Hyperon and Ω^- nonleptonic weak decays

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Nonleptonic hyperon and Ω^- decays are analyzed in the framework of the Weinberg-Salam model of weak interactions. Various effective weak nonleptonic Hamiltonians are tested with and without quantum-chromodynamic corrections, including extensive variations of all the parameters in question. The influence of quark models is also investigated using both the MIT bag model and the harmonic-oscillator model. Various dynamical assumptions connected with particular selections of current-commutator and pole terms are also considered. Irrespective of all variations and uncertainties, it seems that orders of magnitude and relative signs for various amplitudes are always correctly reproduced.

I. INTRODUCTION

Recently, numerous efforts[1-4] have been made claiming to lead to a reasonable estimate of hyperon and Ω^- nonleptonic decays in the Weinberg Salam model. An important ingredient in those attempts was the effective weak Hamiltonian renormalized by quantum chromodynamics (QCD). SU(4)-flavor-symmetry breaking was included in the calculation of QCD corrections, as suggested by Ref. 1. This has the effect of increasing the $\Delta I=\frac{1}{2}$ content of the weak Hamiltonian.

S-wave hyperon-decay amplitudes (A) were approximated through a commutator term, while for p-wave decay amplitudes (B), ground-state pole terms were used. Both amplitudes receive either separable contributions[2] or meson-pole terms.[3] The same type of approximation was used in calculating Ω^- decay amplitudes.[4] In all these cases, baryon matrix elements of the effective weak Hamiltonian were determined using the MIT bag model.[5]

The aim of this paper is to test some of these approaches and to investigate the stability of theoretical results. The calculational scheme used is actually the most commonly adopted "textbook" approach to the problem.[6] We intend, therefore, to keep it as a basis and consider only some variations within this general scheme. This means selecting either separable terms or meson-pole terms and varying their parameters within the boundaries given by experimental[7] and/or theoretical uncertainties. It turns out that theoretical results are not too sensitive to such variations.

Quark models for baryons are still relatively simple semiempirical devices. It would not be convincing if the theoretical results were crucially dependent on the particular model. It is more satisfactory if it turns out that the decisive role is played by SU(6) symmetry underlying any quark model of hadrons and by the consequent Pati-Woo theorem.[8] Such a conclusion seems to follow from the comparison between the results[2-4] based on the MIT bag model[5] with those obtained using an entirely nonrelativistic harmonic-oscillator (HO) quark model.[9,10]

The form of the effective weak Hamiltonian can be varied changing the enhancement factors c_i appearing in front of its operators (see Sec. II). When certain variations were tried, including also QCD corrections, the straightforward theoretical c_i values[1] seemed to lead to satisfactory results. It might be useful to know that there is a relatively stable theoretical scheme, capable of producing reasonable results,[11] as it will be demonstrated below.

II. THE EFFECTIVE WEAK HAMILTONIAN

The effective weak Hamiltonian which includes QCD corrections and SU(4)-flavor-symmetry breaking is of the form

$$H_w^{\text{eff}}(\Delta S=1)=\sqrt{2}\,G_F\sin\theta_C\cos\theta_C\sum_{i=1}^{6}c_iO_i\,. \tag{2.1}$$

It contains the following four-quark operators:

$$\begin{aligned}
O_1&=:(\bar d_L s_L)(\bar u_L u_L)-(\bar d_L u_L)(\bar u_L s_L):\ \ (\underline{20}'',\underline{8},\Delta I=\tfrac{1}{2}),\\
O_2&=:(\bar d_L s_L)(\bar u_L u_L)+(\bar d_L u_L)(\bar u_L s_L)+2(\bar d_L s_L)(\bar u_L u_L)\\
&\qquad+2(\bar d_L s_L)(\bar s_L s_L):\ \ (\underline{84},\underline{8},\Delta I=\tfrac{1}{2}),\\
O_3&=:(\bar d_L s_L)(\bar u_L u_L)+(\bar d_L u_L)(\bar u_L s_L)+2(\bar d_L s_L)(\bar d_L d_L)\\
&\qquad-3(\bar d_L s_L)(\bar s_L s_L):\ \ (\underline{84},\underline{27},\Delta I=\tfrac{1}{2}),\\
O_4&=:(\bar d_L s_L)(\bar u_L u_L)+(\bar d_L u_L)(\bar u_L s_L)\\
&\qquad-(\bar d_L s_L)(\bar d_L d_L):\ \ (\underline{84},\underline{27},\Delta I=\tfrac{3}{2}),\\
O_5&=:(\bar d_L\lambda_\alpha s_L)(\bar q_R^i\lambda^\alpha q_R^i):\ \ (\underline{15},\underline{8},\Delta I=\tfrac{1}{2}),\\
O_6&=:(\bar d_L s_L)(\bar q_R^i q_R^i):\ \ (\underline{15},\underline{8},\Delta I=\tfrac{1}{2}).
\end{aligned} \tag{2.2}$$

We indicate the SU(4) and the SU(3) flavor content of the particular operator. Otherwise the notations mean

$$\bar{q}_L q_L = \tfrac{1}{2}\bar{q}\gamma_\mu(1-\gamma_5)q \,. \qquad (2.3)$$

The coefficients appearing in formula (2.1) were derived in Ref. 1. They are functions of the charmed-quark mass m_c [which breaks SU(4) flavor symmetry], the renormalization mass μ, the intermediate-vector-boson mass m_W, and the running coupling constant $g^2(\mu^2)/4\pi$. Choosing as an example the values

$$m_W = 100 \text{ GeV}, \quad m_c = 2 \text{ GeV}, \quad \mu = 0.5 \text{ GeV}, \quad \frac{g^2(\mu^2)}{4\pi} = 1 , \qquad (2.4)$$

one obtains

$$c_1 = -2.358, \quad c_2 = 0.080, \quad c_3 = 0.082, \quad c_4 = 0.411, \quad c_5 = -0.080, \quad c_6 = -0.021 . \qquad (2.5)$$

The operators O_5 and O_6 appear because of the SU(4)-flavor-symmetry breaking.[1] All other four-quark operators appear already in the QCD-unrenormalized weak Hamiltonian. Such a Hamiltonian is recovered for the following values of the coefficients:

$$c_1 = -1, \; c_2 = \tfrac{1}{5}, \; c_3 = \tfrac{2}{15}, \; c_4 = \tfrac{2}{3}, \; c_5 = c_6 = 0 . \qquad (2.6)$$

The semiempirical coefficients c_+, c_- of Ref. 12 can be introduced by the substitution

$$c_1 = -c_-, \quad c_2 = \tfrac{1}{5}c_+, \quad c_3 = \tfrac{2}{15}c_+, \quad c_4 = \tfrac{2}{3}c_+ . \qquad (2.7)$$

By making this choice (2.7), we have suppressed SU(4)-symmetry breaking. In this way, we were able to establish a straightforward connection between the $\Delta s = 1$ Hamiltonian (2.1) and the $\Delta s = \Delta c = 1$ Hamiltonian of Ref. 12. With SU(4)-symmetry breaking included, operators such as O_5, O_6 appear only in the $\Delta s = 1$ case, where there are closed $\bar{c}c$ quark loops.[13]

III. DYNAMICS: HYPERON DECAYS

In our calculation we do not consider closed hadronic loops, i.e., those containing baryons or measons. Such an approximation is fairly standard[6] and/or has been discussed in detail elsewhere.[1-4]

Parity-violating hyperon-nonleptonic-decay amplitudes A receive contributions A^c from current-algebra commutator terms. For example,[14]

$$A^c\begin{pmatrix}\Xi_-^-\\ \Sigma_0^+\end{pmatrix} = \frac{-1}{f_\pi}\tilde{G}_F\begin{pmatrix} a_{\Xi^0\Lambda} \\ \frac{1}{\sqrt{2}}a_{\Sigma^+p}\end{pmatrix}, \quad \tilde{G}_F = \frac{G_F}{2\sqrt{2}}\cos\theta_C\sin\theta_C ,$$

$$a_{\Xi^0\Lambda} = (\tfrac{2}{3})^{1/2}\{-12c_1(a+b) + (c_6 - \tfrac{8}{3}c_5)[3(a-a') - (9b-b')]\} , \qquad (3.1)$$

$$a_{\Sigma^+p} = \tfrac{2}{3}[-18c_1(a+b) + (c_6 - \tfrac{8}{3}c_5)(3a - 13b)] .$$

Both A and B amplitudes also receive contributions from separable terms in which two quark fields are sandwiched between baryon states, while the other two quarks are responsible for pion emission:

$$A^s\begin{pmatrix}\Xi_-^-\\ \Sigma_0^+\end{pmatrix} = (\tfrac{2}{3})^{1/2}\tilde{G}_F f_\pi\begin{pmatrix}(m_\Xi - m_\Lambda)a^s_{\Xi^-\Lambda} \\ \frac{-1}{\sqrt{3}}(m_\Sigma - m_p)a^s_{\Sigma^+p}\end{pmatrix} ,$$

$$a^s_{\Xi^-\Lambda} = \left[c_1 - 2(c_2+c_3+c_4) - (c_6 + \tfrac{16}{3}c_5)\frac{m_\pi^2}{(m_s - m_u)(m_d + m_u)}\right] , \qquad (3.2)$$

$$a^s_{\Sigma^+p} = \left[-c_1 + 2(c_2+c_3-2c_4) + (c_6 + \tfrac{16}{3}c_5)\frac{m_\pi^2}{(m_s - m_d)(2m_d)}\right] ,$$

$$B^s\begin{pmatrix}\Xi_-^-\\ \Sigma_0^+\end{pmatrix} = \frac{\sqrt{2}}{3}\tilde{G}_F f_\pi\begin{pmatrix}\frac{1}{\sqrt{3}}(\tilde{D} - 3\tilde{F})(m_\Xi + m_\Lambda)b^s_{\Xi^-\Lambda} \\ (\tilde{D} - \tilde{F})(m_\Sigma + m_p)b^s_{\Sigma^+p}\end{pmatrix} ,$$

$$b^s_{\Xi^-\Lambda} = \left[-c_1 + 2(c_2+c_3+c_4) - (c_6 + \tfrac{16}{3}c_5)\frac{m_\pi^2}{(m_s + m_u)(m_d + m_u)}\right] , \qquad (3.3)$$

$$b^s_{\Sigma^+p} = \left[c_1 - 2(c_2+c_3-2c_4) + (c_6 + \tfrac{16}{3}c_5)\frac{m_\pi^2}{(m_s + m_d)(2m_d)}\right] .$$

Here we have used the current quark masses[15]

$$m_u = 4.2 \text{ MeV}, \quad m_d = 7.5 \text{ MeV}, \text{ and } m_s = 150 \text{ MeV}.$$

Alternatively, B^s terms can be replaced[3] by K-meson pole terms B^K as follows:

$$B^K\begin{pmatrix}\Xi_-^-\\ \Sigma_0^+\end{pmatrix} = 4g\tilde{G}_F\frac{2E_\pi}{m_K^2 - m_\pi^2}\begin{pmatrix}\frac{-1}{\sqrt{3}}(d-3f)a_{K^-\pi^-} \\ (d-f)a_{K^0\pi^0}\end{pmatrix} ,$$

$$a_{K^-\pi^-} = \{[-c_1 + 2(c_2+c_3+c_4)](a-3b) - (c_6 + \tfrac{16}{3}c_5)(a+b)\} , \qquad (3.4)$$

$$a_{K^0\pi^0} = \{[c_1 - 2(c_2+c_3-2c_4)](a-3b) + (c_6 + \tfrac{16}{3}c_5)(a+b)\} ,$$

$$2E_\pi m_i = m_i^2 - m_f^2 + m_\pi^2 .$$

It is easy to see that the weak matrix elements $a_{K\pi}$ (3.4) have the same SU(4) structure as the b^s_{if}

terms in (3.3). The strong baryon-kaon vertex is parametrized through f and d in the usual way. Here m_i and m_f correspond to the initial and final baryon mass, respectively.

A^s terms can also be replaced by K^*-pole terms. The weak vertex in such a term can be readily estimated in the separable approximation.[16] However, the result is then practically equivalent to (3.2).

The main contributions to the B amplitudes come from baryon pole terms; for example,

$$B^P(\Xi^-_-) = -g\tilde{G}_F(m_\Xi + m_\Lambda)\left[\frac{a_{\Xi^0\Lambda}}{\sqrt{2}}\frac{f-d}{m_\Xi(m_\Xi - m_\Lambda)} + \frac{a_{\Xi^-\Sigma^-}}{\sqrt{3}}\frac{2d}{(m_\Xi - m_\Sigma)(m_\Sigma + m_\Lambda)}\right], \tag{3.5}$$

$$B^P(\Sigma^+_0) = -g\tilde{G}_F \frac{m_\Sigma + m_p}{m_\Sigma - m_p} a_{\Sigma^+ p}\left(\frac{1}{2m_p} - \frac{f}{m_\Sigma}\right).$$

The constants c_i appearing in formulas (3.1)–(3.5) are defined in Sec. II. The weak-interaction constants are

$$G_F = 1.026\times 10^{-5}\, m_p^{-2} = 1.166\times 10^{-2}\ \text{GeV}^{-2},$$
$$\sin\theta_C = 0.23, \tag{3.6}$$
$$f_\pi = 0.945 \quad m_\pi = 0.128\ \text{GeV}.$$

In any model employing valence quarks only, the amplitudes $a_{\Sigma^+ p}$, etc., receive contributions only from the operators O_1, O_5, and O_6. This is a consequence of the symmetry properties of the baryon states and operators.[8] The quantities $a_{\Sigma^+ p}$, etc., can be expressed through combinations of integrals over quark-model wave functions. This can be performed with the help of Table I using the definitions (2.2) of the four-quark operators. The integrals a, I_{01}, etc., appearing in Table I will be defined in the next section, where quark models are discussed. Other quantities are current form-factor parameters $\tilde{F}$ and $\tilde{D}$, and the hadron coupling parameters f and d. For the parameters $\tilde{F}$ and $\tilde{D}$ one can select

$$\tilde{F} + \tilde{D} = g_A = 1.25, \quad \tilde{D}/g_A = 0.65. \tag{3.7}$$

The baryon coupling parameters f and d are not too well determined experimentally.[7] A possible choice can be based on the relation[6]

$$\frac{g_{KN\Lambda}}{g_{KN\Sigma^0}} = \frac{g_A^\Lambda}{g_A^\Sigma}\frac{m_\Lambda + m_N}{m_\Sigma + m_N}, \tag{3.8}$$

which gives

$$f + d = 1, \quad f = 0.345. \tag{3.9}$$

Ω^- decays

It is well known[4,17,18] that the Ω^- decay probability is determined, to a large extent, by the parity-

TABLE I. Matrix elements of the four-quark operators as functions of the MIT-bag-model and the harmonic-oscillator-model integrals.

$\langle f\,\vert O\vert\, i\rangle$ \ O	$(\bar d s)(\bar u u)$ VV	$(\bar d s)(\bar u u)$ AA	$(\bar d u)(\bar u s)$ VV	$(\bar d u)(\bar u s)$ AA	$(\bar d s)(\bar d d)$ VV	$(\bar d s)(\bar d d)$ AA	$(\bar d s)(\bar s s)$ VV	$(\bar d s)(\bar s s)$ AA	HO quark model
$\langle\pi^-\vert O\vert K^-\rangle$	$-(a+5b)$	$(3a-b)$	0	$6(a-3b)$	0	0	0	0	$a=I'_{01}$, $b=0$
$\langle K^0\vert O\vert\pi^0\rangle$	0	$3\sqrt{2}(a-3b)$	$\frac{-1}{\sqrt{2}}(a+5b)$	$\frac{1}{\sqrt{2}}(3a-b)$	$\frac{1}{\sqrt{2}}(a+5b)$	$\frac{-1}{\sqrt{2}}(9a-19b)$	0	0	
$\langle n\vert O\vert\Lambda^0\rangle$	$\frac{-1}{\sqrt{6}}(3a+7b)$	$\frac{-1}{\sqrt{6}}(3a-b)$	$\frac{1}{\sqrt{6}}(8b)$	$\frac{2}{\sqrt{6}}(3a-b)$	$\frac{-1}{\sqrt{6}}(3a-b)$	$\frac{1}{\sqrt{6}}(3a-b)$	0	0	
$\langle n\vert O\vert\Sigma^0\rangle$	$\frac{-1}{3\sqrt{2}}(3a+23b)$	$\frac{-5}{3\sqrt{2}}(3a-b)$	$\frac{2}{3\sqrt{2}}(3a+11b)$	$\frac{4}{3\sqrt{2}}(3a-b)$	$\frac{-1}{3\sqrt{2}}(3a-b)$	$\frac{1}{3\sqrt{2}}(3a-b)$	0	0	$a=I_{01}$, $b=0$
$\langle p\vert O\vert\Sigma^+\rangle$	$\frac{-2}{3}(3a+11b)$	$\frac{-4}{3}(3a-b)$	$\frac{1}{3}(3a+23b)$	$\frac{5}{3}(3a-b)$	0	0	0	0	
$\langle\Lambda^0\vert O\vert\Xi^0\rangle$	$\frac{-3}{\sqrt{6}}(a+5b)$	$\frac{-3}{\sqrt{6}}(3a-b)$	$\frac{3}{\sqrt{6}}(a+5b)$	$\frac{3}{\sqrt{6}}(3a-b)$	0	0	$\frac{-1}{\sqrt{6}}(3a'-b')$	$\frac{1}{\sqrt{6}}(3a'-b')$	$a=I_{12}$
$\langle\Sigma^-\vert O\vert\Xi^-\rangle$	0	0	0	0	$\frac{-1}{3}(3a-b)$	$\frac{1}{3}(3a-b)$	$\frac{-1}{3}(3a'-b')$	$\frac{1}{3}(3a'-b')$	$a'=I_{12}$, $b=0$
$\langle\Sigma^0\vert O\vert\Xi^0\rangle$	$\frac{-1}{3\sqrt{2}}(3a-b)$	$\frac{1}{3\sqrt{2}}(3a-b)$	$\frac{-1}{3\sqrt{2}}(3a-b)$	$\frac{1}{3\sqrt{2}}(3a-b)$	0	0	$\frac{-1}{3\sqrt{2}}(3a'-b')$	$\frac{1}{3\sqrt{2}}(3a'-b')$	$b'=0$
$\langle\Xi^{*-}\vert O\vert\Omega^-\rangle$	0	0	0	0	0	0	$-4\sqrt{3}(3a'-b')$	$4\sqrt{3}(3a'-b')$	$a'=I_{23}$, $b'=0$

conserving p-wave decay amplitude B. They receive contributions from the pole diagrams

$$B^P(\Omega_K)=\tilde{G}_F\left[g_{\Lambda\Xi^{*-}K^-}\frac{a_{\Xi^{*-}\Omega^-}}{(m_{\Omega^-}-m_{\Xi^{*-}})}-g_{\Xi^0\Omega^-K^-}\frac{a_{\Xi^0\Lambda}}{(m_{\Xi^0}-m_\Lambda)}\right],$$

$$B^P(\Omega_-^-)=g_{\Xi^0\Xi^{*-}\pi^-}\tilde{G}_F\frac{a_{\Xi^{*-}\Omega^-}}{m_{\Omega^-}-m_{\Xi^{*-}}}, \quad (3.10)$$

$$B^P(\Omega_0^-)=g_{\Xi^-\Xi^{*-}\pi^0}\tilde{G}_F\frac{a_{\Xi^{*-}\Omega^-}}{m_{\Omega^-}-m_{\Xi^{*-}}}.$$

Here we have used the notations Ω_K for the $\Omega^-\to\Lambda^0K^-$ mode Ω_-^- for the $\Omega^-\to\Xi^0\pi^-$ mode and Ω_0^- for the $\Omega^-\to\Xi^-\pi^0$ mode. The matrix elements of the weak Hamiltonian determine the quantities $a_{\Xi^0\Lambda}$ (3.1) and $a_{\Xi^{*-}\Omega^-}$ as follows:

$$a_{\Xi^{*-}\Omega^-}=-8\sqrt{3}\,(c_6-\tfrac{8}{3}c_5)(3a'-b'). \quad (3.11)$$

These formulas can be easily found from Table I.

Strong meson coupling constants can be determined using SU(3) relations[7]

$$g_{p\Delta^{++}\pi^+}=g_{\Xi^0\Omega^-K^-}=-g_{\Xi^-\Omega^-\pi^0}=\sqrt{2}\,g_{\Lambda\Xi^{*-}K^-}=\sqrt{3}\,g_{\Xi^0\Xi^{*-}\pi^-}=\sqrt{6}\,g_{\Xi^-\Xi^{*-}\pi^0}=13.01\ \text{GeV}^{-1}. \quad (3.12)$$

The processed Ω_-^- and Ω_0^- also receive separable contributions:

$$B^s\begin{pmatrix}\Omega_-^-\\ \Omega_0^-\end{pmatrix}=\sqrt{2}f_\pi\tilde{G}_F F_-\begin{pmatrix}b^s_{\Omega^-\Xi^0}\\ \frac{-1}{\sqrt{2}}b^s_{\Omega^-\Xi^-}\end{pmatrix}, \quad (3.13)$$

$$b^s_{\Omega^-\Xi^0}=b^s_{\Xi^-\Lambda},\quad b^s_{\Omega^-\Xi^-}=b^s_{\Sigma^+p}.$$

The quantity F_- is connected with the octet-decuplet matrix element of the axial vector current as follows:

$$\langle\Xi^0|A^\mu|\Omega^-\rangle=F_-\bar{u}_{\Xi^0}(u\epsilon^\mu)_{\Omega^-}. \quad (3.14)$$

F_- is the integral over the quark-model wave functions. It is listed explicitly in Sec. IV. The spirit of the approximation defined by (3.10) and (3.12) for Ω_K and (3.13) for Ω_-^- and Ω_0^- is close to the one of Ref. 17. However, our approach involves well-defined quark models and includes systematically all possible separable and nonseparable contributions (octet-pole and decuplet-pole diagrams).[19]

The Ω^- decay probability can be calculated from the following formula:

$$\Gamma=\frac{|\vec{p}\,'|^3}{12\pi m_\Omega}\left[(E'+m_f)|B|^2+(E'-m_f)|C|^2\right]. \quad (3.15)$$

The contribution of the parity-violating d-wave amplitudes C to (3.15) is multiplied by an unfavorable factor $(E'-m_f)$. In our approach the amplitude C can receive contributions only from pole diagrams involving negative-parity $\frac{1}{2}^-$ and $\frac{3}{2}^-$ baryon resonances. This means that the sums of of baryon masses appear in the denominators instead of the differences as in (3.10). These two small factors (provided that the vertices in both B and C pole contributions are of the same order of magnitude) make Ω^- decay almost parity conserving.

The separable contributions (3.13) can also be replaced by K-pole terms as was done in formulas (3.3) and (3.4). For Ω^- decay, one finds

$$B^K\begin{pmatrix}\Omega_-^-\\ \Omega_0^-\end{pmatrix}=g_{\Xi^0\Omega^-K^-}\tilde{G}_F\frac{2E_\pi}{(m_K^2-m_\pi^2)}\begin{pmatrix}a_{K^-\pi^-}\\ \frac{-1}{\sqrt{2}}a_{K^0\pi^0}\end{pmatrix}. \quad (3.16)$$

Here $a_{K\pi}$ are given by (3.4).

IV. QUARK MODELS FOR HADRONS

The MIT version of the bag model (MIT) has been extensively used and explained in the literature.[2-5] Therefore it is sufficient here just to list integrals appearing in Table I. They are determined by the bag model wave functions, with the symbols u_q and v_q denoting large and small components, respectively[5]:

$$a=\int_0^R d^3r[u_u^3(r)u_s(r)+v_u^3(r)v_s(r)],$$

$$b=\int_0^R d^3r[u_u^2(r)v_u(r)v_s(r)+v_u^2(r)u_u(r)u_s(r)],$$

$$a'=\int_0^R d^3r[u_s^3(r)u_u(r)+v_s^3(r)v_u(r)], \quad (4.1)$$

$$b'=\int_0^R d^3r[u_s^2(r)v_s(r)v_u(r)+v_s^2(r)u_s(r)u_u(r)],$$

$$F_-=2\sqrt{6}\int_0^R d^3r[u_u(r)u_s(r)-\tfrac{1}{3}v_u(r)v_s(r)].$$

Some more details will be given about the harmonic-oscillator (HO) quark model in order to illustrate how the HO model compares with the MIT model. Both models have an identical spin-flavor-color structure. The existing dynamical difference is felt through the radial dependence. The HO model is usually treated in the configurational space.[9,10] However, as both pictures are essentially equivalent, the HO model can be easily transformed into the N representation, if desired. In the configuration representation, the baryon states are, for example, for the Σ^+ particle:

$$\langle\vec{r}_1,\vec{r}_2,\vec{r}_3|\Sigma^+\rangle=\frac{1}{\sqrt{2}}G(1,2,3)(\Phi^\rho_{\Sigma^+}\chi^\rho_++\Phi^\lambda_{\Sigma^+}\chi^\lambda_+)c. \quad (4.2)$$

The spin χ^ρ, χ^λ and flavor Φ^ρ, Φ^λ states are defined in Refs. 9 and 10. The spatial part has to be properly symmetrized as follows[20]:

$$G(1,2,3)=N\{g(12,3)+g(13,2)+g(23,1)\},$$

$$g(12,3)=\left(\frac{\alpha_\rho\alpha_\lambda}{\pi}\right)^{3/2}\exp[-\tfrac{1}{2}(\alpha_\rho^2\rho^2+\alpha_\lambda^2\lambda^2)],$$

$$g(23,1)=\left(\frac{\alpha_\rho\alpha_\lambda}{\pi}\right)^{3/2}\exp\{-\tfrac{1}{8}[(\alpha_\rho^2+3\alpha_\lambda^2)\rho^2+(\alpha_\lambda^2+3\alpha_\rho^2)\lambda^2+2\sqrt{3}(\alpha_\lambda^2-\alpha_\rho^2)\vec{\rho}\vec{\lambda}]\},$$

$$g(13,2)=\left(\frac{\alpha_\rho\alpha_\lambda}{\pi}\right)^{3/2}\exp\{-\tfrac{1}{8}[(\alpha_\rho^2+3\alpha_\lambda^2)\rho^2+(\alpha_\lambda^2+3\alpha_\rho^2)\lambda^2-2\sqrt{3}(\alpha_\lambda^2-\alpha_\rho^2)\vec{\rho}\vec{\lambda}]\},\quad(4.3)$$

$$N^{-2}=3+2\left(\frac{16\alpha_\rho^2\alpha_\lambda^2}{(\alpha_\rho^2+3\alpha_\lambda^2)(\alpha_\lambda^2+3\alpha_\rho^2)}\right)^{3/2}+4\left(\frac{64\alpha_\rho^2\alpha_\lambda^2}{(5\alpha_\rho^2+3\alpha_\lambda^2)(5\alpha_\lambda^2+3\alpha_\rho^2)-3(\alpha_\lambda^2-\alpha_\rho^2)^2}\right)^{3/2},$$

$$\vec{\rho}=\frac{1}{\sqrt{2}}(\vec{r}_1-\vec{r}_2),\quad \vec{\lambda}=\frac{1}{\sqrt{6}}(\vec{r}_1+\vec{r}_2-2\vec{r}_3),$$

$$\alpha_\rho^2=m_u\omega_\rho=m_u\left(\frac{3K}{m_u}\right)^{1/2},\qquad \alpha_\lambda^2=m_\lambda\omega_\lambda=m_\lambda\left(\frac{3K}{m_\lambda}\right)^{1/2},$$

$$m_\lambda=\frac{3m_um_s}{2m_u+m_s}.$$

Mass differences force us to distinguish between u, d, and s quarks in the spatial part of the wave function. This is achieved by the following self-explanatory substitution, which is important when Ξ and Ω particles are concerned:

$$\tilde{g}=g(qqs\to qqq),\quad q=u,d,$$
$$\hat{G}=G(qqs\to ssq)=\hat{N}\Sigma\hat{g}(ssq),\qquad(4.4)$$
$$\hat{\tilde{g}}=g(qqs\to sss)=\tilde{g}(qqq\to sss)=\hat{g}(ssq\to sss).$$

The integrals appearing in Table I are

$$I'_{01}=\int d^3r\,\delta^{(3)}(\vec{r})g(1,2)\tilde{g}(1,2)=\left(\frac{\beta\beta'}{\pi}\right)^{3/2},$$

$$I_{01}=\tfrac{1}{6}\sum_{i,j=1}^{3}\int d^3\rho\,d^3\lambda\,\delta^{(3)}(\vec{r}_i-\vec{r}_j)\tilde{g}(12,3)G(1,2,3)$$
$$=\frac{N}{2\sqrt{2}}\left(\frac{2}{\pi}\alpha_\rho^3\alpha_\lambda\right)^{3/2}[(\alpha_\rho^2+\alpha_\lambda^2)^{-3/2}+2(\alpha_\rho^2+\alpha_\lambda'^2)^{-3/2}],$$

$$I_{12}=\tfrac{1}{6}\sum_{i,j=1}^{3}\int d^3\rho\,d^3\lambda\,\delta^{(3)}(\vec{r}_i-\vec{r}_j)G(1,2,3)\hat{G}(1,2,3)$$
$$=\frac{N\hat{N}}{2\sqrt{2}}\left(\frac{2}{\pi}\alpha_\rho\alpha_\lambda\hat{\alpha}_\rho\hat{\alpha}_\lambda\right)^{3/2}[(\alpha_\rho^2+\hat{\alpha}_\lambda^2)^{-3/2}+2(\alpha_\lambda^2+\hat{\alpha}_\lambda'^2)^{-3/2}+2(\alpha_\lambda'^2+\hat{\alpha}_\lambda^2)^{-3/2}+4(\alpha_\lambda'^2+\hat{\alpha}_\lambda'^2)^{-3/2}],\qquad(4.5)$$

$$I_{23}=\tfrac{1}{6}\sum_{i,j=1}^{3}\int d^3\rho\,d^3\lambda\,\delta^{(3)}(\vec{r}_i-\vec{r}_j)\hat{G}(1,2,3)\hat{\tilde{g}}(12,3)$$
$$=\frac{\hat{N}}{2\sqrt{2}}\left(\frac{2}{\pi}\hat{\alpha}_\rho^3\hat{\alpha}_\lambda\right)^{3/2}\left[(\hat{\alpha}_\rho^2+\hat{\alpha}_\lambda^2)^{-3/2}+2(\hat{\alpha}_\rho^2+\hat{\alpha}_\lambda'^2)^{-3/2}\right]\quad(i\neq j),$$

$$F_-=2\sqrt{6}\int d^3\rho\,d^3\lambda\,\hat{\tilde{g}}(12,3)\hat{G}(1,2,3)$$
$$=2\sqrt{6}\,\hat{N}(2\hat{\alpha}_\rho^3\hat{\alpha}_\lambda)^{3/2}\{[\hat{\alpha}_\rho^2(\hat{\alpha}_\rho^2+\hat{\alpha}_\lambda^2)]^{-3/2}+4[(\hat{\alpha}_\rho^2+\hat{\alpha}_\rho'^2)(\hat{\alpha}_\rho^2+\hat{\alpha}_\lambda'^2)-\tfrac{3}{4}(\hat{\alpha}_\lambda^2-\hat{\alpha}_\rho^2)^2]^{-3/2}\}.$$

The parameters are connected with quark mass ratio x $(=m_u/m_s)$ and the harmonic oscillator shape parameters α_ρ (for baryons) and β (for mesons) by simple relations

$$\alpha_\lambda^2=\left(\frac{3}{2x+1}\right)^{1/2}\alpha_\rho^2,\quad \alpha_\lambda'^2=\tfrac{1}{4}\left[1+3\left(\frac{3}{2x+1}\right)^{1/2}\right]\alpha_\rho^2,$$

$$\hat{\alpha}_\lambda^2=\left(\frac{3}{x+2}\right)^{1/2}\alpha_\rho^2,\quad \hat{\alpha}_\lambda'^2=\tfrac{1}{4}\left[1+3\left(\frac{3}{x+2}\right)^{1/2}\right]\alpha_\rho^2,\qquad(4.6)$$

$$\hat{\alpha}_\rho^2=\frac{1}{\sqrt{x}}\alpha_\rho^2,\quad \beta^2=\frac{1}{\sqrt{3}}\alpha_\rho^2,\quad \beta'^2=\left(\frac{2}{x+1}\right)^{1/2}\beta^2.$$

The harmonic-oscillator quark-model integral I'_{01} (4.5) has the value comparable with the MIT-bag integrals for mesons.[21]

Wave functions analogous to (4.2) correspond to pure ground-state configurations. All admixtures of excited-state configurations[22] have been omitted. This corresponds exactly to the approach based on the MIT model. Numerical values of the integrals (4.1) and (4.5) depend strongly on the model parameters. It seems natural to use parameters which correspond to the ground-state baryons. In the literature, one can find two sets of ground-state parameters for the MIT bag model. The set of Ref. 5 is the only one that leads to a satisfactory prediction on nonleptonic decay amplitudes. It predicts the charge radius of the proton to be $\langle r_p^2\rangle_{\rm th}^{1/2}=0.73$ fm, which is somewhat smaller than the experimental value $\langle r_p^2\rangle^{1/2}_{\rm exp}=0.88$ fm.[23] The other set of ground-state MIT parameters[24] predicts the charge radius to be $\langle r_p^2\rangle^{1/2}_{\rm th}=1.03$ fm, which is too large. The HO model parameters corresponding to the ground-state baryons[10] also predict the charge radius of the proton to be $\langle r_p^2\rangle_{\rm th}^{1/2}=0.63$ fm, which is too small. Therefore, there is some consistency between the parameters used for the two models. Their respective values are the following.

MIT bag model:

baryons:

$$R_u \simeq R_s = R = 4.96 \text{ GeV}^{-1},$$

$$m(u)=m(d)=0, \quad m(s)=0.279 \text{ GeV}, \tag{4.7a}$$

$$\omega_{1,-1}(u)=2.0428, \quad \omega_{1,-1}(s)=2.8448,$$

$$a=1.5475\times10^{-3} \text{ GeV}^3,$$

$$a'=1.9107\times10^{-3} \text{ GeV}^3,$$

$$b=0.5936\times10^{-3} \text{ GeV}^3, \tag{4.7b}$$

$$b'=0.4995\times10^{-3} \text{ GeV}^3,$$

$$F_-=2\sqrt{6}\times0.7091.$$

mesons:

$$R_u \simeq R_s = R = 3.26 \text{ GeV}^{-1},$$

$$m(u)=m(d)=0, \quad m(s)=0.279 \text{ GeV}, \tag{4.8a}$$

$$\omega_{1,-1}(u)=2.0428, \quad \omega_{1,-1}(s)=2.5407,$$

$$a=5.3436\times10^{-3} \text{ GeV}^3,$$

$$b=2.1762\times10^{-3} \text{ GeV}^3. \tag{4.8b}$$

HO quark model:

mesons and baryons:

$$m(u)=m(d)=0.33 \text{ GeV}, \quad m(s)=0.55 \text{ GeV},$$

$$x=0.6, \quad \alpha_\rho=0.32 \text{ GeV}, \tag{4.9a}$$

$$I'_{01}=2.8069\times10^{-3} \text{ GeV}^3,$$

$$I_{01}=2.1168\times10^{-3} \text{ GeV}^3,$$

$$I_{12}=2.5188\times10^{-3} \text{ GeV}^3, \tag{4.9b}$$

$$I_{23}=2.9801\times10^{-3} \text{ GeV}^3,$$

$$F_-=2\sqrt{6}\times0.8022.$$

Alternative sets of the HO-model parameters,[9] more appropriate for the excited states of baryons, lead to unsatisfactory magnitudes of the integrals (4.5).

V. RESULTS AND DISCUSSION

In the present investigation, we have tried to test systematically various inputs entering the calculation of nonleptonic hyperon and Ω^- decays. The points tested are the following.

(1) The effective weak nonleptonic Hamiltonian. This includes variations of parameters such as μ, m_w, current quark masses, etc., which determine QCD renormalization (or enhancement) coefficients. The Hamiltonian without QCD renormalization was also considered, as well as the Hamiltonian containing empirically determined enhancement coefficients.

(2) Quark models of hadrons. Calculations were made using both the MIT bag model and the harmonic-oscillator (Isgur-Karl) quark model. Various sets of model parameters, such as quark masses, bag radii, harmonic-oscillator strengths, are discussed.

(3) Dynamical assumptions underlying calculations, for example, which current commutator, pole terms, and/or separable terms to take and which semiempirical and/or calculated parameters to use to determine strong vertices or weak-current form factors.

Permutations and combinations of possible selections listed under 1, 2, and 3 are numerous. We have tried to make some selections which would give some feeling of this complex situation.

First, we have used QCD enhancement factors as listed in (2.5), and calculated nonleptonic hyperon amplitudes through the current commutator terms (3.1), the baryon pole terms (3.5), and the separable terms (3.2) and (3.3). This seems to be the combination suggested in the original paper.[1] Calculations were carried out for both the MIT bag model and the harmonic-oscillator model with parameters as defined by (4.7a), (4.8a), and (4.9a). Numerical results are presented in full detail in Table II. It is immediately obvious that there is not much to choose between the two models and that it is fair to say that the results for hyperon nonleptonic decays are truly quark-model independent.

Secondly, for the QCD enhancement factors c_i, we have used analytical expressions as given by Ref. 1. The masses appearing in c_i were varied in the range $60< m_w\,(\text{GeV})< 120$, $1< m_c(\text{GeV})< 2$, $0.3< \mu(\text{GeV})< 0.7$. Calculations were also performed for the following spread of the vertex parameters:

$$0.3< f< 0.375, \quad 0.4< \tilde{F}< 0.5.$$

They were carried out for both MIT and HO quark models using the dynamical scheme defined by formulas (3.1), (3.2), (3.3), and (3.5). The computer output was studied for numerous combinations of parameters, with the conclusion that the values are fairly stable, extreme changes in ab-

TABLE II. Hyperon nonleptonic decay amplitudes. Experimental values are from Ref. 25. The values of the parameters are: $m_W=100$ GeV, $m_c=2$ GeV, $\mu=0.5$ GeV, $f=0.315$, $\tilde{F}=0.5$.

	MIT bag model						HO quark model						
	Nonseparable		Separable				Nonseparable		Separable				
Amplitude	O_{1-4}	$O_{5,6}$	O_{1-3}	O_4	$O_{5,6}$	Total	O_{1-4}	$O_{5,6}$	O_{1-3}	O_4	$O_{5,6}$	Total	Expt.
$10^6A(\Lambda_-^0)$	0.19	−0.01	−0.05	−0.01	0.09	0.21	0.19	−0.01	−0.05	−0.01	0.09	0.21	0.32
$10^6A(\Xi_-^-)$	−0.38	0.01	0.06	−0.02	−0.10	−0.43	−0.45	0.00	0.06	−0.02	−0.01	−0.42	−0.44
$10^6A(\Sigma_0^+)$	−0.33	0.00	0.04	−0.02	−0.07	−0.38	−0.33	−0.00	0.04	−0.02	−0.07	−0.38	−0.32
$10^6A(\Sigma_-^-)$	0.47	−0.00	−0.06	−0.02	0.10	0.49	0.46	0.01	−0.06	−0.02	0.10	0.49	0.42
$10^6A(\Sigma_+^+)$	0	0	0	0	0	0	0	0	0	0	0	0	0.01
$10^6B(\Lambda_-^0)$	0.65	−0.10	0.42	0.12	0.72	1.81	0.69	−0.14	0.42	0.12	0.72	1.81	2.16
$10^6B(\Xi_-^-)$	1.56	−0.12	−0.17	−0.05	−0.28	0.94	1.84	−0.16	−0.17	−0.05	−0.28	1.18	1.45
$10^6B(\Sigma_0^+)$	1.83	−0.01	0.08	−0.04	0.14	2.00	1.73	0.02	0.08	−0.04	0.14	1.93	2.60
$10^6B(\Sigma_-^-)$	−0.10	0.04	−0.12	−0.03	−0.20	−0.41	−0.01	0.07	−0.12	−0.03	−0.20	−0.29	−0.14
$10^6B(\Sigma_+^+)$	2.50	0.03	0	0	0	2.53	2.44	0.10	0	0	0	2.54	4.13

solute values differing by about 80% at most. As an example, the extreme values of Ξ_-^- amplitudes in the two models are the following.

MIT bag model:

$$-0.48 < A(\Xi_-^-)\times 10^6 < -0.36\,,$$

$$0.33 < B(\Xi_-^-)\times 10^6 < 1.54\,,$$

HO quark model:

$$-0.56 < A(\Xi_-^-)\times 10^6 < -0.44\,,$$

$$0.47 < B(\Xi_-^-)\times 10^6 < 1.84\,.$$

In all these parametric variations, the relative signs of the hyperon-decay amplitudes were always reproduced correctly. This also holds for the relative magnitudes of the various amplitudes; a few interesting examples shown in Fig. 1.

A completely analogous approach, using the same effective weak Hamiltonian (2.1)–(2.5) and an analogous mixture of pole and separable terms (3.10) and (3.13) was then used to calculate B amplitudes for Ω^- decays. Numerical values are summarized in Table III. It seems that the results based on the MIT bag model, which have already been published,[4] agree with experimental values somewhat better than the results based on the HO model. However, this advantage is very slight and probably not significant at all. Both quark models reproduce one experimental amplitude almost exactly. For the MIT model, it is $B(\Omega_K)$, while for the HO model, it is $B(\Omega_-^-)$.

It is necessary to point out that all the above results follow for a particular choice of parameters, either as in Ref. 5 or as in Ref. 10. An alternative set of MIT parameters[24] leads to much poorer predictions of the decay amplitudes. The results based on the HO model also depend on the selection of the parameter α_ρ. If we had used the parameter α_ρ of Ref. 9, we would have obtained

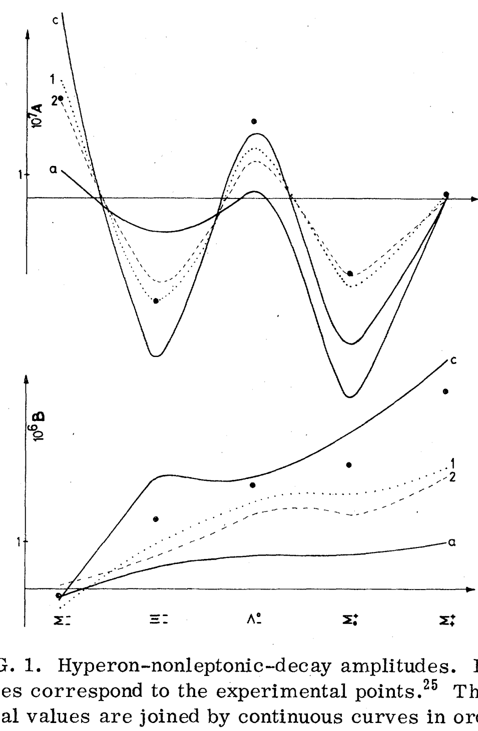

FIG. 1. Hyperon-nonleptonic-decay amplitudes. Full circles correspond to the experimental points.[25] Theoretical values are joined by continuous curves in order to facilitate visual display. Curve 1 corresponds to the values given in Table II. Curve 2 was found by using the same dynamics, but changing the parameters to $m_W=60$ GeV, $m_c=1$ GeV, $\mu=0.7$ GeV, $f=0.375$, and $\tilde{F}=0.5$. Curves a and c correspond to the columns in Table IV labeled by the same letters.

TABLE III. Ω^- decay amplitudes in GeV^{-1}. Experimental values are from Ref. 26. The values of the parameters are the same as in Table II.

	MIT bag model						HO quark model						
	Nonseparable		Separable				Nonseparable		Separable				
Amplitude	O_{1-4}	$O_{5,6}$	O_{1-3}	O_4	$O_{5,6}$	Total	O_{1-4}	$O_{5,6}$	O_{1-3}	O_4	$O_{5,6}$	Total	Expt.
$10^6 B(\Omega_K)$	3.21	0.80	0	0	0	4.01	3.78	1.46	0	0	0	5.24	4.03
$10^6 B(\Omega^-_-)$	0	−0.70	0.78	0.23	1.29	1.60	0	−1.19	0.85	0.25	1.36	1.27	1.32
$10^6 B(\Omega^-_0)$	0	−0.49	0.55	−0.32	0.90	0.64	0	−0.84	0.63	−0.36	1.01	0.44	0.78

$I_{01} = 6.6 \times 10^{-3}$ instead of the value quoted in (4.9b). In some earlier papers[27,28] in which nonleptonic decays were treated within the context of the HO model, one can also find large values for I_{01}. In those papers, I_{01} was determined by fitting some other experimental quantity—electromagnetic and strong hyperfine splittings, for example.[28] The result based on Ref. 9, where the spectrum and mixing angles of negative-parity baryons were discussed, is close to the I_{01} values of Refs. 27 and 28. The HO-model parameters[10] which we have selected are those which are determined in a way consistent with the determination of the MIT-model parameters.[5]

Keeping to an identical dynamical scheme, i.e., based on (2.1), (3.1), (3.3), (3.5), (3.10), and (3.13), we look at what happens if there are no important QCD-renormalization effects. This is performed by using the values (2.6) for the coefficients c_i. Such a possibility has been suggested and discussed recently.[29] As shown in Tables IV and V, the agreement with experiment, especially where B amplitudes are concerned, is poorer than with QCD corrections included. This disagreement is not strong because orders of magnitude and signs still come out correctly. It is possible that alternative dynamical assumptions concerning both the dominant Feynman diagrams and the quark models could lead to noticable improvement.

A recent study of D-meson decay has shown that enhancement coefficients should be actually much larger than found in the usual QCD calculations.[1] This corresponds to the choice (2.7). Calculations carried out for the dynamical scheme as outlined above are summarized in Tables IV and V. Taking into account certain freedom one has in selecting dynamical parameters (see above and Ref. 3), one can claim no clear-cut preference for either (2.5) or (2.7).

It is interesting that very good results can be obtained by weakening artifically the $\Delta I = \frac{3}{2}$ piece of the weak Hamiltonian (2.1). This can be achieved

TABLE IV. Hyperon nonleptonic decay amplitudes calculated for various effective weak Hamiltonians and for various pole diagrams.

Amplitude	Total[a] without QCD	Total[b] including K pole $c_5=-0.08$	$c_5=+0.1$	$c_5=-/+0.5$	Total[c] for $c_-=5$ $c_+=0.45$	Expt.
$10^6 A(\Lambda^0_-)$	0.03	0.21	0.20	0.12/0.25	0.28	0.32
$10^6 A(\Xi^-_-)$	−0.15	−0.44	−0.39	−0.31/−0.44	−0.68	−0.44
$10^6 A(\Sigma^+_0)$	−0.85	−0.38	−0.33	−0.33	−0.62	−0.32
$10^6 A(\Sigma^-_-)$	0.12	0.49	0.47	0.37	0.78	0.42
$10^6 A(\Sigma^+_+)$	0	0	0	0	0	0.01
$10^6 B(\Lambda^0_-)$	0.70	0.62	1.39	−1.17/3.10	2.30	2.16
$10^6 B(\Xi^-_-)$	0.44	1.43	1.63	0.95/2.09	2.36	1.45
$10^6 B(\Sigma^+_0)$	0.70	1.75	2.04	1.06/2.70	3.46	2.60
$10^6 B(\Sigma^-_-)$	−0.16	−0.14	−0.80	1.42/−2.29	−0.21	−0.14
$10^6 B(\Sigma^+_+)$	0.99	2.53	2.46	2.71/2.29	4.81	4.13

[a] (B pole)+(separable).
[b] (B pole)+(K pole).
[c] (B pole)+(separable).

TABLE V. Ω^- decay amplitudes in GeV^{-1} calculated for the same variations as displayed in Table IV.

Amplitude	Total[a] without QCD	Total[b] including K pole $c_5=-0.08$	$c_5=0.1$	$c_5=-/+0.5$	Total[c] for $c_-=5$ $c_+=0.45$	Expt.
$10^6\lvert B(\Omega_K)\rvert$	1.27	4.01	2.03	8.00/2.40	6.33	4.03
$10^6\lvert B(\Omega_-^-)\rvert$	0.82	0.84	0.34	1.50/1.53	1.61	1.32
$10^6\lvert B(\Omega_0^-)\rvert$	0.19	0.31	0.07	1.40/1.50	0.79	0.78

[a] (B pole)+(separable).
[b] (B pole)+(K pole).
[c] (B pole)+(separable).

making the replacement

$$c_4 \to \tfrac{1}{3}c_4$$

in the theoretical set of values for c_i (2.5). Such a replacement does not affect nonleptonic hyperon-amplitudes very much. However, it leads to a marked improvement in Ω^- decay. Agreement between theory and experiment is almost ideal, i.e., better than 10%.

Reference 3 used a slightly different dynamical scheme. Its main point is in replacing separable terms (contributing to the B amplitudes) by K-meson pole terms. For hyperon nonleptonic decays, such a choice has a theoretical foundation.[30] We have also introduced K-meson pole terms in the Ω^--decay calculations. In this scheme, nonleptonic hyperon-decay amplitudes A retain the separable terms (3.7). In Ref. 3, they were deduced from the current-algebra commutator term.

The contribution of the K-pole term is numerically smaller than the corresponding separable terms, which contain small current-quark masses in the denomiantor. Thus, the values obtained using the theoretical enhancement coefficients (columns for $c_5=-0.08$ in Tables IV and V) contain some amplitudes which are too small. However, the c_5 and c_6 coefficients are products of SU(4) symmetry breaking, and cannot be estimated theoretically with great accuracy.[1,3]

It is thus permissible to vary them within certain limits. We used $c_5=0.1$, ∓ 0.5, leaving the theoretical value (2.5) for c_6.[31]

All these combinations work badly for the $B(\Sigma_-^-)$ and $B(\Omega_K)$ amplitudes. Again, this cannot be regarded as an argument against K-pole dominance. The calculation of the K-pole term depends actually much more on the quark-model dynamics than the estimate of the separable terms. It is possible that the dynamics, as depicted by the simple quark models we have been using, is not sophisticated enough for an accurate estimate of K-pole terms. It might be that K-pole contributions are dominated by separable terms (which are one of the contributions[3]) for which the semiempirical estimates (1.2) work well.

As shown in Secs. III and IV, the estimate of the K-pole term based on the HO correponds roughly to the estimate based on the MIT bag model.[21]

With the dynamical scheme we have been employing, some QCD enhancement seems to be absolutely necessary. However, our approach includes many other dynamical assumptions which have not yet been well understood. The quark models we were using introduce model quark masses for their valence quarks, which should be understood as suitable parameters used in fitting certain experimental data, such as masses, magnetic moments, charge distributions, etc. The MIT bag model employs very light valence quarks and the relativistic dynamics inside the bag, while the HO quark-model valence quarks are heavy and their dynamics are strictly nonrelativistic. However, hadron states in both models are classified according to the same spin-flavor symmetry. The fact that both models lead to similar results means that the spin-flavor structure of the hadron states is decisive for correct treatment of the dynamics.

From one point of view it might still be surprising that both models lead to similar results. What one is actually calculating is a product of the QCD enhancement factor c_j and the matrix element of the four-quark operator O_j between the quark-model states (i,f):

$$c_j(\mu^2)\langle f|O_j|i\rangle(\mu^2)\,.$$

The operators O_j (Sec. II) are made out of current quarks. Without any transformation they are used to act as operators in the space of valence quarks constituting baryon or meson states $|i\rangle$. It is not at all obvious that such matching should be equally good (or bad) for any quark model. Furthermore, the product $c_j\langle f|O_j|i\rangle$ should be independent of the renormalization mass μ.[32] This can come about only if the renormalization mass is selected in such a way that the matrix element $\langle f|O_j|i\rangle$ carries a proper concealed μ dependence. Again, this can easily depend on the particular quark model for hadrons.

Naturally, this last statement loses its meaning if QCD corrections turn out to be unimportant and if everything depends on the modification of hadron states.[29,33]

Furthermore, in our approach we have not considered the $\pi^0-\eta^0$ mixing effect which could change the implication of the $\Delta I=\frac{1}{2}$ selection rule for A amplitudes.[35]

Irrespective of all these open questions, and possible shortcomings, it seems that the general dynamical scheme, outlined in Secs. II, III, and IV satisfactorily works. Obviously, we have not controlled all details, neither minor dynamical ones, nor really important ones such as QCD corrections. However, it seems that with the Weinberg-Salam model of weak interactions applied in the framework of a relatively simple dynamics, we might be on the right track.

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