the p.c. NLHD amplitudes $B(B_i \to B_f\pi)$ on (1) F/D describing the SU(3) structure of $\pi BB'$ couplings, and (2) f_P/d_P characterizing the SU(3) structure of the matrix elements of the parity-conserving part $H_W^{\text{p.c.}}$ of the weak Hamiltonian is (see e.g. [8,12]):

$$\begin{aligned}
B(\Sigma^+ \to p\pi^0) &\equiv B(\Sigma^+_0) = \frac{1}{\sqrt{2}}\left(\frac{f_P}{d_P} - 1\right)\left(1 - \frac{F}{D}\right)N,\\
B(\Sigma^+ \to n\pi^+) &\equiv B(\Sigma^+_+) = -\frac{4}{3}N\\
B(\Sigma^- \to n\pi^-) &\equiv B(\Sigma^-_-) = \left[\left(\frac{f_P}{d_P} - 1\right)\frac{F}{D} - \frac{1}{3}\left(3\frac{f_P}{d_P} + 1\right)\right]N\\
B(\Lambda \to n\pi^-) &\equiv B(\Lambda^0_-) = -\sqrt{2}B(\Lambda^0_0) = \frac{1}{\sqrt{6}}\left[\frac{f_P}{d_P} + 3 + \left(3\frac{f_P}{d_P} + 1\right)\frac{F}{D}\right]N\\
B(\Xi^- \to \Lambda\pi^-) &\equiv B(\Xi^-_-) = -\sqrt{2}B(\Xi^0_0) = -\frac{1}{\sqrt{6}}\left[3 - \frac{f_P}{d_P} + \left(3\frac{f_P}{d_P} - 1\right)\frac{F}{D}\right]N,
\end{aligned} \tag{8}$$

with the standard notation given in the second column (see also Table IV below).

In writing Eqs. (8) we assumed as in [8,12] that all pole denominators are equal i.e. that

$$\frac{1}{m_\Sigma - m_N} = \frac{1}{m_\Lambda - m_N} = \frac{1}{m_\Xi - m_\Sigma} = \frac{1}{m_\Xi - m_\Lambda} \equiv \frac{1}{\Delta m_s}, \tag{9}$$

with $\Delta m_s = 190$ MeV, and absorbed them into an overall normalization factor N

$$N = \frac{2m_8}{F_\pi}\frac{Dd_P}{\Delta m_s}, \tag{10}$$

where $F_\pi = 94$ MeV, m_8 is some average value of baryon ground-state octet masses, taken as $m_8 = (m_N + m_\Xi)/2 \approx 1130$ MeV, and d_P together with f_P describe the SU(3) structure of the parity-conserving weak Hamiltonian. The form of SU(3) breaking specified in Eq. (9) was used in previous papers on the subject, and specifically in Ref. [11], and does not constitute the novelty. The difference with respect to Ref. [11] lies in the explicit consideration of SU(3) breaking in the parity-violating sector (Sec. IV).

Our assumption of equal pole denominators (i.e. no $\Sigma - \Lambda$ splitting) corresponds to the simplest form of $SU(3)_F$ symmetry breaking one can consider, the whole effect of SU(3) breaking being due to a heavier mass of the strange quark. Other elements of the description (such as the strong $B'BP$ couplings, or the matrix elements of the p.c. weak Hamiltonian) are $SU(3)_F$-symmetric.

As the $\Sigma - \Lambda$ splitting results from spin-spin effects, it follows that taking this splitting into account would require the consideration of the influence of spin-spin SU(3) breaking effects in strong meson-baryon couplings (and possibly in weak transition amplitudes). These are not understood well, however. Consequently, $SU(3)_F$ symmetry was assumed in this paper for the meson-baryon couplings. For reasons of consistency, therefore, we cannot take the spin-interaction-induced $\Sigma - \Lambda$ splitting into consideration. Equations (8) may be also viewed as just a simple parametrization of the $B_i \to B_f\pi$ amplitudes. Transition from these amplitudes to the amplitudes with pion replaced by a U-spin singlet vector meson U^0 (a linear combination of ρ, ω, and ϕ), as needed in the next subsection, is achieved via $SU(6)_W$ symmetry. If the $B(B_i \to B_f\pi)$ amplitudes are well described by Eqs. (8) (and they indeed are, see Table I), then the $B(B_i \to B_fU^0)$ amplitudes should also be well described.

Our normalization of f_P, d_P can be read off from

$$\langle p|H_w^{\mathrm{p.c.}}|\Sigma^+\rangle = \sqrt{2}(d_P - f_P), \tag{11}$$

(compare also Table III).

TABLE I. P-wave NLHD amplitudes $B(B_i \to B_f\pi)$ (in units of 10^{-7}) using Eq. (8) (with $f_P/d_P = -1.9$, $F/D = 0.55$, $N = -31 \times 10^{-7}$) and the data.

Decay	Equation (8)	Data
Σ^+_0	28.6	26.6 ± 1.3
Σ^+_+	41.3	42.4 ± 0.35
Σ^-_-	0.9	-1.44 ± 0.17
Λ^0_-	18.8	22.1 ± 0.5
Ξ^-_-	15.4	16.6 ± 0.8

For $F/D = 0.55$ $(F = 0.44, D = 0.81)$, $N = -31$ (in units of 10^{-7}), and $f_P/d_P = -1.90$ one obtains a very good description of the data (see Table I, also [12]). Note that our scheme satisfies the $\Delta I = 1/2$ rule. Consequently, one cannot expect here a better agreement in view of the violation of this rule: e.g. the $\Delta I = 1/2$ relation $\sqrt{2}B(\Sigma^+_0) = B(\Sigma^+_+) - B(\Sigma^-_-)$ experimentally reads: $37.6 \pm 1.8 = 43.8 \pm 0.4$, indicating that the $\Delta I = 3/2$ effects are of the order of 5–10%.

From Eq. (10) one finds

$$d_P = \frac{F_\pi}{D}\frac{\Delta m_s}{2m_8}N \approx -3.0 \times 10^{-5}\ \mathrm{MeV}$$
$$f_P \approx 5.8 \times 10^{-5}\ \mathrm{MeV}. \tag{12}$$

Our f_P and d_P parameters are related to the ones used in [1] by

$$f_P = -2\sqrt{3}F_\pi f_P([1]) \tag{13}$$

$$d_P = -2\sqrt{3}F_\pi d_P([1]) \tag{14}$$

The values of $f_P([1]) = -1.44 \times 10^{-7}$ and $d_P([1]) = 0.8 \times 10^{-7}$ given in [1] correspond to our

$$d_P \approx -2.6 \times 10^{-5}\ \mathrm{MeV} \qquad f_P \approx 4.7 \times 10^{-5}\ \mathrm{MeV}. \tag{15}$$

The difference between the latter numbers and the estimates of Eq. (12) indicates how large the uncertainty in the extracted values of f_P and d_P might be.[1]

B. Radiative decays

For the WRHD the parity-conserving amplitudes $C(B_i \to B_f\gamma)$, obtained in the ground-state baryon pole model from the p.c. NLHD amplitudes via the $SU(2)_W$ + VMD route, are given by

$$C(B_i \to B_f\gamma) = \left(\frac{e}{g}\right)\frac{1}{(m_i + m_f)\sqrt{2}}B(B_i \to B_fU^0). \tag{16}$$

In the above equation $e/g = 0.0606$ is the VMD factor

[1] Uncertainties of this order might result e.g. from the treatment of kaon poles which were neglected by us but taken into account in [1].

($e^2/4\pi = 1/137$, $g = 5.0$), and $B(B_i \to B_f U^0)$ describe amplitudes for the emission of a linear superposition U^0 of virtual vector mesons ρ^0, ω, ϕ, corresponding to a photon and obtained by the $SU(6)_W$ symmetry from the NLHD amplitudes $B_i \to B_f \pi$ of Eq. (8). For the $B(B_i \to B_f U^0)$ amplitudes one gets

$$\begin{aligned}
B(\Sigma^+ \to pU^0) &= \sqrt{2}\left(\frac{f_P}{d_P} - 1\right)(\mu_{\Sigma^+} - \mu_p)\frac{N}{\mu_p D}\\
B(\Sigma^0 \to nU^0) &= \left[-\left(\frac{f_P}{d_P} - 1\right)(\mu_{\Sigma^0} - \mu_n) + \frac{1}{\sqrt{3}}\left(3\frac{f_P}{d_P} + 1\right)\mu_{\Sigma\Lambda}\right]\frac{N}{\mu_p D}\\
B(\Lambda \to nU^0) &= \left[\frac{1}{\sqrt{3}}\left(3\frac{f_P}{d_P} + 1\right)(\mu_\Lambda - \mu_n) - \left(\frac{f_P}{d_P} - 1\right)\mu_{\Sigma\Lambda}\right]\frac{N}{\mu_p D}\\
B(\Xi^0 \to \Lambda U^0) &= \left[-\frac{1}{\sqrt{3}}\left(3\frac{f_P}{d_P} - 1\right)(\mu_{\Xi^0} - \mu_\Lambda) - \left(\frac{f_P}{d_P} + 1\right)\mu_{\Sigma\Lambda}\right]\frac{N}{\mu_p D}\\
B(\Xi^0 \to \Sigma^0 U^0) &= \left[\left(\frac{f_P}{d_P} + 1\right)(\mu_{\Xi^0} - \mu_{\Sigma^0}) + \frac{1}{\sqrt{3}}\left(3\frac{f_P}{d_P} - 1\right)\mu_{\Sigma\Lambda}\right]\frac{N}{\mu_p D}\\
B(\Xi^- \to \Sigma^- U^0) &= -\sqrt{2}\left(\frac{f_P}{d_P} + 1\right)(\mu_{\Xi^-} - \mu_{\Sigma^-})\frac{N}{\mu_p D},
\end{aligned} \tag{17}$$

where we used equal mass splittings as suggested both by the success of Eqs. (8) when describing the data, and by the analysis of the Lee-Sugawara relations performed in [13]. The appearance of magnetic moments in Eqs. (17) will be explained shortly.

The particular form of the right-hand side of Eqs. (17) was obtained in [13], where the standard expressions for the $B(B_i \to B_f U^0)$ amplitudes [depending on the F and D couplings and similar to Eqs. (8)] were rewritten in terms of the corresponding magnetic moments to which they would be proportional in the SU(6) symmetry limit. The reason for using this representation is that ultimately we want to describe photon couplings which originate from the $\bar{u}_1 \sigma_{\mu\nu} q^\mu u_2 A^\nu$ terms and thus are expressed in terms of the anomalous parts of baryon magnetic moments. Now, the description of baryon magnetic moments provided by SU(6) (or when the strange quark is assumed to be heavier and has a smaller magnetic moment) is not good enough for our purposes. This is because there are substantial cancellations between various terms in Eqs. (17), and the results depend on the detailed values of baryon magnetic moments. In fact, it is known that substantial nonadditivities are observed in the experimental values of baryon magnetic moments: a thorough analysis performed in [14] revealed that the nonstrange quark contributions to baryon magnetic moments in protons and neutrons are significantly larger than in the baryons containing strange quarks. Since ultimately we want to describe photon couplings, the need to use magnetic moments on the left-hand side of Eq. (16) requires that its right-hand side takes them into account as well, as in Eqs. (17), thus modifying the vector-meson couplings accordingly.

In order to describe the photon couplings best, we chose to use in Eqs. (17) the experimental values themselves, i.e. (from [15]): $\mu_p = 2.793$, $\mu_n = -1.913$, $\mu_{\Sigma^+} = 2.46 \pm 0.01$, $\mu_{\Sigma\Lambda} = 1.61 \pm 0.08$, $\mu_{\Sigma^-} = -1.16 \pm 0.025$, $\mu_\Lambda = -0.613 \pm 0.004$, $\mu_{\Xi^0} = -1.25 \pm 0.014$, $\mu_{\Xi^-} = -0.651 \pm 0.003$, $\mu_{\Sigma^0} = (\mu_{\Sigma^+} + \mu_{\Sigma^-})/2$. From Eq. (16), using the fit of Table I, one can predict the p.c. WRHD amplitudes $C(B_i \to B_f \gamma)$. The relevant numbers for the related $B(B_i \to B_f U^0)$ amplitudes are given in Table II.

The numbers given in Table II result from cancellations between various terms in Eqs. (17). Such cancellations are strongest for the $B(\Lambda \to nU^0)$ amplitude. Specifically, for $B(\Lambda \to nU^0)$ the three terms seen in Eqs. (17) contribute as follows: the μ_Λ term gives -22.9, the μ_n term: $+71.4$, and the $\mu_{\Sigma\Lambda}$ term: -64.2. A change of $\mu_{\Sigma\Lambda}$ by 1 standard deviation from 1.61 to 1.69 leads to the absolute value of $B(\Lambda \to nU^0)$ larger than the value given in Table II by 20%. For the remaining amplitudes the uncertainty in $\mu_{\Sigma\Lambda}$ leads to errors of the order of a few percent only.

One should keep also in mind that our approach leads to an additional uncertainty in the size of the p.v. amplitudes for the $\Sigma^+ \to p\gamma$ and $\Xi^- \to \Sigma^-\gamma$ decays. This is because within our treatment of SU(3) breaking the differences of the anomalous parts of baryon magnetic moments reduce to the differences of baryon magnetic moments themselves.

TABLE II. P-wave amplitudes $B(B_i \to B_f U^0)$ (in units of 10^{-7}) as obtained from Eqs. (17) with $f_P/d_P = -1.9$, $d_P = -3.0 \times 10^{-5}$ MeV, $N = -31 \times 10^{-7}$.

Decay	Equation (17)
$\Sigma^+ \to pU^0$	-18.8
$\Sigma^0 \to nU^0$	-42.1
$\Lambda \to nU^0$	-15.7
$\Xi^0 \to \Lambda U^0$	$+13.9$
$\Xi^0 \to \Sigma^0 U^0$	$+62.1$
$\Xi^- \to \Sigma^- U^0$	-8.9

TABLE III. Parity-violating amplitudes in NLHD.

Amplitude	Parametrization		Values (in units of 10^{-7})	
	Commutator	Diagram decomposition	Experiment	Description ($b_S = -5$, $c_S = 12$)
$A(\Lambda^0_-)$	$\frac{1}{2F_\pi}\sqrt{\frac{2}{3}}(3f_P + d_P)$	$-\frac{1}{2\sqrt{6}}(b_S - c_S)$	3.25	3.47
$A(\Lambda^0_0)$	$-\frac{1}{2F_\pi}\frac{1}{\sqrt{3}}(3f_P + d_P)$	$\frac{1}{4\sqrt{3}}(b_S - c_S)$	−2.37	−2.46
$A(\Sigma^+_+)$	0	0	0.13	0
$A(\Sigma^+_0)$	$-\frac{1}{2F_\pi}\sqrt{2}(f_P - d_P)$	$\frac{1}{2\sqrt{2}}(b_S - \frac{c_S}{3})$	−3.27	−3.18
$A(\Sigma^-_-)$	$\frac{1}{2F_\pi}2(f_P - d_P)$	$-\frac{1}{2}(b_S - \frac{c_S}{3})$	4.27	4.50
$A(\Xi^0_0)$	$-\frac{1}{2F_\pi}\frac{1}{\sqrt{3}}(-3f_P + d_P)$	$-\frac{1}{2\sqrt{3}}(b_S - \frac{c_S}{2})$	3.43	3.18
$A(\Xi^-_-)$	$\frac{1}{2F_\pi}\sqrt{\frac{2}{3}}(-3f_P + d_P)$	$\frac{1}{\sqrt{6}}(b_S - \frac{c_S}{2})$	−4.51	−4.49

IV. PARITY-VIOLATING AMPLITUDES

A. Nonleptonic decays

In the first approximation the parity-violating amplitudes $A(\alpha \to \beta\pi)$ are given by the soft-meson estimate [16,17]:

$$\langle \pi_a \beta | H_W^{\text{p.v.}} | \alpha \rangle = \frac{-i}{F_\pi} \langle \beta | [F_a^5, H_W^{\text{p.v.}}] | \alpha \rangle + q_\mu M_a^\mu \xrightarrow{q \to 0} \frac{-i}{F_\pi} \langle \beta | [F_a^5, H_W^{\text{p.v.}}] | \alpha \rangle, \tag{18}$$

where F_5^a is the axial charge, $F_\pi = 94$ MeV, and $H_W^{\text{p.v.}}$ is the p.v. part of the weak Hamiltonian. Since

$$[F_a^5, H_W^{\text{p.v.}}] = -[F_a, H_W^{\text{p.c.}}], \tag{19}$$

the commutator term $\langle \beta | [F_a^5, H_W^{\text{p.v.}}] | \alpha \rangle$ may be expressed in terms of $\langle \beta | [F_a, H_W^{\text{p.c.}}] | \alpha \rangle$, with F_a the generator of ordinary flavor symmetry. Consequently, the commutator term may be expressed in terms of the matrix elements of $H_W^{\text{p.c.}}$ between appropriate baryon states.

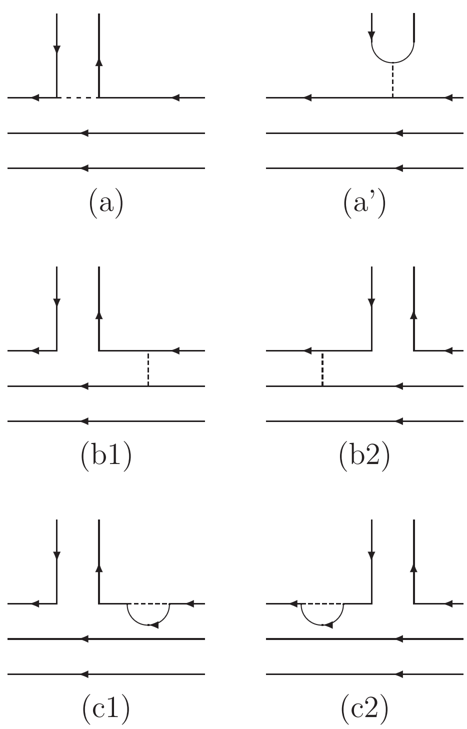

FIG. 1. Quark diagrams for parity-violating weak transitions $B_i \to B_f M$.

In Table III we gather expressions for the p.v. NLHD amplitudes given in terms of the commutator parameters f_P and d_P as well as in terms of the observed p.v. amplitudes b_S and c_S corresponding to the contributions of the topology of the W-exchange and penguin diagrams, respectively. The relevant diagrams are shown in Fig. 1. As shown, the comparison with the data requires

$$b_S \approx -5 \times 10^{-7} \qquad c_S \approx +12 \times 10^{-7}. \tag{20}$$

[Diagrams (a) and (a') of Fig. 1 do not contribute in NLHD.]

When the SU(3) amplitudes f and d describing the p.c. transitions are extracted from b_S and c_S one obtains [2]:

$$f_S \equiv -\frac{F_\pi}{4}\left(b_S - \frac{2}{3}c_S\right) \approx 3.1 \times 10^{-5} \text{ MeV}$$
$$d_S \equiv \frac{F_\pi}{4} b_S \approx -1.2 \times 10^{-5} \text{ MeV}. \tag{21}$$

Comparing with Eq. (12) we see that

$$d_P \approx 2.6 d_S \qquad f_P \approx 1.9 f_S. \tag{22}$$

If one introduces

$$b_P \equiv \frac{4}{F_\pi} d_P \qquad c_P \equiv \frac{6}{F_\pi}(f_P + d_P), \tag{23}$$

one further finds

$$b_P \approx -12.9 \times 10^{-7} \qquad c_P \approx 17.5 \times 10^{-7}. \tag{24}$$

If instead of Eq. (12) one uses the estimates of f_P and d_P given in Eq. (15) one obtains

[2]The values in Eq. (21) are in full agreement with those given in Eq. (6.12) of [1], with the relative relation being $f_S = -2\sqrt{3}F_\pi f_S([1])$, $d_S = -2\sqrt{3}F_\pi d_S([1])$, where $f_S([1]) = -0.92 \times 10^{-7}$, $d_S([1]) = 0.38 \times 10^{-7}$.

$$d_P \approx 2.2 d_S \qquad f_P \approx 1.5 f_S, \tag{25}$$

and

$$b_P \approx -11.1 \times 10^{-7} \qquad c_P \approx 13.4 \times 10^{-7}. \tag{26}$$

Thus, the description of NLHD amplitudes in terms of the leading order terms: the ground-state-baryon pole contribution for the P-waves, and the current-algebra commutator term for the S-waves presents a problem. While soft-meson theorems lead to $f_P = f_S$ and $d_P = d_S$, the values of f_P and d_P as extracted from the parity-conserving amplitudes are larger by a factor of around 2 than those (i.e. f_S, d_S) needed for the description of the parity-violating amplitudes (the S:P problem). In addition, their ratios differ significantly: $f_P/d_P \approx -1.8$ to -1.9, while $f_S/d_S \approx -2.6$. When the problem is expressed in terms of amplitudes b and c corresponding to W-exchange-mediated and single-quark transitions we observe from Eqs. (20), (24), and (26) that the reduction of b_P to b_S is much stronger than a similar reduction of c_P to c_S. This indicates that the dominant correction to the soft-meson formula originates from the W-exchange diagrams.

Following the success of the ground-state baryon pole model in the p.c. sector it was proposed in [18] that the soft-meson expression for the p.v. amplitudes should be supplemented with a substantial correction term $R = q_\mu M^\mu$ originating from a pole model contribution of the negative-parity intermediate excited states [$(70, 1/2^-)$ in the quark model].

The pole model contribution contains two terms [diagrams (1) and (2) in Fig. 1—for both (b) and (c) type transitions], differing in the order of the action of weak and strong transitions. The parity-violating weak transition is described by (see e.g. [19])

$$a_{kl} \bar{u}_k u_l, \tag{27}$$

where the pair of indices k, l describes a pair of $1/2^+$ and $1/2^-$ baryons (B, B^*), i.e. $(k, l) \leftrightarrow (B_k^*, B_l)$ or (B_k, B_l^*). Hermiticity and CP invariance require a to be purely imaginary and antisymmetric [19,20]:

$$a_{kl} = -a_{lk}. \tag{28}$$

The (parity-conserving) strong transition is described by a gradient coupling of the pion. For simplicity, we shall consider π^0 only as its C-parity is well-defined:

$$f_{kl} \bar{u}_k \not{q} u_l \pi^0, \tag{29}$$

with k, l describing as before a pair of $1/2^+$ and $1/2^-$ baryons. Hermiticity and CP invariance require f to be real and antisymmetric:

$$f_{kl} = -f_{lk}. \tag{30}$$

Diagrams (1) [i.e. (b1) and (c1)) and (2) (i.e. (b2) and (c2)] lead to the following total contribution of the $1/2^-$ poles to the $B_i \rightarrow B_f \pi^0$ decay:

$$\left\{ \frac{f_{fk^*} a_{k^*i}}{m_i - m_{k^*}} + \frac{a_{fk^*} f_{k^*i}}{m_f - m_{k^*}} \right\} \bar{u}_f \not{q} u_i \pi^0, \tag{31}$$

where the first (second) term originates from diagram (1) [respectively (2)] and the subscripts k^* label intermediate excited $1/2^-$ baryons. When $m_f = m_i$ one finds that the term in braces is symmetric under $i \leftrightarrow f$ interchange.

Using symmetry properties of a_{k^*l} and f_{k^*l} one can see from Eq. (31) that for $m_i \approx m_f$ the contribution from diagram (1) [alternatively diagram (2)] in $B_i \rightarrow B_f \pi^0$ transition must be equal to the contribution from diagram (2) [alternatively diagram (1)] in $B_f \rightarrow B_i \pi^0$ transition.

In the present paper the sums over intermediate $(70, 1/2^-)$ baryons k^* are not actually performed since our approach deals with their end results only. Thus, only the $SU(6)_W$ structure of the latter is important. Originally, calculations of this structure were carried out in [8,10]. The only difference of this paper with respect to [8,10] is a different relative sign of contributions from diagrams (1) and (2). The presence of this difference is understable as $SU(6)_W$ relates all amplitudes corresponding to diagram (1) [alternatively diagram (2)], but does not relate amplitudes of diagrams (1) to those of diagrams (2). Relation between amplitudes corresponding to diagrams (1) and (2) is dictated by the considerations above and, in particular, by the gradient form of pion coupling. In the $SU(6)_W$ symmetric scheme supplied with the above $i \leftrightarrow f$ symmetry condition the relevant expressions may be therefore readily copied from [12,19] with appropriate sign adjustments. These amplitudes, expressed in terms of amplitudes b_R and c_R, corresponding to W-exchange and penguin diagrams, respectively, are gathered in Table IV. In order to show in an explicit way the $i \leftrightarrow f$ symmetry property required by the gradient coupling, in addition to the amplitudes for the observed decays (Σ_0^+, Σ_+^+ etc.) we also listed the amplitudes for the kinematically forbidden transitions $p \rightarrow \Sigma^+ \pi^0$, $p \rightarrow p \pi^0$, and $\Sigma^+ \rightarrow p U_P^0$ (with $U_P^0 = (\sqrt{3}\pi^0 + \eta_8)/\sqrt{2}$).

When expressed in the language of b and c amplitudes the correction term R leads to

$$b_S = b_P + b_R \qquad c_S = c_P - c_R, \tag{32}$$

with b_R, c_R representing the corrections. The terms b_R and c_R are proportional to $m_i - m_f$ [in Eq. (31) this originates from the $\bar{u}_f \not{q} u_i$ factor] and therefore they vanish in the limit when $q^0 = m_i - m_f \rightarrow 0$. The above formulas are quite general as they follow from SU(3) and the gradient-coupling form only. Later we shall consider SU(3) breaking in the propagators of the intermediate $1/2^-$ states. The size of the $1/2^-$-induced correction terms may be estimated in a quark model [18] and is sizable. Still, the error of such an estimate may be substantial (in [18] it is judged to be of the order of 50%). Consequently, we shall try a different route and estimate the size and sign of the cor-

TABLE IV. Contributions of diagrams (b1), (b2) and (c1), (c2) of Fig. 1 to NLHD amplitudes using $SU(6)_W$ with gradient pion coupling (using [12,19], $U_P^0 = (\sqrt{3}\pi^0 + \eta_8)/\sqrt{2}$).

	Transition	(b1)	(b2)	(c1)	(c2)
Σ_0^+	$\Sigma^+ \to p\pi^0$	0	$\frac{1}{2\sqrt{2}} b_R$	$\frac{1}{6\sqrt{2}} c_R$	0
Σ_+^+	$\Sigma^+ \to n\pi^+$	0	0	0	0
Σ_-^-	$\Sigma^- \to n\pi^-$	0	$-\frac{1}{2} b_R$	$-\frac{1}{6} c_R$	0
Λ_-^0	$\Lambda \to p\pi^-$	0	$-\frac{1}{2\sqrt{6}} b_R$	$-\frac{1}{2\sqrt{6}} c_R$	0
Λ_0^0	$\Lambda \to n\pi^0$	0	$\frac{1}{4\sqrt{3}} b_R$	$\frac{1}{4\sqrt{3}} c_R$	0
Ξ_-^-	$\Xi^- \to \Lambda\pi^-$	0	$\frac{1}{\sqrt{6}} b_R$	$\frac{1}{2\sqrt{6}} c_R$	0
Ξ_0^0	$\Xi^0 \to \Lambda\pi^0$	0	$-\frac{1}{2\sqrt{3}} b_R$	$-\frac{1}{4\sqrt{3}} c_R$	0
	$p \to \Sigma^+\pi^0$	$\frac{1}{2\sqrt{2}} b_R$	0	0	$\frac{1}{6\sqrt{2}} c_R$
	$p \to p\pi^0$	$\frac{1}{2\sqrt{2}} b_R \cot\theta_C$	$\frac{1}{2\sqrt{2}} b_R \cot\theta_C$	$\frac{1}{6\sqrt{2}} c_R \cot\theta_C$	$\frac{1}{6\sqrt{2}} c_R \cot\theta_C$
	$\Sigma^+ \to pU_P^0$	$\frac{1}{2\sqrt{6}} b_R$	$\frac{1}{2\sqrt{6}} b_R$	$\frac{1}{6\sqrt{6}} c_R$	$\frac{1}{6\sqrt{6}} c_R$

rection term from weak radiative hyperon decays using VMD and $SU(2)_W$ ($SU(6)_W$) spin symmetry.

B. Radiative decays

The parity-violating WRHD amplitudes are obtained from those of NLHD amplitudes by replacing the emission of a pion with that of a photon. The connection between the two sets of amplitudes may be obtained via VMD and the symmetry of $SU(6)_W$ ($SU(2)_W$).

When VMD and $SU(6)_W$ are together applied to the commutator term of NLHD they lead to the violation of Hara's theorem [4] in the SU(3) limit [8]. Although the simple quark model, the bag model and early experiments also hinted at the violation of Hara's theorem in that limit, thus suggesting that some assumption of the theorem is violated, the question of Hara's theorem violation is now experimentally settled in the negative by the measurement of the $\Xi^0 \to \Lambda\gamma$ asymmetry [6] as discussed in [11]. Thus, in agreement with the general theoretical expectations (cf. the argument presented in [11]), the soft-meson commutator term present in the p.v. NLHD amplitudes has no $SU(2)_W$-generated counterpart in the WRHD sector.

Consequently, up to an appropriate VMD factor, the parity-violating WRHD amplitudes are the $SU(2)_W$-generated counterparts of the $q_\mu M^\mu$ term in NLHD. As discussed in the previous subsection and in [18], in NLHD this term originates from the pole-model contribution of the intermediate $1/2^-$ excited baryons. Following [18], it was therefore proposed in [21] that the parity-violating WRHD transitions are generated in an analogous manner.

Thus, in the considerations of the previous section one has to replace the strong vertices of Eq. (29) by the electromagnetic ones:

$$\mu_{kl}\bar{u}_k i\sigma^{\mu\nu}\gamma_5 q_\nu u_l A_\mu, \tag{33}$$

with k, l describing as before a pair of $1/2^+$ and $1/2^-$ baryons. Hermiticity and CP invariance require μ to be purely imaginary and symmetric [19,20]:

$$\mu_{kl} = \mu_{lk}. \tag{34}$$

Diagrams (1) and (2) lead to the following total contribution of the $1/2^-$ poles to the $B_i \to B_f\gamma$ decay:

$$\left\{\frac{\mu_{fk^*}a_{k^*i}}{m_i - m_{k^*}} + \frac{a_{fk^*}\mu_{k^*i}}{m_f - m_{k^*}}\right\}\bar{u}_f i\sigma^{\mu\nu}\gamma_5 q_\nu u_i A_\mu. \tag{35}$$

When $m_i \approx m_f$ one finds that the term in braces in Eq. (35) is antisymmetric under $i \leftrightarrow f$ interchange.

Using symmetry properties of a_{k^*l} and μ_{k^*l} one can see from Eq. (35) that for $m_i \approx m_f$ the contribution from diagram (1) [alternatively diagram (2)] in $B_i \to B_f\gamma$ transition must be opposite in sign and equal in absolute magnitude to the contribution from diagram (2) [alternatively diagram (1)] in $B_f \to B_i\gamma$ transition. For $i = \Sigma^+$ and $f = p$ it then follows from the symmetry of the weak interaction Hamiltonian under $s \leftrightarrow d$ (and thus $\Sigma^+ \leftrightarrow p$) that the term in braces in Eq. (35) has to be both antisymmetric and symmetric under $i \leftrightarrow f$ interchange. This means that the p.v. $\Sigma^+ \to p\gamma$ amplitude must vanish in the SU(3) limit (Hara's theorem [4]).

As in the case of NLHD, in the $SU(6)_W$ approach to WRHD the sums over intermediate $(70, 1/2^-)$ baryons k^* are not actually performed, since we deal with the end results only. Up to an appropriate normalization and the VMD factor the $SU(6)_W$ scheme relates then the correction terms to the p.v. NLHD amplitudes and the p.v. WRHD amplitudes for diagrams (1). Similar connection exists (separately) for diagrams (2). The obtained amplitudes $A(B_i \to B_f U^0)$, copied from [2,12] with appropriate sign adjustment (as in Table IV), are gathered in Table V. Contributions from the coupling of the U^0 vector meson (later photon) to the strange quark are described by parameter ϵ [$= 1$ in SU(3)]. All single-quark contributions may be lumped into a single unknown parameter s_R (which includes c_R and the amplitudes corresponding to

TABLE V. Contributions of diagrams (b1), (b2) of Fig. 1 and single-quark transitions to $A(B_i \to B_f U^0)$ amplitudes using $SU(6)_W$ and Table IV (from [2,12]).

Process	diagram (b1)	diagram (b2)	single-quark
$\Sigma^+ \to pU^0$	$\frac{1}{3\sqrt{2}} b_R$	$-\frac{2+\epsilon}{9\sqrt{2}} b_R$	$\frac{1}{3\sqrt{2}} s_R$
$\Sigma^0 \to nU^0$	$\frac{1}{6} b_R$	$\frac{2+\epsilon}{18} b_R$	$-\frac{1}{6} s_R$
$\Lambda \to nU^0$	$-\frac{1}{6\sqrt{3}} b_R$	$\frac{2+\epsilon}{6\sqrt{3}} b_R$	$-\frac{\sqrt{3}}{2} s_R$
$\Xi^0 \to \Lambda U^0$	0	$-\frac{2+\epsilon}{9\sqrt{3}} b_R$	$\frac{1}{2\sqrt{3}} s_R$
$\Xi^0 \to \Sigma^0 U^0$	$-\frac{1}{3} b_R$	0	$-\frac{5}{6} s_R$
$\Xi^- \to \Sigma^- U^0$	0	0	$\frac{5}{3\sqrt{2}} s_R$

diagrams (a), (a') of Fig. 1, see [8,10]). Any SU(3)-breaking effects may be absorbed into its definition. The relevant contributions are also gathered in Table V.

In the pole model with broken SU(3) ($\Delta m_s = 190$ MeV) the contribution from diagrams (1) and (2) given in Tables IV and V will be somewhat modified. Namely, diagrams (1) are associated with the presence of mass denominators $\Delta\omega - \Delta m_s$, while for diagrams (2) these mass denominators contain $\Delta\omega + \Delta m_s$, where $\Delta\omega \approx 570$ MeV is the average splitting between the $(56, 1/2^+)$ and $(70, 1/2^-)$ multiplets. Assuming that all of the SU(3) breaking originates from mass differences (plus possibly from a reduced coupling of U^0/photon to the strange quark, i.e. $\epsilon < 1$), we may take it into account by multiplying the contributions of diagrams (1) by $\Delta\omega/(\Delta\omega - \Delta m_s) \equiv 1/(1-x)$ with $x \approx 1/3$. For diagrams (2) the relevant factor is $1/(1+x)$. It is mainly through the presence of these SU(3) breaking effects that the present paper differs from Ref. [11].

Using the above considerations one obtains the following expressions for the parity-violating WRHD amplitudes:

$$D(B_i \to B_f \gamma) = \left(\frac{e}{g}\right) \frac{1}{(m_i - m_f)\sqrt{2}} A(B_i \to B_f U^0), \quad (36)$$

where amplitudes A are related by $SU(2)_W$ to the (vanishing in the soft-meson limit) correction terms in NLHD:

$$A(\Sigma^+ \to pU^0) = \frac{1}{9\sqrt{2}} \frac{6x + (1-\epsilon)(1-x)}{1-x^2} b_R + \frac{1}{3\sqrt{2}} s_R$$
$$A(\Sigma^0 \to nU^0) = \frac{1}{18} \frac{6-(1-\epsilon)(1-x)}{1-x^2} b_R - \frac{1}{6} s_R$$
$$A(\Lambda \to nU^0) = -\frac{1}{6\sqrt{3}} \frac{4x - 2 + (1-\epsilon)(1-x)}{1-x^2} b_R - \frac{\sqrt{3}}{2} s_R$$
$$A(\Xi^0 \to \Lambda U^0) = -\frac{2+\epsilon}{9\sqrt{3}} \frac{1-x}{1-x^2} b_R + \frac{1}{2\sqrt{3}} s_R$$
$$A(\Xi^0 \to \Sigma^0 U^0) = -\frac{1}{3} \frac{1+x}{1-x^2} b_R - \frac{5}{6} s_R$$
$$A(\Xi^- \to \Sigma^- U^0) = \frac{5}{3\sqrt{2}} s_R. \quad (37)$$

TABLE VI. Numerical values of coefficients at b_R in Eqs. (37).

Process	$x = 0$, $\epsilon = 1$	$x = 1/3$, $\epsilon = 2/3$
$\Sigma^+ \to p\gamma$	0	0.196
$\Lambda \to n\gamma$	0.192	0.048
$\Xi^0 \to \Lambda\gamma$	−0.192	−0.128
$\Xi^0 \to \Sigma^0\gamma$	−0.333	−0.5

The factor $1/(m_i - m_f)$ in Eq. (36) is canceled by factors $m_i - m_f$ contained in amplitudes $A(B_i \to B_f U^0)$ since b_R and s_R vanish in the limit $m_i - m_f \to 0$, as discussed after Eq. (32).

In Table VI we show the sizes of the coefficients at the b_R terms in Eqs. (37) [and hence, up to a factor, in Eq. (36)] for the SU(3)-symmetric case ($x = 0$, $\epsilon = 1$), and the SU(3)-breaking case ($x = 1/3$, $\epsilon = 2/3$). Please note that with growing x the $\Sigma^+ \to p\gamma$ coefficient increases from zero very fast, so that at $x = 1/3$ it becomes larger than the absolute value of the corresponding coefficient for $\Xi^0 \to \Lambda\gamma$ [for $\Xi^0 \to \Lambda(\Sigma^0)\gamma$ the relevant change is of the order of 30%, as naively expected for SU(3) breaking effects]. Thus, for the $\Sigma^+ \to p\gamma$ p.v. amplitude the SU(3) breaking effect is very large indeed.

V. FROM RADIATIVE TO NONLEPTONIC DECAYS

Since the parity-conserving WRHD amplitudes are known via the symmetry connection to the parity-conserving NLHD amplitudes (see Table II), the branching ratios and asymmetries of WRHD provide information on the size of parity-violating WRHD amplitudes, and, consequently, on parameters b_R and s_R.

Present data on WRHD are gathered in Table VII. In order to get information on the size of b_R and s_R we performed fits to the five known branching ratios (given in Table VII) and the three well-known asymmetries (as in Table VII with the exception of $\Xi^- \to \Sigma^-\gamma$). Since only a rough description of the data can be achieved in this way, we decided not to use the experimental errors in the fitting procedures. Still, the fits yield a fairly well-defined value of s_R (around −0.75). For the fixed value $s_R = -0.75$ one can study then how the branching ratios and asymmetries depend on the value of b_R. In Table VII we present results of such calculations for three values of b_R, i.e. for $b_R = +4.2$, +5.3, and +6.5. One can see that in this range of b_R theory is in a reasonable agreement with the data.

Quantification of this agreement in terms of a χ^2-like function depends on the details of how the errors are treated. A reasonable requirement to impose is to admit equal deviations from unity of the ratios of $x_i \equiv \mathcal{B}_i(\text{the})/\mathcal{B}_i(\text{exp})$ and $y_k \equiv \alpha_k(\text{the})/\alpha_k(\text{exp})$ with $\mathcal{B}_i$ (α_k) denoting the branching ratios (asymmetries) in question. Since the position of the minimum depends somewhat on whether one uses $\sum_i (x_i - 1)^2 + \sum_k (y_k - 1)^2$ or a similar

TABLE VII. Fit to branching ratios and asymmetries of weak radiative hyperon decays; data from [15] and from [6] (marked with *).

Process	Data	Fit $b_R = +4.2$, $s_R = -0.75$	Fit $b_R \approx +5.3$, $s_R \approx -0.75$	Fit $b_R = +6.5$, $s_R = -0.75$	Ref. [21]
		Branching ratio (in units of 10^{-3})			
$\Sigma^+ \to p\gamma$	1.23 ± 0.05	0.68	0.72	0.78	$0.92^{+0.32}_{-0.14}$
$\Lambda \to n\gamma$	1.75 ± 0.15	0.74	0.77	0.80	0.62
$\Xi^0 \to \Lambda\gamma$	$1.16 \pm 0.08^{*)}$	0.91	1.02	1.17	3.0
$\Xi^0 \to \Sigma^0\gamma$	3.33 ± 0.10	3.80	4.42	5.33	7.2
$\Xi^- \to \Sigma^-\gamma$	0.127 ± 0.023	0.16	0.16	0.16	
		Asymmetry			
$\Sigma^+ \to p\gamma$	-0.76 ± 0.08	−0.54	−0.67	−0.79	$-0.80^{+0.32}_{-0.19}$
$\Lambda \to n\gamma$		−0.90	−0.93	−0.95	−0.49
$\Xi^0 \to \Lambda\gamma$	$-0.78 \pm 0.19^{*)}$	−0.92	−0.97	−0.99	−0.78
$\Xi^0 \to \Sigma^0\gamma$	-0.63 ± 0.09	−0.78	−0.92	−0.99	−0.96
$\Xi^- \to \Sigma^-\gamma$	$+1.0 \pm 1.3$	+0.8	+0.8	+0.8	

function with $x_i \to 1/x_i$ and $y_k \to 1/y_k$, we decided to consider the minimization of the function

$$\sum_{i=1}^{5}\left(\frac{\mathcal{B}_i(\text{the}) - \mathcal{B}_i(\text{exp})}{\mathcal{B}_i(\text{the}) + \mathcal{B}_i(\text{exp})}\right)^2 + \sum_{k=1}^{3}\left(\frac{\alpha_k(\text{the}) - \alpha_k(\text{exp})}{\alpha_k(\text{the}) + \alpha_k(\text{exp})}\right)^2, \tag{38}$$

which embodies such requirements in a more symmetric way (i.e. it treats the theoretical and the experimental entries in the same way). The fitted values of b_R and s_R are then (in units of 10^{-7})

$$b_R \approx +5.3 \tag{39}$$

$$s_R \approx -0.75. \tag{40}$$

Putting aside the $\Xi^- \to \Sigma^-\gamma$ branching ratio which depends on s_R only, we observe from Table VII that the $\Xi^0 \to \Sigma^0\gamma$ branching ratio is overestimated while the branching ratios of $\Sigma^+ \to p\gamma$, and $\Lambda \to n\gamma$ are underestimated. This suggests that there might be a problem with $\Xi^0 \to \Sigma^0\gamma$. Consequently, it seems more likely that b_R is somewhat larger than 5.3.

For the $\Sigma^+ \to p\gamma$ the discrepancy between the model and experiment is about 20% at the amplitude level. For the $\Xi^0 \to \Sigma^0\gamma$ the discrepancy is larger. There seems to be an even larger discrepancy for the $\Lambda \to n\gamma$ branching ratio, but—as already discussed—this is the decay for which strong cancellations occur in the parity-conserving amplitude with the result depending quite substantially on the precise value of the transition moment $\mu_{\Sigma\Lambda}$. The overall description of the data is rough but fairly satisfactory and indicates that $b_R \approx -b_S$ as discussed in [11]. The fits have a clear tendency to choose a small negative value for s_R, thus predicting a substantial positive asymmetry for the $\Xi^- \to \Sigma^-\gamma$ decay. For comparison, in Table VII we also quoted the branching ratios and asymmetries calculated in [21].

When one inserts the value $b_R \approx 6.0 \times 10^{-7}$ into Eq. (32) one obtains (in units of 10^{-7})

$$-5 \approx -12.9 + \frac{6.0}{1 + x} = -8.4, \tag{41}$$

or, if the estimate of b_P (Eqs. (15, 26)) from [1] is used,

$$-5 \approx -11.1 + \frac{6.0}{1 + x} = -6.6, \tag{42}$$

where the factor $1/(1 + x)$ takes into account the SU(3) breaking in the propagators of amplitudes (b2). [This is consistent with the analogous factors used in the derivation of Eqs. (37).] The discrepancy between the P- and S-waves is now significantly smaller, especially for the values of f_P and d_P extracted in [1].

If one accepts that the small size of s_R suggests the smallness of c_R [s_R contains contributions from c_R and diagrams (a), (a') in Fig. 1, and therefore one cannot determine c_R uniquely], one concludes that one should have

$$c_S \approx c_P. \tag{43}$$

This is indeed true for the parameters of [1] for which Eq. (43) reads:

$$12 \times 10^{-7} \approx 13.4 \times 10^{-7}. \tag{44}$$

In conclusion, we have shown that the argument of Ref. [11] works fairly well also when SU(3) is broken in the p.v. sector as well. Still, some room for unaccounted corrections is obviously present.

VI. CONCLUSIONS

The aim of this paper was to provide both an explanation of the S:P puzzle in NLHD and a successful description of the gross structure of the observed pattern of asymmetries and branching ratios in WRHD in the SU(3) breaking case, the two explanations being related as discussed in [11]. The scheme maintains an intimate connection between NLHD and WRHD, uses VMD, and yet it does not lead to the Hara's-theorem-violating results of both the constituent quark model [7] and the original VMD approach of [8,13].

The resolution of the problem of weak hyperon decays given in the present paper was originally suggested in [18,21]. However, complex quark model calculations of [18,21] did not make it easy to see the simple $SU(2)_W$ symmetry connection existing between the p.v. WRHD amplitudes and the correction term in the p.v. NLHD amplitudes. In Ref. [18] the correction to the CA commutator term in the p.v. NLHD amplitudes is due to the $(70, 1/2^-)$ intermediate states. In our approach explicit calculations of the contributions from the individual intermediate states and the subsequent summation are not performed. Instead, we work at the level of the total resulting contribution. Still, the symmetry properties of the correction term in Ref. [18] and in this paper are identical in the appropriate limit. The difference is that in our paper, instead of calculating the overall size of the correction in a quark model as the authors of Ref. [18] do, we extract both its size and sign from WRHD [via $SU(2)_W$ and VMD], thus bypassing many quark model uncertainties.

Our identification of how symmetry should be applied for a successful joint description of nonleptonic and radiative weak hyperon decays leads to problems elsewhere, however. Namely, present understanding of nuclear parity violation (cf. Ref. [10]) is based on symmetry between the *full* p.v. weak amplitudes $B' \rightarrow BP$ and $B' \rightarrow BV$. According to [10,22] the explanation of data on nuclear parity violation can be obtained through the dominance of the weak ρ-nucleon coupling of the form $\bar{u}_N \gamma_\mu \gamma_5 u_N \rho^\mu$. Via vector-meson dominance this leads to photon-nucleon coupling $\bar{u}_N \gamma_\mu \gamma_5 u_N A^\mu$ which entails the violation of Hara's theorem [2,8]. Since Hara's theorem is satisfied, it follows that either the standard form of VMD is not universal or our present understanding of nuclear parity violation is not fully correct.

ACKNOWLEDGMENTS

This work was supported in part by the Polish State Committee for Scientific Research Grant No. 2 P03B 046 25.

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