WEAK RADIATIVE DECAYS OF HYPERONS: QUARK MODEL AND NONLOCALITY

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It is proved that symmetry structure of the parity-violating amplitudes of weak radiative hyperon decays in the Vector-Meson Dominance (VMD) approach, and the violation of Hara's theorem in particular, are also obtained when direct coupling $e_q\ \bar{q}\gamma_\mu q A^\mu$ of photon to quarks is used in place of VMD (with calculations performed in the limit of static quarks). Thus, violation of Hara's theorem in VMD-based models does not result from the lack of gauge invariance. It is further shown that, in the static limit of the quark model, the current-algebra commutator term in the parity-violating amplitudes of nonleptonic hyperon decays and the parity-violating $\Sigma^+ \to p\gamma$ decay amplitude are proportional to each other. As a result, Hara's theorem may be satisfied in this limit if and only if the contribution from the current-algebra commutator in nonleptonic hyperon decays is zero. Violation of Hara's theorem is traced back to the nonlocality of quark model states in the static limit. It is argued that the ensuing intrinsic baryon nonlocality does not have to be unphysical. It is stressed that the measurement of the $\Xi^0 \to \Lambda\gamma$ asymmetry will provide very important information concerning the presence or absence of nonlocal features in parity-violating photon coupling to baryons at vanishing photon momentum. If the $\Xi^0 \to \Lambda\gamma$ asymmetry is found negative, Hara's theorem is satisfied but the gauge-invariant quark model machinery predicting its violation must miss some contribution, or be modified. If experiment confirms positive $\Xi^0 \to \Lambda\gamma$ asymmetry, then, most likely, Hara's theorem is violated. Although positive $\Xi^0 \to \Lambda\gamma$ asymmetry admits of the possibility that Hara's theorem is satisfied, this alternative is in disagreement with hints suggested by the similarity of photon and vector-meson couplings and the observed size of parity-violating nuclear forces.

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1. Introduction

Thanks to the experimental programmes pursued at Fermilab and CERN, important data on weak radiative hyperon decays (WRHDs) should be available within the next year or two. These data will provide crucial input, which will direct further attempts to understand WRHDs.

Theoretical approaches to WRHDs may be divided into two classes according to whether they (1) satisfy or (2) violate the theorem due to Hara [1], which states that the parity-violating (p.v.) amplitude of the $\Sigma^+ \to p\gamma$ decay should vanish in the SU(3) limit. Since the assumptions of Hara's theorem are fundamental (*i.e.* CP conservation, gauge-invariant local field theory at hadron level, and the absence of exactly massless hadrons), one may be tempted to discard class (2), and claim that any obtained violation of the theorem must result from an unjustified assumption or an erroneous calculation. Yet, despite the fundamental nature of the assumptions of Hara's theorem, the data seem to favour its violation [2]. Clearly, it may happen that some of these data turn out erroneous, and that the combined set of WRHD data finally agrees with the theorem. However, irrespectively of whether one rejects or admits of the possibility that Hara's theorem may be violated in Nature, it is important to identify the origins of the violation in any given calculation. Only mathematically precise analyses may provide us with a deeper understanding of the problem, and its resolution when crucial experimental data become available.

In the quark model of Kamal and Riazuddin (KR) [3], Hara's theorem is violated in the SU(3) limit. It has been claimed that this result is due to KR calculation being not gauge-invariant (see [4] and references therein). As pointed out in Ref. [5] (see also [6]), such claims are based on logically incorrect inferences. Technically speaking, violation of Hara's theorem in KR originates from a contribution in which the intermediate quark enters its mass shell, thus satisfying the condition for a regular free particle [7]. Since quarks are not such particles, the KR calculation must be considered unphysical, *if taken literally.* On the other side, however, it indicates what would be needed for the violation of Hara's theorem to occur in Nature. Since entering mass shell is a feature of asymptotic states, the KR calculation hints that the violation of Hara's theorem requires some kind of nonlocality. Connection with nonlocality can also be anticipated from the example of Ref. [9], in which it is proved at hadron level that the violation of Hara's theorem may occur only if the (conserved) electromagnetic axial baryon current exhibits some degree of nonlocality (see also [5]).

It may be argued that one should replace the KR approach with a model in which quarks are confined to a small region of space. Hadrons should be then well described by an effective local field theory, and Hara's theorem *must* be satisfied as its violation then requires the presence of unobserved exactly massless hadron [2]. The problem is, however, that quark unobservability may be taken care of only using models based on what is *expected* of confinement, not on its calculable properties: one cannot trace if and how quark unobservability modifies KR results. This means that the correctness of these expectations cannot be proved or disproved.

Thus, we are stuck in a stalemate: in order to resolve the puzzle, we have to take quark unobservability into account. Yet, we do not know how to do that properly. A possible way out consists in finding a description in which use of free or confined intermediate quarks is altogether avoided. Such reasoning led to the idea of the $\mathrm{SU}(6)_W\times$ VMD approach of Refs. [2, 10, 11], in which photons couple to hadrons always through intermediate vector mesons (VMD is known to work extremely well). Then, all unknowns related to quark-level problems are hidden in the p.v. meson–baryon couplings. The merit of this approach is that the latter couplings are experimentally accessible in nuclear p.v. processes. On the other hand, the use of VMD may be considered questionable: Ref. [4] attributes violation of Hara's theorem in [2, 10] to the lack of gauge invariance. The question of gauge invariance arises when VMD is understood in a dynamical sense with vector mesons mediating the coupling. Still, even if one rejects the KLZ scheme [12,13] ensuring gauge invariance of VMD, the explanation of the violation of Hara's theorem in Refs. [2,10] by gauge noninvariance of the underlying calculations is incorrect. Proving this is one of the goals of the present paper.

In this paper (Sections 2,3) it is shown that the basic results of Refs. [2,10] are *independent* of the above dynamical understanding of VMD and hold also for manifestly gauge-invariant direct photon–quark coupling. Consequently, the violation of Hara's theorem obtained in [2, 10] has nothing to do with gauge noninvariance.

In Section 4 the whole scheme and, in particular, the pattern of WRHD asymmetries are cross-checked against the PCAC approach to NLHDs. It is shown that if the Current Algebra (CA) commutator in NLHDs is nonzero, the approach leads to the violation of Hara's theorem, while predicting large positive asymmetry of the $\Xi^0 \to \Lambda\gamma$ decay. It is stressed that negative $\Xi^0 \to \Lambda\gamma$ asymmetry (automatically consistent with Hara's theorem) would present a serious problem for the quark model. The decisive role of the sign of the $\Xi^0 \to \Lambda\gamma$ asymmetry as far as the issue of the violation of Hara's theorem is concerned can also be seen from Table I where present experimental branching ratios and asymmetries of WRHDs are gathered together with the predictions of two typical approaches to WRHDs.

TABLE I

Asymmetries and branching ratios of four most important WRHDs: HS — Hara's theorem Satisfied, Ref. [14]; HV — Hara's theorem Violated, Ref. [2].

Decay	Asymmetries			Branching ratios$\times 10^3$		
	experiment	HS	HV	experiment	HS	HV
$\Sigma^+ \to p\gamma$	-0.76 ± 0.08	$-0.80^{+0.32}_{-0.19}$	-0.95	1.23 ± 0.06	$0.92^{+0.26}_{-0.14}$	$1.3 - 1.4$
$\Lambda \to n\gamma$		-0.49	$+0.8$	1.75 ± 0.15	0.62	$1.4 - 1.7$
$\Xi^0 \to \Lambda\gamma$	$+0.43 \pm 0.44$	-0.78	$+0.8$	1.06 ± 0.16	3.0	$0.9 - 1.0$
$\Xi^0 \to \Sigma^0\gamma$	-0.63 ± 0.09	-0.96	-0.45	3.34 ± 0.10	7.2	$4.0 - 4.1$

Providing an interpretation for the origin of the violation of Hara's theorem in the quark model constitutes our another aim. In Section 5 we reflect on the concept of quark position. This reveals that for static quarks (relevant for the calculations in the SU(3) limit), the composite quark-model states possess nonlocal quantum properties. Thus, the violation of Hara's theorem is traced to intrinsic nonlocality of baryons. The question whether this property of the quark model constitutes a drawback or a virtue is discussed and it is argued that in the limit of zero photon momentum such nonlocality does not have to be unphysical. It is stressed that the sign of the $\Xi^0 \to \Lambda\gamma$ asymmetry will provide crucial information on whether the difficulties with WRHDs are due to a serious problem in the quark model or to intrinsic baryon nonlocality.

In Section 6 we consider the issue whether a positive sign of the $\Xi^0 \to \Lambda\gamma$ asymmetry requires violation of Hara's theorem. It is pointed out that the connection between WRHDs and the (observed) p.v. effects in NN interactions constitutes an additional hint against Hara's theorem.

Our conclusions are given in Section 7.

2. Quark model in the static limit

In this section we discuss the essential steps of the derivation of our results (as well as those of [10] and of paper [15] by Desplanques, Donoghue and Holstein (DDH)). It is known that explanation of the observed pattern of the WRHD branching ratios requires that the dominant contribution should come from $us \to du\gamma$ processes shown in Fig. 1.

Other possible diagrams have been estimated in various papers as negligible [2]. We want to evaluate the joint effects of the process of W-exchange and the direct coupling of photon to quarks in an approach in which explicit treatment of intermediate quarks as free particles (present in KR) is avoided or at least confirmed in an independent way. This was also the original mo-

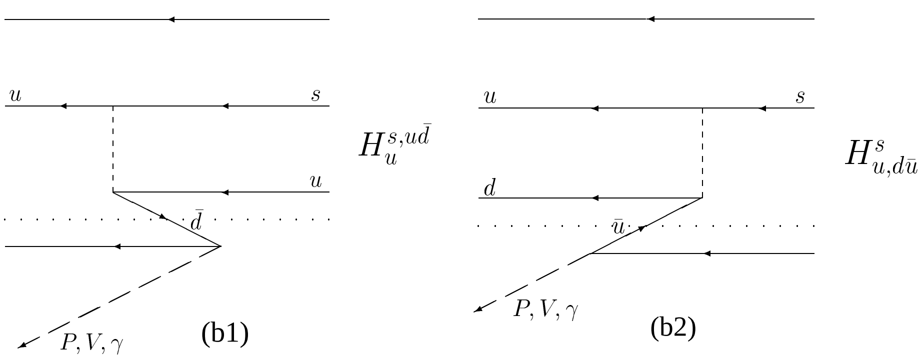

Fig. 1. Quark diagrams for W-exchange contributions to parity-violating amplitudes of WRHDs.

tivation for the VMD approach of Ref. [10]. All complications from strong interactions reside either in the wave functions of external baryons, or are taken into account by the completeness of the set of states intermediate between photon emission and W-exchange. Couplings of pseudoscalar and vector mesons are described by similar diagrams with the photon replaced by a meson and a suitable choice of current in the coupling. Although one may criticise the latter assumption on the grounds that mesons are not pointlike, symmetry structure of these couplings suffices to derive the results of [15]. Since ultimately we are interested in the limit of zero mass difference between initial and final baryons, our calculations shall be done in the extreme nonrelativistic (static) limit, with momenta of both photon and individual quarks approaching zero. Taking this limit does not violate gauge invariance: one might keep small terms of higher order in momenta and drop them at the end. The validity of the assumption of static quarks shall be discussed later. Below we give a schematic presentation of quark model calculations. Practical equivalence of SU(6)$_W$ and the quark model technique was discussed in [15].

2.1. Strong and electromagnetic couplings

Parity-conserving (p.c.) interactions of fields with quark currents are given by:
for pseudoscalar mesons

$$g_P \ \bar{q}_l \gamma_5 q_m \ P^{lm} \, , \tag{1}$$

for vector mesons

$$g_V \ \bar{q}_l \gamma_\mu q_m \ V^{lm,\mu} \, , \tag{2}$$

for photons

$$e_l \; \bar{q}_l \gamma_\mu q_l \; A^\mu \,, \tag{3}$$

where $l, m = u, d, s$ label quark–antiquark operator fields q_l. Colour indices are suppressed.

In the static limit, neglecting inessential factors, one may rewrite the relevant structures in Eqs. (1)–(3) as follows (the direction of the 3rd axis is defined as the direction of meson/photon momentum $\boldsymbol{k}$ before taking the limit $k_\mu \to 0$):

for pseudoscalar mesons

$$[a^+(\bar{m}l(\uparrow\downarrow - \downarrow\uparrow)) + a(\bar{l}m(\uparrow\downarrow - \downarrow\uparrow))] \; a^+(P(m\bar{l}))\,, \tag{4}$$

(+ terms not involving $a^+(P(m\bar{l}))$);

for vector mesons

$$-[a^+(\bar{m}l(\uparrow\uparrow)) + a(\bar{l}m(\downarrow\downarrow))] \; a^+(V(m\bar{l},-1))\,, \tag{5}$$

(+ terms not involving $a^+(V(m\bar{l},-1))$);

for photons

$$-[a^+(\bar{l}l(\uparrow\uparrow)) + a(\bar{l}l(\downarrow\downarrow))] \; a^+(A(-1)) \tag{6}$$

(+ terms not involving $a^+(A(-1))$).

Here $a^+(\bar{m}l(\uparrow\downarrow - \downarrow\uparrow))$ denotes difference of two terms consisting of products of creation operators of quark l and antiquark $\bar{m}$ in spin states described by arrows:

$$a^+(\bar{m}l(\uparrow\downarrow - \downarrow\uparrow)) = a^+(\bar{m}\uparrow)a^+(l\downarrow) - a^+(\bar{m}\downarrow)a^+(l\uparrow)\,. \tag{7}$$

In $a(\bar{l}m(\uparrow\downarrow - \downarrow\uparrow)$ we have annihilation operators of antiquark $\bar{l}$ and quark m.

The ordering of indices corresponds to the ordering of quark creation (annihilation) operators, which satisfy standard anticommutation relations. Furthermore, $a^+(P(m\bar{l}))$ is a creation operator of pseudoscalar field describing meson composed of quark m and antiquark $\bar{l}$, while $a^+(V(m\bar{l},-1))$ corresponds to vector meson with spin projection down. Similarly, $a^+(A(-1))$ describes creation of photon with its spin directed along the negative axis. For future discussion, we note that for the coupling

$$g'_V \; \bar{q}_l i\sigma_{\mu\nu} k^\nu q_m \; V^{lm,\mu} \tag{8}$$

at small momentum $k^\nu, k^2 \approx 0$ one obtains the following structure ($k = k^0 \approx k^3$)

$$-k[a^+(\bar{m}l(\uparrow\uparrow)) - a(\bar{l}m(\downarrow\downarrow))] \; a^+(V(m\bar{l},-1))\,. \tag{9}$$

The plus sign in Eq. (5) is here replaced by the minus sign. Corresponding formulas for photons are obtained from Eqs. (8),(9) by taking $l = m$ and performing substitutions $V \to A$, $g'_V \to \kappa_l$, with κ_l being the anomalous magnetic moment of quark q_l. Formulas (4-6,9) may be conveniently represented in terms of quark diagrams shown below the dotted lines in Fig. 1. The ordering of quark lines (from top to bottom) corresponds to the ordering of quark creation (annihilation) operators.

2.2. Weak interactions

Starting from the standard $(V - A) \times (V - A)$ weak interaction, after expressing axial and vector weak currents relevant for the transition $us \to du$ in a way analogous to that given in the previous subsection, the contribution from the p.v. part of W-exchange is proportional in the static limit to

$$H_W^{\text{p.v.}} = H_u^{s,u\bar{d}} + H_d^{u,s\bar{u}} + H^s_{u,d\bar{u}} + H^u_{d,u\bar{s}}\,, \tag{10}$$

where

$$\begin{aligned} H_u^{s,u\bar{d}} &= a^+(u\downarrow)a(s\downarrow)a(u\uparrow)a(\bar{d}\downarrow) - a^+(u\uparrow)a(s\uparrow)a(u\downarrow)a(\bar{d}\uparrow) \\ &+ a^+(u\uparrow)a(s\downarrow)a(u\uparrow)a(\bar{d}\uparrow) - a^+(u\downarrow)a(s\uparrow)a(u\downarrow)a(\bar{d}\downarrow) \end{aligned} \tag{11}$$

and

$$\begin{aligned} H^s_{u,d\bar{u}} &= a^+(u\uparrow)a(s\uparrow)a^+(d\downarrow)a^+(\bar{u}\uparrow) - a^+(u\downarrow)a(s\downarrow)a^+(d\uparrow)a^+(\bar{u}\downarrow) \\ &+ a^+(u\uparrow)a(s\downarrow)a^+(d\downarrow)a^+(\bar{u}\downarrow) - a^+(u\downarrow)a(s\uparrow)a^+(d\uparrow)a^+(\bar{u}\uparrow) \end{aligned} \tag{12}$$

with $a^+(u)a(s)$ describing $s \to u$ transition of one of the quarks. The above two expressions may be conveniently represented diagrammatically as shown above the dotted lines in Fig. 1. The ordering of three creation (or three annihilation) operators corresponds to the ordering of relevant quark lines (from top to bottom). Note that in this language the description of p.v. processes is made possible by the appearance of (negative parity) antiquarks [15]. Additional four terms with $d \leftrightarrow s$ are to be added to Eq. (10) if $ud \to su$ processes are to be described as well. This makes the whole interaction symmetric under $d \leftrightarrow s$.

Starting from the $(V - A) \times (V - A)$ interaction, we have also found by explicit calculation that the corresponding expression for $\bar{u}\bar{s} \to \bar{d}\bar{u}$ is obtained from Eq. (10) by replacing quarks with antiquarks and vice versa (without acting on the spin degrees of freedom), and by changing the overall sign of the Hamiltonian, as expected. Similarly, explicit calculation has shown that the corresponding p.c. Hamiltonian does not change its sign after charge conjugation. These checks have confirmed that the calculation is performed correctly and does not involve any artificial CP-violating effects, forbidden by the assumptions of Hara's theorem.

2.3. Final prescription

The final calculation to be performed consists in the evaluation of matrix elements of quark vector or pseudoscalar current by sandwiching it in between external baryonic states, described by standard spin-flavour wave functions of ground-state baryons and modified by weak interaction of Eq. (10).

Calculations of the p.v. NLHD and WRHD amplitudes proceed in two steps:

(1) evaluation of the admixture of $q\bar{q}$ pairs in qqq baryons generated by Eq. (10), and

(2) calculation of the matrix elements of currents $\bar{q}\gamma_5 q$ and $\bar{q}\gamma_\mu q$ in between states with these admixtures.

For p.v. WRHD amplitudes the matrix elements to be evaluated are of the form

$$\begin{aligned} &\langle q_1'q_2'q_3'|\ H_W^{\rm p.v.}\ \bar{q}\gamma_\mu q\ |q_1q_2q_3\rangle \qquad &(\mathrm{b1})\,, \\ &\langle q_1'q_2'q_3'|\ \bar{q}\gamma_\mu q\ H_W^{\rm p.v.}\ |q_1q_2q_3\rangle \qquad &(\mathrm{b2})\,. \end{aligned} \tag{13}$$

For NLHDs, replace the vector current $\bar{q}\gamma_\mu q$ with the pseudoscalar current $\bar{q}\gamma_5 q$ (thus, NLHDs and WRHDs are related through the *spin properties of the quark model*). Energy denominators have been suppressed: as they correspond to energy difference between qqq and $qqq\bar{q}q$ states, they are identical for diagrams (b1) and (b2). The diagrammatic notation of Fig. 1 helps in the calculations which are a little tedious but straightforward. They may be simplified by exploiting total symmetry of ground-state baryon spin-flavour wave function. Thanks to the symmetry of the wave functions, it is sufficient to evaluate the contribution when photon (meson) emission proceeds from quark creation or annihilation operator (from the $\bar{q}\gamma_\mu q$ current) contracted with the third (by definition) quark in a baryon. The third quark is also one of quarks between which W-exchange occurs. Because of wave-function symmetry it is sufficient to consider contributions in which the other quark undergoing weak interactions is quark number two. The resulting diagrams to be evaluated are precisely the diagrams of Fig. 1, in which quarks number 1,2,3 are ordered from top to bottom. When the actions of weak Hamiltonian on external states are worked out, one obtains:

$$\begin{aligned} &\langle q_1q_2q'\bar{q}'q_3'|\ \bar{q}\gamma_\mu q\ |q_1q_2q_3\rangle \qquad &(\mathrm{b1})\,, \\ &\langle q_1q_2q_3'|\ \bar{q}\gamma_\mu q\ |q_1q_2q'\bar{q}'q_3\rangle \qquad &(\mathrm{b2}) \end{aligned} \tag{14}$$

with $\bar{q}\gamma_\mu q$ acting on the bottom quark line, as shown in Fig. 1.

The above equations are completely analogous to those appearing in the standard quark-model calculation of baryon magnetic moments: the latter are evaluated from the nonrelativistic reduction of

$$\langle q_1q_2q_3|\ \bar{q}\gamma_\mu q\ |q_1q_2q_3\rangle\,, \tag{15}$$

where, because of the spin-flavour symmetry of external three-quark states, contribution from the third quark only needs to be calculated.

3. Pattern of parity-violating amplitudes and asymmetries

Using direct photon–quark coupling of Eqs. (3),(6) and the weak Hamiltonian of Eq. (10), it is straightforward to evaluate the p.v. WRHD amplitudes. In this calculation, relative signs of various contributions are fixed by the employed group-theoretic structure. In particular, for the $\Sigma^+ \to p\gamma$ p.v. amplitude one obtains

$$A(\Sigma^+ \to p\gamma) = -\frac{1}{3\sqrt{2}}b_\gamma - \frac{1}{3\sqrt{2}}b_\gamma\,, \tag{16}$$

where the first term comes from diagram (b1) and the other one from diagram (b2). The overall size of the amplitude is taken care of by b_γ, which is proportional to electric charge e and Fermi coupling constant $G_{\rm F}$. The subscript γ stresses that the calculation has been performed using direct photon–quark coupling, *i.e.* without any intermediate vector meson. Normalisation of numerical factors in front of b_γ has been chosen so that when the subscript $_\gamma$ is omitted, one recovers formulas of Table 7.2 in Ref. [2] (Table 3 in Ref. [7]). These formulas are repeated here in Table II. Relative

TABLE II

Amplitudes b_1 and b_2.

Decay	Diagram (b1)	Diagram (b2)
$\Sigma^+ \to p\gamma$	$-\frac{1}{3\sqrt{2}}b_\gamma$	$-\frac{1}{3\sqrt{2}}b_\gamma$
$\Lambda \to n\gamma$	$+\frac{1}{6\sqrt{3}}b_\gamma$	$+\frac{1}{2\sqrt{3}}b_\gamma$
$\Xi^0 \to \Lambda\gamma$	0	$-\frac{1}{3\sqrt{3}}b_\gamma$
$\Xi^0 \to \Sigma^0\gamma$	$\frac{1}{3}b_\gamma$	0

sizes and signs of all contributions in these tables are a consequence of the group-theoretic assumption of the quark model and *not* of the assumption of VMD in its dynamical sense. This should be obvious from the comparison of interactions of vector meson and photons with quarks in Eqs. (5),(6), where vector currents $\sum_l e_l \bar{q}_l \gamma_\mu q_l$ and $\sum_{l,m} g_V \bar{q}_l \gamma_\mu (\lambda_3 + \lambda_8/\sqrt{3})_{lm} q_m$ are proportional. Clearly, the evaluation of matrix elements of a current in between hadronic states is completely oblivious to the nature of the field coupled to it. With quark currents for photon and vector-meson couplings identical, the calculated symmetry structure of p.v. couplings of vector mesons and photons must be the same. The only difference is the size of the coupling. Thus, VMD may be also understood as a two-step *merely technical* prescription of substitution:

(i) evaluate the couplings of vector mesons to hadrons in the quark model,

(ii) in order to see what direct photon–quark coupling would give, perform the *substitution* (here for the ρ meson): $\rho_\mu \to \frac{e}{g_\rho} A_\mu$.

In this sense, the results of VMD calculations of Refs. [2,10,11] cannot be any less gauge-invariant than those obtained from minimal direct photon–quark coupling.

From Eq. (16) it follows that the p.v. amplitude for the $\Sigma^+ \to p\gamma$ decay is equal to $-\frac{2}{3\sqrt{2}} b_\gamma$. Hara's theorem might be satisfied in the SU(3) limit provided b_γ vanishes. The latter is not true in our quark model calculations in the static limit.

p.v. amplitudes of the remaining WRHDs (Table II) are also proportional to b_γ. This means that in order for Hara's theorem to be satisfied, p. v. amplitudes of *all* WRHDs must vanish in the SU(3) limit [16]. Modifications of the predictions of the static limit so as to take into account SU(3) breaking shall be discussed in Section 6. In Sections 3–5 we accept that experimental asymmetries of WRHDs at $k_\mu \neq 0$ can be well approximated by the static quark model prescription. A similar assumption was used in Ref. [15] when considering p.v. couplings of mesons to baryons. Further discussion of this assumption shall be given later.

With the coupling of photon to quark vector current proportional to the coupling of U-spin-singlet vector meson to that current, it is obvious that all the relative signs of p.v. amplitudes calculated in the present scheme with direct photon–quark coupling must be proportional to the amplitudes of the $SU(6)_W$ + VMD approach of [2,10], *i.e.* they are given by the sums of entries in columns (b1) and (b2) in Table II. (For clarity and in order to exhibit some of the $s \leftrightarrow d$ symmetry properties, the contributions from vector currents with definite quark content are given in Table III for the $\Sigma^+ \to p$ transitions with p.v. weak interaction in the initial or final baryon.)

Description of WRHDs requires knowledge of p.c. amplitudes as well. These amplitudes have been evaluated in many approaches, such as different versions of the pole or quark models. In the pole model the dominant contribution comes from intermediate ground states. One may also identify

TABLE III

Weights of amplitudes b_1 and b_2 for $\Sigma^+ \stackrel{\bar{q}\gamma_\mu q}{\longrightarrow} p$ in the presence of strangeness-changing p.v. weak interaction (W-exchange) in initial or final baryon for vector currents with well-defined $q\bar{q}$ content.

Current $\bar{q}\gamma_\mu q$	Diagram (b1) $\langle\Sigma^+\| \bar{q}\gamma_\mu q\, H_W^{\text{p.v.}} \|p\rangle$	Diagram (b2) $\langle\Sigma^+\| H_W^{\text{p.v.}}\, \bar{q}\gamma_\mu q \|p\rangle$
$\bar{u}\gamma_\mu u$	$-\frac{1}{3\sqrt{2}}$	$-\frac{1}{3\sqrt{2}}$
$\bar{d}\gamma_\mu d$	$+\frac{1}{3\sqrt{2}}$	0
$\bar{s}\gamma_\mu s$	0	$+\frac{1}{3\sqrt{2}}$
$\frac{2\bar{u}\gamma_\mu u-\bar{d}\gamma_\mu d-\bar{s}\gamma_\mu s}{\sqrt{6}}$	$-\frac{1}{2\sqrt{3}}$	$-\frac{1}{2\sqrt{3}}$

a correspondence between the pole and quark model prescriptions: up to some details, the symmetry structure of the p.c. WRHD amplitudes is similar in both approaches. Therefore, p.c. WRHD amplitudes may be safely described in terms of the pole model. Reliability of the model is confirmed by its successful description of p.c. NLHD amplitudes (*cf.* [2]).

From the pole model we know the approximate sizes and signs of the p.c. WRHD amplitudes. In our conventions, these signs are $-,-,+,+$ for $\Sigma^+ \to p\gamma$, $\Lambda \to n\gamma$, $\Xi^0 \to \Lambda\gamma$, and $\Xi^0 \to \Sigma^0\gamma$ respectively, see *e.g.* [2]. Upon taking the sum of contributions from diagrams (b1) and (b2) (Table II), the asymmetries of the $\Sigma^+ \to p\gamma$ and $\Xi^0 \to \Sigma^0\gamma$ decays turn out to be of the same sign, while for $\Xi^0 \to \Lambda\gamma$ the asymmetry is of the opposite sign [2, 10]. Relative sizes of p.v. and p.c. amplitudes are such that if one of the asymmetries is large, all of them are large. In particular, in Ref. [2] the asymmetry in the $\Xi^0 \to \Sigma^0\gamma$ decay was predicted to be around -0.45 (or slightly more negative, see [17]), away from its first measurement of $+0.20 \pm 0.32$ [18]. A recent experiment [19] gave $\alpha(\Xi^0 \to \Sigma^0\gamma) = -0.63 \pm 0.09$, confirming that the $\Xi^0 \to \Sigma^0\gamma$ asymmetry is significantly negative (and pointing out an error in Ref. [18]). With substantially negative asymmetry of $\Xi^0 \to \Sigma^0\gamma$, using the experimental branching ratio and the pole model prediction for the p.c. amplitude of this decay, one can determine b_γ and predict the asymmetry of the $\Xi^0 \to \Lambda\gamma$ to be large and positive. Although in Ref. [2] this prediction was obtained in the framework of a VMD approach, *no dynamical VMD is needed to obtain this result.* Here this conclusion follows from direct photon–quark coupling in the static limit of the quark model.

In fact, the sizes and relative signs of p.v. and p.c. WRHD amplitudes can be fixed without any recourse to experimental data on WRHDs or vector meson couplings to baryons. One needs to know the data on NLHDs only.

Namely, in both p.v. and p.c. NLHD amplitudes there occurs the same interaction of quark pseudoscalar current with an emitted pseudoscalar meson, and, similarly, in both p.v. and p.c. WRHD amplitudes there occurs the same interaction of quark vector current with photon. With spin-flavour symmetry providing the connection between the two currents, the NLHD and WRHD W-exchange contributions to amplitudes become related (separately in p.c. and p.v. sectors) and the prediction for WRHD asymmetries and branching ratios becomes absolute. Roughly speaking, one has to remove the factor of g_P from the NLHD amplitudes, perform the appropriate symmetry transformation (*i.e.* replace the pseudoscalar current with the vector one), and multiply the results by the electric charge. Up to some details, this prediction is numerically the one given in Refs. [2, 10, 11]. Further discussion of the connection with NLHDs is given in the next section.

At this point one may ask what part of the results of [2, 10, 11] really depends on the assumption of VMD understood in its dynamical sense. The answer is: not much. Since formulas for the p.v. amplitudes in the present scheme and in Refs. [2,10,11] are identical in the SU(3) limit, the only place where dynamical VMD effects may be somewhat important is the description of p.c. amplitudes. In fact, in Refs. [2,11] these amplitudes are assumed in a form obtained from the experimental p.c. NLHD amplitudes by symmetry transformation from pseudoscalar to vector mesons. This amounts to some fine tuning of the p.c. WRHD amplitudes as compared to calculations avoiding that route, thus helping a little with the fits. The signs and approximate size of asymmetries in [2, 11] do not depend on this fine tuning, however.

4. Relative sign of b_1 and b_2 amplitudes and PCAC

Although the direct photon–quark coupling approach gives absolute predictions for the pattern of asymmetries in the static limit, an independent check on its predictions would be welcome, especially in view of the fact that the pattern depends on one single relative sign between the contributions from diagrams (b1) and (b2). One may wonder whether this sign has been fixed correctly or whether the intermediate states have been treated in a way consistent with experiment elsewhere. To answer these questions, we turn to the PCAC reduction of p.v. NLHD amplitudes in the limit when pion momentum k_μ goes to zero.

Direct calculation of the contribution from W-exchange processes to the p.v. NLHD amplitudes along the lines of Section 2 gives the results shown in Table IV. The amplitude b in Table IV is proportional to b_γ used in WRHDs in the previous section, with the proportionality factor including g_P/e (compare Eqs. (1),(3)).

TABLE IV

W-exchange-induced ($H_W^{\rm p.v.}$) contributions to p.v. amplitudes and the corresponding expressions obtained through PCAC reduction.

	$\langle pP\|H_W^{\rm p.v.}\|\Sigma^+\rangle$			$\langle \Sigma^+ P\|H_W^{\rm p.v.}\|p\rangle$		
Meson $P\,(q\bar{q})$	(b1)	(b2)	PCAC	(b1)	(b2)	PCAC
$u\bar{u}$	$-\frac{1}{2}b$	$\frac{1}{2}b$		$-\frac{1}{2}b$	$\frac{1}{2}b$	
$d\bar{d}$	$-\frac{1}{2}b$	0		0	$\frac{1}{2}b$	
$s\bar{s}$	0	$\frac{1}{2}b$		$-\frac{1}{2}b$	0	
$\frac{1}{\sqrt{2}}(u\bar{u}-d\bar{d})$	0	$\frac{1}{2\sqrt{2}}b$	$\frac{i}{2F_\pi}\langle p\|H^{\rm p.c.}\|\Sigma^+\rangle$	$-\frac{1}{2\sqrt{2}}b$	0	$-\frac{i}{2F_\pi}\langle \Sigma^+\|H^{\rm p.c.}\|p\rangle$
$\frac{1}{\sqrt{6}}(2u\bar{u}-d\bar{d}-s\bar{s})$	$-\frac{1}{2\sqrt{6}}b$	$\frac{1}{2\sqrt{6}}b$	0	$-\frac{1}{2\sqrt{6}}b$	$\frac{1}{2\sqrt{6}}b$	0

From the comparison of the prediction of the static quark model with the results of the PCAC calculation (also given in Table IV), we see that the relative sign of contributions from diagrams (b1) and (b2) has to be positive. Indeed, on account of $d \leftrightarrow s$ symmetry of weak Hamiltonian we have

$$\langle p|H^{\rm p.c.}|\Sigma^+\rangle = \langle \Sigma^+|H^{\rm p.c.}|p\rangle \tag{17}$$

and a negative sign between contributions (b1) and (b2) would lead to contradiction with the PCAC description of p.v. NLHD amplitudes. This is also seen when P is a U-spin singlet $\frac{1}{2}[\pi^0+\sqrt{3}\eta_8]$. Please note that the formulas obtained by PCAC reduction do *not* depend on any intermediate states. Thus, for NLHDs the relative positive sign between the two contributions has been confirmed for $k_\mu \to 0$ in a way that *completely avoids using any intermediate states.* Since within the calculational scheme employed, the difference between NLHDs and WRHDs consists solely in the replacement of the quark pseudoscalar current by the vector one, the relative contributions from diagrams (b1) and (b2) in the p.v. WRHD amplitudes (Table II) must be added. In other words, consistency with formulas obtained by PCAC reduction of NLHD amplitudes in the $k_\mu \to 0$ limit requires that in this limit the contributions from diagrams (b1) and (b2) have to be added for the $\Sigma^+ \to p\gamma$ p.v. amplitude as well (as in Eq. (16)): the $k^\mu \to 0$ limit is relevant for both Current-Algebra (CA) soft-pion estimates and for Hara's theorem. It follows that consistency with the PCAC-based description of NLHDs requires in particular that the $\Xi^0 \to \Lambda\gamma$ asymmetry be positive, in disagreement with Ref. [14] and chiral approaches [20].

At this point it is appropriate to discuss the connection between the framework considered in this paper so far and the standard CA prescription supplemented with SU(3)-breaking resonance-induced corrections. Namely, the p.v. NLHD amplitudes $A(B_i \to B_f\pi)$ are usually described in terms of the contribution from the CA commutator plus terms proportional to pion momentum k_μ:

$$A(B_i \to B_f \pi^{\rm a}) = -\frac{i}{F_\pi} \langle B_f | [F_5^{\rm a}, H^{\rm p.v.}] | B_i \rangle + k_\mu M^\mu \,. \tag{18}$$

The second term on the right vanishes in the SU(3) limit $k_\mu \to 0$. This term is thought to be dominated by the SU(3)-breaking corrections from the intermediate $(\mathbf{70}, 1^-)$ resonances [21]. Using the properties of the Cabibbo Hamiltonian to replace $[F_5^{\rm a}, H^{\rm p.v.}]$ with $[F^{\rm a}, H^{\rm p.c.}]$, and evaluating the contribution from resonances, one obtains:

$$A(B_i \to B_f \pi^{\rm a}) \propto \langle \tilde{B}_f | H^{\rm p.c.} | \tilde{B}_i \rangle + {\rm const}(B_f, B_i) \times (m_s - m_d) \,, \tag{19}$$

where $\tilde{B}_{i(f)}$ denote baryon states, of which one is obtained from $B_{i(f)}$ by the action of isospin generator $I^{\rm a}$ and the other is left unchanged. The leading term (in SU(3) breaking) in Eq. (19) is $\langle \tilde{B}_f | H^{\rm p.c.} | \tilde{B}_i \rangle$. It contains the W-exchange-induced term $\langle \tilde{B}_f | H^{\rm p.c.}_W | \tilde{B}_i \rangle$. It must correspond to the $b_1 + b_2$ term obtained in the scheme of this paper at $k^\mu \to 0$: one might have assumed that the states $|B_i\rangle, |B_f\rangle$ in between which the CA commutator is to be evaluated are the states of the static quark model.

Alternatively, one might *saturate* the CA commutator with this part of contribution from resonances which does *not* vanish in the SU(3) limit. That is, one may replace the r.h.s of Eq. (18) with the pole model prescription. This amounts to considering energy denominators for the intermediate states in the amplitudes of diagrams (b1) and (b2). When one restricts to the $(\mathbf{70}, 1^-)$ (B^*) resonances in intermediate states, for $\Sigma \to N$ transitions these energy denominators are $m_\Sigma - m_{N^*}$ for diagram (b1) , and $m_N - m_{\Sigma^*}$ for diagram (b2). Assuming

$$\begin{aligned} m_\Sigma &\approx m_N + (m_s - m_d) \,, \\ m_{\Sigma^*} &\approx m_{N^*} + (m_s - m_d) \,, \\ m_{N^*} - m_N = \; &\omega \; = m_{\Sigma^*} - m_\Sigma \end{aligned} \tag{20}$$

we see that $(m_s - m_d)$ enters with opposite signs into the two denominators. Thus SU(3) corrections to diagrams (b1) and (b2) are of opposite signs:

$$\begin{aligned} b_1 + b_2 \;\xrightarrow{\rm SU(3)\ breaking}\; & \frac{b_1 \omega}{\omega - (m_s - m_d)} + \frac{b_2 \omega}{\omega + (m_s - m_d)} \\ & \approx b_1 + b_2 + (b_1 - b_2) \frac{m_s - m_d}{\omega} \,. \end{aligned} \tag{21}$$

As a result, the symmetry structure of the SU(3)-breaking correction in Eq. (19) is different from the SU(3)-preserving CA term (b_1–b_2 *versus* b_1+b_2). Please compare the form of the right-hand sides of Eqs. (19) and (21).

In the approach to WRHDs discussed in [14] and in chiral perturbation theory (ChPT), only the WRHD counterpart of the $b_1 - b_2$ term above is considered. In that approach there is no counterpart to the CA commutator term of NLHDs: the $b_1 + b_2$ term in WRHDs is set to zero by hand because it is thought to violate gauge invariance. On the other hand, the gauge-invariant scheme of this paper indicates that such a term (generated by the replacement of current $\bar{q}\gamma_5 q$ for NLHDs with current $\bar{q}\gamma_\mu q$ for WRHDs) exists and has symmetry properties corresponding to the sums of b_1 and b_2 (Table II). Therein lies the difference between the genuine quark model approach and the standard approach of Ref. [14] or more recent ChPT attempts.

With nonzero b_γ it follows that Hara's theorem must be violated, unless a good reason for rejecting the WRHD counterpart to the CA commutator is given. One cannot claim that this counterpart to the CA commutator violates gauge invariance: the input into the calculation is gauge-invariant direct photon–quark coupling and all the steps in the calculation are correct. The (nonzero) $b_1 + b_2$ term must be present in NLHDs because it corresponds to nonzero value of $\langle p|H^{\mathrm{p.c.}}|\Sigma^+\rangle = \langle \Sigma^+|H^{\mathrm{p.c.}}|p\rangle$ in the quark model in the SU(3) limit. This is in agreement with the present descriptions of NLHDs, in which the contribution from the commutator is nonzero and large. Subtracting the $b_1 + b_2$ term in WRHDs while keeping it in NLHDs is completely arbitrary. Thus, the $b_1 + b_2$ term should be present in both NLHDs (the CA commutator) and WRHDs (the counterpart to the CA commutator).

As the p.v. $\Xi^0 \to \Lambda\gamma$ amplitude is due to amplitude b_2 only (Table II), a negative sign of the $\Xi^0 \to \Lambda\gamma$ experimental asymmetry would mean that amplitudes b_2 must enter with an additional negative sign, thus signalling the $b_1 - b_2$ structure, and the cancellation of b_1 and b_2 terms for $\Sigma^+ \to p\gamma$. However, what is more important, a large negative $\Xi^0 \to \Lambda\gamma$ asymmetry would also mean that either something is badly wrong with the quark-model machinery which connects the calculations for quark pseudoscalar and vector currents (I reject this possibility because it concerns *spin* properties of the quark model), or a completely new contribution reveals itself in WRHDs *only*, not only cancelling the counterpart to the CA commutator, but effectively reversing its sign. There is no hint of what such a contribution could be.

As mentioned in Section 2, calculations of the p.v. amplitudes proceed in two steps: (1) evaluation of the admixture of $q\bar{q}$ pairs in qqq baryons, and (2) calculation of the matrix elements of currents $\bar{q}\gamma_5 q$ and $\bar{q}\gamma_\mu q$ in between states with these admixtures. As the latter do not depend on the current, there is no reason why the scheme should fail for the vector current while being correct for the pseudoscalar one. Thus, if the CA commutator is dominant in NLHDs (the $k_\mu \to 0$ term dominates), positive $\Xi^0 \to \Lambda\gamma$ asymmetry is a *must*. Consequently, negative $\Xi^0 \to \Lambda\gamma$ experimental asymmetry would signal deep trouble for the quark model. Note also that the violation of Hara's theorem must be related to step (1) above, *i.e.* to the properties of *states* in the quark model (the interaction $\bar{q}\gamma_\mu q A^\mu$ of step (2) is gauge invariant).

5. Origin of the violation of Hara's theorem in the quark model

As discussed, the static gauge-invariant quark-model calculation unavoidably violates Hara's theorem. This is in direct conflict with what is easily proved in hadron-level approach. Before we jump to the conclusion that there is something wrong with the static quark model, let us discuss the origin of our result from a different perspective.

We start with the standard hadron-level picture. In this picture hadrons are assumed to be well described by an effective local field theory. We are interested in what happens at point $k_\mu = 0$. This point (as any other point in momentum space) corresponds to particles (baryons, photons) described by plane waves. The positions of corresponding particles are not well defined: the particles may be found anywhere. It might help to think of particles as "potentially being everywhere".

The static-limit calculations in the quark model involve juggling spin, flavour and parity (quark–antiquark) indices only (see Eqs. (6),(10),(13),(14), Fig. 1): the quarks themselves are in states of vanishing but definite momenta. This prescription, when *interpreted* in position space, corresponds to the situation in which positions of quarks "are" arbitrary, *i.e.* now it is *individual* quarks that "are potentially everywhere". In particular, this includes configurations with baryon quarks "potentially being" arbitrarily far away from each other. Such nonlocal quark configurations are forbidden in an effective hadron-level local field theory, when quarks are assumed to be close to one another. Thus, the static quark model includes configurations which are not and cannot be taken into account in the language of effective hadron-level local field theory. Hadron- and quark-level prescriptions are generically *different*. Quark-model violation of Hara's theorem must come from these nonlocal configurations.

The nonlocal configurations of baryon quarks may be considered unacceptable: their presence is in conflict with the expectations based on the idea of confinement. We know, however, that in general composite quantum states may exhibit nonlocal features. Since baryons are quantum states made out of quarks, the real question is whether we can exclude such (admittedly weird) configurations on the grounds of either experiment or of general ideas of quantum physics (as opposed to those based on certain theoretical expectations)? Consider therefore the situation from the point of view of what is measurable. The closer we approach the $k_\mu = 0$ limit (assuming we can manipulate $m_s - m_d$), the worse spatial resolution we have. Consequently, we cannot experimentally exclude that in the $k_\mu = 0$ limit, quarks may be thought of as "being" arbitrarily far away from each other, *i.e.* that the photon–baryon coupling is intrinsically nonlocal. Note that if experiment is set up to measure momenta it is not meaningful to talk about hadron position in any other way than with the help of a theory. The same applies to the positions of quarks. In order to see small distances and check what quark positions "really are", one has to look "deep into hadrons". This requires high, not low momenta and amounts to asking a completely different experimental question. Consequently, it is hard to see how one can reject the violation of Hara's theorem on the grounds of general principles. This interpretation of quark-model results is clearly more general than the model considered: the composite quark state may couple to zero-energy photon in a genuinely nonlocal way.

Note that the interpretation in terms of quarks "being" arbitrarily far away naturally follows also when baryon magnetic moments are calculated in the original quark model (*cf.* Eq. (15)). A possible resolution to this problem of free quarks, corresponding to the *expected* effective local theory at hadron level, was proposed a long time ago in the form of the confinement idea. When confinement is imposed in a way consistent with the standard expectations, no clear-cut conflict between naive quark model calculations and hadron-level expectations emerges for baryon magnetic moments. For WRHDs, however, the predictions of the static quark model, and those of the hadron-level effective local theory, are different. Thus, it appears that WRHDs probe in a subtle way the original question of apparent quark freedom and unobservability at low energies. For this reason, I think that the issue of the violation of Hara's theorem is extremely interesting and very important.

There is another way to see that the violation of Hara's theorem requires some kind of nonlocality. Namely, it has been shown in Ref. [9] that when the violation of Hara's theorem with built-in current conservation is *forced* into the "corset" of effective hadron-level local description, the electromagnetic axial baryonic current cannot be strongly suppressed at infinity: the fall-off must be as slow as $1/r^3$. This should be contrasted with models in which the assumed exponential suppression of the current leads to Hara's theorem. Since there is no massless hadron, the nonlocality discussed in Ref. [9] should be thought of as simulating some other kind of nonlocality.

One may argue that the static calculation is inadequate and that small components of Dirac spinors should be taken into account. Indeed, small components necessarily emerge when quarks are confined to a restricted volume of space. One may hope that contributions from these components could cancel the static limit term, thus restoring Hara's theorem. This may be so (see, however, Ref. [22]). The real problem is that one should include such terms in other calculations as well. In particular, they should be taken into account in the calculations of p.v. meson-baryon couplings. The relevant calculations (equivalent to those of Section 2) were performed by DDH [15] in the nonrelativistic (actually static) limit. Thus, the celebrated DDH results stem from an approach that violates standard ideas about confinement.

As the photon and vector-meson p.v. couplings to baryons are calculated using the same quark-level current, the Dirac structures of photon and vector-meson couplings at hadron level should be identical. The only gauge-invariant coupling acceptable in local hadron-level framework is that involving $\bar{B}_f\sigma_{\mu\nu}\gamma_5 k^\nu B_i$ (with $B_{i(f)}$ being baryon bispinors). If Hara's theorem is satisfied, also for vector mesons this structure only should appear. However, the interaction $\bar{B}_f\sigma_{\mu\nu}\gamma_5 k^\nu B_i V^\mu$ is inconsistent with the data on p.v. effects in NN interactions (see Section 6).

The above arguments suggest that the language of effective local theories at hadron level might sometimes constitute an insufficient approximation to reality. In other words, the assumption that spin-$1/2$ baryon may *always* be well approximated by a Dirac spinor field depending on a well-defined single point x might be too strong[1]. Although one could speculate about possible deeper origins and implications of the suspected limitations of the standard language, such speculations will be warranted only if large positive value of the $\Xi^0 \to \Lambda\gamma$ experimental asymmetry is established.

Of course, any limitation of the standard hadron-level language does not mean that the quark-level theory cannot be local or that individual quarks have to exhibit intrinsic nonlocality. The calculation of this paper constitutes an explicit counterexample. Point-like interactions of the standard model underly the whole picture herein considered. It is the question of proper description of baryons as composite states that is being discussed here. Clearly, an effective local field theory at hadron level should constitute a sufficient approximation for most of the present practical purposes.

[1] Field theories of nonlocal type have been considered as a possible theoretical vehicle for describing hadrons since the times of Yukawa [23].

6. Dependence of b_γ on SU(3) breaking parameter $m_s - m_d$

So far, the (mutually proportional) amplitudes b_γ and b have been nonzero constants. As the calculation was performed in the SU(3) limit, these constants were independent of the SU(3) breaking parameter $m_s - m_d$. This was also the case in Ref. [15], where the $\Delta S = 1$ p.v. amplitudes extracted from NLHDs (and in particular, their scale) were used to the description of $\Delta S = 0$ nuclear p.v. processes and their scale.

Let us assume that experiment confirms the pattern corresponding to the sum of b_1 and b_2, thus indicating spin symmetry between NLHDs and WRHDs at $m_s \neq m_d$. This does not mean yet that the preceding sections describe the physical situation in a qualitatively correct way. The static quark model gives predictions at $k_\mu = 0$, while in the real world $k_\mu \neq 0$. Still, one suspects that spin symmetry would survive should the masses of s and d quarks be different. Thus, in a theoretical description in which quark masses are free parameters, let $b \propto b_\gamma$ be a function of $\delta \equiv m_s - m_d$ (*i.e.* $b \to b(\delta)$), which can be expanded into a series in the vicinity of $\delta = 0$ (thus, the form $b \propto b(0) + b'|m_s - m_d|$ is not allowed). Hara's theorem is satisfied if $b(0) = 0$. From symmetry properties under $s \leftrightarrow d$ interchange (Table IV) we see that only nonzero even powers of $m_s - m_d$ may appear in the expansion of $b(\delta)$ and $b_\gamma(\delta)$. With $b(0) = 0$, the lowest order term is proportional to $(m_s - m_d)^2$. If $b(\delta) = b''\delta^2$, the relevant hadron-level structure of photon–hadron p.v. interaction may be written as

$$(m_s - m_d)(b''_1 + b''_2)\bar{B}_f i\sigma_{\mu\nu}\gamma_5 k^\nu B_i A^\mu \,, \tag{22}$$

where $b''_1(b''_2)$ denotes amplitude for diagram (b1) (diagram (b2)) with the $(m_s - m_d)^2$ factor removed. In Ref. [16], a pole model leading to the above solution was discussed. Clearly, this way of making positive $\Xi^0 \to \Lambda\gamma$ asymmetry consistent with Hara's theorem is in conflict with the nonzero value of the CA commutator and with the quark model results for $\langle p|H^{\text{p.c.}}|\Sigma^+\rangle$, according to which the contribution of W-exchange between quarks does *not* vanish for $m_s = m_d$. Internal consistency requires that for $m_s = m_d$ both $\langle p|H_W^{\text{p.c.}}|\Sigma^+\rangle$ and W-exchange contributions to $A(\Sigma^+ \to p\gamma)$ (Eq. (16)) behave in the same way in the static limit: they are either both zero or both nonzero.

Note that if Hara's theorem is to be satisfied in the SU(3) limit for positive $\Xi^0 \to \Lambda\gamma$ asymmetry, *all* p.v. amplitudes must vanish in this limit (Eq. (22)). This is not the resolution proposed in [14] and chiral approaches [4], where only the $\Sigma^+ \to p\gamma$ p.v. amplitude vanishes, while the p.v. amplitudes of the remaining three WRHDs stay nonzero. The symmetry structure of the latter approach yields the pattern $(-,-,-,-)$ for the asymmetries of $\Sigma^+ \to p\gamma, \Lambda \to n\gamma, \Sigma^0 \to \Lambda\gamma$, and $\Sigma^+ \to \Sigma^+\gamma$ respectively.

Such symmetry structure might be obtained in the static limit with direct photon–quark coupling if the photon were to couple to quarks through the nonminimal coupling of Eqs. (8),(9). This is not acceptable.

Extracting information on $b_\gamma(\delta)$ for small δ from the analysis of NLHDs and WRHDs is impossible: with m_s fixed, there is no way to check whether the transition amplitudes are proportional to $(m_s - m_d)^2$ or not. Despite that, it is still possible to get some experimental indication as to which of the two theoretical solutions: $b_\gamma(0) \approx b_\gamma(m_s - m_d)$ or $b_\gamma(0) = 0$, is favoured. In order to see what these indications are, we recall that the coupling of photon to baryon results from γ coupling to quark vector current. As discussed in Section 2, the coupling of U-spin-singlet vector meson to baryons results from its coupling to the *same* current. Since the whole quark model machinery does not depend on the nature of the field this current is coupled to, we conclude that one should be able to get some hints about the possible dependence of b_γ on $m_s - m_d$ by looking at the $\Delta S = 0$ p.v. couplings of vector mesons to nucleons. If the factor of $(m_f - m_i)^2$ is present in b and b_γ, it follows that *all* $\Delta S = 0$ p.v. weak couplings of mesons to baryons should be very small. Thus, we should observe nearly vanishing p.v. effects in the interactions of nucleons. With respect to NLHD-based SU(3) estimates, they should be down by a factor of the order of at least $(m_n - m_p)^2/(m_N - m_\Sigma)^2 \approx (m_u - m_d)^2/(m_s - m_d)^2 \approx 10^{-3}$ or less depending on the process and the precise mass values used. This is not what is experimentally observed. The present data on p.v. nuclear forces indicate that the p.v. couplings of mesons to nucleons are roughly of the order expected from NLHDs by an SU(3)-symmetric extrapolation, *i.e.* as if there was no $(m_f - m_i)^2$ suppression. This is in particular supported by p.v. effects observed in pp scattering, which constitute the best experimentally established effect [15,24,25]. These data require the presence of a $\overline{N}\gamma_\mu\gamma_5 N V^\mu$ term, and in particular, nonvanishing contributions from b diagrams. If these b (and the so-called c, *cf.* [15]) diagrams were negligible and only the factorisation a terms (see [15]) were allowed, one could not describe the p.v. effect observed in pp scattering [25, 26]. A $\overline{N} i\sigma_{\mu\nu}k^\nu\gamma_5 N V^\mu$ term cannot describe the data [25] either. This suggests that the coupling of photons need not vanish, unless some additional cancellation takes place. This reasoning is correct, provided p.v. couplings of vector mesons to baryons can be correctly described by their evaluation in the static limit of the quark model. This is how they were estimated in [15], where they were also linked through symmetry to the soft-pion contribution to p.v. NLHD amplitudes.

It may be that the effects observed in nuclear parity violation mean that the $\Delta S = 0$ p.v. couplings of vector mesons to baryons cannot be reliably estimated in this way. For example it may be argued that in the p.v. nuclear force the couplings of vector mesons at $k^2 = m_\rho^2$ should be used, while the scheme of this paper yields them at $k^2 = 0$. Clearly, experimental results on p.v. nuclear forces may be considered as a hint only. In the meantime, in order to learn more about WRHDs we have to wait for the results of the KTeV and NA48 experiments on $\Xi^0 \to \Lambda\gamma$ decays [27, 28].

7. Conclusions

In this paper we have analysed the assumptions needed to generate the formulas of the VMD approach to WRHDs [2, 10]. We have shown that all their characteristic features are obtained also in an explicitly gauge-invariant quark-level scheme in the limit of static quarks with direct photon–quark coupling. Consequently, the claim of Ref. [4] that the results of the VMD $\times$ $SU(6)_W$ scheme of Refs. [2, 10] (and the violation of Hara's theorem) follow from a lack of gauge invariance, is incorrect.

Our analysis shows that the standard pole-model approach of [14] and ChPT miss a quark-level contribution, which we call the WRHD counterpart to the CA commutator in NLHDs (the lowest-order term in ChPT). This contribution is obtained by replacing the interaction of pseudoscalar meson with the quark pseudoscalar current by the interaction of photon with the quark vector current, *i.e.* through the use of *spin properties of the quark model*. The obtained term is gauge invariant and proportional to matrix elements of the CA commutator in NLHDs. Consequently, the vanishing of the latter matrix elements in between two states of strangeness differing by $\Delta S = 1$ should hold in the $m_s \to m_d$ limit if Hara's theorem is to be satisfied. Because these matrix elements are proportional to the elements of the p.c. weak Hamiltonian in between similar two baryonic states, the latter matrix elements should also vanish in this limit. However, in quark model calculations they are non-zero.

Detailed analysis of the direct photon–quark approach shows that in the quark model the pattern of asymmetries predicted in Refs. [2, 11] is unavoidable. In particular, large positive $\Xi^0 \to \Lambda\gamma$ asymmetry is necessarily predicted. Any other experimental value for this asymmetry constitutes *a serious problem for the quark model*.

Negative $\Xi^0 \to \Lambda\gamma$ asymmetry of the (consistent with Hara's theorem) standard pole model [14] and ChPT, cannot be obtained in calculations performed with static quarks. Such calculations, when applied to NLHDs, agree with the expectations based on the dominance of CA commutator. If experiments find negative $\Xi^0 \to \Lambda\gamma$ asymmetry, the discrepancy with this paper cannot be blamed on the violation of gauge invariance, and would constitute a serious problem *for the quark model*. A close-to-zero experimental value for this asymmetry would constitute another problem.

Positive $\Xi^0 \to \Lambda\gamma$ asymmetry does not mean yet that Hara's theorem must be violated, but that this is likely. If Hara's theorem is satisfied, *all* WRHD amplitudes must vanish in the SU(3) limit. Although we cannot vary $m_s - m_d$, getting experimental information on the $m_s - m_d$ dependence of theoretical amplitudes is to some extent possible. The experiments in question deal with the sector of weak p.v. $\Delta S = 0$ nuclear transitions and indicate a lack of the SU(3) suppression for p.v. weak couplings of vector mesons to baryons. This hints that there is no such suppression for photons either, unless some additional cancellations occur.

We have argued that although the violation of Hara's theorem is in conflict with standard theoretical Ansätze at hadron level, it does not have to be unphysical. An interpretation of the quark-model result in terms of an intrinsic baryon nonlocality was presented. It was stressed that the issue of Hara's theorem probes the original quark-model question of apparent quark freedom and additivity at low energy. It may be that the quark model is an abstraction that went too far in assigning particle-like properties to quarks, thus leading us into conflict with the hadron-level description. However, it may also be that quark model results should be treated seriously. Measurement of the $\Xi^0 \to \Lambda\gamma$ asymmetry will provide crucial information, pointing out either to *a serious problem for the quark model* or to *intrinsic nonlocality of baryons* when these are probed by low energy photons. Whatever the resolution, WRHDs should teach us a lot about the quark model in the low-energy domain, *i.e.* about the way quarks combine to form hadrons.

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