PHYSICAL REVIEW D VOLUME 50, NUMBER 5 1 SEPTEMBER 1994

Weak hyperon decays: Quark sea and SU(3) symmetry breaking

P. Żenczykowski
Department of Theoretical Physics, H. Niewodniczański Institute of Nuclear Physics, Radzikowskiego 152, Kraków, Poland
(Received 16 November 1993)

An explanation of the difference in the values of the apparent f/d ratios for the S- and P-wave amplitudes of nonleptonic hyperon decays is proposed. The argument is formulated in the framework of the standard pole model with $(56,0^+)$ ground-state and $(70,1^-)$ excited baryons as intermediate states for the P and S waves, respectively. Under the assumption that the dominant part of the deviation of $(f/d)_{P\text{ wave}}$ from -1 is due to large quark sea effects, SU(3) symmetry breaking in energy denominators is shown to lead to a prediction for $(f/d)_{S\text{ wave}}$ which is in excellent agreement with experiment. This corroborates our previous unitarity calculations which indicated that the matrix elements $\langle B|H^{\rm PC}_{\rm weak}|B'\rangle$ of the parity-conserving weak Hamiltonian between the ground-state baryons are characterized by $f_0/d_0 \approx -1.6$ or more. A brief discussion of the problem of the relative size of S- and P-wave amplitudes is given. Finally, implications for weak radiative hyperon decays are also discussed.

PACS number(s): 13.30.Eg, 11.40.Ha, 14.20.Jn

I. INTRODUCTION

Despite several decades of theoretical inquiry, our understanding of weak hyperon decays has remained elusive and controversial [1]. Dominantly, hyperons decay weakly into two-body pion+baryon channels. Various models proposed for a theoretical description of these nonleptonic processes always relate to an approach based on PCAC (partial conservation of axial vector current) and current algebra (CA) [2]. One of the reasons for such a pronounced role of that approach is that it is theoretically attractive: It allows a parallel treatment of the S and P waves, expressing both of these as functions of the transition matrix elements $\langle B'|H^{\rm PC}_{\rm weak}|B\rangle$ of the parity-conserving part of the weak Hamiltonian.

Unfortunately, this PCAC-CA approach is less appealing when confronted with experiment as it presents us with two serious difficulties. The first concerns the relative size of the S and P waves: Current algebra overestimates the S:P ratio by a factor of around 2. The second is related to the SU(3) structure of the decays. The quark model prediction for the two SU(3)-invariant couplings f_0, d_0 describing the SU(3) structure of the $\langle B'|H^{\rm PC}_{\rm weak}|B\rangle$ matrix elements is $f_0/d_0 = -1$, while the experimental S waves require $f/d \approx -2.5$. Similarly, the value of the f/d ratio extracted from the P waves is different from -1. Its exact value is sensitive to the way one treats SU(3) breaking in energy denominators and couplings. When SU(3)-symmetric $\pi BB'$ couplings and equal spacing of ground-state octet baryons are used, one infers from the P-wave amplitudes that $f/d \approx -1.8$ or -1.9 [1,3,4].

As yet, there is no general consensus as to what a full resolution of the above problems might be. On one side, it is rather generally acknowledged that an important correction to the CA results stems from a more realistic treatment of the contribution from the intermediate $(70,1^-)$ baryons. In particular, SU(3) breaking in energy denominators generates corrections which subtract from the standard soft pion contribution [5]. The correction is of order $\delta s/\Delta\omega_s \approx 0.3-0.4$ relative to that of the commutator [δs is the SU(3)-breaking parameter (≈ 190 MeV) and $\Delta\omega_s$ is the mean spacing of $(56,0^+)$ and $(70,1^-)$ baryons]. On the other side, however, no such consensus has been reached so far on the question of the f/d ratio. In fact, several different explanations of the deviation of f/d from -1 have been proposed.

In their original paper [5], LeYaouanc *et al.* have suggested that f/d is larger in parity-violating amplitudes because for different decays such as $\Lambda\to N\pi, \Sigma\to N\pi, \ldots$ the corrections due to $(70,1^-)$ baryons appear to be proportional to different mass differences of ground-state baryons $(\Lambda-N, \Sigma-N, \ldots)$. With $\Sigma-\Lambda\neq 0$ one obtains then an increase of the *effective* f/d ratio. The problem with this explanation is that Σ-Λ splitting is a second order effect due to spin-spin interactions which were neglected in the intermediate $(70,1^-)$ baryons in Ref. [5]. If spin-spin interactions are also neglected for ground-state baryons, one recovers for the $(70,1^-)$ correction the canonical quark model value $f/d = -1$.

Another possible and at first sight natural explanation is to attribute the departure of f/d from -1 to a contribution of diagrams with weak Hamiltonian acting in the meson leg. Such diagrams are characterized by $d_{\rm mes}/f_{\rm mes} = 0$, and thus they might provide the much needed enhancement of f. For the S waves, they were invoked by Gronau [6], who introduced the contribution of the K^* intermediate meson. The contribution of such diagrams has been later discussed in various papers by Bonvin [7], Nardulli [8], Stech and Xu [9], and others. The main problem with this line of reasoning is that one expects such contributions to be small on general grounds. Indeed, for the P waves, the K-pole contribution is proportional to $p_\pi \cdot p_K \sim m_\pi^2$ as a result of chiral symmetry [1] and it should vanish for $m_\pi^2 \to 0$. For the S waves, one can show that in the limit of exact SU(3) symmetry such diagrams should give a vanishing contribution as well (see, e.g., Ref. [10]). In the case of broken SU(3), one might expect corrections to the quark model value of -1 of order δs/(hadron mass scale)$\approx$20–30 %, but not 100–150 %.

The third possibility discussed in the literature consists of a large departure from the assignment of the canonical value of $f_0/d_0=-1$ to the (directly not measurable) matrix elements $\langle B|H^{\rm PC}_{\rm weak}|B'\rangle$ of the parity-conserving part of the weak Hamiltonian between the ground-state baryons. This departure is attributed to the contribution from the sea quarks [10,11]. In quantum chromodynamics this corresponds to the consideration of penguin diagrams. On one side, direct evaluation of these diagrams leads to a small increase of f/d only [12]. On the other side, if one estimates the contribution of the penguins by relating them to the gluon-induced Δ-N splitting, one obtains [11] a substantial increase of f_0/d_0 to -1.6. Although the size of this renormalization of f/d is determined by the experimentally observed Δ-N splitting, it corresponds to a large value of the QCD coupling constant, believed by many to be unrealistic (see, however, Ref. [13]). A different origin for a large contribution from sea quarks has been proposed recently in Ref. [14]. It has been shown there that the interference of strong and parity-conserving weak (P-wave) amplitudes leads to a substantial increase of the f_0/d_0 ratio characterizing the $\langle B'|H^{\rm PC}_{\rm weak}|B\rangle$ matrix elements. When the size of hadronic loops thus generated by unitarity is estimated by comparison with hadron mass splittings, one finds that f_0/d_0 is shifted by such hadronic penguins to around -1.6 or more. The exact value depends slightly on how much of the Δ-N splitting is attributed to hadron-level (unitarity) effects. Even with a moderate (around 80 MeV) pion-induced contribution to the Δ-N splitting, one obtains $f_0/d_0=-1.5$ [14]. For larger contributions of this type as in the unitarized quark model [15,16], one gets f_0/d_0 around -1.6 or more. Thus one can have both a smaller QCD coupling governing the short distance effects and large (hadron-level-induced) sea effects.

In this paper we study in more detail how these sea effects manifest themselves in S- and P-wave amplitudes. We work in the framework of a kind of "skeleton" pole model which both includes the essential SU(3)-breaking effects of the pole model and, at the same time, retains much of the simplicity of the PCAC-CA approach by bypassing the need to use a detailed information on the $\frac{1}{2}^-$ baryons in the intermediate states.

We find that the model thus constructed explains the f/d structure of both the P- and S-wave amplitudes very naturally. In fact, joint consideration of large quark sea effects and SU(3) breaking in energy denominators leads, *without any new parameters,* to the following approximate relationship between the deviations from -1 of the observed[1] f/d ratios in S- and P-wave amplitudes:

$$\frac{(f/d+1)_{S\ \rm wave}}{(f/d+1)_{P\ \rm wave}}=\frac{1+x}{1-x}\ , \tag{1}$$

where $x=\delta s/\Delta\omega_s\approx 0.3$–$0.4$.

Using the experimental values for the corresponding f/d ratios (-2.6 for S waves, -1.85 to -1.9 for P waves), Eq. (1) reads 1.8–$1.9=2.1\pm0.25$. The experimentally observed deviation of $(f/d)_{P\ \rm wave}$ from -1 is in agreement with the unitarity-based calculation [14] of the SU(3) structure of the $\langle B'|H^{\rm PC}_{\rm weak}|B\rangle$ matrix elements: $(f/d)_{P\ \rm weak}\approx f_0/d_0$, or -1.8 to -1.9 ≈ -1.6 to -1.7. This is consistent with general hadron-level arguments permitting only a small correction from meson-leg diagrams to $(f/d)_{P\ \rm wave}$.

The paper is organized as follows. In the next section we exhibit the basic SU(3)-symmetric connections between the quark diagrams, the pole model, and the PCAC-CA approach for the S-wave amplitudes. In Sec. III the standard description of the P-wave amplitudes and the assignment of the dominant part of the deviation of $(f/d)_{P\ \rm wave}$ from -1 to quark sea effects is discussed in some detail. Section IV contains the analysis of the SU(3)-symmetry-breaking effects in the energy denominators of the pole model for the S-wave amplitudes. Equation (1) is derived there. It is also shown there that the S-wave reduction mechanism of LeYaouanc *et al.* becomes unimportant for $f_0/d_0\approx -1.7$. In an attempt to deal with this reappearing S:P problem, in Sec. V we briefly consider the contribution from the radially excited $(56,0^+)^*\frac{1}{2}^+$ baryons. We find that, if the relevant f^*/d^* ratio is equal to that of ground-state baryons, the contribution of radially excited states cannot cure the S:P problem. We argue then that the smallness of the experimental S:P ratio may be related to the departure of the ratio $g_{B(1/2^+)B^*(1/2^-)P}/g_{B(1/2^+)B'(1/2^+)P}$ of strong hadron couplings from quark model predictions. In Sec. VI a brief discussion is given of the modifications to the combined symmetry–vector-meson-dominance approach to weak radiative hyperon decays that originate from the effect considered in this paper. Finally, in Sec. VII we reiterate the main points of our paper.

II. PARITY-VIOLATING AMPLITUDES

All quark-line diagrams that may in principal contribute to weak hyperon decays are shown in Fig. 1. Diagrams (a) and (a′) correspond to the meson-leg topology, while diagrams (b), (c), (d), and (e) admit intermediate baryons in between the action of the weak Hamiltonian and the strong (meson-emission) vertex.

For the parity-violating amplitudes, the contributions from diagrams (a) and (a′) vanish in the SU(3)-symmetry limit [10]. Similarly, the Lee-Swift theorem [17] requires the vanishing of diagrams (d) and (e). Diagrams (b) are the familiar W-exchange processes that lead to $f/d=-1$, while diagrams (c) are the sea diagrams (with $d=0$). In an SU(6)$_W$-symmetric approach, the contributions from diagrams (b1), (b2), (c1), and (c2) can be calculated using the quark model technique of Desplanques, Donoghue, and Holstein [10] and are gathered in Table I. For completeness, the weights for the kinematically forbidden transitions are also given.

In terms of the reduced matrix elements b and c corresponding to diagrams (b1),(b2) and (c1),(c2), respectively, one obtains from Table I the following expressions for the

[1] If meson-leg contributions to f are small.

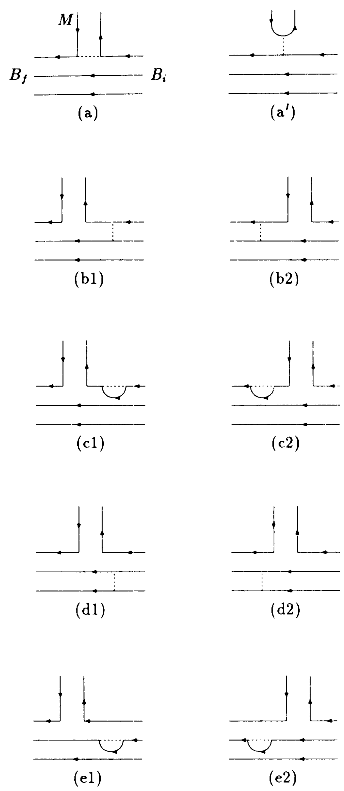

FIG. 1. Quark diagrams for weak decays.

parity-violating amplitudes:

$$\begin{aligned}
&A(\Sigma_0^+)=\frac{1}{2\sqrt{2}}b-\frac{1}{6\sqrt{2}}c\ ,\\
&A(\Sigma_+^+)=0\ ,\\
&A(\Sigma_-^-)=-\tfrac{1}{2}b+\tfrac{1}{6}c\ ,\\
&A(\Lambda_-^0)=-\sqrt{2}A(\Lambda_0^0)=-\frac{1}{2\sqrt{6}}b+\frac{1}{2\sqrt{6}}c\ ,\\
&A(\Xi_-^-)=-\sqrt{2}A(\Xi_0^0)=\frac{1}{\sqrt{6}}b-\frac{1}{2\sqrt{6}}c\ .
\end{aligned} \tag{2}$$

For the kinematically forbidden amplitudes, one gets, similarly,

$$A(\Sigma^+\to p\eta_8)=\left[-\frac{1}{\sqrt{6}}-\frac{1}{2\sqrt{6}}\right]b+\left[\frac{1}{6\sqrt{6}}+\frac{1}{3\sqrt{6}}\right]c\ , \tag{3}$$

etc. (i.e., the entries from Table I that correspond to diagrams (b1),(b2) [(c1),(c2)] are to be *added*). Experiment fixes then $b=-5$, $c=+12$ (in units of 10^{-7}; see Ref. [3]), $f/d=-1+(2c)/(3b)\approx-2.6$.

Let us discuss how formulas (2) are related to the pole model and the PCAC-CA approach. For the sake of definiteness, consider the $\Sigma^+\to p\pi^0$ decay. Upon using the PCAC relation between the pion field and the divergence of the axial vector current, the calculation of the *S*-wave amplitude $A(\Sigma^+\to p\pi^0)$ in the pole model in-

TABLE I. Weights of quark diagrams (b) and (c) for the *S*-wave (PV) amplitudes.

	Transition	(b1)	(b2)	(c1)	(c2)
Σ_0^+	$\Sigma^+\to p\pi^0$	0	$\frac{1}{2\sqrt{2}}$	$-\frac{1}{6\sqrt{2}}$	0
Σ_+^+	$\Sigma^+\to n\pi^+$	0	0	0	0
	$\Sigma^+\to p\eta_8$	$-\frac{1}{\sqrt{6}}$	$-\frac{1}{2\sqrt{6}}$	$\frac{1}{6\sqrt{6}}$	$\frac{1}{3\sqrt{6}}$
Σ_-^-	$\Sigma^-\to n\pi^-$	0	$-\frac{1}{2}$	$\frac{1}{6}$	0
Λ_-^0	$\Lambda\to p\pi^-$	0	$-\frac{1}{2\sqrt{6}}$	$\frac{1}{2\sqrt{6}}$	0
Λ_0^0	$\Lambda\to n\pi^0$	0	$\frac{1}{4\sqrt{3}}$	$-\frac{1}{4\sqrt{3}}$	0
	$\Lambda\to n\eta_8$	$-\frac{1}{6}$	$-\frac{1}{12}$	$\frac{1}{12}$	$\frac{1}{6}$
Ξ_-^-	$\Xi^-\to\Lambda\pi^-$	0	$\frac{1}{\sqrt{6}}$	$-\frac{1}{2\sqrt{6}}$	0
Ξ_0^0	$\Xi^0\to\Lambda\pi^0$	0	$-\frac{1}{2\sqrt{3}}$	$\frac{1}{4\sqrt{3}}$	0
	$\Xi^0\to\Lambda\eta_8$	$\frac{1}{3}$	$\frac{1}{6}$	$-\frac{1}{12}$	$-\frac{1}{6}$
	$\Xi^-\to\Sigma^-\pi^0$	0	0	$\frac{1}{6\sqrt{2}}$	0
	$\Xi^-\to\Sigma^-\eta_8$	0	0	$-\frac{1}{6\sqrt{6}}$	$-\frac{1}{3\sqrt{6}}$
	$p\to K^0p$	0	$-\frac{1}{2}$	0	$\frac{1}{6}$

volves consideration of the expressions

$$A_{(1)}(\Sigma^+ \to p\pi^0) = \frac{\langle p|\partial_\mu A_\mu^{(0)}|N^*\rangle\langle N^*|H_{\rm weak}^{\rm PV}|\Sigma^+\rangle}{\Delta\omega_{W_1}} \quad (4)$$

and

$$A_{(2)}(\Sigma^+ \to p\pi^0) = \frac{\langle p|H_{\rm weak}^{\rm PV}|\Sigma^*\rangle\langle \Sigma^*|\partial_\mu A_\mu^{(0)}|\Sigma^+\rangle}{\Delta\omega_{W_2}} , \quad (5)$$

corresponding to diagrams (b1),(c1) and (b2),(c2), respectively. [We have ignored uninteresting factors such as $1/f_\pi$ on the right-hand side (RHS) of Eqs. (4) and (5).] In Eqs. (4) and (5), the dominant contribution is expected to arise from the N^* and Σ^* $(70,1^-)\frac{1}{2}^-$ intermediate states. The energy denominators $\Delta\omega_{W_1},\Delta\omega_{W_2}$ have subscripts W_1,W_2 since they correspond to the energy difference "across" the *weak* interaction:

$$\Delta\omega_{W_1} = N^* - \Sigma, \quad \Delta\omega_{W_2} = \Sigma^* - p . \quad (6)$$

Since the matrix elements of the spatial components A_k of the axial vector current between $\langle p|$ and $|N^*\rangle$ ($\langle\Sigma^*|$ and $|\Sigma^+\rangle$) vanish (see Ref. [5]), we have

$$\begin{aligned} &\frac{1}{i}\langle p|\partial_\mu A_\mu|N^*\rangle = \Delta\omega_s\langle p|A_0|N^*\rangle , \\ &\frac{1}{i}\langle \Sigma^*|\partial_\mu A_\mu|\Sigma^+\rangle = -\Delta\omega_s\langle \Sigma^*|A_0|\Sigma^+\rangle , \end{aligned} \quad (7)$$

where we have used the subscript s to denote the baryon energy difference "across" the *strong* vertex:

$$\Delta\omega_s = N^* - p = \Sigma^* - \Sigma . \quad (8)$$

In the SU(3) limit, we have $\Delta\omega_s = \Delta\omega_{W_1} = \Delta\omega_{W_2}$ and one obtains, from Eqs. (4) and (5),

$$\begin{aligned} &A_{(1)}(\Sigma^+ \to p\pi^0) = \langle p|A_0|N^*\rangle\langle N^*|H_{\rm weak}^{\rm PV}|\Sigma^+\rangle , \\ &A_{(2)}(\Sigma^+ \to p\pi^0) = -\langle p|H_{\rm weak}^{\rm PV}|\Sigma^*\rangle\langle \Sigma^*|A_0|\Sigma^+\rangle ; \end{aligned} \quad (9)$$

i.e., we recover the standard commutator prescription of current algebra,

$$A = A_{(1)} + A_{(2)} = \langle p|[A_0, H_{\rm weak}^{\rm PV}]|\Sigma^+\rangle , \quad (10)$$

which, upon using the commutation relation

$$[A_0, H_{\rm weak}^{\rm PV}] = [V_0, H_{\rm weak}^{\rm PC}] , \quad (11)$$

enables us to express $A(\Sigma_0^+)$ in terms of the matrix element $\langle p|H_{\rm weak}^{\rm PC}|\Sigma^+\rangle$.

III. PARITY-CONSERVING AMPLITUDES

The SU(6) structure of the parity-conserving amplitudes corresponding to the diagrams of Fig. 1 may be calculated using, as before, the quark model technique of Refs. [10,3]. This time, however, the dominant contribution is expected to come from the ground-state baryons as intermediate states. This introduces the energy denominators (here for $Y \to N\pi$ processes)

$$\frac{1}{N-Y} \left[\frac{1}{Y-N} \right] \quad (12)$$

for diagrams (b1),(c1) [(b2),(c2)], respectively. On account of the sign difference between these energy denominators, the SU(6) factors corresponding to diagrams (b1),(b2) should be subtracted [and similarly for diagrams (c1),(c2)]. For diagrams (d1),(d2) and (e1),(e2), this subtraction procedure leads to the total cancellation of their contributions. The relevant SU(6) factors are gathered in Table II, where, for completeness, the factors corresponding to the separate diagrams (b1), (b2), (c1), and (c2) are given. Contributions from the individual diagrams (d1), (d2), (e1), and (e2), though nonzero in general, are not shown. The entries in Table II correspond to the F/D ratio of SU(6), i.e., equal to $\frac{2}{3}$. Phenomenologically more successful fits are obtained in the pole models in which F/D differs slightly from its SU(6) value: $F/D \approx 0.56$ or 0.58. The explicit dependence on F/D of the ground-state baryon pole model formulas is given in Eq. (13). (See also Table III where weights of individual baryon pole contributions corresponding to the two [(1) and (2)] different orderings of the strong and weak transitions (see Fig. 2) are exhibited. Table III includes the effects of all the quark diagrams (b), (c), (d), and (e).)

In Eq. (13), f_0/d_0 characterizes the $\langle B'|H_{\rm weak}^{\rm PC}|B\rangle$ matrix elements, while all energy denominators $\pm 1/(N-Y)$ (we use $\Sigma - N = \Lambda - N = \Xi - \Sigma$) are contained in the overall normalization factor $C = -33$ (see also Ref. [3]):

$$\begin{aligned} B(\Sigma_0^+) &= \frac{1}{\sqrt{2}}\left[\frac{f_0}{d_0} - 1\right]\left[1 - \frac{F}{D}\right] C , \\ B(\Sigma_+^+) &= -\tfrac{4}{3} C , \\ B(\Sigma_-^-) &= \left[\left[\frac{f_0}{d_0} - 1\right]\frac{F}{D} - \frac{1}{3}\left[3\frac{f_0}{d_0} + 1\right]\right] C , \\ B(\Lambda_-^0) &= -\sqrt{2} B(\Lambda_0^0) \\ &= \frac{1}{\sqrt{6}}\left[\frac{f_0}{d_0} + 3 + \left[3\frac{f_0}{d_0} + 1\right]\frac{F}{D}\right] C , \\ B(\Xi_-^-) &= -\sqrt{2} B(\Xi_0^0) \\ &= -\frac{1}{\sqrt{6}}\left[3 - \frac{f_0}{d_0} + \left[3\frac{f_0}{d_0} - 1\right]\frac{F}{D}\right] C . \end{aligned} \quad (13)$$

The correspondence between the expressions resulting from the use of Table II through

$$B(\Sigma_0^+) = -\frac{1}{6\sqrt{2}}\beta - \frac{1}{9\sqrt{2}}\gamma \quad (14)$$

and Table III through Eq. (13) is given by taking, in Eq. (13), $F/D = \frac{2}{3}$ and identifying

$$\beta = 4C, \quad \gamma = -3\left[1 + \frac{f_0}{d_0}\right] C . \quad (15)$$

In Eq. (14), β and γ are the reduced matrix elements corresponding to diagrams (b1),(b2) and (c1),(c2), respective-

TABLE II. Weights of quark diagrams (b) and (c) for the P-wave (PC) amplitudes.

	Transition	(b1)	(b2)	(c1)	(c2)
Σ_0^+	$\Sigma^+\to p\pi^0$	$-\frac{1}{3\sqrt{2}}$	$-\frac{1}{6\sqrt{2}}$	$-\frac{1}{9\sqrt{2}}$	0
Σ_+^+	$\Sigma^+\to n\pi^+$	$-\frac{1}{3}$	0	0	0
	$\Sigma^+\to p\eta_8$	0	$-\frac{1}{2\sqrt{6}}$	$\frac{1}{9\sqrt{6}}$	$-\frac{2}{9\sqrt{6}}$
Σ_-^-	$\Sigma^-\to n\pi^-$	0	$\frac{1}{6}$	$\frac{1}{9}$	0
Λ_-^0	$\Lambda\to p\pi^-$	$-\frac{1}{3\sqrt{6}}$	$-\frac{1}{2\sqrt{6}}$	$-\frac{1}{\sqrt{6}}$	0
Λ_0^0	$\Lambda\to n\pi^0$	$\frac{1}{6\sqrt{3}}$	$\frac{1}{4\sqrt{3}}$	$\frac{1}{2\sqrt{3}}$	0
	$\Lambda\to n\eta_8$	0	$\frac{1}{4}$	$-\frac{1}{6}$	$\frac{1}{3}$
Ξ_-^-	$\Xi^-\to\Lambda\pi^-$	0	$\frac{1}{3\sqrt{6}}$	$\frac{1}{3\sqrt{6}}$	0
Ξ_0^0	$\Xi^0\to\Lambda\pi^0$	0	$-\frac{1}{6\sqrt{3}}$	$-\frac{1}{6\sqrt{3}}$	0
	$\Xi^0\to\Lambda\eta_8$	0	$-\frac{1}{6}$	$\frac{1}{18}$	$-\frac{1}{9}$
	$\Xi^-\to\Sigma^-\pi^0$	0	0	$-\frac{5}{9\sqrt{2}}$	0
	$\Xi^-\to\Sigma^-\eta_8$	0	0	$\frac{5}{9\sqrt{6}}$	$-\frac{10}{9\sqrt{6}}$
	$p\to K^0p$	0	$-\frac{1}{6}$	0	$-\frac{1}{9}$

ly. As is clearly seen from Eq. (15), in the ground-state baryon pole model the deviation of the experimentally observed f/d from its canonical value of -1 is attributed to a substantial contribution from diagrams (c), which modifies the f_0/d_0 structure of the $\langle B'|H_{\rm weak}^{\rm PC}|B\rangle$ matrix elements. Equation (13) describes the P-wave data very well (see Table IV). Note that one cannot expect here a better agreement in view of the violation of the $\Delta I=\frac{1}{2}$ rules by the data. For example, the $\Delta I=\frac{1}{2}$ rule $\sqrt{2}\Sigma_0^+=\Sigma_+^+-\Sigma_-^-$ experimentally reads $37.6\pm1.8=43.8\pm0.4$. The not-well-understood $\Delta I=\frac{3}{2}$ amplitudes are of the order of a few percent.

From Table IV we see that the data seem to require $(f/d)_{P\ \rm wave}\approx-1.85$ to -1.9. The ground-state baryon pole model identifies this f/d as the f_0/d_0 ratio characterizing the $\langle B'|H_{\rm weak}^{\rm PC}|B\rangle$ matrix elements. Although it is hard to make a fully reliable calculation of sea quark effects, the estimates of f_0/d_0 performed by Donoghue and Golowich [11] and by the author [14] lead to $f_0/d_0\approx-1.6$ or more. In Ref. [11] quark sea effects are due to short-distance QCD interactions, while in Ref. [14] hadron-level unitarity plays the dominant role in boosting the value of f_0/d_0 away from -1. The precise division of how much of this shift is due to short- and

TABLE III. Weights of baryon pole contributions for the P-wave amplitudes.

	Fig. 2(1)	Fig. 2(2)
Σ_0^+	$\frac{1}{\sqrt{2}}\left[1+\frac{F}{D}\right]\left[1-\frac{f}{d}\right]$	$\frac{1}{\sqrt{2}}\left[1+\frac{F}{D}\right]\left[1-\frac{f}{d}\right]-\frac{1}{\sqrt{2}}\left[1-\frac{F}{D}\right]\left[1-\frac{f}{d}\right]$
Σ_+^+	$\left[1+\frac{F}{D}\right]\left[1-\frac{f}{d}\right]$	$\left[1+\frac{F}{D}\right]\left[1-\frac{f}{d}\right]-\frac{4}{3}$
Σ_-^-	0	$\left[1-\frac{F}{D}\right]\left[1-\frac{f}{d}\right]-\frac{4}{3}$
Λ_-^0	$-\frac{1}{\sqrt{6}}\left[1+\frac{F}{D}\right]\left[1+3\frac{f}{d}\right]$	$\frac{2}{\sqrt{6}}\left[1-\frac{f}{d}\right]$
Ξ_-^-	$\frac{2}{\sqrt{6}}\left[1+\frac{f}{d}\right]$	$\frac{1}{\sqrt{6}}\left[1-\frac{F}{D}\right]\left[3\frac{f}{d}-1\right]$

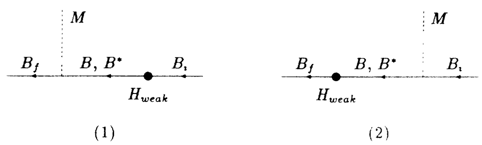

FIG. 2. Baryon pole diagrams for weak decays.

long-distance effects is not important here: The size of the quark sea contribution is determined by the total of these effects. It is this total that can be directly linked with the experimental value of the Δ-N splitting. In this way the large deviation of f_0/d_0 from -1 is correlated with the size of the Δ-N splitting.

In conclusion, it is natural to expect that the dominant part of the deviation of $(f/d)_{P\ \mathrm{wave}}$ from -1 is due to quark sea effects as identified in Eq. (15) and that f_0/d_0 is close to (say) -1.7. The remaining small enhancement of $(f/d)_{P\ \mathrm{wave}}$ may come from the meson-leg diagrams. For example, Stech and Xu [9] estimate the contribution to f arising from nonpenguin factorization diagrams to be around $f_{\mathrm{nonpenguin}}/d_0 \approx -0.15$ to -0.2.

IV. QUARK SEA EFFECTS IN *S* WAVES

To reconcile the value of the f/d ratio observed in the P-wave amplitudes with the one needed for a proper description of the S-wave amplitudes, we shall consider SU(3) symmetry breaking in the energy denominators of the latter. This effect was originally discussed by LeYaouanc *et al.* [5], who have shown how its inclusion works toward reducing the discrepancy in size between the CA estimate of the $S{:}P$ ratio and experiment. What LeYaouanc *et al.* did not consider was the presence of SU(3) symmetry breaking in denominators *in conjunction* with large quark sea effects. When SU(3) breaking is taken into account, the RHS's of Eqs. (9) are modified and one obtains

$$A_{(1)}(\Sigma^+\to p\pi^0) = \frac{\Delta\omega_s}{\Delta\omega_s-\delta s}\langle p|A_0|N^*\rangle\langle N^*|H^{\mathrm{PV}}_{\mathrm{weak}}|\Sigma^+\rangle ,$$
$$A_{(2)}(\Sigma^+\to p\pi^0) = -\frac{\Delta\omega_s}{\Delta\omega_s+\delta s}\langle p|H^{\mathrm{PV}}_{\mathrm{weak}}|\Sigma^*\rangle\langle \Sigma^*|A_0|\Sigma^+\rangle . \tag{16}$$

In Eqs. (16) we have put

$$\Delta\omega_{W_1} = N^*-\Sigma = \Delta\omega_s-\delta s ,$$
$$\Delta\omega_{W_2} = \Sigma^*-p = \Delta\omega_s+\delta s , \tag{17}$$

with $\Delta\omega_s \approx 570$ MeV being the average splitting between the $(56,0^+)\frac{1}{2}^+$ and $(70,1^-)\frac{1}{2}^-$ multiplets and $\delta s \approx 190$ MeV being the mass difference associated with a change of strangeness by -1. The sums over intermediate states N^*,Σ^* on the RHS's of Eqs. (16) are implicit in the weights of Table I. These weights are in turn proportional to the numerators of the pole model amplitudes. Using Table I, the sums in Eqs. (16) may be expressed therefore as

$$\langle p|A_0|N^*\rangle\langle N^*|H^{\mathrm{PV}}_{\mathrm{weak}}|\Sigma^+\rangle = -\frac{1}{6\sqrt{2}}kc_0 ,$$
$$-\langle p|H^{\mathrm{PV}}_{\mathrm{weak}}|\Sigma^*\rangle\langle\Sigma^*|A_0|\Sigma^+\rangle = +\frac{1}{2\sqrt{2}}kb_0 , \tag{18}$$

where k is some proportionality constant and b_0, c_0 are the SU(3)-invariant couplings characterizing the $\langle B|H^{\mathrm{PC}}_{\mathrm{weak}}|B'\rangle$ matrix elements. Indeed, if there is no SU(3) breaking in energy denominators, from Eqs. (16) and (18) we obtain the $A(\Sigma_0^+)$ amplitude (which is proportional to the $\langle p|H^{\mathrm{PC}}_{\mathrm{weak}}|\Sigma^+\rangle$ matrix element) of Eq. (2) with

$$b = kb_0, \quad c = kc_0 . \tag{19}$$

Thus, in the limit of exact SU(3), the f/d ratio for the S-wave amplitudes $[(f/d)_{S\ \mathrm{wave}} = -1+2c/3b]$ is the same as the f_0/d_0 ratio for the $\langle B|H^{\mathrm{PC}}_{\mathrm{weak}}|B'\rangle$ matrix elements $(f_0/d_0 = -1+2c_0/3b_0)$.

When $\delta s \neq 0$ we obtain from Eq. (16)

$$A(\Sigma_0^+) = \frac{1}{2\sqrt{2}}\frac{1}{1+x}kb_0 - \frac{1}{6\sqrt{2}}\frac{1}{1-x}kc_0 , \tag{20}$$

where $x = \delta s/\Delta\omega_s$. All the other *pion*-emission amplitudes of Eq. (2) are modified in the same way as $A(\Sigma_0^+)$, i.e.,

$$kb_0 \to = kb_0/(1+x) ,$$
$$kc_0 \to = kc_0/(1-x) . \tag{21}$$

The above simple prescription does not apply to the *nonpion*-emission amplitudes, which are, however, kinematically forbidden. Using $(f/d)_{S\ \mathrm{wave}} = -1+2c/3b$ and Eq. (21), one immediately obtains Eq. (1).

Inclusion of large quark sea effects explains the

TABLE IV. P-wave amplitudes (in units of 10^{-7}) from Eq. (13).

	$F/D=0.56,\ C=-33$				
	Baryon legs			Meson leg	
Process	$f_0/d_0=-1.7$	-1.85	-1.9	$f_0'/d_0=-0.15$	Data
Σ_0^+	27.7	29.3	29.8	1.6	$26.6\ \pm 1.3$
Σ_+^+	44.0	44.0	44.0	0	$42.4\ \pm 0.35$
Σ_-^-	4.8	2.6	1.9	-2.2	-1.44 ± 0.17
Λ_-^0	13.4	18.8	20.6	5.4	$22.1\ \pm 0.5$
Ξ_-^-	17.3	15.9	15.5	-1.4	$16.6\ \pm 0.8$

TABLE V. S-wave amplitudes (in units of 10^{-7}) as calculated from P-wave amplitudes.

	$x=0$		$x=1/3$		
f_0/d_0	-1	-1.7	-1	-1.7	Data
Σ_0^+	-5.4	-7.3	-4.1	-6.8	-3.27
Σ_+^+	0	0	0	0	0.13
Σ_-^-	7.7	10.3	5.8	9.7	4.27
Λ_-^0	3.1	6.4	2.3	7.1	3.23
Ξ_-^-	-6.3	-9.5	-4.7	-9.5	-4.50

difference in the size of the (f/d) ratios of S- and P-wave amplitudes in a very natural way. At the same time, however, the S-wave reduction mechanism proposed by LeYaouanc *et al.* [5] to bring the $S:P$ ratio into agreement with experiment becomes essentially unimportant. Reference [5] corresponds to $c_0=0$ and leads to a reduction of the S-wave amplitudes by 25–30% for $x=0.3$–0.4 [see Eq. (21) and Table V]. With $f_0/d_0\approx-1.7$ ($c_0\approx-b_0$), this reduction is, however, negligible. We shall discuss this reappearing question of the $S:P$ ratio in the next section. Below, for completeness and possible future use, we rewrite the $B_i\to B_fP$ parity-violating $\Delta S=1$ amplitudes in an explicit SU(3) language.

The relevant amplitudes are given by, for the (b1) diagrams,

$$\frac{1}{2}\mathrm{Tr}(SP^\dagger B_f^\dagger B_i)\frac{kb_0}{1-x}, \tag{22}$$

for the (b2) diagrams,

$$-\frac{1}{2}\mathrm{Tr}(P^\dagger SB_f^\dagger B_i)\frac{kb_0}{1+x}, \tag{23}$$

for the (c1) diagrams,

$$\frac{1}{6}[\mathrm{Tr}(P^\dagger S[B_f^\dagger,B_i])-\mathrm{Tr}(P^*S)\mathrm{Tr}(B_f^\dagger B_i)]\frac{kc_0}{1-x}, \tag{24}$$

and for the (c2) diagrams,

$$-\frac{1}{6}[\mathrm{Tr}(SP^\dagger[B_f^\dagger,B_i])-\mathrm{Tr}(P^\dagger S)\mathrm{Tr}(B_f^\dagger B_i)]\frac{kc_0}{1+x}. \tag{25}$$

In Eqs. (22)–(25),

$$S=\lambda_6=\begin{bmatrix}0&0&0\\0&0&1\\0&1&0\end{bmatrix} \tag{26}$$

is the spurion representing the weak Hamiltonian and B_i,B_f,P are the standard 3×3 matrices corresponding to the hadrons in question.

For the pions ($P=P_\pi$), only the $\mathrm{Tr}(P^\dagger SB_f^\dagger B_i)$ and $\mathrm{Tr}(P^\dagger SB_iB_f^\dagger)$ traces in Eqs. (22)–(25) are nonzero. Consequently, the pion-emission amplitudes are

$$\begin{aligned}A(B_i\to B_f\pi)&=-\frac{1}{2}\mathrm{Tr}(P_\pi^\dagger SB_f^\dagger B_i)\frac{kb_0}{1+x}\\&\quad+\frac{1}{6}\mathrm{Tr}(P_\pi^\dagger S[B_f^\dagger,B_i])\frac{kc_0}{1-x}\\&=d\,\mathrm{Tr}(P_\pi^\dagger S\{B_i,B_f^\dagger\})+f\,\mathrm{Tr}(P_\pi^\dagger S[B_i,B_f^\dagger]),\end{aligned} \tag{27}$$

with

$$d=-\frac{1}{4}\frac{kb_0}{1+x},\quad f=\frac{1}{4}\frac{kb_0}{1+x}-\frac{1}{6}\frac{kc_0}{1-x}, \tag{28}$$

and the apparent (i.e., applicable to pions amplitudes only) $(f/d)_{S\ \mathrm{wave}}$ ratio is given by

$$(f/d)_{S\ \mathrm{wave}}=-1+\frac{2}{3}\frac{c_0}{b_0}\frac{1+x}{1-x}. \tag{29}$$

V. PROBLEM OF THE $S:P$ RATIO

Large quark sea effects constitute an attractive explanation of the deviation of f/d from -1 because (1) their large size is consistent with unitarity-based calculations with the scale provided by Δ-N splitting [14], and (2) they explain in a nice way the difference in the apparent f/d ratios of the S and P waves.

However, when sea effects are large, the S-wave reduction mechanism induced by the SU(3)-breaking effects ceases to be significant and the problem of the $S:P$ ratio reappears. A possible way to deal with the latter has been discussed by Milosević, Tadić, and Trampetić [18], Bonvin [7], and Nardulli [8]. These authors considered the radially excited $(56,0^+)^*\frac{1}{2}^+$ baryons B^* in the intermediate states of the P-wave amplitudes and found that their contribution has the same sign and order of magnitude as the contribution from the ground-state baryons. The details of the decomposition of the P-wave amplitudes into various contributions differed in these papers substantially even though the f_0^*/d_0^* ratio for the radially excited baryons was assumed equal to -1 in all these papers. Since in this paper we argue that for ground-state baryons f_0/d_0 deviates from -1 significantly, it is natural to expect the same of f_0^*/d_0^*. In fact, it is natural to expect that $f_0^*/d_0^*=f_0/d_0$: The relative size of the contributions from the W exchange and sea pieces of the weak Hamiltonian in $\langle B|H_{\mathrm{weak}}^{\mathrm{PC}}|B'\rangle$ should be independent of whether the external state $|B'\rangle$ is a ground-state or radially excited baryon. The contributions from the radially excited $\frac{1}{2}^+$ baryons can be read off from the weights of Table III. Assuming that radial excitations are heavier than the ground states by $\Delta\omega^* \approx 450$ MeV, the weights corresponding to diagrams (1) and (2) of Fig. 2 have to be added leading to

$$\begin{aligned}
B(\Sigma_0^+) &= \frac{1}{\sqrt{2}}\left\{\left[2\left(1+\frac{F^*}{D^*}\right)\right.\right.\\
&\quad\left.\left.-\left(1-\frac{F^*}{D^*}\right)\right]\left(1-\frac{f_0^*}{d_0^*}\right)\right\}G\,,\\
B(\Sigma_+^+) &= \left\{2\left(1+\frac{F^*}{D^*}\right)\left(1-\frac{f_0^*}{d_0^*}\right)-\frac{4}{3}\right\}G\,,\\
B(\Sigma_-^-) &= \left\{\left(1-\frac{F^*}{D^*}\right)\left(1-\frac{f_0^*}{d_0^*}\right)-\frac{4}{3}\right\}G\,,\\
B(\Lambda_-^0) &= \frac{1}{\sqrt{6}}\left\{-\left(1+\frac{F^*}{D^*}\right)\left(3\frac{f_0^*}{d_0^*}+1\right)\right.\\
&\quad\left.+2\left(1-\frac{f_0^*}{d_0^*}\right)\right\}G\,,\\
B(\Xi_-^-) &= \frac{1}{\sqrt{6}}\left\{2\left(1+\frac{f_0^*}{d_0^*}\right)\right.\\
&\quad\left.+\left(1-\frac{F^*}{D^*}\right)\left(3\frac{f_0^*}{d_0^*}-1\right)\right\}G\,,
\end{aligned} \tag{30}$$

with

$$G = C\frac{\delta s}{\Delta\omega^*}\frac{g^*}{g}\,. \tag{31}$$

In Eq. (30), F^*/D^* (=0.56) is the F/D ratio for the B^*BP couplings and g^*/g describes the relative size and sign of the B^*BP and BBP couplings. (The ratio g^*/g may be considered as including the relative size of d_0^*/d_0, which in Ref. [7] was found to be close to 1, however.) In the quark model, the ratio g^*/g is calculable and turns out to be *negative* and small [see, e.g., Eq. (10) of Ref. [7]]:

$$\frac{g^*}{g} \approx -0.1 \text{ to } -0.2\,. \tag{32}$$

In writing Eq. (30), we have neglected SU(3) breaking in energy denominators. Inclusion of this effect generates an additional contribution whose symmetry structure is identical to that of the intermediate ground-state baryons. It adds constructively to the latter one, though with a small relative size of $-(\delta s/\Delta\omega^*)^2 g^*/g \le 3\%$ only.

The contribution of the radially excited states [Eq. (30)] violates the Lee-Sugawara (LS) sum rule [19]

$$2\Xi_-^- + \Lambda_-^0 = \sqrt{3}\Sigma_0^+ \tag{33}$$

for the P waves. With the inclusion of radially excited states, Eq. (33) reads

$$\begin{aligned}
&\frac{1}{\sqrt{6}}\left\{3\left(1-\frac{f_0}{d_0}\right)\left(\frac{F}{D}-1\right)+\frac{\delta s}{\Delta\omega^*}\frac{g^*}{g}3\left(1+3\frac{F^*}{D^*}\right)\left(1-\frac{f_0^*}{d_0^*}\right)\right\}C+\frac{1}{\sqrt{6}}\frac{\delta s}{\Delta\omega^*}\frac{g^*}{g}\left[8\left(\frac{f_0^*}{d_0^*}-\frac{F^*}{D^*}\right)\right]C\\
&\qquad=\frac{1}{\sqrt{6}}\left\{3\left(1-\frac{f_0}{d_0}\right)\left(\frac{F}{D}-1\right)+\frac{\delta s}{\Delta\omega^*}\frac{g^*}{g}3\left(1+3\frac{F^*}{D^*}\right)\left(1-\frac{f_0^*}{d_0^*}\right)\right\}C\,.
\end{aligned} \tag{34}$$

Using experimental numbers, Eq. (33) reads

$$55.3 = 46.1\,. \tag{35}$$

The negative sign of g^*/g leads to the violation of the LS rule in the direction opposite to the experimental one. This violation comes about as follows. For (all) Σ and Λ decays, the contribution of radially excited states has the same sign as that of ground states and thus seems to help in the explanation of the $S{:}P$ ratio. However, for Ξ decays this relative sign is negative. If $f_0^*/d_0^* = -1$ is used as in Refs. [18,7,8], the size of the contribution of radially excited states to Ξ decays is small. For $f_0^*/d_0^* = -1.7$, however, this contribution is larger by a factor of 2.5 [see Eq. (30)] and it reduces the Ξ_-^- amplitudes [and the LHS of Eq. (33)] very strongly. Consequently, only a small contribution (characterized by $g^*/g \le -0.05$) of the radially excited states can be tolerated if $f_0^*/d_0^* \approx -1.7$. Inspection of Eqs. (13) and (30) shows then that radially excited states may increase the Λ,Σ amplitudes by $\approx 15\%$ only. Thus, if $f_0^*/d_0^* \approx -1.7$, the radially excited states cannot be held responsible for the experimentally observed large size of P-wave amplitudes (or small size of S-wave amplitudes).

In search for an explanation of the experimentally observed suppression of the $S{:}P$ ratio, let us note that in the preceding sections we have pointed at *SU(3) symmetry breaking* as the possible origin of different deviations of apparent f/d from -1. Thus it was essentially proposed that the quark model as used in the PCAC-CA approach has too much built-in symmetry. Similarly, the relative size of various hadron couplings does not have to follow the quark model predictions closely. For example, it is well known that the $\Delta \to N$ magnetic transition is misjudged in the quark model by 30% if the magnetic moment of the proton is used to set the scale of quark-level couplings. Now Δ and N are still members of the same $(56,0^+)\frac{1}{2}^+$ SU(6)×O(3) multiplet. It is therefore conceivable that similar or larger deviations from quark model predictions may appear when one attempts to estimate the $B(\frac{1}{2}^+)B^*(\frac{1}{2}^-)P$ couplings from the knowledge of familiar couplings of ground-state baryons to pseudoscalar mesons. After all, we are dealing now with two different SU(6)×O(3) multiplets: $(56,0^+)$ and $(70,1^-)$. A 30% reduction in the overall size of the $g_{B^*(1/2^-)BP}$ and

$\langle B|H^{\mathrm{PV}}_{\mathrm{weak}}|B^*(\frac{1}{2}^-)\rangle$ couplings with respect to those calculated from $g_{BB'P}$ and $\langle B|H^{\mathrm{PC}}_{\mathrm{weak}}|B'\rangle$ by the quark model route is a totally plausible possibility. It would provide the missing factor of 2 by reducing k in Eq. (18) *without affecting* the relationship of Eq. (1) between the f/d ratios of S- and P-wave amplitudes. Clearly, the above argument constitutes a suggestion only. It would require a thorough investigation, which, for obvious reasons, is beyond the scope of this paper: At the moment we do not know how to modify the quark model to improve its predictions for couplings.

VI. WEAK RADIATIVE HYPERON DECAYS

As already discussed in the preceding sections, in the literature on weak nonleptonic hyperon decays there is no consensus on the origin of (1) the suppression of the $S:P$ ratio and (2) the deviations of f/d from -1. The general theoretical framework is not disputed, however. This is not the case for weak radiative hyperon decays, which, for the last 25 years, have constituted a real puzzle that has even been termed "the last low-q^2 frontier of weak interaction physics." For a thorough presentation of this highly controversial topic, see the upcoming review in Ref. [20]. At present, there is only one approach that seems capable of describing fairly well the existing experimental data on asymmetries and branching ratios of these decays. This approach, developed recently by the author [3,4], is based on a combination of the arguments of symmetry with the idea of vector meson dominance (VDM) [21]. Although joint consideration of weak interactions, symmetry, and vector meson dominance looks innocent, it is possible that it is intricately linked with very deep issues (see Ref. [20]). Now the SU(3)-symmetry-breaking effects discussed in the present paper have not been considered within that approach as yet. Therefore it is of great importance to see how the results of Ref. [4] might be changed if SU(3) symmetry breaking in energy denominators is taken into account.

Calculation of the relevant weights is straightforward and leads to Table VI. In this table only the weights corresponding to diagrams (b1) and (b2) have been given. Contributions from diagrams (c1) and (c2) add up to the

TABLE VI. Weights of quark diagrams (b) for the radiative S-wave amplitudes.

Transition	(b1)	(b2)
$\Sigma^+\to p\gamma$	$-\frac{1}{3\sqrt{2}}$	$-\frac{2+\epsilon}{9\sqrt{2}}$
$\Sigma^0\to n\gamma$	$-\frac{1}{6}$	$\frac{2+\epsilon}{18}$
$\Lambda\to n\gamma$	$\frac{1}{6\sqrt{3}}$	$\frac{2+\epsilon}{6\sqrt{3}}$
$\Xi^0\to\Lambda\gamma$	0	$-\frac{2+\epsilon}{9\sqrt{3}}$
$\Xi^0\to\Sigma^0\gamma$	$\frac{1}{3}$	0
$\Xi^-\to\Sigma^-\gamma$	0	0

same general SU(3) structure irrespectively of whether or not SU(3) is broken in energy denominators. This general structure has been treated in Ref. [4] with the help of a parameter [d' in Eq. (30) below]. The SU(3)-symmetry-breaking effects of the type considered in this paper do affect the size of this parameter. However, they do not affect the *relative* sizes of the single-quark contributions to various radiative decays. Since in the VDM$\times$symmetry approach of Ref. [4] d' is treated as a parameter to be fitted, SU(3) symmetry breaking in energy denominators is phenomenologically discernible in the contributions from the (b)-type processes only.

The parity-violating amplitudes due to (b)-type diagrams can be read off from Table V, and together with the single-quark contributions, they give [up to an overall VDM factor of e/g ($e^2/4\pi=\frac{1}{137}$, $g=5.0$)]

$$\begin{aligned}
A(\Sigma^+\to p\gamma)&=-\frac{b}{9\sqrt{2}}\left\{2+\epsilon+3\frac{1+x}{1-x}\right\}+\frac{1}{\sqrt{2}}d'\,,\\
A(\Sigma^0\to n\gamma)&=-\frac{b}{18}\left\{3\frac{1+x}{1-x}-2-\epsilon\right\}-\frac{1}{2}d'\,,\\
A(\Lambda\to n\gamma)&=\frac{b}{6\sqrt{3}}\left\{2+\epsilon+\frac{1+x}{1-x}\right\}-\frac{3\sqrt{3}}{2}d'\,,\\
A(\Xi^0\to\Lambda\gamma)&=-\frac{2+\epsilon}{9\sqrt{3}}b+\frac{\sqrt{3}}{2}d'\,,\\
A(\Xi^0\to\Sigma^0\gamma)&=-\frac{1}{3}b\frac{1+x}{1-x}-\frac{5}{2}d'\,,\\
A(\Xi^-\to\Sigma^-\gamma)&=\frac{5}{\sqrt{2}}d'\,,
\end{aligned}\tag{36}$$

with $b=kb_0/(1+x)=-5$ (in units of 10^{-7}) and small negative d'.

From Eq. (36) it follows that, when compared to the SU(3) symmetric case ($x=0$), the SU(3) symmetry breaking in energy denominators (1) increases the parity-violating amplitudes in the $\Sigma^+\to p\gamma$, $\Sigma^0\to n\gamma$, $\Lambda\to n\gamma$, and $\Xi^0\to\Sigma^0\gamma$ decays, and (2) leaves the $\Xi^0\to\Lambda\gamma$ parity-violating amplitude unchanged.

No change of sign of the (b)-type two-quark contribution is observed. Since the contribution of the single-quark parity-violating amplitudes [terms proportional to d' in Eq. (36)] is strongly limited from above by the recently measured branching ratio of the $\Xi^-\to\Sigma^-\gamma$ decay [22], we conclude that the basic expectations of the SU(3)-symmetric approach of Ref. [4] (such as signs and approximate size of asymmetries) cannot change much when the effect of SU(3) symmetry breaking in the denominators is included. However, the slight increase in the value of $A(\Lambda\to n\gamma)$ would make it easier to fit the observed $\Lambda\to n\gamma$ branching ratio [23]. At the same time, the increase of $A(\Xi^0\to\Sigma^0\gamma)$ would manifest itself mostly in a more negative asymmetry of the $\Xi^0\to\Sigma^0\gamma$ decay. The only experiment performed so far [24] yields a slightly positive (albeit with a large error) value for this asymmetry ($+0.2\pm0.32$). The calculations of this paper confirm therefore that it is very important to measure the $\Xi^0\to\Sigma^0\gamma$ asymmetry precisely. Should this asymmetry stay significantly positive, it would add yet another question mark to the long-standing enigma of weak radiative hyperon decays.

VII. SUMMARY

In this paper we have carried out an analysis of the joint influence of large quark sea and SU(3)-symmetry-breaking effects in weak hyperon decays. An explanation of the difference between the values of the apparent f/d ratios for the S- and P-wave amplitudes of nonleptonic decays has been proposed. It was pointed out that quark sea effects in the matrix elements of the parity-conserving part of the weak Hamiltonian between the ground-state baryons are additionally enhanced in the S-wave amplitudes by the presence of the SU(3)-symmetry-breaking effects in energy denominators. A formula for this enhancement has been derived and shown to agree with the data extremely well if the dominant part of the derivation of $(f/d)_{P\ \mathrm{wave}}$ from -1 is due to sea quarks. This corroborates our earlier calculations which indicated that large deviations of $(f/d)_{\mathrm{soft\ pion}}$ from its naive quark model value of -1 are to be expected when the quark model is properly unitarized. Thus the commonly used quark model value of $(f/d)_{\mathrm{soft\ pion}}=-1$ should be replaced by a value close to -1.7. We suggest that a possible way to resolve the $S:P$ problem is to break the naive quark model predictions relating the values of matrix elements involving the $(56,0^+)$ and $(70,1^-)$ baryons. In view of unsolved difficulties existing elsewhere in similar problems involving baryon couplings, this possible route of explaining the $S:P$ size problem cannot be properly handled at the moment: We do not know how to modify the (oversimplified) naive quark model predictions for couplings. Finally, implications of this paper for the weak radiative hyperon decays have been briefly discussed. It was shown that the signs of the asymmetries previously calculated in the SU(3)-symmetric approach are unchanged by the inclusion of the SU(3)-symmetry-breaking effects in energy denominators.

ACKNOWLEDGMENTS

This research has been supported in part by Polish Committee for Scientific Research Grant No. 2 0198 91 01.

[1] J. F. Donoghue, F. Golowich, and B. Holstein, Phys. Rep. **131**, 319 (1986).
[2] M. Gell-Mann, Physics **1**, 63 (1964); S. Weinberg, Phys. Rev. Lett. **16**, 879 (1966); R. Dashen, Phys. Rev. **183**, 1245 (1969).
[3] P. Żenczykowski, Phys. Rev. D **40**, 2290 (1989).
[4] P. Żenczykowski, Phys. Rev. D **44**, 1485 (1991).
[5] A. LeYaouanc, O. Péne, J.-C. Raynal, and L. Oliver, Nucl. Phys. **B149**, 321 (1979).
[6] M. Gronau, Phys. Rev. D **5**, 118 (1972).
[7] M. Bonvin, Nucl. Phys. **B238**, 241 (1984).
[8] G. Nardulli, Nuovo Cimento A **100**, 485 (1988).
[9] B. Stech and Q. P. Xu, Z. Phys. C **49**, 491 (1991).
[10] B. Desplanques, J. F. Donoghue, and B. R. Holstein, Ann. Phys. (N.Y.) **124**, 449 (1980).
[11] J. F. Donoghue and E. Golowich, Phys. Lett. **69B**, 437 (1977).
[12] J. F. Donoghue, E. Golowich, W. A. Ponce, and B. Holstein, Phys. Rev. D **21**, 186 (1980).
[13] J. O. Eeg, Phys. Lett. **155B**, 115 (1985).
[14] P. Żenczykowski, Phys. Rev. D **45**, 3153 (1992).
[15] N. A. Törnqvist, Acta Phys. Pol. B **16**, 503 (1985); N. A. Törnqvist, in *The Hadron Mass Spectrum,* Proceedings of the Rheinfels Workshop, St. Goar, Germany, 1990, edited by E. Klempt and K. Peters [Nucl. Phys. B (Proc. Suppl.) **21** (1991)].
[16] P. Żenczykowski, Ann. Phys. (N.Y.) **169**, 453 (1986).
[17] B. W. Lee and A. R. Swift, Phys. Rev. **136**, B228 (1964).
[18] M. Milosević, D. Tadić, and J. Trampetić, Nucl. Phys. **B207**, 461 (1982).
[19] B. W. Lee, Phys. Rev. Lett. **12**, 83 (1964); H. Sugawara, Prog. Theor. Phys. (Kyoto) **31**, 213 (1964).
[20] J. Lach and P. Żenczykowski (unpublished).
[21] J. J. Sakurai, *Currents and Mesons* (University of Chicago Press, Chicago, 1969), pp. 65–71.
[22] T. Dubbs *et al.*, Phys. Rev. Lett. **72**, 808 (1994).
[23] A. J. Noble *et al.*, Phys. Rev. Lett. **69**, 414 (1992); K. D. Larson *et al.*, Phys. Rev. D **47**, 799 (1993).
[24] S. Teige *et al.*, Phys. Rev. Lett. **63**, 2717 (1989).