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J. Lach

Fermi National Accelerator Laboratory
P.O. Box 500, Batavia, Illinois 60510

P. Zenczykowski

Insitute of Nuclear Physics
Radzikowskiego 152, 31-342 Krakow, Poland

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J. LACH
Fermilab, PO Box 500 MS 220
Batavia, IL 60510, USA

P. ŻENCZYKOWSKI
Institute of Nuclear Physics
Radzikowskiego 152
31-342 Krakow, Poland

ABSTRACT

We review the experimental and theoretical status of weak radiative hyperon decays. Our discussion centers around a controversy over the validity of Hara's theorem originally expected to be respected by these decays. After presenting the hadron-level theorem we describe experiments that have provided convincing evidence against its applicability to these decays. In the theoretical part we juxtapose the hadron-level and quark-level approaches and discuss the violation of Hara's theorem in the latter. We review quark-model phenomenology which offers a promising description of experimental data. Measurements that should be done to settle the theoretical controversy are pinpointed. The importance of radiative hyperon decays in understanding the nonlocal composite nature of hadrons is stressed.

1. Introduction

Hyperon radiative decays exhibit the full interplay of the electromagnetic, weak, and strong interactions. One would think that because of their simple two body kinematics: just the decay of one baryon into another with the emission of a photon, they should be amenable to insightful theoretical analysis and clean experimental probing. In fact, they have proved to be a challenge to both the theorist and experimenter.

These strangeness changing decays are induced by the weak interactions but their final state photon ensures that the electromagnetic forces are also involved. Since baryons are strongly interacting particles, the strong force is also important.

The baryon octet provides us with multiple reactions of this class with varying quark content of the initial and final state baryons. These are the decays

$$
\begin{array}{ll}
\Sigma^+ \to p\ \gamma & \Xi^\circ \to \Sigma^\circ\ \gamma \\
\Sigma^\circ \to n\ \gamma & \Xi^\circ \to \Lambda^\circ\ \gamma \\
\Lambda^\circ \to n\ \gamma & \Xi^- \to \Sigma^-\ \gamma
\end{array}
\qquad (1)
$$

Because the weak decay, $\Sigma^\circ \to n\ \gamma$, is completely overwhelmed by the simpler strangeness conserving electromagnetic decay, $\Sigma^\circ \to \Lambda^\circ\ \gamma$, it has not been observed. Except for this decay, all decays of the baryon octet have received major attention and been observed.

Although decays from the baryon decuplet are also of great interest, the only member of the decuplet with a sufficiently long lifetime to make it accessible to experimental study is the Ω^-. Its decay to a member of the octet

$$\Omega^- \rightarrow \Xi^- \gamma$$

is expected to be dominant.

So far, however, only upper limits have been put on its branching fraction (R). The Ω^- decay to a member of the decuplet

$$\Omega^- \rightarrow \Xi^*(1530)\gamma$$

is expected to have a branching fraction an order of magnitude smaller than the octet mode.[1,2]

The weak radiative hyperon decays (WRHD) pose significant experimental challenges. They have small branching fractions, $\approx 10^{-3}$, and copious photon backgrounds due to their more abundant decay modes involving $\pi^0 \rightarrow \gamma\gamma$. The most sensitive tests require the use of polarized hyperons. Modern hyperon beams have provided effective tools for overcoming these difficulties.

Theoretical difficulties manifest themselves in a long history of unsuccessful attempts to describe data in hadron level approaches and in the appearance of a basic conflict between these approaches and the quark model. For these reasons WRHD have been regarded as "the last low q^2 frontier of weak interaction physics[3]", "unsolved puzzle[4]", "the long-standing $\Sigma^+ \rightarrow p\gamma$ puzzle[5]", "a puzzle which has so far defied a simple and widely accepted solution[2]", and "a long standing discrepancy[6]". Clearly, there are still many unsolved questions in the domain of low q^2 weak interaction physics (for example the origin of the $\Delta I=1/2$ rule, etc.). The problem of WRHD seems to be of a more fundamental nature, however.

Yet, significant experimental and theoretical progress has been made in the last ten years. The aim of this paper is to review the present status of both the theory and experiments, discuss what we believe are existing problems and point to promising future directions.

We follow this introduction with Section 2, a brief discussion of the experimental techniques, and present the basic measurements in Table 2.1. This is followed by Section 3 devoted to Hara's theorem in which we try to crystallize the nature of the problem. Before embarking on a discussion of the experimental results, we present theoretical lower bounds on the WRHD branching fractions in Section 4. These bounds are imposed by unitarity and therefore they are very reliable.

In sections 3 and 4 we develop general theoretical arguments and explore the nature of the controversy. In Section 5, we return for a more detailed description of the experimental measurements. Section 6 develops the hadron level formalism for nonleptonic and radiative hyperon decays, Section 7 explores the phenomenology of the standard approaches, Section 8 looks at the single quark processes, Section 9 considers other approaches and, finally, we present our conclusions in Section 10.

2. Experimental Techniques and Data Summary

Hyperon radiative decay measurements consist of branching fractions and asymmetry parameters. The measurement of branching fractions requires identification of the hyperons and of the unique radiative decay final state. Although limited statistically, reliable measurements came from early bubble chamber experiments. More recent measurements have employed electronic techniques and most have relied on high energy hyperon beams at Fermilab.

Hyperon beams can provide substantial fluxes of hyperons; furthermore, they can be produced with significant polarization. The direction and magnitude of these polarizations can be controlled thus providing an important tool for the evaluation of systematic uncertainties. These high energy hyperon beams with their easily controlled polarizations have also allowed us to make precision measurements of hyperon static properties.

They have allowed us to study polarization effects in Σ^- beta decay,[7] high statistics weak radiative decays,[8] and to make precision measurements of hyperon magnetic moments.[9-11] Hyperon polarization has provided an extremely useful tool for the study of hyperon fundamental properties, although the production mechanism which produces these polarizations is not well understood. A number of reviews describing hyperon beams and the physics programs that have utilized them are available.[12-15]

In recent years it has become clear that hyperon polarization itself is a complex process whose energy and P_t dependence is different[16-19] for each of the hyperons. This has provided significant challenges to our theoretical understanding of polarization mechanisms.

Table 2.1 shows the present experimental status of weak radiative hyperon decays. Not included are some early experiments which have presented upper limits that have been superseded by more recent experiments which have observed the decay. The newer measurements are consistent with the previously measured limits.

In Table 2.1, we list the experimental branching fractions (R), asymmetry parameters (α), number of events, year and place of measurement, and refer to the experimental group by the first author. We quote both the statistical and systematic uncertainty (in that order) for each measurement if available. For those decays where more than one measurement exists, we first combine the statistical and systematic uncertainties quadratically and then form the weighted mean for each set of measurements.

The $\Sigma^+ \to p\gamma$ reaction was the first WRHD to be observed and stimulated the controversy that is still with us. In Table 2.1 we have not included $\Sigma^+ \to p\gamma$ measurements which contain less than 25 events. These are early emulsion[20] and bubble

Table 2.1 Hyperon Radiative Decay Measurements

	R (x10^3)	α	Events	Year	Group
$\Sigma^+ \to p\gamma$	1.20±0.06±0.05		31901	1994	Fermilab, Timm[36]
		-0.720±0.086±0.045	34754	1992	Fermilab, Foucher[8]
	$1.45 \pm 0.20\,^{+0.11}_{-0.22}$		408	1989	BNL, Hessey[37]
	1.30±0.15	-0.86±0.13±0.04	190	1987	KEK, Kobayashi[38]
	1.27±0.17		155	1985	CERN, Biagi[39]
	1.09±0.20	$-0.53\,^{+0.38}_{-0.36}$	30 (R) 46 (α)	1980	CERN, Manz[25]
	1.42±0.26	s $-1.03\,^{+0.52}_{-0.42}$	31 (R) 61 (α)	1969	LBL, Gershwin[24]
	1.23±0.06	**-0.76±0.08**	**Combined**	**Weighted**	**Mean**
$\Lambda \to n\gamma$	1.75±0.15		1816	1993	BNL, Larsen[27]
	1.02±0.33		24	1986	CERN, Biagi[29]
	1.63±0.14		**Combined**	**Weighted**	**Mean**
$\Xi^\circ \to \Lambda\gamma$	1.06±0.12±0.11	0.43±0.44	116 (R) 87 (α)	1990	Fermilab, James[32]
	5±5		1	1974	BNL, [34]
	1.06±0.16	**0.43±0.44**	**Combined**	**Weighted**	**Mean**
$\Xi^\circ \to \Sigma^\circ\gamma$	3.56±0.42±0.10	0.20±0.32±0.05	85	1989	Fermilab, Teige[33]
	<8.0			1988	BNL, Bensinger[40]
	3.56±0.43	**0.20±0.32**	**Combined**	**Weighted**	**Mean**
$\Xi^- \to \Sigma^-\gamma$	0.122±0.023±0.006	1.0±1.3	211	1994	Fermilab, Dubbs[30]
	0.23±0.10		~10	1987	CERN Biagi[31]
	0.128±0.023	**1.0±1.3**	**Combined**	**Weighted**	**Mean**
$\Omega^- \to \Xi^-\gamma$	<0.46			1994	Fermilab, Albuquerque[35]
	<2.2			1984	CERN, Bourquin[41]

chamber[21-23] measurements. Some of the branching fraction measurements[24,25] quote their results as the ratio ($\Sigma^+ \rightarrow p\gamma / \Sigma^+ \rightarrow p\pi^\circ$). We have converted[26] these to absolute branching fractions. Considering the difficulty of these experiments, the measurements are in remarkably good agreement.

The decay rate for $\Lambda^\circ \rightarrow n\gamma$ has been observed by two groups, one working at Brookhaven National Laboratory (BNL) and the other working at CERN. The results from the BNL group[27,28] are contained in two papers, the more recent[27] includes the data from the earlier.[28] Results from the CERN[29] and BNL experiments differ by about 2.0 σ.

The decay $\Xi^- \rightarrow \Sigma^- \gamma$ has now been observed by two groups.[30,31] The more recent measurement[30] has over two hundred events and exhibits a clear signal. Unfortunately, the only asymmetry measurement[30] has a large statistical uncertainty and is only able to provide weak evidence for the sign.

For the Ξ° decays the two modes, $\Xi^\circ \rightarrow \Lambda^\circ \gamma$ and $\Xi^\circ \rightarrow \Sigma^\circ \gamma$ have both been observed.[32-34] The limited statistics (≈ 100 events) in each final state severely limit the precision of the asymmetry and branching fraction measurements. It is very important that the measurements of these branching fractions and asymmetries be repeated with higher precision.

None of the Ω^- radiative decays have been observed although a recent experiment[35] has reduced the limit on the $\Omega^- \rightarrow \Xi^- \gamma$ branching fraction significantly.

3. Hara's Theorem

Weak radiative hyperon decays (WRHD) are a puzzle because of Hara's theorem[42] which states that the parity violating amplitude of the decay $\Sigma^+ \rightarrow p\gamma$ (as well as that of $\Xi^- \rightarrow \Sigma^- \gamma$) should vanish in the limit of SU(3) flavor symmetry. This theorem is crucial to understanding the theoretical implications for the weak radiative hyperon decays. We thus start the theoretical part of this review with the presentation of Hara's theorem.

In his original paper Hara[42] assumed octet dominance of the nonleptonic weak interactions. This assumption is experimentally well verified in all strangeness changing weak decays involving hadrons. It states that the weak interaction Hamiltonian transforms like a member of the octet of flavor SU(3). Contributions from other representations (i.e. the 27-plet) contained in the product $3 \otimes 3 \otimes \bar{3} \otimes \bar{3}$ (describing possible SU(3) transformation properties of the Fermi interactions of four quarks) are assumed negligible. Subsequently, proofs of Hara's theorem have been given by Lo[43] and Gourdin.[44] Insightful comments have also been made by others.[45,46] In the formulation of Gourdin, octet dominance was not used and the requirement of SU(3) symmetry was replaced by the weaker requirement of U-spin symmetry (basically, simple interchange of s and d quarks). Our approach below is similar to that of Gourdin. Later, in Section 6 we shall place Hara's theorem in a more elaborate theoretical framework.

3.1. *Gauge invariance and U-spin arguments*

Let us start our discussion by writing the most general parity violating coupling of photon to hadrons in the standard hadron-level language:

$$\overline{\psi}_1\left[\gamma_\mu F_1(q^2)+q_\mu F_2(q^2)+i\sigma_{\mu\nu}q^\nu F_3(q^2)\right]\gamma_5\psi_2 A^\mu . \tag{1}$$

Invariance under the gauge transformation

$$A^\mu \rightarrow A^\mu + q^\mu \chi \tag{2}$$

requires the vanishing of the additional term

$$\overline{\psi}_1\left[\not{q}F_1 + q^2 F_2\right]\gamma_5\psi_2\chi \tag{3}$$

generated by this transformation.

Consequently (as required also by current conservation), we must have

$$\overline{\psi}_1\left[(m_1+m_2)F_1 + q^2 F_2\right]\gamma_5\psi_2 = 0. \tag{4}$$

From (4) it follows that

$$F_1 = -\frac{q^2}{m_1+m_2}F_2. \tag{5}$$

Since F_2 cannot have a pole at $q^2=0$ (no massless hadrons exist), F_1 must vanish at $q^2=0$, i.e., for real photons.

For real photons $\left(q^2 = q^\mu \cdot \varepsilon_\mu = 0\right)$, only the third term, $F_3(0)$, in (1) may be nonvanishing. Since weak radiative hyperon decays are CP-conserving processes, we must deduce what restrictions the requirement of CP-invariance imposes upon this term. Under the operations of charge conjugation we have

$$\overline{\psi}_1 i\sigma_{\mu\nu}\gamma_5\psi_2 q^\mu \xrightarrow[C]{} -\overline{\psi}_2 i\sigma_{\mu\nu}\gamma_5\psi_1 q^\mu \tag{6}$$

corresponding to C-parity equal to -1 for the diagonal term

$$\overline{\psi}_1 i\sigma_{\mu\nu}\gamma_5\psi_1 q^\mu \tag{7}$$

describing the situation when the incoming and outgoing baryons are identical ($2\rightarrow 1$). Since the parity of expression (7) is +1, it cannot be coupled to the photon $A^\nu\left(C_\gamma = P_\gamma = -1\right)$ if CP is to be conserved.

In general, however, the incoming and outgoing particles are not identical, as in (1), and one can write an expression that is *antisymmetric* under the interchange of particles labels ($1 \rightleftarrows 2$) in the initial (and final) state

$$\overline{\psi}_1 i\sigma_{\mu\nu}\gamma_5\psi_2 q^{\mu} - \overline{\psi}_2 i\sigma_{\mu\nu}\gamma_5\psi_1 q^{\mu} \tag{8}$$

and which goes into itself (with + sign) under the operation (6) of charge conjugation. Thus, expression (8) has *even* C-parity and, consequently, a coupling of the form

$$\left[\overline{\psi}_1 i\sigma_{\mu\nu}\gamma_5\psi_2 - \overline{\psi}_2 i\sigma_{\mu\nu}\gamma_5\psi_1\right] q^{\mu} A^{\nu} \tag{9}$$

is permitted by CP-conservation.

For $1 \to p$, $2 \to \Sigma^+$ we get the term considered by Hara:

$$\left[\overline{\psi}_p i\sigma_{\mu\nu}\gamma_5\psi_{\Sigma^+} - \overline{\psi}_{\Sigma^+} i\sigma_{\mu\nu}\gamma_5\psi_p\right] q^{\mu} A^{\nu} \tag{10}$$

as the only nonvanishing parity violating $\Sigma^+ p\gamma$ coupling permitted in the standard hadron level language.

Hara's theorem immediately follows from (10) as explained below. Indeed, since the weak $\Delta S = 1$ Hamiltonian is symmetric under the $s \rightleftarrows d$ interchange one concludes that only the $s \rightleftarrows d$ symmetric part of (10) may be non zero. However, $s \rightleftarrows d$ corresponds to $\Sigma^+(uus) \rightleftarrows p(uud)$, and under the $\Sigma^+ \rightleftarrows p$ interchange, expression (10) is *anti*-symmetric. Consequently, the $s \rightleftarrows d$ symmetric part of (10) is zero. Thus, in a U-spin symmetric world the $\Sigma^+ \to p\gamma$ parity violating amplitude should vanish. Similar considerations apply to the $\Xi^- \to \Sigma^-\gamma$ process since, under the $s \rightleftarrows d$ interchange, $\Xi^-(ssd) \rightleftarrows \Sigma^-(dds)$.

3.2. *Abandoning exact SU(3)*

In the real world, the strange quark is heavier than the down quark. Consequently, with U-spin symmetry broken, nonvanishing parity violating amplitudes (and therefore also asymmetries) are expected for the $\Sigma^+ \to p\gamma$ and $\Xi^- \to \Sigma^-\gamma$ processes.

This situation has been discussed by Vasanti.[47] To get a prediction for the sign of the resulting asymmetry let us consider his argument for the effective $s \to d\gamma$ transition. This transition is described by

$$M \propto \bar{d}\sigma_{\mu\nu}\left(a + b\gamma_5\right)s\, q^{\mu} A^{\nu}. \tag{11}$$

In Eq. (11) a and b depend on the masses m_s, m_d and b must vanish for $m_s = m_d$ as required by Hara's theorem.

Since the theory is invariant under the following transformations of fields and masses:

$$
\begin{aligned}
q &\to -\gamma_5 q \\
\bar{q} &\to \bar{q}\gamma_5 \\
m_q &\to -m_q
\end{aligned}
\tag{12}
$$

corresponding to

$$
\begin{aligned}
m_q\bar{q}q &\to \left(-m_q\right)\bar{q}\gamma_5\left(-\gamma_5 q\right) = m_q\bar{q}q \\
\bar{q}'\gamma_\mu\left(1-\gamma_5\right)q &\to \bar{q}'\gamma_\mu\left(1-\gamma_5\right)\left(-\gamma_5 q\right) = \bar{q}'\gamma_\mu\left(1-\gamma_5\right)q
\end{aligned}
$$

etc.
and (12) holds separately for each flavor we may perform transformation (12) in Eq. (11) for the strange quark *only*. Assuming for the moment that a and b are odd in quark masses:

$$
\begin{aligned}
a &= \varsigma m_d + \beta m_s \\
b &= \gamma m_d + \delta m_s
\end{aligned}
\tag{13}
$$

(and $\varsigma, \beta, \gamma, \delta$ are even)

we obtain then

$$
\begin{aligned}
&\bar{d}\sigma_{\mu\nu}\left[\left(\varsigma m_d + \beta m_s\right) + \left(\gamma m_d + \delta m_s\right)\gamma_5\right]s \to \\
&\to \bar{d}\sigma_{\mu\nu}\left[-\gamma m_d + \delta m_s + \left(-\varsigma m_d + \beta m_s\right)\gamma_5\right]s.
\end{aligned}
\tag{14}
$$

Invariance under (12) requires then $\varsigma = -\gamma, \beta = \delta$ and Hara's theorem itself (i.e., vanishing of b for $m_s = m_d$) fixes $\beta = \varsigma$.

As far as the terms even in quark masses are concerned, one can similarly show that they must vanish. To this end one has to consider separately two transformations: $s \to -\gamma_5 s,\ m_s \to -m_s$ and $\bar{d} \to \bar{d}\gamma_5, m_d \to -m_d$.

Thus, from (14) and (11) we obtain

$$
M \;\propto\; \bar{d}\sigma_{\mu\nu}\left(1 + \frac{m_s - m_d}{m_s + m_d}\gamma_5\right)s\, q^\mu A^\nu. \tag{15}
$$

From (15) it follows that the asymmetry parameter is positive:

$$
\alpha = \frac{m_s^2 - m_d^2}{m_s^2 + m_d^2} \tag{16}
$$

and close to +1 if current quark masses are used $\left(m_s \gg m_d\right)$. If, on the other hand, one uses constituent quark masses ($m_s \approx 500$ MeV, $m_d \approx 330$ MeV) one gets α around +0.4 to +0.5.

One might hope that a similar argument may be applied to all hadron level transitions $B_1 \to B_2 \gamma$ when baryons B_1, B_2 are members of a U-spin doublet (as are s and d in (11)) That is, one would expect positive asymmetries for the $\Sigma^+ \to p\gamma$ and $\Xi^- \to \Sigma^- \gamma$ decays. As will be discussed at length in this review this expectation is not confirmed when one uses the quark model perscription for the structure of baryons. Yet, if single-quark transitions are dominant, the argument of Vasanti[47] is valid and leads to positive asymmetries for *all* weak radiative decays. If these transitions are not dominant, the above arguments still apply to the $\Xi^- \to \Sigma^- \gamma$ decay since single-quark transitions are the only ones that may contribute to this decay (Section 7).

3.3. *The controversy*

It came, therefore, as a great surprise[48,49] when the first measurements[24,25] indicated a *large negative* asymmetry in the $\Sigma^+ \to p\gamma$ decay. As discussed in Section 2 the most recent high statistics experiment performed at Fermilab[8] confirms these findings and leaves no doubts as to the sign and size of the $\Sigma^+ \to p\gamma$ asymmetry:

$$\alpha\left(\Sigma^+ \to p\gamma\right) = -0.72 \pm 0.086 \pm 0.045 \tag{17}$$

where the quoted errors are statistical and systematic respectively. Thus, standard hadron-level arguments appear to be at gross variance with experiment.

The theoretical situation became muddled in 1983 when Kamal and Riazuddin[50] (KR) reconsidered the question within the framework of the quark model. The astonishing result of their simple, explicitly gauge-invariant calculation was that in the quark model, Hara's theorem is not satisfied in the SU(3) limit. Since the remaining assumptions upon which this theorem rests seem to be unshakable, their result has been considered by some workers as revealing a kind of pathology of the quark model.

Others, nonetheless, tried to find a place for it in the existing hadron-level formalism. We come back to the KR paper in Section 7 where we discuss if (and how) it is possible to fit this paper into the existing standard hadron-level theoretical framework of Section 6 as well as which of the assumptions of Hara's theorem appears to conflict with the quark model.

A resolution of the problem has been proposed[51] but may be regarded by some as itself controversial. It is therefore of paramount importance to have a sound experimental input against which theoretical ideas may be tested. We shall review the actual status of our experimental knowledge on weak radiative hyperon decays in Section 5. Before embarking on a tour of the experimental side of the studies of weak radiative hyperon decays, in the next brief section we shall present lower bounds on the WRHD branching fractions. These bounds are imposed by unitarity and therefore they are very reliable.

4. Unitarity Bounds

Presentation and discussion of the predictions of specific models of WRHD will be given in Sections 7-9. Here we gather the most important and essentially model-independent lower bounds on the branching fractions of WRHD, that follow from unitarity. These bounds result from the nonvanishing of the contribution of πB intermediate states as shown in Fig. 4.1

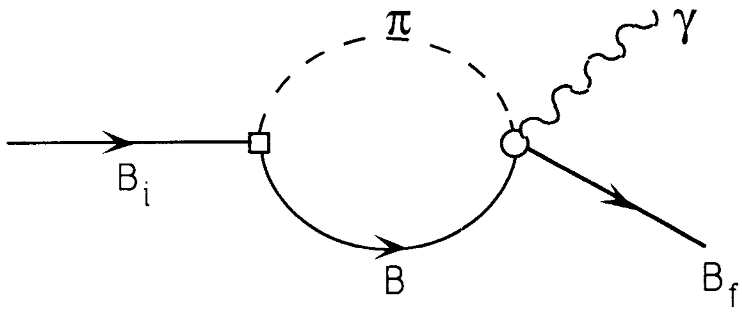

Fig. 4.1. Unitarity-induced contribution of the πB intermediate state to the WRHD $B_i \to B_f \gamma$ (□-weak nonleptonic decay; O-pion photoproduction).

The first estimate of this contribution has been made by Zakharov and Kaidalov.[52] In their paper they considered $\Sigma^+ \to p\gamma$, $\Lambda \to n\gamma$ and $\Xi^- \to \Sigma^-\gamma$ decays.

For the $\Sigma^+ \to p\gamma$ decay, Im $M(\Sigma^+ \to p\gamma)$ can be expressed in terms of the amplitudes of the $\Sigma \to N\pi$ nonleptonic decays and of those of pion photoproduction on nucleons $(\gamma p \to \pi^+ n)$. Using the results of a phase-shift analysis of the photoproduction of pions on protons Zakharov and Kaidalov[52] estimated that the branching fraction R for the $\Sigma^+ \to p\gamma$ decay satisfies

$$R(\Sigma^+ \to p\gamma) \geq (0.69 \pm 0.40) * 10^{-4}. \tag{1}$$

The above number corresponds to case (1) of Zakharov and Kaidalov (i.e., to the domination of the p-wave in the decay $\Sigma^+ \to n\pi^+$, as it has been experimentally determined after their publication). For the $\Lambda \to n\gamma$ and $\Xi^- \to \Sigma^-\gamma$ decays the necessary experimental input in the form of the relevant phase-shift analyses was not available. Using perturbation theory estimates of the s- wave pion photoproduction amplitudes (the s-wave constitutes the dominant amplitude in relevant nonleptonic decays), Zakharov and Kaidalov concluded that the following lower bounds for the $\Lambda \to n\gamma$ and $\Xi^- \to \Sigma^-\gamma$ branching fractions should hold:

$$R(\Lambda \to n\gamma) \geq 0.83 \times 10^{-3} \quad (2)$$

$$R(\Xi^- \to \Sigma^-\gamma) \geq 0.13 \times 10^{-3}. \quad (3)$$

The same number for the lower bound on the $\Lambda \to n\gamma$ branching fraction has been independently obtained by Farrar.[53] Adding an estimate of the real part she concluded that

$$R(\Lambda \to n\gamma) \approx (1.9 \pm 0.8) \times 10^{-3} \quad (4)$$

and that the corresponding asymmetry is likely to be positive.

For the $\Sigma^+ \to p\gamma$ lower bound Farrar found a value smaller by an order of magnitude from the one given by Zakharov and Kaidalov[52] in Eq.1.

An independent estimate of the branching fraction for the $\Xi^- \to \Sigma^-\gamma$ decay has been made by Kogan and Shifman.[54] Their calculation of the diagram of Fig. 4.1 gives

$$R(\Xi^- \to \Sigma^-\gamma) \geq 0.10 \times 10^{-3}. \quad (5)$$

Taking the real part into account they estimate

$$R(\Xi^- \to \Sigma^-\gamma) \geq 0.17 \times 10^{-3}. \quad (6)$$

Finally, a thorough study of the contribution of πN intermediate states to the weak radiative decays of Σ and Λ hyperons has been carried out by Reid and Trofimenkoff.[55,56] Their approach contains some technical and phenomenological improvements over that of Farrar.[53] Reid and Trofimenkoff[55] also contain references to earlier papers on the contribution of the πB intermediate states.

We have gathered all these lower bounds determined by the imaginary parts of the amplitudes corresponding to Fig. 4.1 and the full predictions (which include, fairly uncertain, estimates of real parts) in Table 4.1.

Table 4.1 Comparison of estimates of πB contributions to the branching fractions of WRHD. (in units of 10^{-3})

	Zakharov[52]	Farrar[53]		Kogan[54]		Reid[55,56]
Process	**lower bound**	**lower bound**	**full estimate**	**lower bound**	**full estimate**	**full estimate**
$\Sigma^+ \to p\gamma$	0.07±0.04	0.007	0.3±1.2			$0.77^{+1.29}_{-0.49}$
$\Lambda \to n\gamma$	0.83	0.85	1.9±0.8			$1.20^{+0.46}_{-0.04}$
$\Xi^- \to \Sigma^-\gamma$	0.13			0.10	0.17	
$\Omega^- \to \Xi^-\gamma$				0.008	0.01	

5. Specific Measurements and Techniques

5.1 $\Sigma^+ \to p\gamma$

Measured for the first time twenty five years ago, the large negative asymmetry in this reaction spurred further work in WRHDs. Figure 5.1 shows the history of this asymmetry measurement. In these plots we have combined in quadrature the statistical and systematic uncertainties for each of the measurements.

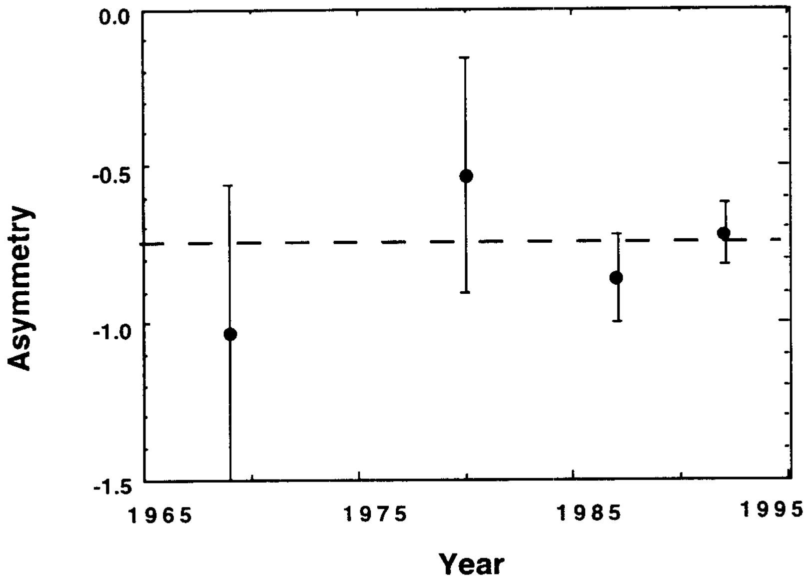

Figure 5.1 History of measurement of $\Sigma^+ \to p\gamma$ asymmetry parameter

From Fig. 5.1 we see a steady reduction in the uncertainties of the individual measurements. Combining these measurements we find for the $\Sigma^+ \to p\gamma$ asymmetry parameter, α=-0.76±0.08. This is shown on Fig 5.1 by the dashed line. Crucial to these experiments is the ability to produce a Σ^+ with well known and controllable polarization. The primary measurement is of the decay asymmetry which is the product of the polarization and the intrinsic asymmetry parameter, α. Knowledge of the polarization comes from the measurement of a decay which has a known α parameter in the same beam or from a reliance on some other method of the determination of the polarization. Knowledge of the production polarization through known phase shifts has provided this for the low energy experiments.

These measurements utilize a variety of techniques to produce the polarized Σ^+ needed for asymmetry measurement which is illustrated in Table 5.1. The fact that the polarizations are derived from different reactions and are of differing magnitudes but that the experiments still give consistent values of α lead one to the inescapable conclusion that α is large and negative in the decay $\Sigma^+ \to p\gamma$.

Table 5.1 Properties of $\Sigma^+ \to p\gamma$ asymmetry experiments

Experiment	**Laboratory**	**Reaction**	**Σ^+ Momentum GeV/c**	**Polarization %**
Foucher[8]	Fermilab	$p\ Cu \to X\ \Sigma^+$	375.	12.
Kobayashi[38]	KEK	$\pi^+ p \to K^+ \Sigma^+$	1.7	87.
Manz[25]	CERN	$\pi^- p \to K^- \Sigma^+$	0.42-0.50	≈10-90
Gershwin[24]	LBL	$K^- p \to \pi^- \Sigma^+$	0.5	40.

The branching fraction measurements are plotted in Fig. 5.2 and give a similarly consistent picture. Again these experiments use different techniques and have differing systematic uncertainties. Their weighted mean and standard deviation is

$$(\Sigma^+ \to p\gamma)/(\Sigma^+ \to \text{all}) = (1.23 \pm 0.06) * 10^{-3}$$

and is represented by a dashed line in Fig. 5.2.

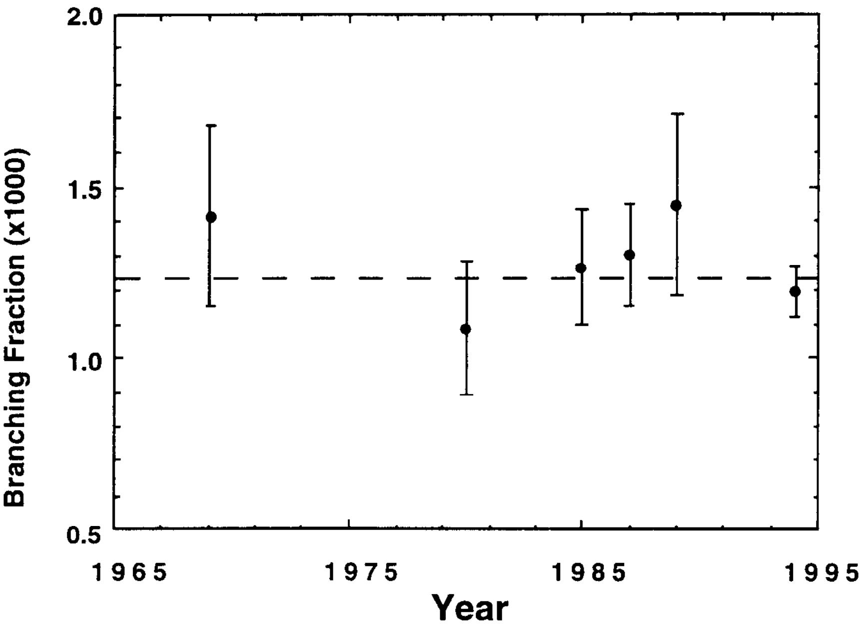

Figure 5.2 History of measurement of $\Sigma^+ \to p\gamma$ branching fraction

Needed in these experiments is both the ability to produce sufficient samples of polarized hyperons as well as careful control of systematic uncertainties. We illustrate how this is done by looking in some detail at one of these experiments.

The experiment of Foucher et. al.[8] is shown in Fig. 5.3. In this classic high energy charged hyperon beam experiment, the Σ^+ are produced by 800 GeV protons incident on a small Cu target at the entrance of a large "hyperon" magnet. The latter serves as a magnetic channel selecting particles within a narrow momentum and angle range thus defining the transverse momentum, p_t, and Feynman x, x_f, of the produced hyperons. Reversing the sign of the targeting angle reverses the sign of the hyperon polarization. Since this can be done by changing currents in magnets upstream of the "hyperon" magnet, the resolutions and backgrounds in the spectrometers are not affected. This is a powerful technique for controlling systematic uncertainties.

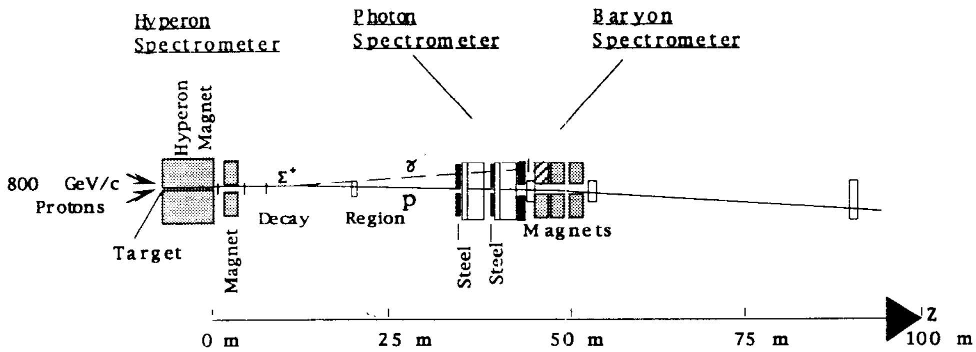

Fig. 5.3 $\Sigma^+ \rightarrow p\gamma$ Apparatus of Foucher et. al.[8]

The charged beam had a mean momentum of 375 GeV/c and provided a large flux ($\approx$2000 Σ^+ per second) of Σ^+ at the decay region indicated in Figure 5.3. The momentum and direction of the beam particles were measured by the magnets and detectors of the hyperon spectrometer. The decay products of the Σ^+ were measured by the photon and baryon spectrometers.

High spatial resolution detectors in the hyperon and baryon spectrometers of Fig. 5.3 allowed excellent mass resolution. The required trigger was simple in that it only required the conversion of a neutral photon into a charged electromagnetic shower in a set of steel plates. This means that Σ^+ decaying through the $\Sigma^+ \rightarrow p\gamma$ were recorded at the same time as $\Sigma^+ \rightarrow p\pi^\circ$ decays, thus providing a measurement of the beam polarization from the well known decay properties of the $\Sigma^+ \rightarrow p\pi^\circ$ Fig. 5.4 shows a mass squared distribution (M_x^2) of the missing neutral particle (X°) for the hypothesis $\Sigma^+ \rightarrow pX^\circ$ where we assume the hyperon track is a Σ^+ and the baryon track is a proton.

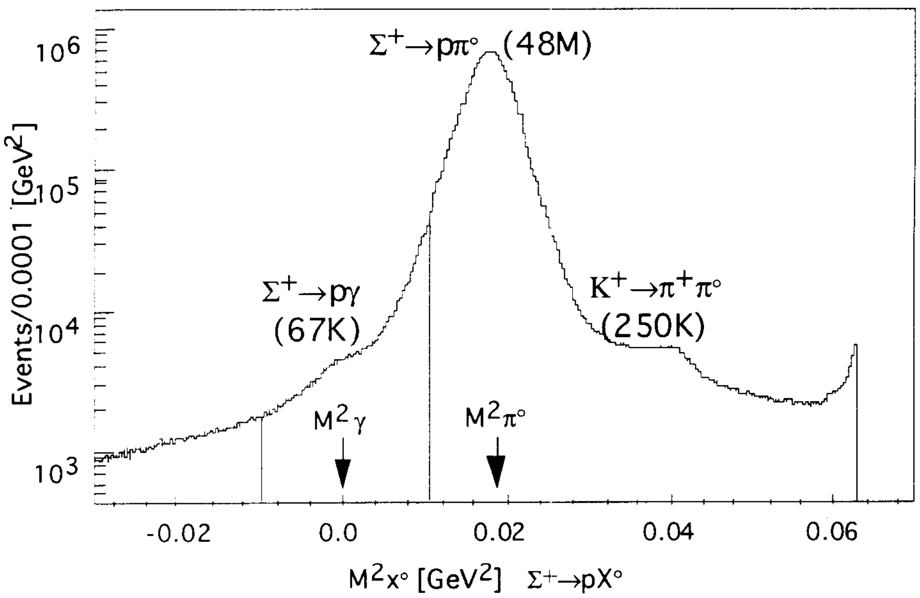

Fig. 5.4 Event distributions of the mass squared of the missing neutral particle (X°) for the hypothesis $\Sigma^+ \to p X^\circ$

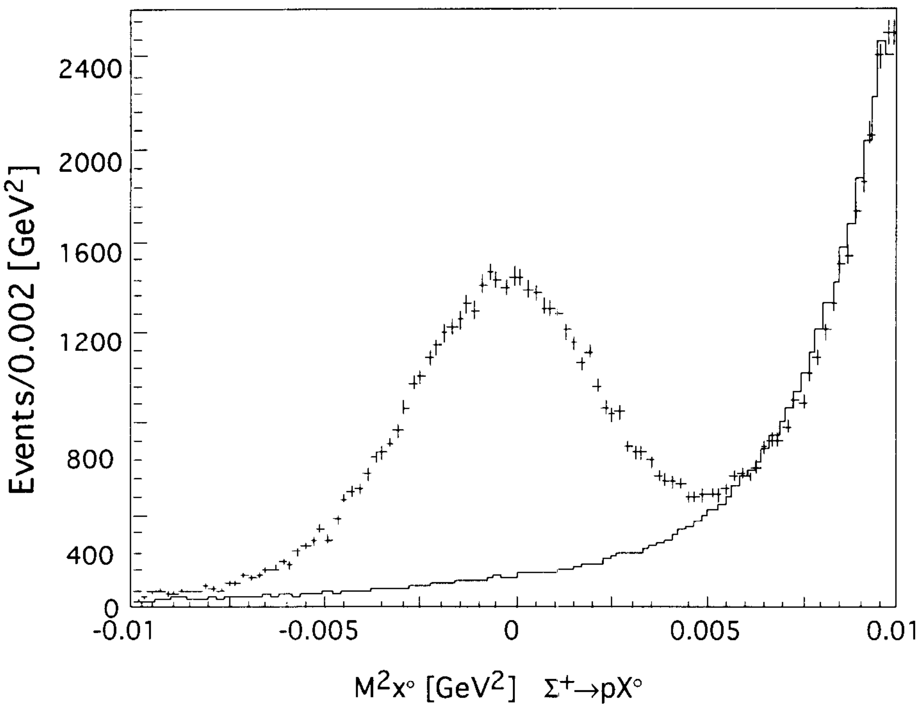

Fig. 5.5 Same as Fig. 5.4 after selection criteria are imposed. Note that one can estimate the background from the $\Sigma^+ \to p X^\circ$ peak.

Note the size of the event sample, the small shoulder corresponding to the radiative decay, and the peak corresponding to the decay of beam kaons. Imposition of selection criteria on the energy and angle of the neutral particle[8] results in the event distribution of Fig. 5.5. Here the radiative decays are clearly seen above a relatively small background. The larger event sample and the ability to change the direction of polarization allowed this experiment to compete favorably with Kobayashi et. al[38] even though their Σ^+ polarization was much larger as indicated in Table 5.1.

The measurements of Kobayashi et. al[38] and Foucher et.al[8] used very different experimental techniques. Not only were their Σ^+ produced with different energies and polarizations, but also different methods of identifying the radiative decays were employed. Yet the fact that both experiments give similar and unambiguous results as shown in Figures 5.1 and 5.2 should reassure the reader that there are no hidden sources of systematic uncertainty. The $\Sigma^+ \rightarrow p\,\gamma$ branching fraction and asymmetry parameter are the most precisely measured of any of the radiative decays. There is no way of escaping the fact that the asymmetry is *large* and *negative*. Statistically, it is almost *ten standard deviations from zero*.

By reversing the currents in the magnets shown in the experiment of Fig. 5.5, a measurement[16,57] has been made of a WRHD of an antibaryon, $\overline{\Sigma}^- \rightarrow \bar{p}\gamma$. Its measured decay parameters are consistent with CPT invariance.

5.2 $\Lambda^\circ \rightarrow n\gamma$

This decay presents special problems to the experimenter since both the initial and final states are neutral. Although the asymmetry has not been measured so far, two measurements have been made of the branching fraction.[27,29] The first measurement[29] utilized Λ° from the decay $\Xi^- \rightarrow \Lambda^\circ\pi^-$ in the CERN charged hyperon beam. The momentum and the direction of the Λ° were determined from the momentum and direction of the Ξ^- and π^-. Although the Λ° resulting from the Ξ^- decay are polarized, the small event sample from this experiment (31 events) allowed for a measurement of the branching fraction only.

The second experiment[3,27,28] utilized a very different technique. A stopping beam of K^- produces Λ°s through the reaction $K^- p \rightarrow \Lambda^\circ \pi^\circ$. Measurement of the energy and direction of the two photons from the π° decay fixes the kinematics of the Λ°. In this case the Λ° is unpolarized. As can be seen from Table 2.1 the two experiments are in poor agreement differing by about $2\,\sigma$. High intensity charged hyperon beams are available at Fermilab which could produce large fluxes of polarized Λ° from Ξ^- decays. Definitive measurements of both the branching fraction and asymmetry could be made at Fermilab.

5.3 $\Xi^\circ \to \Lambda^\circ \gamma$ and $\Xi^\circ \to \Sigma^\circ \gamma$

The identification of these all neutral topologies relies on the observation of the decay $\Lambda^\circ \to p\pi^-$ for the $\Xi^\circ \to \Lambda^\circ \gamma$ or the electromagnetic decay $\Sigma^\circ \to \Lambda^\circ \gamma$ for the $\Xi^\circ \to \Sigma^\circ \gamma$. Both of these Ξ° WRHDs have been measured[32,33] in Proton Center neutral beams at Fermilab. The geometry of these experiments was similar and is illustrated in Fig. 5.6. In each case a high energy proton beam impinged on a small target. A large high field magnet served to deflect charged particles and produce a collimated neutral beam containing Ξ° as well as more copious amounts of Λ°, K°, and neutrons. The Ξ° component can be identified because it is the only source of Λ° which do not originate in the target. (Since the Σ° lifetime is very short, Λ° produced by its decay appear to come from the target.) Because the Ξ° were inclusively produced, one does not have direct measurement of their momenta. However, from the reconstruction of the direction and momenta of the Λ° from its decay $\Lambda^\circ \to p\pi^-$, one can combine this with photons in the same event to determine the Ξ° direction and momentum.

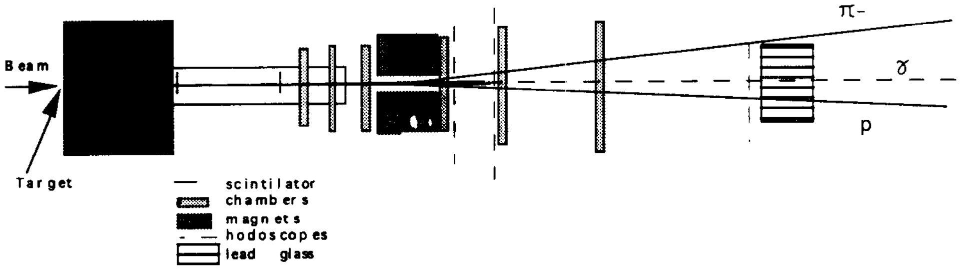

Fig. 5.6 Simplified version of the apparatus of Teige et al.[33]

We note that these experiments were done as subsidiary measurements in existing experimental configurations and each contained less than 100 events. These pioneering measurements demonstrated the versatility of high energy neutral hyperon beams in extracting the parameters of the Ξ° WRHDs. Enhanced statistical precision is clearly needed, however. This is particularly important for the determination of the asymmetry parameters since at this time one is not even sure of their signs.

The Fermilab neutral kaon facility now under construction may offer the best possibility for new information on the Ξ° WRHDs.[58] This experiment has as its primary goal the measurement of CP violating parameters in the kaon system and has excellent photon detection capabilities. Utilization of the Ξ° component of this beam has the capabilities of increasing the statistics of the Ξ° WRHDs by one to two orders of magnitude.

5.4 $\Xi^- \to \Sigma^- \gamma$

This decay presents experimental challenges because of its small branching ratio ($\approx 10^{-4}$), the small polarization of the Ξ^- ($\approx 10\%$), and the need to identify the Σ^- either directly or through it major decay mode $\Sigma^- \to n\pi^-$ in order to suppress backgrounds. The branching fraction was first measured in the CERN hyperon beam[31] with a sample of 11 events. More recently the Fermilab group[30] with about 200 events was able to improve on its value as well as present weak evidence that the asymmetry parameter is positive. The measurement of the asymmetry parameter is of particular importance since this is the most accessible WRHD that cannot proceed by a two quark diagram (Fig. 7.1). Consequently, an improved measurement would help shed light on the other processes.

5.5 $\Omega^- \to \Xi^- \gamma$ and $\Omega^- \to \Xi^*(1530)\, \gamma$

Neither of these decays have been seen. The considerably lower fluxes of Ω^- in hyperon beams (compared to Σ^- and Ξ^-) coupled with the observation that Ω^- are produced unpolarized[59,60] make these branching fractions and asymmetry parameters particularly difficult to measure. Tertiary beams of polarized Ω^- have been produced and used to measure the Ω^- magnetic moment.[61] However since they involve using polarized hyperon interactions to produce polarized Ω^- in a spin transfer mechanism, there is a further reduction in Ω^- beam rate. While a new experiment might be expected to push the $\Omega^- \to \Xi^- \gamma$ branching fraction to a level where it might be seen, a measurement of the asymmetry is not on the near horizon.

6. General Theoretical Framework

Great interest in weak radiative hyperon decays was stimulated both by the apparent disagreement between Hara's theorem and experiment and by the argument that nonetheless these hyperon decays should appear simpler and more susceptible to theoretical description than the nonleptonic ones. In the latter case the presence of two strongly interacting particles in the final state requires consideration of all complications due to final-state strong interactions while in WRHD one of the two outgoing particles is a strong-interaction-blind photon. Thus, final-state strong interactions appear to be absent in WRHD whose description may consequently be expected to be less dependent upon unknown details of strong interaction dynamics.

However, as the problems with Hara's theorem indicate, this expectation is misleading. Proper description of weak radiative hyperon decays is, most probably, at least as difficult as that of nonleptonic ones where there is no consensus as to the relative size (and sometimes also sign) of the contributions from various physical mechanisms. The standard general theoretical framework used in the description of nonleptonic decays is not disputed, however. Most attempts at a description of WRHD fit into a similar standard framework. These two frameworks constitute two parts of a single theoretical scheme that unites weak couplings of pseudoscalar and vector particles to baryons. This general theoretical scheme has been known for years.[62] It has been reviewed recently anew in an updated form[63-65] which takes into account the development of the quark model in the intervening years. In this section, we describe this general scheme of the weak couplings of pseudoscalar and vector particles to baryons. We shall consider here the requirements imposed by the conditions of gauge invariance only very briefly. Their further discussion is shifted to appropriate sections of this review where we show how calculations of various papers fit into the general scheme presented here.

Let us therefore consider the couplings $B_i B_f M$ (where $B_{i(f)}$ denotes the initial (final) baryon and M is pseudoscalar (P) or vector (V) particle) in the presence of weak interactions.

In general, the weak interaction Hamiltonian may act in any one of the three legs of the $B_i B_f M$ coupling as shown in Fig. 6.1.

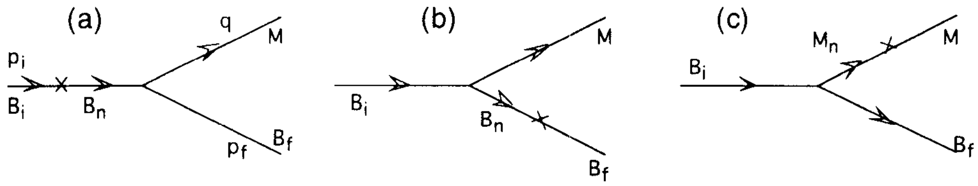

Fig. 6.1abc. Diagrams for the weak $B_i B_f M$ coupling.
The cross denotes the action of the weak Hamiltonian.

We postpone the discussion of the boson-leg contribution (c) for the moment and focus on diagrams (a) and (b). To stress similarities between the weak couplings of pseudoscalar and vector particles, we shall consider both of them alongside each other starting from a fairly extensive discussion of the troublesome parity violating amplitudes. This will be followed by a brief presentation of the standard approach to the parity conserving amplitudes.

6.1 *The contribution of baryon poles*

6.1.1 *Parity violating amplitudes*

Let us consider the action of the parity violating part $H^{p.v.}$ of the weak Hamiltonian in the baryon legs of Fig. 6.1. The intermediate states B_n may be the ground $B_n\left(\tfrac{1}{2}^+\right)$ and the excited $B_{n^*}\left(\tfrac{1}{2}^-\right)$, $B_{n^*}\left(\tfrac{1}{2}^+\right)$,... baryon states. The Lee-Swift theorem[66] requires that in the SU(3) symmetry limit the matrix elements of $H^{p.v.}$ between the ground states vanish. Therefore, the contribution of the intermediate ground states $B_n\left(\tfrac{1}{2}^+\right)$ is generally neglected. The dominant contribution is then expected to come from the excited $B_{n^*}\left(\tfrac{1}{2}^-\right)$ baryons.

6.1.1.1 *Pseudoscalar particles – nonleptonic decays*

General expressions for the parity-violating (s-wave) $B_i \to B_f M$ nonleptonic hyperon decay (NLHD) amplitudes are

$$
\begin{aligned}
&M^{p.v.}\left(B_i \to B_f P\right) = A\bar{u}_f u_i \\
&M^{p.v.}\left(B_i \to B_f V\right) = \varepsilon^{*\mu}\bar{u}_f\left[A_1\gamma_\mu + A_2 q_\mu + A_3 i\sigma_{\mu\nu}q_\nu\right]\gamma^5 u_i \\
&\left(\sigma_{\mu\nu} = \tfrac{i}{2}\left[\gamma_\mu \cdot \gamma_\nu\right]\right).
\end{aligned}
\tag{1}
$$

It is straightforward to show that for the pseudoscalar mesons, the excited $B_{n^*}\left(\tfrac{1}{2}^-\right)$ baryons contribute to the A amplitude as follows

$$
A = \sum_{B_{n^*}\left(\frac{1}{2}^-\right)}\left[\frac{g_{B_f B_{n^*} P}\, b_{n^* i}}{m_i - m_{n^*}} + \frac{b_{fn^*}\, g_{B_{n^*} B_i P}}{m_f - m_{n^*}}\right] \tag{2}
$$

where

$$
\begin{aligned}
&\left\langle B_{n^*}\left(\frac{1}{2}^-\right)\middle| H^{p.v.} \middle| B_i\left(\frac{1}{2}^+\right)\right\rangle = b_{n^* i}\bar{u}_{n^*}u_i \\
&\left\langle B_f\left(\frac{1}{2}^+\right)\middle| H^{p.v.} \middle| B_{n^*}\left(\frac{1}{2}^-\right)\right\rangle = b_{fn^*}\bar{u}_f u_{n^*}
\end{aligned}
\tag{3}
$$

with $b_{in^*} = b_{n^* i}$ (from hermiticity and CP invariance). In Eq. 2. $g_{BB^*P}\left(g_{B^*BP}\right)$ are strong (parity conserving) couplings of pseudoscalar mesons to the $\left(\tfrac{1}{2}^+, \tfrac{1}{2}^-\right)$ pair of baryons

$$
\begin{aligned}
&M^{Strong}\left(B_i \to B_{n^*}P\right) = g_{B_{n^*}B_i P}\bar{u}_{n^*}u_i \\
&M^{Strong}\left(B_{n^*} \to B_f P\right) = g_{B_f B_{n^*} P}\bar{u}_f u_{n^*}.
\end{aligned}
\tag{4}
$$

The standard current algebra form[67] is obtained from (2) through the use of generalized Goldberger-Treiman relations[65,68]

$$g_{B^*BP} = \frac{\sqrt{2}}{f_P}\left(m_{B^*} - m_B\right)g_A^{B^*B}$$
$$= g_{BB^*P} = \frac{\sqrt{2}}{f_P}\left(m_B - m_{B^*}\right)g_A^{BB^*} \tag{5}$$

where f_P is the decay constant of the pseudoscalar meson P and $g_A^{B^*B} = -g_A^{BB^*}$ is the axial vector coupling constant.

In the soft meson limit $\left(q^\mu \to 0,\ m_i \to m_f\right)$ the use of Eq. 5 reduces Eq. 2 to the standard commutator relation of the current algebra approximation A^{CA} to A:

$$A^{CA} = \lim_{q \to 0} A = \frac{\sqrt{2}}{f_{pa}} \sum_{B_{n^*}\left(\frac{1}{2}^-\right)} \left(g_{Aa}^{B_f B_{n^*}} b_{n^*i} - b_{fn^*} g_{Aa}^{B_{n^*}B_i} \right)$$
$$= +\frac{\sqrt{2}}{f_{pa}} \left\langle B_f \left| \left[Q_5^a , H^{p.v.} \right] \right| B_i \right\rangle \tag{6}$$

where Q_5^a is the axial charge. Away from the soft meson limit, we have

$$A = A^{CA} + \frac{\sqrt{2}}{f_{pa}}\left(m_f - m_i\right) \sum_{B_{n^*}\left(\frac{1}{2}^-\right)} \left[\frac{g_A^{B_f B_{n^*}} b_{n^*i}}{m_i - m_{n^*}} + \frac{b_{fn^*} g_A^{B_{n^*}B_i}}{m_f - m_{n^*}} \right] \tag{7}$$

where the second term describes this part of the contribution from the excited intermediate states $B_{n^*}\left(\tfrac{1}{2}^-\right)$ that vanishes in the $q^\mu \to 0\left(m_i = m_f\right)$ limit and therefore cannot be absorbed into the standard commutator term of current algebra. The advantage of current algebra approach over that of the pole model appears in the limit of exact SU(3), when the second term in Eq. 7 vanishes and, consequently, no information on the $\tfrac{1}{2}^-$ poles is needed.

For further discussions of the relationship of the general scheme to the quark model calculations we need to establish a connection (if any) between the above considerations and the quark model. That such a connection exists has been observed by Körner and Gudehus.[69] In 1979 Körner, Kramer, and Willrodt[70] proved that the soft meson approach and the quark model are totally equivalent in a group theoretical sense. The question has been discussed also by Desplanques, Donoghue, and Holstein.[64] We shall come to the questions of dynamics in the quark model after completing our presentation of the standard general scheme.

6.1.1.2 *Vector particles – radiative decays*

For the vector particles, we have two possible (vector and tensor) strong (parity-conserving) couplings of vector mesons to the $\left(\tfrac{1}{2}^{+},\tfrac{1}{2}^{-}\right)$ pair of baryons: $\left[\text{with } \varepsilon^{*\mu}q_{\mu}=0\right]$.

$$
\begin{aligned}
M^{Strong}\left(B_i \to B_{n^*}V^{\mu}\right) &= \varepsilon^{*\mu}\bar{u}_{n^*}\left[g_{B_{n^*}B_iV}\gamma_{\mu} + f_{B_{n^*}B_iV}i\sigma_{\mu\nu}q_{\nu}\right]\gamma_5 u_i \\
M^{Strong}\left(B_{n^*} \to B_f V^{\mu}\right) &= \varepsilon^{*\mu}\bar{u}_f\left[g_{B_fB_{n^*}V}\gamma_{\mu} - f_{B_fB_{n^*}V}i\sigma_{\mu\nu}q_{\nu}\right]\gamma_5 u_{n^*}
\end{aligned}
\tag{8}
$$

with

$$
\begin{aligned}
g_{B_{n^*}B_iV} &= g_{B_iB_{n^*}V} \\
f_{B_{n^*}B_iV} &= f_{B_iB_{n^*}V}.
\end{aligned}
\tag{9}
$$

From Eq. 3 and 8 it follows that: (1) the contribution from the vector coupling g_{B^*BV} to the parity violating $B_i \to B_f V$ amplitude is

$$
\sum_{B_{n^*}\left(\frac{1}{2}^-\right)}\left[\frac{b_{fn^*}g_{B_{n^*}B_iV}}{m_f - m_{n^*}} + \frac{g_{B_fB_{n^*}V}b_{n^*i}}{m_i - m_{n^*}}\right]\varepsilon^{*\mu}\bar{u}_f\gamma_{\mu}\gamma_5 u_i \tag{10}
$$

and thus it determines A_1 in Eq. 1 while, (2) the contribution from the tensor coupling f_{B^*BV} is

$$
\sum_{B_{n^*}\left(\frac{1}{2}^-\right)}\left[\frac{b_{fn^*}f_{B_{n^*}B_iV}}{m_f - m_{n^*}} - \frac{f_{B_fB_{n^*}V}b_{n^*i}}{m_i - m_{n^*}}\right]\varepsilon^{*\mu}\bar{u}_f i\sigma_{\mu\nu}q_{\nu}\gamma_5 u_i \tag{11}
$$

and thus it determines A_3 in Eq. 1.

When the vector particle under consideration is a photon, standard application of the requirement of gauge invariance to the $g_{B^*B\gamma}$ and $f_{B^*B\gamma}$ couplings implies (as in the Section on Hara's theorem) that

$$
g_{B^*B\gamma} = 0. \tag{12}
$$

Only the contribution from the tensor coupling $f_{B^*B\gamma}$ survives then leading (through Eq. 11) to the standard gauge invariant form of the $B_iB_f\gamma$ parity violating coupling used in the derivation of Hara's theorem.

6.1.2 *Parity conserving amplitudes*

Let us now consider the parity conserving amplitudes. Since the matrix elements $\left\langle B_{n*}\left(\frac{1}{2}^-\right)\middle| H^{p.c.}\middle| B\left(\frac{1}{2}^+\right)\right\rangle$ vanish in the SU(3) limit (just as in the Lee-Swift theorem), only the ground $B_n\left(\frac{1}{2}^+\right)$ and the excited $B_{n*}\left(\frac{1}{2}^+\right)$ states may contribute significantly to the parity conserving amplitudes. In simple models the contribution from the ground states is assumed to be dominant. Consequently, with

$$\left\langle B_m\left(\tfrac{1}{2}^+\right)\middle| H^{p.c.}\middle| B_n\left(\tfrac{1}{2}^+\right)\right\rangle = a_{mn}\bar{u}_m u_n. \tag{13}$$

we have the following expressions for the parity conserving (p-wave) $B_i \to B_f M$ amplitudes.

$$M^{p.c.}\left(B_i \to B_f P\right) = \bar{u}_f \gamma_5 u_i \cdot \sum_{B_n\left(\frac{1}{2}^+\right)} \left[\frac{g_{B_f B_n P} a_{ni}}{m_i - m_n} + \frac{a_{fn} g_{B_n B_i P}}{m_f - m_n}\right] \tag{14}$$

and

$$M^{p.c.}\left(B_i \to B_f V\right) = \varepsilon^{*\mu} \bar{u}_f \gamma_\mu u_i \cdot \sum_{B_n\left(\frac{1}{2}^+\right)} \left[\frac{g_{B_f B_n V} a_{ni}}{m_i - m_n} + \frac{a_{fn} g_{B_n B_i V}}{m_f - m_n}\right]$$

$$+\varepsilon^{*\mu} \bar{u}_f i\sigma_{\mu\nu} q_\nu u_i \sum_{B_n\left(\frac{1}{2}^+\right)} \left[\frac{f_{B_f B_n V} a_{ni}}{m_i - m_n} + \frac{a_{fn} g_{B_n B_i V}}{m_f - m_n}\right] \tag{15}$$

where $g_{BBV}\left(f_{BBV}\right)$ are vector (tensor) parts of the $B_i B_n V$ or $B_n B_f V$ coupling constants.

6.2 *The quark model and QCD*

As mentioned before, for the parity violating NLHD amplitudes the soft meson approximation and the quark model results were shown to be equivalent in a group-theoretical sense. Such an equivalence leaves plenty of room for the dynamics. The modern way of supplementing the quark model with the dynamics involves introduction of quantum chromodynamics and its subsequent treatment through the application of the operator-product expansion and the renormalization-group techniques.[71]

The effective operator employed to analyze $\Delta S = 1$ weak interactions has the form[72-74]

$$H_{\Delta S=1} = \frac{G_F \sin\theta_c \cos\theta_c}{2\sqrt{2}} \sum_i c_i O_i + h.c. \tag{16}$$

where the four-quark operators $\{O_i\}$ are the lowest-dimension operators appearing in the operator-product expansion and are defined by

$$\begin{aligned}
O_1 &= H_A - H_B \\
O_2 &= H_A + H_B + 2H_c + 2H_D \\
O_3 &= H_A + H_B + 2H_c - 3H_D \\
O_4 &= H_A + H_B - H_c \\
O_5 &= \bar{d}\Gamma_L^\mu \lambda^A s \sum_Q \left(\bar{Q}\Gamma_{R\mu}\lambda^A Q\right) \\
O_6 &= \bar{d}\Gamma_L^\mu s \sum_Q \left(\bar{Q}\Gamma_{R\mu} Q\right)
\end{aligned} \tag{17}$$

with

$$\begin{aligned}
H_A &= \bar{d}\ \Gamma_L^\mu u\ \bar{u}\Gamma_{L\mu} s \\
H_B &= \bar{u}\ \Gamma_L^\mu u\ \bar{d}\Gamma_{L\mu} s \\
H_C &= \bar{d}\ \Gamma_L^\mu d\ \bar{d}\Gamma_{L\mu} s \\
H_D &= \bar{s}\ \Gamma_L^\mu s\ \bar{d}\Gamma_{L\mu} s
\end{aligned} \tag{18}$$

and

$$\Gamma_L^\mu \equiv \gamma^\mu(1+\gamma_5),\ \Gamma_R^\mu \equiv \gamma^\mu(1-\gamma_5) \tag{19}$$

The long-distance physics resides in the matrix elements of the O_i operators. The "penguin" operators O_5, O_6 (see Fig. 6.2.) have a $(V-A)(V+A)$ chiral structure whereas the remaining O_i operators are $(V-A)(V-A)$.

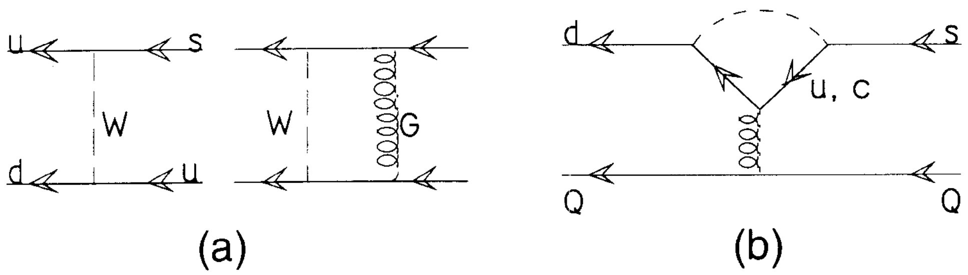

Fig. 6.2. Diagrams corresponding to the effective $\Delta S = 1$ Hamiltonian: (a) nonpenguin (b) penguin.

Operators O_1, O_2, O_5 and O_6 transform like SU(3) octets and are $\Delta I = \frac{1}{2}$ operators. O_3 and O_4 are 27-plets carrying $\Delta I = \frac{1}{2}$ and $\Delta I = \frac{3}{2}$ respectively. The coefficients c_i contain the short-distance effects of hard gluons and are calculated by studying the QCD renormalization group equations. Without QCD evolution one has $c_1 = 1$, $c_2 = \frac{1}{5}$, $c_3 = \frac{2}{15}$, $c_4 = \frac{2}{3}$, $c_5 = c_6 = 0$. For the case under consideration a typical set[74-76] of these coefficients is

$$\begin{array}{cccc} c_1 \sim 2.5 & c_2 \sim 0.08 & c_3 \sim 0.08 & c_4 \sim 0.40 \\ & c_5 \sim -0.05 \rightarrow -0.1 & c_6 \sim -0.01 \rightarrow -0.05. & \end{array} \tag{20}$$

Thus, radiative QCD corrections result in the enhancement of the octet $\Delta I = \frac{1}{2}$ O_1 operator and the suppression of the $\Delta I = \frac{3}{2}$ O_4 operator.[73,77] This dynamical argument goes some way towards the explanation of the $\Delta I = \frac{1}{2}$ rule. Furthermore, with penguin operators being flavor octet $\Delta I = \frac{1}{2}$ objects, a further enhancement of the $\Delta I = \frac{1}{2}$ amplitudes is predicted. The standard values of the penguin coefficients c_5, c_6 are small: they vanish in the limit of $m_c = m_u$ because of GIM cancellation. They are too small by a factor of order 5 to provide a satisfactory explanation of the $\Delta I = \frac{1}{2}$ rule.[76,78]

In the baryon sector, however, this $\Delta I = \frac{1}{2}$ rule is readily explained as an automatic consequence of color symmetry. The relevant argument, known as the Pati-Woo theorem[79,80] implies the vanishing of the matrix elements of the $\Delta I = \frac{3}{2}$ operator O_4 between the baryonic states. Actually, one can show[76] that

$$\left\langle B' \middle| O_i \middle| B \right\rangle = 0 \quad \text{for } i = 2, 3, 4. \tag{21}$$

Thus, the net effect of the QCD-enhancement factors is just to change the overall size of the Pati-Woo-allowed baryon-to-baryon weak matrix elements of diagrams (a) and (b) in Fig. 6.1. It *cannot* change such qualitative characteristics of quark model calculations as the violation of Hara's theorem in this model, that is the basis of the controversy regarding the WRHD. The QCD considerations of this section might, however, be more important in the boson-leg diagrams (Fig. 6.1c) where the Pati-Woo theorem does not apply.

6.3 *The boson-leg contribution*

At the quark level the boson-leg contribution (Fig. 6.1c) is often identified with the factorizable amplitude. The factorization prescription corresponds to the insertion of the vacuum state between the quark bilinears of the 4-quark Fermi interaction. In the case of nonleptonic hyperon decays, upon invoking PCAC, the use of factorization prescription gives a contribution that vanishes in the SU(3) limit.[81] At the hadron level the strength of the kaon-pole diagrams relevant in the parity conserving amplitudes is governed by the $K\pi$ transition matrix elements. It has been stressed[63,76] that quark model estimates of $\langle \pi | H_{weak} | K \rangle$ involve substantial cancellations and cannot be reliably computed. Application of chiral Lagrangians to provide phenomenological estimates of these matrix elements indicates that kaon pole terms are small in comparison to the baryon pole terms.[63] For a thorough review of the meson sector relevant here see a recent paper by Cheng.[78] In other phenomenological studies of NLHD the kaon pole contribution is substantial, however.[82,83] This disagreement constitutes just one example of the lack of general consensus concerning the relative magnitudes of various contributions in the nonleptonic hyperon decays. We shall discuss other such disagreements in the next section.

Since the decays of hyperons to other ground-state baryons and vector mesons are kinematically forbidden, we know even less about the boson-leg contribution for vector particles. Consequently, the contribution of this type of diagram in the weak radiative hyperon decays is often treated with the help of free parameters. Indeed, as Gilman and Wise[75] put it

> "while sometimes disguised in the language of the operator-product expansion, much of the short-distance analysis boils down in the end to finding the local operators which correspond in a particular model to the amplitude for the transition of an s-quark to a d-quark plus photon".

7. Phenomenology of the Standard Approach

7.1 *Pole models*

As it has been discussed in the previous section, the standard schemes for the description of the nonleptonic and weak radiative hyperon decays belong to the same general theoretical framework. Thus, it is quite plausible that our present phenomenological knowledge of NLHD might be useful in providing not only a background but also some important input needed for an understanding of WRHD. Accordingly, we must present first a brief overview of the present phenomenological situation in the NLHD sector.

7.1.1 *Nonleptonic hyperon decays*

As mentioned before, there is no consensus as to what are the relative sizes of various contributions to the NLHD amplitudes. Before we present the conflicting theoretical views, let us therefore recall the model-independent characteristics of these decays. These are given by SU(3) fits to the relevant experimental amplitudes, obtained as follows.

For the parity violating amplitudes, the soft-pion approximation (Eq. 6.6) can further be reduced through the use of the commutation relation

$$\left[Q_5^a, H^{p.v.}\right] = \left[Q^a, H^{p.c.}\right] \tag{1}$$

where Q^a is an SU(3) generator.

After working out the action of Q^a on the baryon states $\left(Q^a\left|B_i\right\rangle = \left|B_1\right\rangle, \left\langle B_f\right|Q^a = \left\langle B_2\right|\right)$ one finds that the current algebra approximation A^{CA} is given in terms of the matrix elements of the parity-conserving part $H^{p.c.}$ of the weak Hamiltonian between some baryon octet states B, B'

$$\langle B|H^{p.c.}|B'\rangle \tag{2}$$

The SU(3) parameterization of this matrix element is

$$\langle B|H^{p.c.}|B'\rangle = fTr\left(S\left[B', B^+\right]\right) + d\left(S\left\{B', B^+\right\}\right) \tag{3}$$

where B', B on the right-hand side are standard 3×3 matrices corresponding to the baryons in question and $S = \lambda_6$ is the ($s \rightleftarrows d$ symmetric) octet spurion representing the weak Hamiltonian. SU(3) fits to the s-wave amplitudes[84] give $f/d = -2.5$. Similar fits[51,63,85-87] to the p-waves yield $f/d = -1.8$ to -1.9. These experimental numbers still constitute a problem for the *valence* quark model (current algebra)[88] in which one obtains $f/d = -1$.

Existence of relatively good SU(3) fits to the NLHD amplitudes indicates that in these decays SU(3) is a fairly good approximate symmetry. Thus, the Lee-Swift theorem should be satisfied fairly well in the real world. Still, one may wonder how big the $\langle B|H^{p.v.}|B'\rangle$ matrix elements could be when SU(3) is weakly broken. Theoretical estimates of the ratios $\langle B|H^{p.v.}|B'\rangle / \langle B|H^{p.c.}|B'\rangle$ performed by Golowich and Holstein[89] in the context of the bag model indicate that they are of order of 1% in NLHD (5% in WRHD). As far as WRHD are concerned, the above considerations support, therefore, the assumption of the overall SU(3) symmetry used in the proof of Hara's theorem.

The origin of the discrepancy between the *valence* quark model and the experimental values for the f/d ratio has not been yet agreed upon. LeYaouanc et al. proposed[90] that the departure of the phenomenologically determined f/d ratio from its valence quark model/current algebra value is due to a nonvanishing contribution of the second term in Eq. 6.7 (which vanishes for degenerate octet baryons). In explicit quark model calculations, they have found that this term is of the order of 50% and negative with respect to the commutator term. Using the pole model fit of Gronau[88] for the p-waves (which, when experimental baryon masses are used, needs $f/d \approx -1.18$) they were able to explain both the value of f/d observed in the s-wave amplitudes and the relative size of s-wave and p-wave amplitudes (the s:p ratio).

General theoretical scheme of the previous section admits, however, contributions from other intermediate states besides the $(56, 0^+)\frac{1}{2}^+$ and $(70, 1^-)\frac{1}{2}^-$ baryons discussed in LeYaouanc et. al.[90] A thorough study of the contributions from radially excited $(56, 0^+)^* \frac{1}{2}^+$ baryons as well as K^* and K mesons has been carried out by Bonvin.[82] His calculations confirm general qualitative features of LeYaouanc et. al[90] i.e., that the SU(3) symmetry breaking part of the contribution of $(70, 1^-)\frac{1}{2}^-$ is significant and interferes destructively with the commutator contribution, thus partially curing the problem of the s:p ratio. However, he also finds that the contributions from the meson-leg diagrams and the $(56, 0^+)^* \frac{1}{2}^+$ radially excited baryons are far from being negligible. As a result, his decomposition's of the amplitudes A and B into different contributions are totally different from those of LeYaouanc et. al. [90]

Another approach for alleviating the f/d problem has been proposed by Donoghue and Golowich[91] who considered the effects of quark sea on the soft-meson approximation A^{CA}. Both the QCD sea (corresponding to the *enhanced* penguin contribution)[91] and the sea generated by unitarity on the hadron level[92] increase f/d of the soft pion contribution substantially (to around -1.6). This is close to the experimental value extracted from p-wave amplitudes. For the s-waves the f/d ratio is further enhanced to around -2.2 or even -2.5 by the SU(3) symmetry breaking in energy denominators of the intermediate states.[93]

In fact, under certain assumptions a value of -1.6 for the f/d ratio of the soft pion contribution has been determined phenomenologically by Pham.[94] His determination raises further doubts as to the validity of the previous decompositions[82,90] in which the value $(f/d)_{\text{soft pion}} = -1$ was used. The bigger value of the soft-pion f/d ratio was utilized in a recent update on the pole model by Nardulli.[83] His decomposition of the amplitudes again differs significantly from LeYaouanc et. al[90] and Bonvin. [82] In view of the uncertainties just discussed, a recent claim[95] that nonleptonic hyperon decays can be well understood should be considered as over optimistic.

In conclusion, no generally agreed upon explanation of the f/d problem in NLHD exists. This situation is *one* of the reasons for the proliferation of various results for WRHD – all obtained in the framework of the same general theoretical scheme.

7.1.2 *Weak radiative hyperon decays*

As discussed in Section 6 the tensor coupling contribution from the intermediate $(70, 1^-)\frac{1}{2}^-$ baryons leads to the standard form, Eq. 3.9 of the parity violating coupling of photon to baryons. Estimates of the contribution from the $\frac{1}{2}^-$ baryons were performed by many workers[96-100] at the time when the CERN experiment[25] was being carried out and again, more recently[83,101,102] when a new wave of experimental results became imminent. Originally, the most extensive calculations were performed by Gavela et al.[96] To find out the contribution from the $\frac{1}{2}^-$ baryons, they evaluated the parity-conserving $\frac{1}{2}^- \to \frac{1}{2}^+ \gamma$ s-wave decay amplitudes in the quark model following the old Copley et al. method[103] and identified the results with the phenomenological couplings $f_{BB*\gamma}$ in Eq. 6.8.

For the particular decay $\Sigma^+ \to p\gamma$, nonvanishing contributions come from $N^{*+}(^2 8)$ (Fig. 6.1a) and $\Sigma^{*+}(^2 8)$ (Fig 6.1b) members of the $(70, 1^-)\frac{1}{2}^-$ multiplet.

In the limit of exact SU(3) one has

$$\begin{aligned} f_{pN^{*+}\gamma} &= f_{\Sigma^+\Sigma^{*+}\gamma} \\ b_{N^{*+}\Sigma^+} &= b_{\Sigma^{*+}p} \end{aligned} \tag{4}$$

Using Eq. 6.9 and Eq. 6.3, one obtains then that in the SU(3) limit these two contributions cancel each other in expression Eq. 6.11. In this way, contributions from tensor couplings $f_{BB^*\gamma}$ satisfy Hara's theorem. Since the parity conserving $\Sigma^+ \to p\gamma$ amplitudes in the pole model are proportional to the difference $\mu_{\Sigma^+} - \mu_p$ of baryon magnetic moments (i.e., $f_{BB'\gamma}$ in Eq 6.15), the final result for the asymmetry and the branching fraction of the $\Sigma^+ \to p\gamma$ decay is very sensitive to the value of $\mu_{\Sigma^+} - \mu_p$, a feature already observed by Farrar.[53] Thanks to the fact that the experimental value for $\mu_{\Sigma^+} - \mu_p$ is significantly bigger than the quark model result, Gavela et al.[96] were able to obtain a large negative $\Sigma^+ \to p\gamma$ asymmetry. (This would not have been the case had they used the physical Σ^+, p masses and the additive quark model for the evaluation of μ_{Σ^+}, μ_p.) This and other results of Gavela et al.[96] are compared with the results of later experiments (see Table 2.1) in Table 7.1a (branching fractions) and Table 7.1b (asymmetry parameters). It is seen that their predictions went wrong in several places.

The contributions of the Σ^* and N^* resonances to the parity violating $\Sigma^+ \to p\gamma$ amplitude have also been estimated in the bag model.[100] As in Gavela et al.[96], Hara's theorem was satisfied by the cancellation of the contributions from these two resonances. Despite such similarities, the overall size of the parity violating and parity conserving amplitudes was found to be over an order of magnitude smaller than experimentally observed.

One may wonder if contributions from other intermediate states such as decuplet baryons could not bring theory in agreement with experiment. This does not seem to be the case, however. First, in the quark model such contributions vanish since the decuplet wave function is symmetric.[79,90] Second, explicit considerations of decuplet contribution by Scadron and Thebaud[104] have yielded predictions that totally disagree with the present data on WRHD.

Table 7.1a Branching fractions (in units of 10^{-3}) of weak radiative decays: comparison of theoretical predictions with experiment. Input entries are underlined.

	Gavela[96]	Gilman[75]	Kamal[106]		Verma[107]		Verma[1]	Zenc[108]	VDM Update	experiment
			A	B	(cq)	(ld)				
$\Sigma^+ \to p\gamma$	$0.92^{+0.26}_{-0.14}$	1.24	1.24	1.24	1.20	1.37	1.24	1.26	1.3-1.4	1.23±0.06
$\Lambda \to n\gamma$	0.62	22.0	5.97	1.70	0.95	1.66	1.62	1.00	1.4-1.7	1.63±0.14
$\Xi^\circ \to \Lambda\gamma$	3.0	4.0	1.80	1.36	0.36	0.57	0.5	1.03	0.9-1.0	1.06±0.16
$\Xi^\circ \to \Sigma^\circ\gamma$	7.2	9.1	1.48	0.23	4.43	4.06	3.30	3.83	4.0-4.1	3.56±0.43
$\Xi^- \to \Sigma^-\gamma$		11.0	1.20	1.20	0.23	0.23	0.23	0.29	0.15	0.128±0.023
$\Omega^- \to \Xi^-\gamma$		41.	0.60	0.60	0.52	0.52	0.86		0.55±0.20	<0.46

Table 7.1b Asymmetries of weak radiative decays: comparison of theoretical predictions with experiment. Input entries are underlined.

	Gavela[96]	Kamal[106]		Verma[107]		Verma[1]	Zenc[108]	VDM Update	experiment
		A	B	(cq)	(ld)				
$\Sigma^+ \to p\gamma$	$-0.80^{+0.32}_{-0.19}$	$\underline{-0.5}$	$\underline{-0.5}$	-0.59	-0.55	-0.56	-0.97	-0.95	-0.76±0.08
$\Lambda \to n\gamma$	-0.49	-0.87	+0.25	-0.66	-0.52	-0.54	+0.76	+0.8	
$\Xi^o \to \Lambda\gamma$	-0.78	-0.96	-0.45	+0.87	+0.74	+0.68	+0.65	+0.8	+0.43±0.44
$\Xi^o \to \Sigma^o\gamma$	-0.96	-0.3	-0.99	-0.90*	-0.81*	-0.94	-0.36	-0.45	+0.20±0.32
$\Xi^- \to \Sigma^-\gamma$		-0.87	+0.56	$\underline{\pm 0.44}$	$\underline{-0.44}$	-0.60	+0.63	+0.7	+1.0±1.3
$\Omega^- \to \Xi^-\gamma$		-0.87	+0.56	+0.44	-0.44	-0.60		+0.7	

*Verma[107] erroneously reports a positive sign here, corrected later in Verma[1]

7.2 *Quark models*

Pole model calculations of the parity violating amplitudes $B(\frac{1}{2}^+) \to B'(\frac{1}{2}^+)\gamma$ consist in the explicit evaluation of the contribution from the intermediate $\frac{1}{2}^-$ states. Thus, a good knowledge and understanding of both the parity conserving $B^*(\frac{1}{2}^-) \to B(\frac{1}{2}^+)\gamma$ amplitudes and the parity violating $\langle B^*(\frac{1}{2}^-)|H^{p.v.}|B(\frac{1}{2}^+)\rangle$ matrix elements is needed in such calculations. In connection with this requirement, it should therefore be recalled from Section 6.1.1 that in the related case of nonleptonic hyperon decays and in the limit of exact SU(3), the current algebra/quark model approach enables us to bypass the need to know anything about the $\frac{1}{2}^-$ poles. Thus, the quark model should provide a simple and transparent way of understanding the salient features of pole model calculations. Furthermore, calculations of weak radiative hyperon decays in the quark model may be more reliable than in the pole model. One might then also hope that in the quark model the questions of SU(3) symmetry breaking can be dealt with more easily.

7.2.1 *Single quark processes*

Within the quark model framework, the simplest assumption one can make is that the radiative weak decays originate from a strange quark decaying into a down quark with the emission of a photon (Fig. 7.1a). That is, one may assume that contributions from other

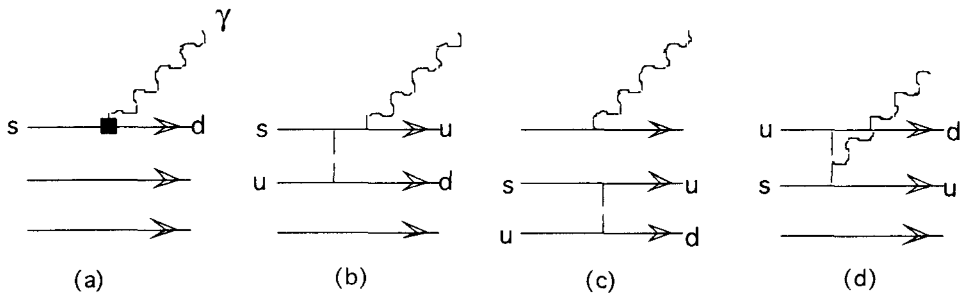

Fig. 7.1a-d Representatives of usually considered quark diagrams. a) single-quark emission b) two-quark bremsstrahlung c) three-quark transition d) internal radiation.

processes such as those shown in Figs. 7.1bc are negligible. The contribution from diagram (d) is suppressed by the presence of two W-propagators. It is a great virtue of the quark model that the general structure of the complicated contributions from baryon poles can be summarized through such simple quark diagrams. In fact, from the presence of the u quark in the structure of diagrams (b, c, d) in Fig 7.1 it follows immediately that the amplitudes of the $\Xi^- \to \Sigma^-\gamma$ and $\Omega^- \to \Xi^-\gamma$ decays may receive contributions from single-quark processes (Fig 7.1a) *only*. Now, assumptions similar to the dominance of diagram (a) worked very well for baryon magnetic moments and so one might think that it should work for the WRHD as well.[105] In 1979, treating the magnitudes of the parity violating and parity conserving amplitudes as free parameters (compare Eq. 6.3), Gilman and Wise[75] found, however, that this assumption is totally incompatible with the existing data. Their estimates of the branching fractions in the single-quark emission model are presented in Table 7.1a.

7.2.2 *W-exchange diagrams and the violation of Hara's Theorem*

First attempts to include diagram (b) from Fig. 7.1 have been made by Kamal and Verma.[106] Considering both diagrams (a) and (b) but neglecting (c), treating the overall sizes of the corresponding amplitudes as free parameters and antisymmetrizing the contributions from various diagrams of type (b) to satisfy Hara's theorem, they carried out a fit to the then-existing data. Their two solutions, shown in Tables 7.1ab are seen to be incompatible with the present data. The antisymmetrization prescription imposed by Kamal and Verma[106] does not belong to the quark model, however. It constitutes an extra condition inconsistent with the standard prescriptions of the quark model (as this model is formulated today). In fact, if one follows the rules of the quark model closely and pays due attention to the question of gauge invariance at the quark level, one finds that the W-exchange diagrams of Fig. 7.1b *violate* Hara's theorem, as observed by Kamal and Riazuddin (KR)[50,109] in 1983.

The KR calculation proceeds as follows. The contribution of the diagram of Fig. 7.1b is proportional to

$$\frac{\bar{u}\left(2\varepsilon\cdot p_u+\not{\varepsilon}\not{q}\right)\Gamma_{L\mu}\,s\bar{d}\,\Gamma_{L\mu}u}{\left(p_u+q\right)^2-m^2} \tag{5}$$

where ε is photon polarization, q its momentum, p_u momentum of the final u quark and m its mass.

The term proportional to $\varepsilon\cdot p_u$ vanishes after the integration over quark momenta in the final baryon.[106] To evaluate the leading term in Eq. 5 it is sufficient to approximate the quark propagator by $1/\left(q\cdot p_u\right)=1/(qm)$ and to use the static approximation for Dirac spinors. Taking the limit of exact SU(3) one then obtains, after the nonrelativistic reduction, the following parity violating Hamiltonian for the sum of all relevant W-exchange diagrams

$$H^{p.v.}\propto\vec{\varepsilon}\cdot\left(u^{+}\vec{\sigma}s\right)\times\left(d^{+}\vec{\sigma}u\right). \tag{6}$$

The expectation value of this operator sandwiched between the SU(6) spin-flavor internal wave functions of Σ^+ hyperon and proton is nonzero. The origin of this surprising result and, in particular, its connection with the pole model is not clear in this calculation, however. Since the assumptions upon which Hara's theorem rests seem unshakable, the result of Kamal and Riazuddin has been dismissed by some workers as exhibiting just a kind of pathology of the quark model. Others, nonetheless, tried to find a place for it in the existing hadron-level formalism.

In their original paper Kamal and Riazuddin claimed that of all assumptions used in the derivation of Hara's theorem, it is the assumption of $\Delta I = \frac{1}{2}$ octet dominance that is not satisfied in the quark model. However, the proof by Gourdin[44] (see Section 3.1) is not based on this assumption. Moreover – as it might have been expected from the related case of NLHD and the Pati-Woo theorem[79,80] – and as it has been checked by explicit calculations[51] -octet dominance is satisfied by the quark diagram (Fig. 7.1b) in question (see next section).

An attempt to obtain the violation of Hara's theorem in the standard pole model through the contribution from the a priori possible infrared singularities[110] has been refuted by M. K. Gaillard.[111] Indeed, such singularities could appear in Eq. 6.11 in the limit of exact SU(3) only if the $B^*\left(\frac{1}{2}^-\right)$ states were degenerate in mass with the ground states (i.e., if $m_\Sigma = m_{\Sigma^*} \neq m_N = m_{N^*}$ before the SU(3) limit is taken). Then, in the limit $q_0 = m_\Sigma - m_N \to 0$ one would obtain from 6.11 a nonvanishing E1 transition amplitude. However, the physical $B^*\left(\frac{1}{2}^-\right)$ states are nondegenerate in mass with the $B\left(\frac{1}{2}^+\right)$ ground states and, consequently, expression in 6.11 must vanish in the SU(3) limit. In a subsequent paper[101] Liu assigns the quark model violation of Hara's theorem to the contribution from multiquark, $M_n = (2q)(2\bar{q})$ 1^{++} intermediate state in the meson leg (Fig. 7.1c). This ad hoc assignment does not help, however, in identifying which of the original assumptions of Hara is violated.

One might wonder if the violation of Hara's theorem exhibited by the naive quark model would not vanish in a more elaborated quark scheme such as the relativistic bag model. However, the calculations performed by Lo[112] in the framework of the MIT bag model (though with broken SU(3)) have yielded the parity violating amplitude much *bigger* than the parity conserving one. If Hara's theorem were satisfied in the bag model, one would expect the opposite inequality.

7.2.3 *Quark model fits*

Whether we do understand the origin of the violation of Hara's theorem in the quark model or not, it is of great interest to try to describe the experimental data in that approach. The relevant calculations have first been carried out by Verma and Sharma.[107] They have studied the combined contribution from the diagrams (a) and (b) of Fig. 7.1 (Contributions from diagram (c) have been found negligible in the bag model calculations.[112] They also vanish in the $SU(6)_W$ symmetry approach discussed in Section 7.3.3). The magnitude of Hara's theorem violating diagrams (b) has been given by direct calculation in terms of the Fermi coupling constant without free parameters. On the other hand, the single-quark emission diagram has been expressed in terms of two parameters (one for the parity violating and one for the parity conserving amplitude, compare Section 6.3). With the experimental value for the $\Xi^- \to \Sigma^- \gamma$ branching fraction setting one constraint for these two parameters, there is only one parameter in their approach.

Alternatively, one may treat the asymmetry of the $\Xi^- \to \Sigma^- \gamma$ decay as an equivalent parameter. Verma and Sharma consider three models for the single-quark transition amplitude. These models yield three different values for the asymmetry in question. Two models give branching fractions which are not in total disagreement with the data. In the first model (cq) constituent quark masses are used in expression (3.15) yielding $\alpha(\Xi^- \to \Sigma^- \gamma) \approx +0.4$. In the second model (ld), motivated by an estimate of long-distance hadron-level contributions to the effective $s \to d\gamma$ vertex performed by Palle[113], they used $\alpha(\Xi^- \to \Sigma^- \gamma) \approx -0.4$. Predictions of the two models are given in Table 7.1ab. Both models yield much too small branching fractions for the $\Xi^\circ \to \Lambda\gamma$ decays. In addition, the branching fraction of the $\Lambda \to n\gamma$ process in model (cq) is in disagreement with the most recent data. As far as asymmetries are concerned, the $\Xi^\circ \to \Sigma^\circ \gamma$ decay asymmetry is predicted in both models to be close to -1 while experiment does not support a large negative value. (Verma[1] has corrected the erroneous positive sign for this asymmetry reported by Verma and Sharma.[107]) Model (ld) is seen to be the best of all models presented so far, but it still strongly disagrees with the data for the $\Xi^\circ \to \Lambda\gamma$ branching fraction and the $\Xi^\circ \to \Sigma^\circ \gamma$ asymmetry. Further work in the framework of quark models[114] has not improved considerably upon the description of Verma and Sharma.[107]

A very important feature of the quark model is that it gives a *positive* asymmetry for the $\Xi^\circ \to \Lambda\gamma$ decay, as suggested by recent experimental findings. This is in disagreement with the predictions of papers[96,106] in which Hara's theorem is enforced. On the whole, it may be seen from Table 7.1ab that the quark model seems to describe the essential features of the data better than previous hadron-level calculations. Thus, it is interesting to see if and how the quark model violations of Hara's theorem can be understood in hadron-level language originally employed in the proof of this theorem. In the next section, such an explanation based on the vector-dominance model (VDM) is given.[51]

7.3 *Vector-meson dominance*

The idea of vector-meson dominance of photon couplings to hadrons (Vector Dominance Model, VDM) is thirty years old. Although there is still no consensus as to whether one should treat VDM as a purely phenomenological prescription or as something more fundamental (though the first point of view is now widely accepted) the usefulness of VDM cannot be disputed since it is generally acknowledged that "vector mesons never miss". As an example supporting this statement, one may cite the calculations by Schwinger[115] who has shown that – when one combines the idea of vector-meson dominance with by now standard assumptions concerning the spin-flavor symmetry of the ground-state baryons – one obtains a parameter-*free* prediction for baryon magnetic moments. In particular, one obtains for the proton $\mu_p = 2m_N/m_\rho \approx 2.5$. (Thus, if the idea of constituent quark masses is accepted, this means that the value of the magnetic moment of the proton is close to 3 just because there are three quarks in a baryon and two in a meson!) Schwinger's model is still one of the best *parameter-free* models of baryon magnetic moments.[116] Another VDM calculation of Schwinger given in the same paper is that of the celebrated $\pi^\circ \to 2\gamma$ decay. His VDM calculation gives $\Gamma_{VDM}(\pi^\circ \to 2\gamma) = 7.4\text{eV}$ to be compared with the experimental value of $7.8 \pm 0.5\text{eV}$. Apparently, for some still not well understood reasons, the quark model and vector-meson dominance are forced to give very similar results. That is, the existence of a connection between VDM and the quark model is very likely. Now, as far as WRHD are concerned, it was precisely the quark model calculation in which Hara's theorem was violated. Since VDM is formulated in terms of hadron-level concepts (which were used in the proof of this theorem), VDM might be expected to shed some light on the origin of the violation of Hara's theorem in the quark model. To be able to pursue this line of thought[51] we shall first present the vector-dominance model and its possible connection with the quark model. The connection in question is known as "the Kroll-Lee-Zumino scheme".

7.3.1 *The Kroll-Lee-Zumino scheme*

On a purely phenomenological level, vector-meson dominance states that the coupling of photon to hadrons, H_1, H_2, may be obtained by (1) evaluating the expression (see Fig. 7.2a)

$$\left\langle H_2 \middle| V_\mu J_\mu^V \middle| H_1 \right\rangle \tag{7}$$

where J_μ^V are currents built out of quarks and V_μ is a vector meson and then (2) performing the substitution

$$V_\mu \to \frac{e}{g_V} A_\mu \tag{8}$$

where $e^2/4\pi = 1/137$ and (for $V = \rho$) $g_\rho = f_{\rho NN} = 5.0$. It is the phenomenological prescription of Eq. 7 and 8 that leads to all VDM successes.

This prescription has been criticized soon after its introduction since it is apparently incompatible with gauge invariance. Indeed, the condition that the vector-meson-mediated coupling of photon to hadrons (Fig. 7.2a) leads to prescription (Eq. 7,8) requires the introduction of the following photon-vector-meson coupling

$$e\frac{m_V^2}{g_V} V^\mu A_\mu . \tag{9}$$

At $q^2 = 0$ the vector-meson propagator in the diagram of Fig. 7.2a cancels m_V^2 from (9) leading to prescription (8). Through the diagrams shown in Fig. 7.2b the iteration of coupling (9) leads then, however, to a nonvanishing photon mass thus violating gauge invariance.

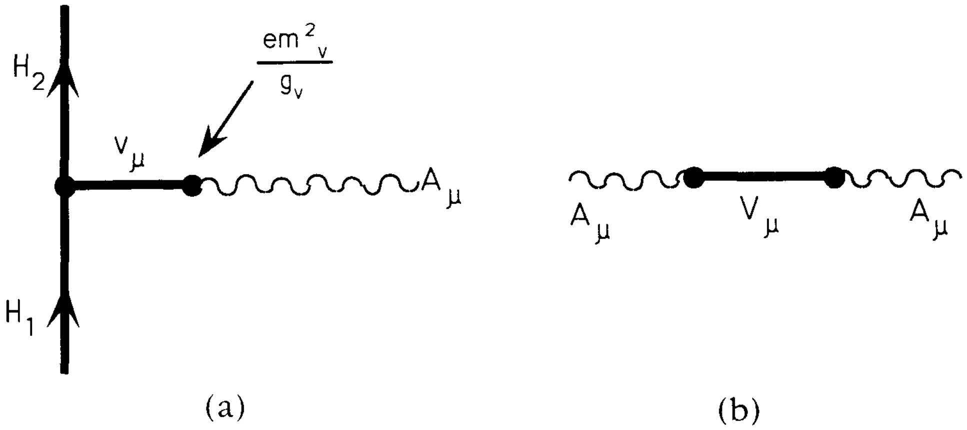

Fig. 7.2 (a) vector-meson-mediated coupling of photon to hadrons.
(b) vector-meson-induced photon mass term.

A possible resolution of this problem was proposed in 1967 by Kroll, Lee, and Zumino[117] (KLZ). A good account of the KLZ idea and VDM itself is given by Sakurai.[118] In brief, the contribution of terms in (9) to the photon mass requires the introduction of a photon mass counterterm

$$-\tfrac{1}{2}\left(\frac{e}{g_V}\right)^2 m_V^2 A_\mu^2 \tag{10}$$

which, together with vector-meson mass term and the photon-vector-meson coupling of Eq. 9 may be combined to form

$$-\tfrac{1}{2} m_V^2 \left[V_\mu - \frac{e}{g_V} A_\mu \right]^2 . \tag{11}$$

At this point, KLZ redefine the vector-meson and photon fields as well as electric charge by

$$V-\frac{e}{g_V}A=V'$$

$$\frac{1}{e^2}+\frac{1}{g_V^2}=\frac{1}{e'^2}$$

$$eA=e'A'$$

$$\left(\text{i.e.}\quad e'=e\left[1-\frac{1}{2}\frac{e^2}{g_V^2}+\ldots\right]\approx e\right). \tag{12}$$

Then, the expression (11) reduces to a mass term for the new vector field V'_μ, while Eq. 7 gives the coupling of the (new) photon field A'_μ to quarks in hadrons

$$e'\left\langle H_2\middle|A'_\mu J^V_\mu\middle|H_1\right\rangle. \tag{13}$$

In addition, the kinetic term

$$-\frac{1}{4}V^2_{\mu\nu}$$

$$\left(V_{\mu\nu}=\partial_\mu V_\nu-\partial_\nu V_\mu\right) \tag{14}$$

from the original vector-meson Lagrangian gives rise to the following explicitly gauge-invariant hadron-level photon-vector meson coupling of the primed fields:

$$-\frac{1}{2}\frac{e}{g_V}\cdot V'_{\mu\nu}F'^{\mu\nu}$$

$$\left(F'_{\mu\nu}=\partial_\mu A'_\nu-\partial_\nu A'_\mu\right). \tag{15}$$

At $q^2=0$ the hadron-level coupling of expression (15) vanishes and, consequently, in the new (primed) formulation the real photon couples to hadrons solely through the explicitly gauge-invariant coupling to quarks in Eq. 13.

In other words, according to KLZ, vector-meson dominance provides *effective means of translation* of the quark-level contribution (Eq. 13) into the standard hadron-level language (Fig. 7.2a). In the process of this translation the explicit gauge invariant formulation in terms of quarks and the primed fields V', A' becomes replaced by the effective VDM prescription of Eq. 7 and 8. This is precisely what is needed for an understanding of the origin of the violation of Hara's theorem in the quark model.[51]

7.3.2 *The quark/VDM approach and the violation of Hara's theorem*

According to Żenczykowski[51] the quark model contribution calculated by Kamal and Riazuddin generates, through the KLZ mechanism, a nonvanishing $F_1(0)\overline{u}_f\gamma_\mu\gamma_5 u_i \cdot A^\mu$ coupling in the *effective* hadron-level theory (see Section 3). Such a parity violating coupling may be non-zero for vector mesons, and hence, by VDM, also for photons. With nonvanishing $F_1(0)$ the assumptions of Hara are not satisfied and the theorem may be violated. Thus, according to the KLZ mechanism and Żenczykowski[51], the origin of the violation of Hara's theorem in the quark model is the neglect of the "contact" photon-quark interaction (Eq. 13) in all standard hadron-level proofs of Hara's theorem. This identification of the quark model amplitudes with the effective $\overline{u}_f\gamma_\mu\gamma_5 u_i A^\mu$ coupling confirms that the violation of Hara's theorem cannot arise in the pole model (see Section 6) from the tensor couplings $f_{B^*B\gamma}$ in Eq. 6.8 as attempted by Liu[101] and refuted by Gaillard.[111] Instead, it indicates that the resulting effective $\overline{u}_f\gamma_\mu\gamma_5 u_i A^\mu$ term arises from nonvanishing *effective* vector couplings $g_{B^*B\gamma}$ in Eq. 6.8 and 6.10. It should be stressed here that the introduction of nonvanishing strong coupling $g_{B^*B\gamma}$ is *not* an independent assumption. It is the only *logically consistent* way to obtain a nonzero effective $F_1(0)\overline{u}_f\gamma_\mu\gamma_5 u_i A^\mu$ coupling in a pole model. The vector couplings for the parity violating $B' \to B\gamma$ and the strong $B^* \to B\gamma$ transitions are either *both* zero as usually thought or *both* nonzero as suggested by VDM.

By the KLZ mechanism, the effective $g_{B^*B\gamma}$ couplings are then equivalent to the couplings obtained in the quark model calculation of the $B^* \to B\gamma$ decays.[103] In principle therefore, using the quark model calculations of the $B^* \to B\gamma$ decays one should obtain violation of Hara's theorem in the pole model. Calculations of Gavela et al.[96] did employ the quark model to evaluate the $B^* \to B\gamma$ amplitudes. However, their identification of the quark model results with the tensor couplings has led to the enforcement of the Hara theorem, and – not surprisingly – to results different from those obtained in genuine quark models (see discussion after Eqs. 7.21 and 7.22 in Section 7.3.4). One also has to bear in mind that in the related case of NLHD the quark model/current algebra approach seems superior to the pole model in the SU(3) limit because no information on the $\frac{1}{2}^-$ poles is then needed. One may expect a similar (purely practical) superiority of the quark model/VDM approach over the $\frac{1}{2}^-$ pole model also in the case of WRHD. Consequently, direct calculation of WRHD in the quark/VDM approach may not only be simpler and more transparent but also more reliable (at least at present) than the direct pole model calculations of the type performed by Gavela et. al.[96]

Depending on one's view concerning the meaning, importance and range of validity of the vector-meson dominance approach, one may accept or reject the above explanation of the violation of Hara's theorem in the quark model. If the KLZ scheme is accepted, there is no doubt that what is being calculated in the quark model is the $F_1(0)\bar{u}_f \gamma_\mu \gamma_5 u_i A^\mu$ coupling in the *effective* hadron-level theory. (In the anapole moment calculations of Musolf and Holstein[119], the vector-meson mediated contribution differs in an essential way from the original VDM.) If, on the other hand, the KLZ scheme is not accepted, we have two independent (quark model, VDM) schemes *both* indicating the violation of Hara's theorem. Their common conclusion is then more credible precisely because of the accepted independence of these schemes.

It should be noted that from the purely technical point of view the origin of the violation of Hara's theorem in the quark model is very simple. It results from the fact that – due to the tensor SU(6) structure of baryon wave functions in the quark model – the connection between the spin and space degrees of freedom in the quark model is different from that in the hadron-level language.[120] It is easy to see this in the naive quark model. Namely, in the most naive quark model the diagram of Fig. 7.1b represents a *scattering* amplitude of *free* quarks. As such it just *cannot* be simply fitted into the standard hadron-level language in which the concept of free quarks does not exist. As demonstrated by the more sophisticated potential[50] or bag model[112] calculations the use of confined quarks does not remedy this situation. This result is expected in all QCD inspired approaches with built-in short distance contributions of *free* quarks.

It was argued by Żenczykowski[120] that this conflict between the quark level calculations and the hadron level arguments seems to be intimately related to the notion of point in space-time. Namely, one of the problems posed by the standard quark model is how the gauge transformations on quark fields located at x_1, x_2, x_3 should be translated into a gauge transformation on the effective hadron field located at the system center of mass ($x = (x_1 + x_2 + x_3)/3$ in the simplest case). Perhaps the KLZ mechanism should be regarded as a (technical?) trick that does just this. Consequently, a deeper understanding of issues involved seems to be related to the composite and nonlocal nature of hadrons.[120] In fact, it is precisely through the description of composite quantum systems that the naive view of underlying continuum-spacetime local reality has been invalidated by Bell's theorem and associated experiments. The character of this review does not permit us to even attempt to pursue these matters in any detail. It is tempting, however, to speculate that the controversy over Hara's theorem may be quite closely related to such fundamental issues.

On the other hand, one might argue here that the quark-level calculation exhibits just a pathology of the quark model which, consequently, should be modified to conform to the standard hadron-level requirements. Indeed, for confined quarks their treatment as free particles is bound to lead to inconsistencies or paradoxial results. Such arguments should be treated with all due respect.

However, given:

- the experimentally observed apparent deviation from the expectations based on Hara's theorem,

- the great success of the simple quark model in providing the basis for our understanding of hadrons in general, and
- the pragmatic success of the VDM idea,

a serious consideration to the strict quark model/VDM approaches should also be given. After all, it used to happen repeatedly in the history of physics that we did not take our theories seriously enough. For all the above reasons we are inclined to think that Hara's theorem is indeed violated in nature.

Extensive studies of the VDM approach to WRHD have been carried out.[51,93,108,121] Let us turn therefore to the phenomenology of the vector-dominance model.

7.3.3 *Parity violating couplings of vector mesons to baryons*

Application of the VDM idea requires a calculation of the relevant parity-violating $\Delta S = 1$ couplings of vector mesons to baryons. The most reliable route is to utilize the $SU(6)_W$ symmetry approach[122,123] and to predict these couplings from those known empirically from NLHD. First calculations along these lines have been carried out by McKellar and Pick.[124] The interrelation of the $SU(6)_W$ and quark model schemes has been discussed thoroughly by Desplanques, Donoghue, and Holstein[64] (DDH). The possible $SU(6)_W$ diagrams are displayed in Fig. 7.3. They correspond precisely to appropriate quark model diagrams.

Diagrams (b1), (b2) and (c1), (c2) (referred to also as ((b) and (c)) of Fig 7.3 in general contribute to all the NLHD. On the other hand, diagram (a) corresponds to the usual factorization contribution and it does not contribute to NLHD in the $SU(3)$ limit. The appearance of diagrams (c) may be understood in the quark model as due to the presence of quark sea in physical hadrons.[91-93] Because of the symmetry of baryon wave functions, contributions of diagram (d) are zero in the $SU(6)_W$ symmetry limit for both pseudoscalar and vector mesons M.

The SU(3) structure of the $B_i \to B_f P$ parity violating amplitudes of Desplanques et. al.[64] has been expressed by Żenczykowski[51] as (see also Lee[125])

$$\frac{1}{4}\left(b-\frac{2}{3}c\right) Tr\left(\left[P^+,S\right]\left[B_i,B_f^+\right]\right)-\frac{1}{4}b\, Tr\left(\left[P^+,S\right]\left\{B_i,B_f^+\right\}\right) \tag{16}$$

where P is the standard 3×3 matrix corresponding to the pseudoscalar mesons and b, c are reduced matrix elements for the $SU(6)_W$ diagrams of Figs. 7.3b and 7.3c respectively. By CP invariance the parameters a, b, c are real. From Eq. 16 one can see that the $SU(6)_W$/quark model approach gives the soft-meson structure of Eq. 3 (as it must in accordance with Körner et. al.[70] and Desplanques et. al.[64]) with

$$\frac{f}{d} = -1 + \frac{2}{3}\frac{c}{b} \tag{17}$$

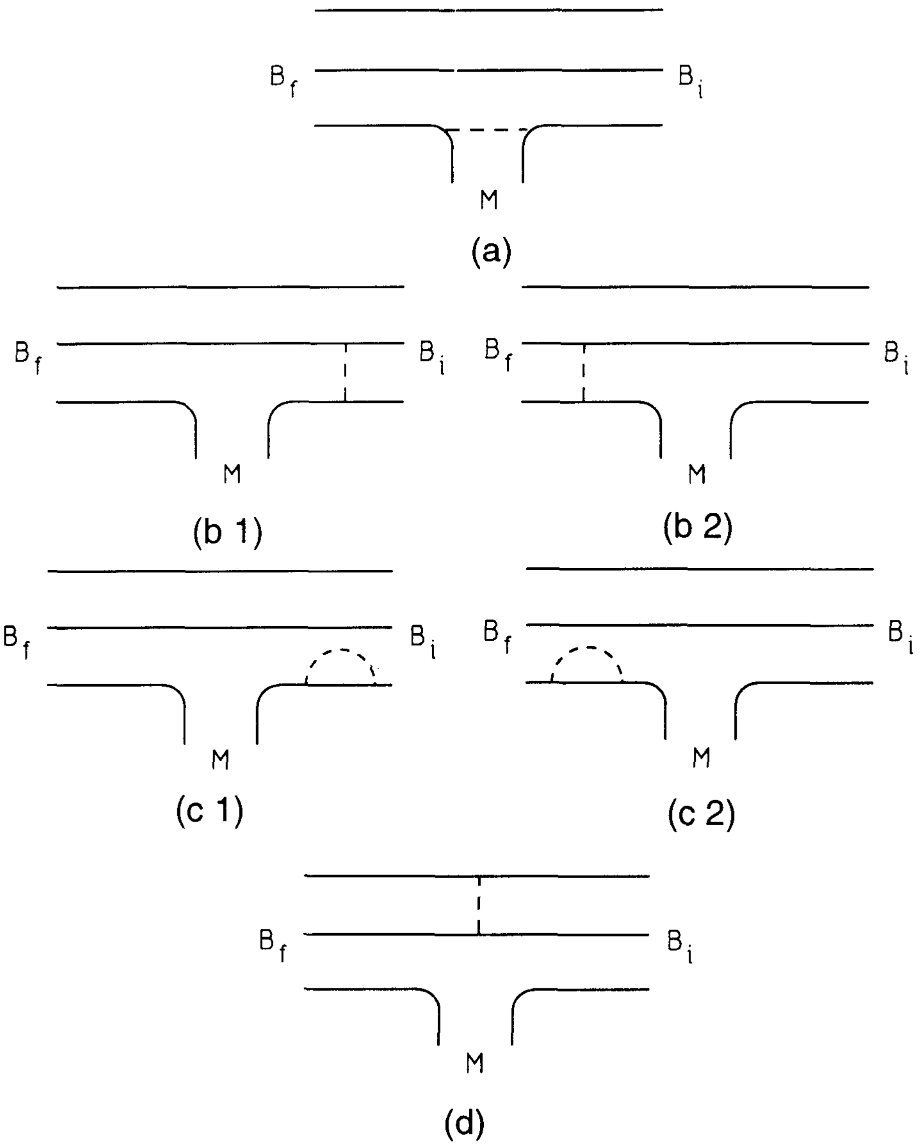

Fig. 7.3a-d SU(6)$_W$ quark-line diagrams.

The values of the reduced matrix elements b, c are determined from experiment to be (in units of 10^{-7})

$$\begin{aligned} b &= -5 \\ c &= +12 \end{aligned} \tag{18}$$

and thus the normalization of the contribution from the diagrams (b) and (c) to the $B_i \rightarrow B_f V$ parity violating amplitudes is fixed.

For (transverse) vector mesons, the contributions from the diagrams of Fig. 7.3 have been identified[64,124] with the vector coupling $\bar{u}_f \gamma_\mu \gamma_5 u_i \varepsilon^{*\mu}$ of Eq. 6.1.

For vector mesons the SU(3) structure of the contributions from the diagrams (b) and (c) is more complicated[51] than in the pseudoscalar case

$$\begin{aligned}
&\frac{1}{3\sqrt{2}} b \cdot \left\{ -4Tr\left(SB_f^+ V^+ B_i\right) - Tr\left(\left\{S, V^+\right\} B_f^+ B_i\right) \right. \\
&\qquad \left. + 2\left[Tr\left(SB_f^+\right) Tr\left(B_i V^+\right) + Tr\left(SB_i\right) Tr\left(B_f^+ V^+\right)\right]\right\} \\
&+ \frac{1}{9\sqrt{2}} c \cdot \left\{ -Tr\left(\left\{S, V^+\right\} B_f^+ B_i\right) - 5Tr\left(\left\{S, V^+\right\} B_i B_f^+\right) \right. \\
&\qquad \left. + 2Tr\left(SV^+\right) Tr\left(B_i B_f^+\right)\right\}.
\end{aligned} \tag{19}$$

Together with the factorization contribution of diagram 7.3a, Eq. 19 constitutes the strict analog in the vector-meson sector of Eq. 16. The latter reduces to the well-known soft meson/current algebra result for nonleptonic hyperon decays (Eq. 3).

As in Eq.16 and 3, the action of the weak Hamiltonian is represented by the $\Delta I = \frac{1}{2}$ octet spurion S. Thus, the octet dominance assumption of Hara is satisfied in the troublesome (b) type (W-exchange) diagrams of the quark model/VDM approach.

Since diagram (a) does not contribute to NLHD in the SU(3) limit, the size of its contribution to $B_i \rightarrow B_f V$ parity violating amplitudes cannot be determined by symmetry arguments. Various theoretical assumptions lead to a very wide range for the size of this term.[108,121,124] It was argued by Żenczykowski,[108,121] that the theoretical determination of its value may be prone to similar uncertainties as the determination of the value of c which in the simple quark model – in total disagreement with experiment – is zero. In fact, as mentioned before, in many approaches the size of the contribution to WRHD from such "single-quark" diagrams is treated as a free parameter. A similar approach was adopted[108,121] in VDM.

7.3.4 *Parity violating couplings of photons to baryons*

Expressing the photon couplings through the standard VDM assumption linking the photon with the linear combination of vector mesons

$$\gamma \leftrightarrow \frac{1}{\sqrt{2}} \rho^\circ + \frac{1}{3\sqrt{2}} \omega^\circ - \frac{1}{3} \varepsilon \varphi \tag{20}$$

($\varepsilon \approx {}^{m_u}\!/_{m_s} \approx {}^{2}\!/_{3}$ is the additive SU(3)-breaking parameter) the resulting parity violating amplitudes are readily calculated[51] from Eq. 19 (up to an overall VDM factor ${}^{e}\!/_{g}$; ${}^{e^2}\!/_{4\pi} = {}^{1}\!/_{137}$, $g = 5.0$):

$$A\left(\Sigma^+ \to p\gamma\right) = -\frac{5+\varepsilon}{9\sqrt{2}}b + \frac{1}{\sqrt{2}}d'$$

$$A\left(\Lambda \to n\gamma\right) = \frac{3+\varepsilon}{6\sqrt{3}}b - \frac{3\sqrt{3}}{2}d'$$

$$A\left(\Xi^\circ \to \Lambda\gamma\right) = -\frac{2+\varepsilon}{9\sqrt{3}}b + \frac{\sqrt{3}}{2}d'$$

$$A\left(\Xi^\circ \to \Sigma^\circ\gamma\right) = \frac{1}{3}b - \frac{5}{2}d'$$

$$A\left(\Xi^- \to \Sigma^-\gamma\right) = \frac{5}{\sqrt{2}}d' \tag{21}$$

where d' is expressed by c and the unknown reduced matrix element a corresponding to Fig. 7.3a through

$$d' = \left((1+\varepsilon)c - 8a\right)/27 \tag{22}$$

The contributions from the SU(6)$_W$ diagrams (b) violate the Hara theorem (as is manifest in Eq. 21) and correspond exactly to the SU(3) limit of the quark model calculations of W-exchange processes, which were carried out by Verma and Sharma.[107]

The W-exchange terms in Eq. 21 correspond to the *sum* of the contributions from diagrams (b1) and (b2) in Fig. 7.3. The weights of the individual contributions of these two diagrams are given in Table 7.2. The entries in Table 7.2 are proportional to the numerators $\Sigma\left\langle B_f\ V\middle|strong\middle|n\right\rangle\left\langle n\middle|H^{\mathrm{p.v.}}\middle|B_i\right\rangle$ (diagram (b1)) and $\Sigma\left\langle B_f\ \middle|H^{\mathrm{p.v.}}\middle|n\right\rangle\left\langle nV\middle|strong\middle|B_i\right\rangle$ (diagram (b2)) of the pole model contributions. As mentioned in Sections 7.1.2 and 7.3.2, Gavela et. al.[96] enforced Hara's theorem by requiring the cancellation of the contributions from the s- and u-channel poles. This procedure corresponds to the *subtraction* of the weights of diagrams (b1) and (b2). It is inconsistent with the quark model prescription of Eq. 21 and with the pole model corresponding to it since energy denominators associated with diagrams (b1) and (b2) are of the same sign.

From Table 7.2 one can see that in the subtraction procedure one obtains the vanishing of the parity violating $\Sigma^+ \to p\gamma$ amplitude in the SU(3) limit $(\varepsilon \to 1)$. From Table 7.2 one can also determine the relative signs of the amplitudes corresponding to the prescription of Gavela et. al.[96] with respect to those given in Eq. 21. To do this one has to express the SU(3)-breaking mechanism of Gavela et al.[96] in the language of quark diagrams. This has been done by Żenczykowski.[93] Basically, the SU(3) breaking of Gavela et al.[96] comes from the energy denominators describing the propagation of the

Table 7.2: Weights of quark diagrams (b1) and (b2) for the s-wave amplitudes.

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

intermediate $\left(70,1^-\right)\frac{1}{2}^-$ states in between the action of the weak Hamiltonian and the photon coupling to baryons. When SU(3) is broken the energy differences in these denominators are different for diagrams (b1) and (b2). They are

$$\Delta\omega - \delta s$$

for diagrams (b1) and

$$\Delta\omega + \delta s$$

for diagrams (b2). Here $\Delta\omega \approx 570 MeV$ is the $\left(56,0^+\right)\frac{1}{2}^+ - (70,1^-)\frac{1}{2}^-$ splitting, and $\delta s \approx 190 MeV$ is the strange-nonstrange quark mass difference. Let us denote $x = \delta s / \Delta\omega$ ($\approx 1/3$). Up to an overall normalization the weights for diagrams (b1) should be multiplied by $1/(1-x)$ while those of diagrams (b2) by $1/(1+x)$. This leads to the formulas for parity violating amplitudes presented in Table 7.3 (for clarity we have put $\varepsilon = 1$, this does not invalidate our discussion below).

We see that the signs of the $\Sigma^+ \to p\gamma$ and $\Xi^\circ \to \Sigma^\circ\gamma$ parity violating amplitudes calculated according to Gavela et. al.[96] are identical with those of the SU(3) breaking VDM/quark model (and those of Eq. 21), while for $\Xi^\circ \to \Lambda\gamma$ they are opposite. For the $\Lambda \to n\gamma$ decay these signs are also opposite since $x \approx 1/3 < 1/2$. If $\varepsilon \approx 2/3$ is reintroduced the corresponding inequality is still satisfied:

$$x \approx 1/3 < \frac{1+\varepsilon}{3+\varepsilon} \approx \frac{5}{11}$$

Table 7.3. Parity violating amplitudes with SU(3) breaking effects.

Process	Hara's theorem satisfied Gavela[96] (b1)-(b2)	Hara's theorem violated VDM/quark model[93,107] (b1)+(b2)
$\Sigma^+ \rightarrow p\gamma$	$-\frac{1}{3\sqrt{2}}\frac{2x}{1-x^2}$	$-\frac{1}{3\sqrt{2}}\frac{2}{1-x^2}$
$\Lambda \rightarrow n\gamma$	$\frac{1}{6\sqrt{3}}\frac{2(2x-1)}{1-x^2}$	$\frac{1}{6\sqrt{3}}\frac{2(2-x)}{1-x^2}$
$\Xi^\circ \rightarrow \Lambda\gamma$	$\frac{1}{3\sqrt{3}}\frac{1}{1+x}$	$-\frac{1}{3\sqrt{3}}\frac{1}{1+x}$
$\Xi^\circ \rightarrow \Sigma^\circ\gamma$	$\frac{1}{3}\frac{1}{1-x}$	$\frac{1}{3}\frac{1}{1-x}$

This comparison of the signs of parity violating amplitudes readily translates into a statement about the signs of the asymmetries (when p.c. amplitudes are treated in a standard manner, see next Section) and explains why – while in Gavela et. al.[96] all asymmetries are negative – in the VDM/quark model approach one obtains positive asymmetries for the $\Lambda \rightarrow n\gamma$ and $\Xi^\circ \rightarrow \Lambda\gamma$ decays (see Section 7.4).

Thus, there is a very important qualitative difference between the results of the two contradictory ways of treating the W-exchange processes in quark-model-*inspired* frameworks.

Namely:

If Hara's theorem is enforced in the SU(3) limit by the cancellation of the s- and u-channel contributions ((b1) and (b2) quark diagrams) all the relevant asymmetries are of the same sign. On the other hand, if the quark model recipe is *strictly* followed, the Hara's theorem is violated and the asymmetries of the $\Lambda \rightarrow n\gamma$ and $\Xi^\circ \rightarrow \Lambda\gamma$ decays are opposite to those of $\Sigma^+ \rightarrow p\gamma$ and $\Xi^\circ \rightarrow \Sigma^\circ\gamma$. Please note that for the $\Xi^\circ \rightarrow \Lambda\gamma$ decay the difference between the two prescriptions is that of a change of sign *only*.

Before proceeding to the discussion of VDM fits, we should mention here that the application of VDM to the description of WRHD was first attempted in Gavroglu et. al.[126] Several essential ingredients of the approach of Żenczykowski[51,108,121] were missing in that paper, however. In particular, only the factorization diagrams (Fig. 7.3a) were treated with the help of VDM. Furthermore, the Lee-Swift theorem was violated and, most important of all, the crucial assumption of SU(6) has not been made.

More recently, an empirical rule correlating the radiative and pionic decays of hyperons (and reminiscent of the VDM-based connections between the two) has been pointed out by Brown and Paschos.[127] In fact, the order of magnitude of the WRHD branching fractions can be determined from the pionic decays of hyperons and the VDM ideas as follows. The NLHD branching fractions into a channel with neutral pion (e.g., $\Sigma^+ \rightarrow p\pi^\circ$) are of order $\frac{1}{2}$. Replacing pseudoscalar meson with vector meson is equivalent to a change in group-theoretical factors which – in both cases – are of the same order. Thus, the whole difference in the branching fractions R of weak radiative and nonleptonic hyperon decays is attributed to the VDM factor $e/_g$, i.e.,

$$R(WRHD) \approx R(NLHD) \cdot \left(\frac{e}{g}\right)^2 \approx \frac{1}{2} \cdot \left(\frac{e}{g}\right)^2 = 1.8 * 10^{-3} \qquad (23)$$

7.3.5 *VDM fits*

The controversy surrounding WRHD concerns the parity violating amplitudes only. Unfortunately, experiments do not provide an unambiguous determination of these and parity conserving amplitudes. From the measured branching fractions and asymmetries one can determine the two amplitudes up to an interchange of their magnitudes only.

Consequently, comparison with experiment of various models of the parity violating amplitudes must also involve other models: those of the parity conserving amplitudes. Although no controversy regarding the general structure of the latter exists, they are very sensitive to various numerical details. Therefore – if a meaningful test of various approaches to the parity violating amplitudes is to be made – it is obvious that one should adopt such a description of the parity conserving amplitudes that is least dependent on such details. The best way would then probably be to exploit somehow our experimental knowledge of the parity conserving amplitudes in the related case of NLHD. The use of specific models for NLHD does not seem to be a good choice since there is no consensus as to the relative sizes of the contributions from the various possible intermediate states (see Section 7.1.1). Even when one considers only the ground-state baryons in the intermediate states, the transition from the NLHD to WRHD still involves two sources of phenomenological uncertainties.

The first of them was noted twenty years ago by Farrar[53]: it is the high sensitivity of the model to the precise values of baryon magnetic moments. Now, the prediction of the standard additive quark model for the *difference* of the magnetic moments of the Σ^+ hyperon and the proton differs from the measured difference by a factor of around 3. This strong nonadditive SU(3) breaking has not been explained in any scheme as yet.[128-130] Since the parity conserving amplitude of the $\Sigma^+ \to p\gamma$ decay is proportional to this difference, it is important that we use here the experimental values and not the additive quark model predictions.[96] Similarly, we should not trust too much the detailed numerical predictions of the quark model for other magnetic moments.

The second uncertainty, first pointed out by Żenczykowski[108], appears when spin symmetry is applied to relate the couplings of pseudoscalar mesons to those of vector mesons. It turns out that the relevant relation – as it emerges in standard pole models – leads to a strong violation of the vector-meson analog of the Lee-Sugawara relation[131,132] For NLHD this relation reads

$$\sqrt{3}\,\Sigma_\circ^+ = \Lambda_-^\circ + 2\Xi_-^- \qquad (24)$$

and it relates decay amplitudes of the Σ^+, Λ and Ξ^- hyperons into other ground-state baryons and π°, π^-, π^- mesons respectively. Since the Lee-Sugawara relation (Eq. 24) is well satisfied for the parity conserving NLHD amplitudes (at the 12% level) the violation of its vector-meson analog is not what one would expect on the basis of spin symmetry.

One can build the Lee-Sugawara relation into the model in such a way that it is preserved also in the vector meson sector (that is, that Eq. 24 is satisfied for rho couplings). This corresponds to assuming that for the parity conserving amplitudes the weak Hamiltonian transforms[63] effectively as λ_7.

In an analysis of these ambiguities[108], fits to the WRHD data were performed with respect to the parameter d' of Section 7.3.3. The fits favor the Lee-Sugawara relation for rho meson couplings. In the favored approach the parity conserving WRHD amplitudes are given by

$$
\begin{aligned}
B\left(\Sigma^+ \to p\gamma\right) &= \sqrt{2}\left(\frac{f}{d}-1\right)\left(\mu_{\Sigma^+}-\mu_p\right)C \\
B(\Lambda \to n\gamma) &= \left\{\frac{1}{\sqrt{3}}\left(1+3\frac{f}{d}\right)\left(\mu_\Lambda-\mu_n\right)-\left(\frac{f}{d}-1\right)\mu_{\Sigma\Lambda}\right\}C \\
B\left(\Xi^\circ \to \Lambda\gamma\right) &= \left\{\frac{1}{\sqrt{3}}\left(1-3\frac{f}{d}\right)\left(\mu_{\Xi^\circ}-\mu_\Lambda\right)-\left(\frac{f}{d}+1\right)\mu_{\Sigma\Lambda}\right\}C \\
B\left(\Xi^\circ \to \Sigma^\circ\gamma\right) &= \left\{\left(\frac{f}{d}+1\right)\left(\mu_{\Xi^\circ}-\mu_{\Sigma^\circ}\right)+\frac{1}{\sqrt{3}}\left(3\frac{f}{d}-1\right)\mu_{\Sigma\Lambda}\right\}C \\
B\left(\Xi^- \to \Sigma^-\gamma\right) &= -\sqrt{2}\left(\frac{f}{d}+1\right)\left(\mu_{\Xi^-}-\mu_{\Sigma^-}\right)C.
\end{aligned} \tag{25}
$$

In Eq. 25 C is an overall normalization factor (negative by convention) and f/d is the ratio of two invariant SU(3) couplings. Both C and f/d can be determined by symmetry from NLHD. Possible uncertainties result from the imprecisely known value of the f/d ratio (which NLHD determines to be around -1.8 to -1.9) and the experimental error in the value of the $\mu_{\Sigma\Lambda}$ transition moment $\left(\mu_{\Sigma\Lambda}^{\exp}=1.61\pm0.08\right)$. Because of strong cancellations, the $\Lambda \to n\gamma$ amplitude is particularly sensitive to such details. In Table 7.1ab we display the original fit of Żenczykowski.[108,121] In this paper we update this fit (VDM Update in Table 7.1) using the most recent experimental data (including the new measurements of the $\Lambda \to n\gamma$ and $\Xi^- \to \Sigma^-\gamma$ branching fractions). Since the fits are sensitive to the precise value of the $\mu_{\Sigma\Lambda}$ transition moment we decided to present their results in a semi-quantitative way. Results shown in Table 7.1ab are obtained for f/d =-1.8 when $\mu_{\Sigma\Lambda}$ deviated by 1.0 or 1.5 standard deviations from the experimental number. The normalization constant C differs by a few percent from its value determined from NLHD. The fits exhibit a good agreement with the experiment with a possible exception of the $\Xi^\circ \to \Sigma^\circ\gamma$ asymmetry. The fitted value of d' is near -0.25 to -0.28.

We should note here that in the related charmed baryon decays the $\Lambda_c^+ \to p\phi$ decay receives contributions only from the factorizable diagram of Fig. 7.3a. Although one has to keep in mind that SU(4) is not such a useful symmetry as SU(3), the measurement of the asymmetry in the $\Lambda_c^+ \to p\phi$ decay could still provide some independent information on the sign (and size) of the reduced matrix element a.

7.4 *Reliability of VDM and quark model fits*

We feel it worthwhile to address now the question of the numerical reliability of the quark/VDM fits of Table 7.1ab. In general, unless there are some substantial cancellations in the amplitudes, the expected errors of VDM fits are 15-20% for the branching fractions and ± 0.15 for the asymmetries. The differences between the predictions of various papers utilizing the quark model route and those where the VDM/symmetry approach[108,121] was employed, originate mainly in the parity conserving amplitudes. The two most important differences in the treatment of the parity conserving amplitudes which lead to the observed diversity of predictions are as follows:

(1) the quark model route implicitly assumes standard additive expressions for baryon magnetic moments, while the VDM/symmetry approach[108,121] takes into account the experimentally observed nonadditivities.
(2) in the quark model approach, the parity conserving single-quark amplitude, Fig. 7.1a, is treated as a free parameter while in the other[108,121] it is fixed by symmetry from experimentally known parity conserving amplitudes of NLHD.

Points (1) and (2) above account for the bulk of differences between the predictions of various versions of the quark model and Żenczykowski.[108,121]

To exhibit the connection between the VDM and the quark model calculations we shall rewrite Eqs. 25 in a form which explicitly displays the contributions from two-quark (W-exchange) and single-quark $(s \to d\gamma)$ processes. Using the additive quark model for the baryon magnetic moments, the formulas (Eqs. 25) reduce to

$$B\left(\Sigma^{+} \to p\gamma\right) = \frac{\sqrt{2}}{9} C'[2(1-\varepsilon) - s]$$
$$B(\Lambda \to n\gamma) = \frac{1}{\sqrt{3}} C'\left[\frac{2}{3}(1+\varepsilon) + s\right]$$
$$B\left(\Xi^{\circ} \to \Lambda\gamma\right) = -\frac{1}{3\sqrt{3}} C'\left[\frac{4}{3}(2+\varepsilon) + s\right]$$
$$B\left(\Xi^{\circ} \to \Sigma^{\circ}\gamma\right) = -\frac{1}{3} C'\left[4 - \frac{5}{3}s\right]$$
$$B\left(\Xi^{-} \to \Sigma^{-}\gamma\right) = -\frac{5\sqrt{2}}{9} C' \cdot s. \tag{29}$$

where: $\varepsilon = {}^{m_u}\!/_{m_s} \approx {}^{2}\!/_{3}$ is the SU(3) breaking parameter, C' is some overall normalization parameter, and s is the single-quark contribution, given in the symmetry approach by:

$$s = \left(\frac{f}{d} + 1\right)(1 - \varepsilon) \tag{30}$$

and thus negative (for ${}^{f}\!/_{d} \approx -1.8$).

The terms independent of s in Eq. 29 correspond to the W-exchange two-quark processes with ${}^{f}\!/_{d} = -1$. Note that for negative s the $B(\Lambda \to n\gamma)$ and $B(\Xi^{\circ} \to \Lambda\gamma)$ amplitudes exhibit cancellations between the single- and the two-quark processes.

Let us now discuss the reliability of the VDM fits of Table 7.1ab for the individual decays as well as their connection with various quark model predictions.

7.4.1 $\Sigma^{+} \to p\gamma$

The branching fraction is correct. The asymmetry is large and negative. Its exact size depends on how one treats SU(3) breaking in the parity conserving amplitude. Strict quark model calculations tend to give smaller parity conserving amplitudes, smaller asymmetries (around -0.6) and smaller branching fractions. If the experimentally observed, but theoretically not understood nonadditive SU(3) breaking in $\mu_{\Sigma^+} - \mu_p$ is used as a guide for what to expect in this amplitude a value of about -0.8 to -1.0 is predicted. (If Hara's theorem is enforced in quark-model-inspired framework,[96] the asymmetry in the SU(3) breaking case is also negative – recall discussion in Section 7.3.5).

7.4.2 $\Lambda \to n\gamma$

This is the case of strong cancellations in the parity conserving amplitude. In the symmetry approach of Żenczykowski[108,121] with ${}^{f}\!/_{d}$ fixed at -1.8, the experimental uncertainty in the value of $\mu_{\Sigma\Lambda}$ magnetic transition moment permits any value of the branching fraction up to about 1.7×10^{-3}. Two-quark (W-exchange) contribution to the parity conserving amplitudes (corresponding to ${}^{f}\!/_{d} = -1$) leads to a positive asymmetry of the decay. The asymmetry remains positive also for ${}^{f}\!/_{d} = -1.8$ when the size of the corresponding single-quark parity conserving amplitude is determined by symmetry from NLHD (see Eq. 29). In quark models, where this amplitude is treated as a free parameter, the only constraint on it is an upper bound resulting from the known $\Xi^{-} \to \Sigma^{-}\gamma$ branching fraction. If the most recent small value[30] of this branching fraction is not used then - due to the cancellations in Eq. 29 - one can obtain a change of sign of the whole parity conserving amplitude. This leads to a wide range of predictions for the $\Lambda \to n\gamma$ asymmetry in the quark model (~ -0.60 in Verma and Sharma[107], 0.0 in Uppal and Verma[114]). In addition, the cancellation in the parity violating amplitudes between the W-exchange and single-quark processes are also non-negligible. A conservative estimate for the asymmetry in the VDM/symmetry approach is +0.4 to +0.9. (On the contrary, in quark-inspired pole models in which Hara's theorem is enforced by the cancellation of the s- and u-channel contributions[96] this asymmetry has a tendency to be negative (c.f. Section 7.3.4)).

7.4.3 $\Xi^\circ \rightarrow \Lambda\gamma$

In the symmetry approach of Żenczykowski[108,121] the cancellation in the parity conserving amplitude is weaker than in $\Lambda \rightarrow n\gamma$ (see Eq. 29). The symmetry approach gives then for the branching fraction a value in the range $0.8 * 10^{-3}$ to $1.0 * 10^{-3}$, while the asymmetry is predicted to be positive and around +0.7. Dependence on the size of the single-quark parity conserving amplitude is fairly weak. Therefore, in quark models where this amplitude is treated as a free parameter, the size of the single-quark contribution is limited to – at the very most – around one third of the two-quark contribution. Thus, no change of sign of the whole parity conserving amplitude is possible. Consequently, the sign and the approximate size of the asymmetry constitute stable predictions of the VDM/quark model approach. Indeed: explicit quark model calculations give positive and substantial asymmetries.[107,114]

Present data are consistent with this prediction. Recall from Section 7.3.5 that if Hara's theorem is enforced upon the quark-inspired pole model[96] the asymmetry for this process is predicted negative. Thus, since the $\Xi^\circ \rightarrow \Lambda\gamma$ asymmetry is experimentally accessible,[58] its precise measurement should settle the *theoretical* question of the breakdown of Hara's theorem in the SU(3) limit.

7.4.4 $\Xi^\circ \rightarrow \Sigma^\circ\gamma$

When the single-quark parity conserving amplitude is determined by symmetry from NLHD, the contributions from the single-quark and W-exchange processes *add* rather than cancel. The parity conserving amplitude is then dominant and the branching fraction is expected to be around 4.0×10^{-3}. Consequently, the asymmetry should be small. The experimental value of the $\Xi^- \rightarrow \Sigma^-\gamma$ branching fraction limits the size of the single-particle parity conserving amplitude so that in those quark models in which this amplitude is treated as a free parameter no change of sign of the whole $\Xi^\circ \rightarrow \Sigma^\circ\gamma$ parity conserving amplitude is possible. As a result, the sign of the $\Xi^\circ \rightarrow \Sigma^\circ\gamma$ asymmetry is predicted in the VDM/quark model approach as negative. In the approach of Żenczykowski[108,121], one expects $\alpha(\Xi^\circ \rightarrow \Sigma^\circ\gamma) \approx -0.3$ to -0.4. The theoretical error on this value could be bigger than ±0.15 because of cancellations in the parity violating amplitudes. The agreement with the experimental value of +0.2±0.32 is satisfactory. It should be noted, however, that a substantial positive experimental value for this asymmetry would add to the puzzle of WRHD, as essentially all models (whether satisfying or violating Hara's theorem) predict negative (and most often large) asymmetries.[114,121] One can see the origin of this result from Table 7.2. Namely, on account of the vanishing contribution from diagram (b2) the sum (Hara's theorem violated) and the difference (Hara's theorem satisfied) of contributions from diagrams (b1) and (b2) are identical. Recent measurements of the $\Xi^- \to \Sigma^-\gamma$ branching fraction limit the size of d' and reduce the cancellation between the W-exchange and single-quark contributions in the parity violating amplitude of this decay. The net effect for a fit is then a more negative (Table 7.1b) asymmetry $\left(\alpha\left(\Xi^\circ \to \Sigma^\circ\gamma\right) \sim -0.4 \text{ to } -0.5\right)$.

7.4.5 $\Xi^- \to \Sigma^-\gamma$

In the symmetry approach, the parity-conserving single-quark $s \to d\gamma$ amplitude is determined by symmetry from NLHD. The measured value of the $\Xi^- \to \Sigma^-\gamma$ branching fraction requires then that the parity violating amplitude is bigger than the parity conserving one. Of the two a priori possible signs for d', the fit rejects the positive value.[108,121] Consequently, the asymmetry was predicted to be large and positive, around +0.6. For a smaller branching fraction like the one reported recently by the E761 collaboration[30] d' must be less negative (≈ -0.25) and an even bigger positive asymmetry is predicted. This constitutes a firm prediction of the VDM/symmetry approach. It is encouraging that it agrees in sign and size with the arguments of Vasanti.[47] In the quark models discussed so far in the literature, our experimental knowledge of NLHD is not used in the determination of the single-quark parity conserving amplitude. As a result, in these models, the $\Xi^- \to \Sigma^-\gamma$ asymmetry and branching fraction are essentially two free parameters.

7.4.6 $\Omega^- \to \Xi^-\gamma$

Since both this decay and the preceding one are due to the same effective $s \to d\gamma$ process, the asymmetry of this decay is predicted in the "symmetry/quark model" approach to be the same as in the $\Xi^- \to \Sigma^-\gamma$ decay. If one relates the hadron-photon coupling in $\Xi^- \to \Sigma^-\gamma$ and $\Omega^- \to \Xi^-\gamma$ decays by the quark model route, the ratio of the branching fractions of the $\Omega^- \to \Xi^-\gamma$ and $\Xi^- \to \Sigma^-\gamma$ decays should be[75] around 3.7. For the fit of Żenczykowski[108,121] this gives for the $\Omega^- \to \Xi^-\gamma$ branching fraction the value of around $1.0*10^{-3}$. On the other hand, if one uses the recent measurement[30] of the $\Xi^- \to \Sigma^-\gamma$ branching fraction as an input, one predicts a smaller value of around $0.45 * 10^{-3}$. Unfortunately, the theoretical error of the "symmetry/quark model" route is here bigger than in the estimates of the other branching fractions: in a similar electromagnetic $\Delta \to N\gamma$ transition there is a 30% discrepancy in amplitude between the symmetry/quark model approach and the experiment. Consequently, a more conservative symmetry expectation for the ratio of the $\Omega^- \to \Xi^-\gamma$ to $\Xi^- \to \Sigma^-\gamma$ branching fractions is 2.5 to 5.2. Taking into account the experimental errors on the $\Xi^- \to \Sigma^-\gamma$ branching fraction one expects the $\Omega^- \to \Xi^-\gamma$ branching fraction to be in the region $(0.25 - 0.75) * 10^{-3}$. This is still in agreement with the experimental upper limit[35] of $0.46 * 10^{-3}$.

In both the quark model and VDM approaches discussed here the SU(3) breaking in parity violating amplitudes was treated additively. On the other hand it is well known that in baryon magnetic moments SU(3) is broken in a nonadditive way.[128,129] A characteristic feature of these nonadditivities is that the contribution from nonstrange quarks is smaller in strange baryons than in nonstrange baryons.[128] If this observation is considered a guide for what may happen in the parity violating WRHD amplitudes, one might expect diminished parity violating amplitudes in Ξ radiative decays leading to smaller asymmetries in $\Xi^\circ \to \Lambda\gamma$ and $\Xi^\circ \to \Sigma^\circ\gamma$ decays.

The above discussion indicates that the combined VDM and symmetry approach does appear quite promising from the phenomenological point of view. Clearly, to make definite statements, we need more precise data on the asymmetries of WRHD. Particularly important here are the $\Xi^\circ \to \Lambda\gamma$ and $\Xi^- \to \Sigma^-\gamma$ asymmetries. A positive sign of the $\Xi^\circ \to \Lambda\gamma$ asymmetry would signify that Hara's theorem is violated (in the SU(3) limit), while a negative sign would indicate that the theorem is satisfied. A remeasurement of the $\Xi^\circ \to \Sigma^\circ\gamma$ asymmetry is also badly needed. If it is indeed positive, as the only measurement[33] seems to indicate, it would add yet another puzzling twist to the long-standing enigma of WRHD.

7.5 *Effective chiral Lagrangians*

Effective chiral Lagrangians and chiral perturbation theory provide a different hadron-level approach to the problem of nonleptonic and weak radiative hyperon decays. However, already in 1971 it has been observed by Holstein[48] that the large negative asymmetry of the $\Sigma^+ \to p\gamma$ decay cannot be explained in a chiral approach as long as it satisfies Hara's theorem in the limit of exact SU(3).

Theoretical work within the framework of chiral perturbation theory has not changed this conclusion. In a recent extensive study along these lines Neufeld[133] gives a couple of numerical predictions obtained when the counterterms of the theory are assumed small (as expected in the approach). He then finds that the asymmetry of the $\Xi^- \to \Sigma^-\gamma$ decay is restricted to the interval (-0.4, +0.3). Using the most recent experimental data on the branching fractions and asymmetries of the $\Xi^\circ \to \Lambda\gamma$ and $\Xi^\circ \to \Sigma^\circ\gamma$ decays he also predicts a negative asymmetry (-0.7 or -0.3) in the $\Lambda \to n\gamma$ decay. Finally and most importantly, he obtains $|\alpha(\Sigma^+ \to p\gamma)| \lesssim 0.2$, in agreement with Hara's theorem and in gross disagreement with data. A similar conclusion has been obtained by Jenkins et. al.[134] It should be stressed that the agreement between theory and experiment claimed by Jenkins et. al.[134] for the remaining (but $\Sigma^+ \to p\gamma$) decays is somewhat illusory: the

8. Analysis of the $s \to d\gamma$ Transition

In the preceding sections two views concerning the violation of Hara's theorem by the quark model have been mentioned. According to the first view, this violation should be regarded as a result specific to the quark model, a result that indicates the need to deepen our understanding of the connection (as provided by VDM) between the quark level description of hadrons and the standard language of the effective hadron-level theory. According to the second view, Hara's theorem must be satisfied and, consequently, its violation in the quark model should be considered as exhibiting a kind of pathology of the quark model that has to be cured. Since the violation originates from the W-exchange diagrams, all those WRHD to which such exchanges can contribute suffer from the above controversy. Fortunately, there does exist a subclass of WRHD to which such processes cannot contribute. These are the $\Xi^- \to \Sigma^-\gamma$, $\Omega^- \to \Xi^-\gamma$ and $\Omega^- \to \Xi^{*-}\gamma$ decays. Due to their quark content (lack of a u quark) diagrams (b), (c) and (d) of Fig 7.1 are absent leaving diagram (a) as the sole contributor. Thus, both experimental and – being free from the above controversy – theoretical studies of these decays are very important. In this Section, we shall present results of the theoretical studies of the $\Xi^- \to \Sigma^-\gamma$ and $\Omega^- \to \Xi^-\gamma$ decays.

As shown in Section 3.2 general arguments of Vasanti[47] lead to the conclusion that the asymmetry in the decays induced by the single-quark $s \to d\gamma$ transition should be positive (Eq.3.16)

$$\alpha = +0.4 \text{ to } +1 \tag{1}$$

(the size depends on whether constituent or current quark masses are used). The VDM fits to all WRHD[108] predict an asymmetry around +0.6 to +0.7 (Table 7.1b) in nice agreement with Eq. 1. Although these fits are fairly successful, they do not address the question of the calculability of the effective $s \to d\gamma$ transition. Rather, they relate one of the corresponding parameters to the experimentally observed parity conserving NLHD amplitudes and thus shift at least part of the question to the NLHD sector where no general consensus exists. Accordingly, the VDM/symmetry approach regards the effective $s \to d\gamma$ transition as incalculable at present: it should become calculable only when a better theoretical understanding of NLHD is achieved.

8.1 *Short distance QCD estimates*

One may be more optimistic with respect to the calculability of the $s \to d\gamma$ decay. The main idea underlying many papers was to calculate the effective $s \to d\gamma$ transition from the first principles of QCD. In QCD, the diagrams of Fig. 7.1. should be understood as indicating the flow of flavor only. When one-gluon exchanges are added, diagram (a) of Fig. 7.1 corresponds then to the sum of diagrams (a1, a2, b, c) shown in Fig. 8.1.

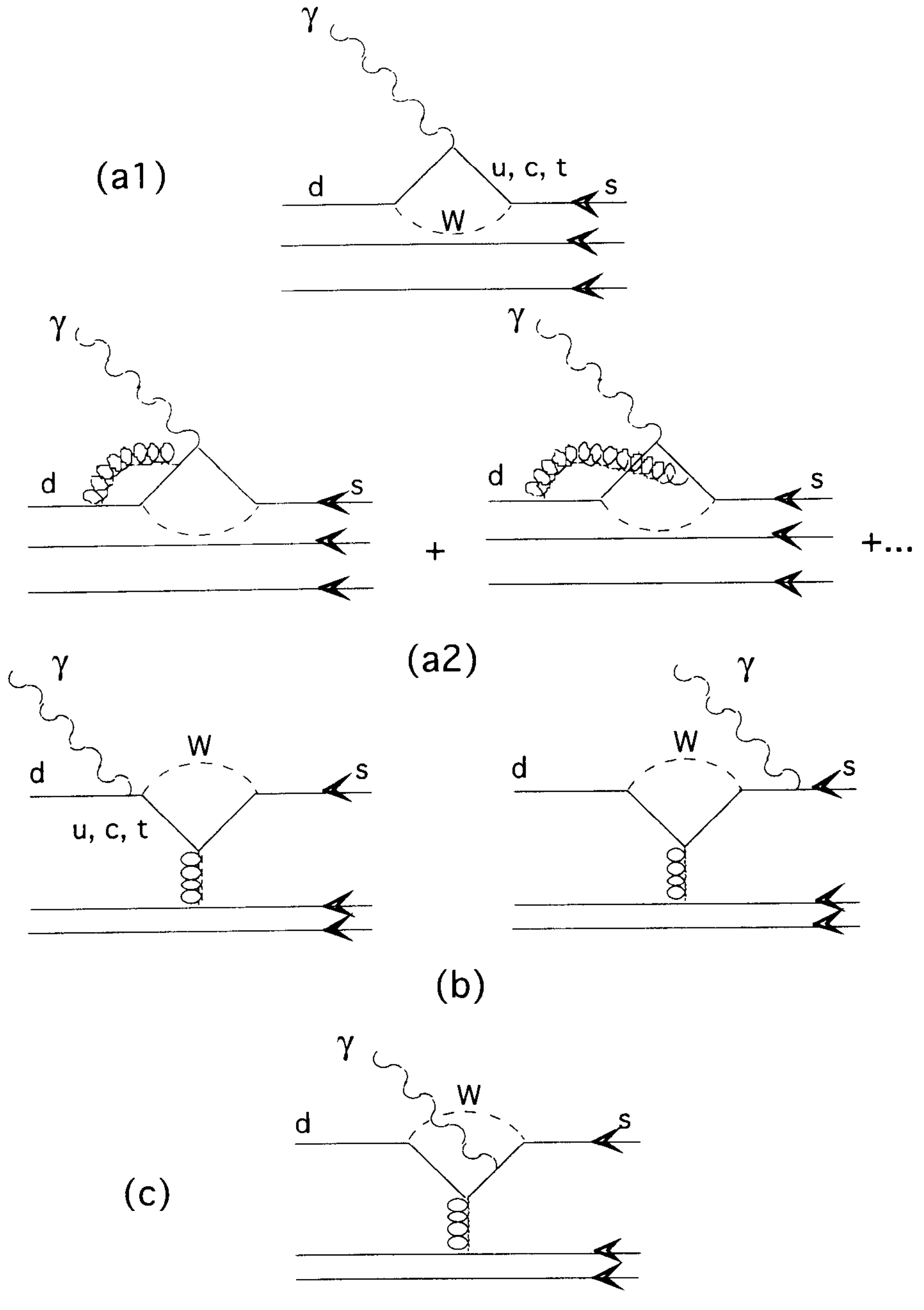

Fig 8.1abc The $s \to d\gamma$ transition in QCD.

The electroweak diagram (a1) of Fig. 8.1 has been widely discussed[72,135] and found totally negligible due to GIM cancellations. Its contribution to the $\Xi^- \to \Sigma^-\gamma$ branching fraction is of order of 10^{-8}. We shall discuss the contribution of the remaining two diagrams (a2) after considering diagrams (b) and (c).

8.1.1 *The penguin*

Analyses of the penguin diagram (b) have been carried out by Kamath[136] and Eeg[137]. Both authors estimated the contribution of the penguin diagrams to the branching fractions of the processes $\Xi^- \to \Sigma^-\gamma$ and $\Omega^- \to \Xi^-\gamma$ as 10^{-8}-10^{-7} and 10^{-5}-10^{-4} respectively if standard values of the penguin coefficients c_5, c_6 (Section 6.2) are used. Even if these contributions were to be underestimated by a factor of 5 or so[91,138] the penguin diagram is totally negligible for the $\Xi^- \to \Sigma^-\gamma$ decay for which the experimental branching fraction is of order of 10^{-4}. Thus, the penguin-induced positive (+0.8) $\Xi^- \to \Sigma^-\gamma$ asymmetry[137] will certainly be swamped by other effects. For the $\Omega^- \to \Xi^-\gamma$ decay the penguin contribution is not that small and might be observable. The predicted asymmetry[137] is -0.9.

8.1.2 *The modified penguin*

A penguin-type diagram with photon emitted from the loop (diagram (c) in Fig. 8.1) has been considered by Eeg[137] and analyzed in more detail by Gaillard, Li and Rudaz.[6] The coefficient function associated with the corresponding operator (not considered in Section 6.2) is not suppressed by GIM cancellations. Furthermore, the use of current quark masses $(m_s \gg m_d)$ appropriate in such a short-distance description ensures comparable strength for the parity violating and parity conserving amplitudes and therefore this operator – if dominant – would yield larger (and negative in explicit calculations) asymmetry parameters for the $\Sigma^+ \to p\gamma$ decay. To estimate if it is really dominant, one needs to know the gluon content of ground state baryons. Using a bag model framework[139], Gaillard, Li, and Rudaz[6] show that the $\Xi^- \to \Sigma^-\gamma$ branching fraction due to this diagram is at most 10^{-6} and thus much too small. A similar result has been found for the $\Omega^- \to \Xi^-\gamma$ decay by Eeg[137] who estimated that diagram (c) contributes at most 10% of the standard penguin contribution. Although explicit calculations indicate that (c) is negligible, it has been argued[6] that such calculations may be very sensitive to the details of the model employed (in bag model large cancellations conspire to give overall small branching fractions). Apart from the question of the overall size of its contribution diagram (c) suffers from yet another drawback, however. This drawback has been identified by Gilman and Wise[75] in the case of a dominant $s \to d\gamma$ effective transition and it persists for diagram (c) even though the spin structure of the 3 quark + gluon wave function is different[139] from ordinary. In their paper, Gilman and Wise have shown that the dominance of the $s \to d\gamma$ effective transition relates the branching fractions R of the $\Sigma^+ \to p\gamma$ and $\Xi^- \to \Sigma^-\gamma$ decays:

$$\left[R\left(\Sigma^+ \to p\gamma\right)/R\left(\Xi^- \to \Sigma^-\gamma\right)\right]_{GW} \sim 10^{-1}. \tag{2}$$

Similar result was found to hold if diagram (c) is dominant:[6]

$$R(\Sigma^+ \to p\gamma)/R(\Xi^- \to \Sigma^-\gamma) \sim 0.2 - 0.4. \qquad (3)$$
(at most 0.7)

This is still a factor of (at least) 10 smaller than the experimental value of around 10. Consequently, diagram (c) - being severely bounded from above by the experimentally known $\Xi^- \to \Sigma^-\gamma$ branching fraction - cannot be dominant in the $\Sigma^+ \to p\gamma$ decay.

8.1.3 *The QCD enhanced electroweak diagram*

We return now to the discussion of the effective single-quark transition of diagrams (a) in Fig. 8.1. Although original calculations[72,135] have indicated that this diagram is totally negligible on account of GIM cancellations, this is no longer so if gluon corrections are considered. The standard GIM power suppression of the type $\left(m_c^2 - m_u^2\right)/m_W^2$ characteristic of diagram (a1) gets replaced by $\ln\left(m_c^2/m_u^2\right)$ factors when two-loop graphs (a2) are considered.[54,140] This enhances the effective $sd\gamma$ vertex by 3 orders of magnitude.[2,54] Phenomenological studies of the implications of this enhancement have been carried out by Singer and collaborators.[2,5,141] The main result of their investigations is that the QCD-enhanced short-distance effective $s \to d\gamma$ transition is still negligible in the $\Xi^- \to \Sigma^-\gamma$ decay, but that its contribution to the $\Omega^- \to \Xi^-\gamma$ decay might be dominant. The branching fraction of the $\Omega^- \to \Xi^-\gamma$ decay is predicted to be around 0.07 * 10^{-3} and it is similar to the estimate (10^{-2} - 10^{-1}) * 10^{-3} of the penguin contribution.[142] Since diagrams (a2) and (b) predict opposite asymmetries for the $\Omega^- \to \Xi^-\gamma$ decay (~+1 for (a2), -0.9 for (b)) the measurement of the characteristics of this decay might discriminate between the two, unless there are other effects that dwarf their contributions.

That this might be so should be obvious from a comparison with the VDM/symmetry prediction which is a factor of at least 5 bigger than either of the contributions of diagrams (a2) and (b).

8.2 *Long distance contributions*

The main difference between the VDM/symmetry approach and the explicit short-distance calculations is that the former does include all the effects (although only implicitly), while the latter may not if the short-distance processes are more severely affected by long-distance effects than we think. In principle, there seems to exist a kind of equivalence between the quark-level and hadron-level approaches. This equivalence has been discussed in the literature.[51,64,65,93,98,112,143]

In practice, hadron-level unitarity-induced lower bounds may exceed short-distance results indicating the failure of our short-distance techniques. This happens, for example, in the $\Xi^- \to \Sigma^-\gamma$ decay. The long-distance, hadron-level unitarity-based prediction for the branching fraction of this process has been calculated by Kogan and Shifman[54] and updated by Singer[5] as $0.18 * 10^{-3}$. This, although a little too big, is still in fair agreement with experiment. This exceeds the purely short-distance effects by a factor of 10. Consequently, the asymmetry of the $\Xi^- \rightarrow \Sigma^-\gamma$ decay should be dominated by hadron-level effects. It has been estimated by Singer[5] as -0.13±0.15. The long-distance calculations of Palle[113] give for it a value of -0.6.

Similar hadron-level unitarity-based calculations of the $\Omega^- \rightarrow \Xi^-\gamma$ branching fraction yield a value of $0.01 * 10^{-3}$, much smaller than the short-distance estimates of the diagrams (a2) and (b). Consequently, it has been argued[2] that in the $\Omega^- \rightarrow \Xi^-\gamma$ decay the short-distance calculation is reliable and that the $\Omega^- \rightarrow \Xi^-\gamma$ branching fraction should be around $0.1 * 10^{-3}$. Unfortunately, however, there is no agreement on the size of long-distance effects in this decay either: calculations of Palle[113] yield a branching fraction of around $0.15 * 10^{-3}$ while the VDM/symmetry approach predicts an even bigger value of at least around $3.7 * R(\Xi^- \rightarrow \Sigma^-\gamma) * 70\% \approx 0.3 * 10^{-3}$ if the recent value for the branching fraction $R(\Xi^- \rightarrow \Sigma^-\gamma)$ is used ($0.7*10^{-3}$ in the original fit of Żenczykowski.[108])

As the above presentation demonstrates, there is no consensus on the $s \rightarrow d\gamma$ single-quark transition. Predictions of various conflicting papers have been gathered in Table 8.1. Comparison data is from Table 2.1. Recall that general arguments of Vasanti[47] predict a significant positive asymmetry for both the $\Xi^- \rightarrow \Sigma^-\gamma$ and $\Omega^- \rightarrow \Xi^-\gamma$ processes. One can see that the measurements of the $\Omega^- \rightarrow \Xi^-\gamma$ branching fraction and of the asymmetries of both the $\Xi^- \rightarrow \Sigma^-\gamma$ and $\Omega^- \rightarrow \Xi^-\gamma$ decays can differentiate between the proposed theoretical approaches.

Table 8.1 Comparison of predictions for the branching fractions (R) (in units of 10^{-3}) and asymmetries of the $\Xi^- \to \Sigma^- \gamma$ and $\Omega^- \to \Xi^- \gamma$ decays.

Process		short-distance: $s \to d\gamma$	short-distance: QCD-corrected [2] $s \to d\gamma$	short-distance: standard penguin[137]	short-distance: modified penguin[6]	long-distance: unitarity[2,5,54]	long-distance: unitarity[113]	VDM/ symmetry	Data
$\Xi^- \to \Sigma^- \gamma$	R	$\sim 10^{-5}$	0.02	$10^{-5} \sim 10^{-4}$	10^{-6}-10^{-3}	0.18	0.04	0.15	0.128 ± 0.023
	α	+1	+1	+0.8	-0.7 +0.47	-0.13	-0.6	+0.7	1.0 ± 1.3
$\Omega^- \to \Xi^- \gamma$	R	$\sim 10^{-4}$	0.07	$10^{-2} \sim 10^{-1}$		0.01	0.15	0.55 ± 0.02	<0.46
	α	+1	+1	-0.9			-0.6	+0.7	

9. Other Approaches

As it has been argued in preceding sections, the problem of WRHD appears to be linked to the question of hadron compositeness and nonlocality. Although most theorists who grappled with the puzzle of WRHD would perhaps agree with this statement, the meanings they assign to the word "nonlocal" may be very different.

In particular, the KLZ mechanism seems to involve the very notion of point in space-time in its translation from the component to the composite level. In other approaches "nonlocality" is usually a synonym for the more traditionally understood "long-distance effects". The term "long-distance effects" does not presuppose the method of their evaluation, either. Thus, various hadron-level calculations are thought of as a phenomenological way of their evaluation. Inclusion of intermediate baryon states is one such way (see comments by Eeg[137] and Palle[113]). Another way is provided by unitarity calculations of Section 4.

In quantum chromodynamics "long-distance effects" translates as "nonperturbative effects". Consequently, apart from lattice simulations, there does not exist any way of actually calculating them from first principles of QCD. A possible phenomenological way to deal with this situation has been proposed in the form of QCD sum rules. This technique has been applied to the WRHD by Khatsymovsky[144-146], Balitsky, Braun and Kolesnichenko[4,147] and Goldman and Escobar.[148]

Since the contribution from the effective $s \to d\gamma$ vertex (Fig. 7.1a) can be estimated in the QCD sum rule approach to be negligible (see also Goldman and Escobar[148]) attention is focused[4,144-147] on the W-exchange processes of Figs 7.1bc. The quarks in the initial and final states of these processes are not (almost) free in the sense of the naive quark (or bag) model (see Section 7.3.2) since they enter the calculations through pointlike baryon currents with proper quantum numbers. Consequently, in the limit of exact SU(3) symmetry, Hara's theorem is satisfied in the QCD sum rule approach. In these papers[4,144,147] it is the contribution of type (b) diagram in Fig. 7.1 that gives rise to large asymmetry in the $\Sigma^+ \to p\gamma$ decay. Unfortunately, the original calculations[144,147] yielded a large *positive* $\alpha(\Sigma^+ \to p\gamma)$ asymmetry (the asymmetry parameter in these papers is defined with an opposite sign to that usually adopted[144,147,149]). Note that in the calculations of Khatsimovsky the asymmetries of the $\Sigma^+ \to p\gamma$ and $\Xi^\circ \to \Lambda\gamma$ decays are of the same sign as it might have been expected. Indeed, Hara's theorem satisfying QCD sum rule approach should be compatible with the standard pole model of Gavela et al.[96] at least as far as the relative sign of these two asymmetries is concerned. As discussed in Section 7.3.4, the equality of these two signs is a general feature of any quark-based approach in which Hara's theorem is enforced by the cancellation of contributions from diagrams (b1) and (b2).

More recent calculations of Balitsky, Braun and Kolesnichenko[4] did give a large negative value for $\alpha(\Sigma^+ \to p\gamma)$. It is not clear, however, what significance should be assigned to this new result.[149,150] In fact, in his last paper Khatsymovsky[145] concluded: "calculations carried out by means of sum rules are very unreliable in the case of complex processes". Predictions of the QCD sum rule approach are compared with the data in Table 9.1. Clearly, to claim real success, calculations[4] should be extended to and explain the remaining WRHD. In particular, their prediction for the asymmetry of $\Xi^\circ \to \Lambda\gamma$ decay is crucial.

Another approach to the problem of nonperturbative effects of quark confinement in QCD has been to connect the large $-N_c$ limit of QCD with the idea of describing baryons as soliton solutions of chiral Lagrangians as originally proposed by Skyrme. An attempt to describe WRHD in the context of the Skyrme model has been made by Kao and Schnitzer.[151] As shown in Table 9.1 the model fails badly when it is compared to the data from Table 2.1.

Table 9.1. Estimates of nonperturbative effects (branching fractions R in units of 10^{-3})

process		QCD W- Khat..[144-146]	sum exchange Balitsky[4]	rules $s \to d\gamma$ Gold..[148]	Skyrme model Kao[151]	Data
$\Sigma^+ \to p\gamma$	R	0.8	0.5 to 1.5	0.047	0.013	1.23±0.06
	$asym$	+1	-0.85±0.15*	-1.0	-0.13	-0.76±0.08
$\Lambda \to n\gamma$	R	2.1-3.1			1.23	1.65±0.12
	$asym$	+0.10 to +0.15			+0.98	
$\Xi^\circ \to \Lambda\gamma$	R	1.1			0.67	1.06±0.16
	$asym$	+0.9			-0.29	+0.43±0.44
$\Xi^\circ \to \Sigma^\circ\gamma$	R				1.04	3.56±0.43
	$asym$				-0.95	+0.2±0.32
$\Xi^- \to \Sigma^-\gamma$	R	0.1 to 0.2		+0.002	0.42	0.128±0.023
	$asym$	+0.4		+0.9	-0.015	1.0±1.3
$\Omega^- \to \Xi^-\gamma$	R	0.2 to 0.4				<0.46

*originally predicted positive[147]

10. Conclusions

Experiments have continued to confirm the large negative value of the asymmetry parameter in $\Sigma^+ \to p\gamma$. This fact, which signaled the controversial nature of WRHD through its violation of Hara's theorem, remains with us and refuses to be ignored. New measurements have improved our knowledge of the branching fractions of $\Lambda \to n\gamma$, and $\Xi^- \to \Sigma^-\gamma$. The measured $\Xi^- \to \Sigma^-\gamma$ branching fraction approaches the unitarity bound while the limit on the $\Omega^- \to \Xi^-\gamma$ branching fraction is within a factor of about 5 of the unitarity bound.

What are the prospects for further experimental progress? The Fermilab charged hyperon beam still offers the best possibility of measuring the $\Lambda \to n\gamma$ asymmetry parameter through the decay chain $\Xi^- \to \Lambda\pi^-$, $\Lambda \to n\gamma$ which provides a polarized Λ°. Improvements on the only existing measurement of the $\Xi^- \to \Sigma^-\gamma$ asymmetry parameter are certainly possible. More challenging but also possible are the measurements of the $\Omega^- \to \Xi^-\gamma$ branching fraction and asymmetry parameter. It is the Fermilab charged hyperon beam that offers the only possibility of measuring these decays in the foreseeable future. The only limitation will be the availability of beam time and the ingenuity of the experimenters.

Present measurements of the decays $\Xi^{\circ} \to \Lambda\gamma$ and $\Xi^{\circ} \to \Sigma^{\circ}\gamma$ are severely limited by statistics, each having only about 100 events. A study of these reactions is being planned as part of the Fermilab program[58] which will investigate CP violation in the neutral kaon system. It is possible that the number of events could be increased by one or two orders of magnitude. As we have argued at length in Sections 7.3.4 and 7.4, precise measurements of the asymmetries of these decays are very important: they could help settle the theoretical controversy concerning the validity of Hara's theorem in the SU(3) limit. In particular, measurement of the $\Xi^{\circ} \to \Lambda\gamma$ asymmetry is absolutely crucial.

At present there is no consensus among theorists as to how the quark-model violation of Hara's theorem should be interpreted. The two contradictory points of view are:

(1) Violation of Hara's theorem is a pathological feature of the quark model. Consequently, the quark-level approaches have to be somehow modified to ensure that in the SU(3) limit this theorem is satisfied. This would presumably require some sort of wave function modification as it is the group-theoretical structure of the composite state of three quarks that directly leads to the result in question. No such proposal has been put forward that would describe experimental data in a satisfactory way.

(2) Violation of Hara's theorem is not an artifact of the quark model, but a feature of nature itself. This would mean that we lack a deep enough understanding of the quark model and its connection to the standard hadron-level language. Through the Kroll-Lee-Zumino scheme this seems to entail that vector meson dominance is more than a mere phenomenological model. The quark/vector-dominance model describes the data better than models in which Hara's theorem is enforced.

Whether one accepts (1) or (2), it is clear that what radiative hyperon decays probe is the composite nature of hadrons.

The problem of radiative hyperon decays bears some resemblance to the story of baryon magnetic moments. In fact, it is precisely through the proton magnetic moment measurement of Stern, Estermann and Frisch[152] that hadron compositeness was revealed to us for the first time. At that time the experimental results were not interpreted as indicative of proton substructure, however. Although the measured value of the magnetic moment was in sharp disagreement with (hadron-level) theoretical expectations,[153] the explanation of baryon magnetic moments in terms of constituents was proposed only thirty years later. The successful explanation was based on the dubious assumption of the additivity of (Dirac) magnetic moments of three free quarks in spin-flavor symmetric state. It is precisely the same set of qualitative features of the three-quark states in the quark model that leads to the violation of Hara's theorem in weak radiative hyperon decays. For this reason we believe that study of these decays should teach us a lot about how quarks combine to form hadrons. The present results and the prospect of new experiments should provide a stimulus for further theoretical progress.

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