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PARITY-VIOLATING NUCLEON-PION AMPLITUDE IN GAUGE-INVARIANT MODELS OF WEAK, ELECTROMAGNETIC AND STRONG INTERACTIONS*)

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Abstract: Nonleptonic effective weak Hamiltonians have been found for unified gauge theories, including asymptotically free strong gluon interactions. Two-quark operators appearing in the short-distance expansion are discussed in detail. Effective Hamiltonians serve as a basis for the deduction of the parity-violating nucleon-nucleon-pion amplitude. Problems connected with this deduction are described and various approximations are compared, including the possibility of the existence of a Zweig-Iizuki rule analogon. There are some indications that certain models predict a parity-violating nucleon-nucleon-pion amplitude which is smaller than 10^{-7} ($\hbar = c = 1$), thus providing a basis for experimental discrimination.

1. Introduction

Ever since the original suggestion[1], the parity-violating pion-nucleon interaction (PV MNπ) has been used to study models of weak Hamiltonians (H_W)[2].

The main purpose of this paper is to reevaluate this problem in the light of recent theoretical developments using gauge-invariant theories of weak and electromagnetic interactions[3-5], which also allow the inclusion of the supposed mediators of strong interactions, coloured vector gluons. It is assumed that strong interactions arise from a renormalizable asymptotically free colour-SU(3) gauge theory[6] and that only colour-singlet states exist in a physical spectrum. Various

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quarks are contained in the models of H_W studied and each is supposed to appear in a colour triplet. The asymptotic freedom of colour SU(3)[7] is used to study the short-distance behaviour of the product of two currents appearing in hyperon--decay amplitudes and the related NNπ amplitudes[8,9]. Assuming that this short--distance contribution is the dominant one, the effective weak Hamiltonians can be found for the weak-interaction models under study. These effective weak Hamiltonians can be used in the sum rules connecting the PV NNπ amplitude to the measured nonleptonic hyperon-decay amplitudes.

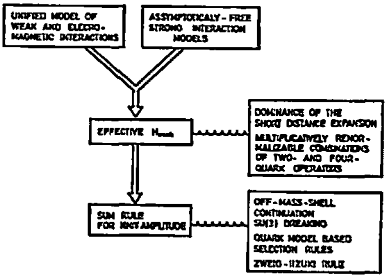

Fig. 1. Outline of the theoretical scheme for the prediction of the PV NNπ amplitude. Inputs are indicated by arrows; difficulties, assumptions and ambiguities by wavy lines.

Unfortunately, none of these steps is free from ambiguities and difficulties. With existing more detailed elementary-particle theories, predictions are more intimately connected with the quark dynamics of elementary-particle models. We feel that the systematic outline of theoretical problems which will be presented here, is an essential step for further research.

In Section 2 we discuss problems connected with the determination of the effective weak Hamiltonians and list preliminary results for three characteristic models of unified weak and clectromagnetic interactions. These are: the recently proposed model containing V + A terms in weak currents[5], the Georgi-Glashow model[4], which has no weak neutral currents, and the Salam-Weinberg model[3], which has been most extensively discussed so far.

Section 3 is concerned with the determination of the PV NNπ amplitude using the effective H_W as an iput. Six possible approaches are defined, discussed and used.

An outline of the scheme for the theoretical prediction of the PV NNπ amplitude is shown in Fig. 1.

In Section 4 the theoretical predictions based on gauge-invariant unified field-theory models[3-5] are compared with the predictions based on an empirical current-current model of weak interactions[10].

2. *Asymptotic freedom and the effective weak Hamiltonian*

The effective weak Hamiltonian H_W^{eff} is found from the second-order contribution in weak interactions

$$H^2 = \int d^4 x \, D_W(x) \, [T(\mathcal{J}_\mu^+(x)\mathcal{J}_-^\mu(0)) + T(\mathcal{J}_\mu^-(x)\mathcal{J}_+^\mu(0))] +$$

$$+ \int d^4 x \, D_Z(x) \, T(\mathcal{J}_\mu^0(x)\mathcal{J}_0^\mu(0)). \tag{2.1}$$

This gauge-independent expression is obtained by neglecting Higgs scalar exchanges[11,12]. Here $D_A(x)$ symbolizes the scalar propagator of a particle of mass M_A, corresponding to intermediate vector bosons (IVB) (i. e. W, Z, etc.) appearing in the unified field theory. The currents $\mathcal{J}^a_\mu$, carrying charge a, contain Heisenberg quark fields with respect to strong interactions, which are thus included in Equ. (2.1) to any order. Asymptotic freedom of strong interactions, allows to calculate strong-interaction corrections to the Wiison expansion of time-ordered products of currents in Equ. (2.1)[13]. The x behaviour of the propagators $D(x)$ suggests that these leading terms for small x might actually give the dominant contribution to (2.1)[8,9,14]. This is certainly the case for free quarks, i. e. the vanishing strong--interaction coupling constants g_S. For $g_s \neq 0$, one cannot exclude the logical possibility that in the region where $D(x)$ is very small the time-ordered product of currents might be very large. In such a case the small-x expansion does not work.

All following conclusions are drawn assuming that this does not happen and that the situation $g_S = 0$ is also reflected in the presence of weak interactions.

The Wilson expansion of the time-ordered product has the general form[13]

$$T(\mathcal{J}(x)\mathcal{J}(0)) = \sum_n C_2^n(x)\,\mathcal{O}_2^{(n)}(0) + \sum_n C_4^{(n)}(x)\,\mathcal{O}_4^n(0) + \ldots . \tag{2.2}$$

Here the indices 2 and 4 denote various two-and four-quark operators, respectively. For asymptotically free theories[6,7],

$$C(x) = C_0(x)\,(1 + b \ln(\mu^2 x^2)). \tag{2.3}$$

Here b has been found by solving the Callan-Symanzik[15] equation or a corresponding one[16]. C_0 corresponds to the case $g_S = 0$ and has to be found by studying Equ. (2.1) for free-quark fields. For four-quark operators, C_0 can be found in a trivial way, corresponding simply to the replacement

$$T(\mathcal{J}(x)\mathcal{J}(0)) \to N(\mathcal{J}(0)\mathcal{J}(0)), \tag{2.4}$$

where N symbolizes the normal product. Simply speaking, four-quark operators are normalized to the well-known local current-current theory.

When dealing with two-quark operators, one has to study expansions in small x for all possible contractions in Equ. (2.1); for example

$$\bar{q}(x)\, C_+\, \Gamma^+_\mu\, S_C(x)\, C_-\, \Gamma^\mu_-\, q(0). \tag{2.5}$$

Here $C_\pm$ are some matrices characterizing the symmetry structure of the model and will be defined later for each model studied. The symbols Γ_μ describe Lorentz properties of weak currents and S_C is the fermion propagator.

We want to stress certain general properties which follow from the structure of unified gauge models[3-5]. Details are given in Appendix A. In general, such models do contain left- and right-handed currents. Contractions indicated by Equ. (2.5) give as leading terms operators containing either two derivatives of quark fields

$$\mathcal{O}_2^{(2)} = \bar{q} \hat{\nabla} \hat{\nabla} N \gamma_5 q \tag{2.6}$$

or three derivatives of quark fields

$$\mathcal{O}_2^{(3)} = \bar{q} \hat{\nabla} \hat{\nabla} \hat{\nabla} L \gamma_5 q. \tag{2.7}$$

The operator (2.6) is obtained by combining left- and right currents, while Equ. (2.7) follows from right-right and left-left combinations. The matrix N in Equ. (2.6) contains the quark mass matrix M, which also contributes to the asymptotic freedom enhancement[16,17]. It is well known that operators without derivatives or with one derivative are absorbed in renormalization constants and subtractions[11].

Even when SU(n) symmetry, corresponding to n-quark flavours, is completely broken by M, $\mathcal{O}_2^{(2)}$ coming from charged currents does not contribute to $\Delta S = 0$ processes to the order $g_S = 0$, as shown in Appendix A. On the contrary, the L matrix coming from charged currents can contribute only to $\Delta S = 0$ processes up to the order $\Delta m_q^2/M_W^2$. Neutral currents can contribute only to $\Delta S = 0$ processes. The suppression of $\Delta S \neq 0$ processes to the order $\Delta m_q^2/M_W^2$ is usually known as GIM-mechanism suppression[18]. The above statements follow from the structure of the L matrix

$$L = \{C^+, C^-\}_+ - \{D^+, D^-\}_+ + 0\left(\frac{\Delta m_q^2}{M_W^2}\right). \tag{2.8}$$

In the weak-interaction models studied here, the anticommutators of the matrices C and D are always diagonal.

At the end of this section we will give L and M matrices for the models under consideration.

We now turn to four-quark operators. The leading contributions of the six-dimensional four-quark operators were calculated in Refs.[8,9]. While the results of these references are exact for $\Delta S = 1$ processes, in the $\Delta S = 0$ case a number of operators appear which are multiplicatively renormalizable only in their linear combinations and not individaully. In these combinations, strong-interaction renormalization mixes two-quark operators with four-quark ones. This was treated in Ref.[9], assuming that certain two-quark operators appearing in the operator mixing are unimportant. If one wants to study two-quark operators (2.7), the approach of Ref.[9] has to be reconsidered.

For orientation, we give enhancement factors for four-quark operators calculated by existing methods[8,9]. The results obtained for the 27-dimensional representation of the standard strong SU(3) symmetry (i. e. $|\Delta \vec{I}| = 3/2, 2$) are exact.

We consider three models.

The Six-Quark de Rujula-Georgi-Glashow (DGG) model. In this model[5] currents transform as a 35-dimensional representation of the SU(6) group. Using the quark spinor $\bar{\psi} = (\bar{p}, \bar{n}, \bar{\lambda}, \bar{p}', \bar{p}'', \bar{n}'')$ for a basis, the parity-violating product of charged currents can be decomposed according to

$$35 \oplus 35 = 405 \oplus 280^{*} \oplus 280 \oplus 189 \oplus 35_S \oplus 35_A \oplus 1. \tag{2.9}$$

The resulting expression is

$$(\mathcal{J}^{+}_{\mu}\mathcal{J}^{\mu}_{-} + \text{h. c.})_{PV} = \frac{G}{2\sqrt{2}}(P^{SS} + P^{AS} + P^{SA} + P^{AA}). \tag{2.10}$$

The model is defined by the six-dimensional matrices

$$D_{+} = (D_{-})^{\dagger} = \begin{bmatrix} \cdot & \cdot & \cdot & \cdot & \cdot & 1 \\ \cdot & \cdot & \cdot & \cdot & \cdot & \cdot \\ \cdot & \cdot & \cdot & \cdot & \cdot & \cdot \\ \cdot & l & r & \cdot & \cdot & \cdot \\ \cdot & -r & l & \cdot & \cdot & \cdot \\ \cdot & \cdot & \cdot & \cdot & \cdot & \cdot \end{bmatrix},$$

$$C_{+} = (C_{-})^{\dagger} = \begin{bmatrix} \cdot & c & s & \cdot & \cdot & \cdot \\ \cdot & \cdot & \cdot & \cdot & \cdot & \cdot \\ \cdot & \cdot & \cdot & \cdot & \cdot & \cdot \\ \cdot & -ds & dc & \cdot & \cdot & t \\ \cdot & +st & -tc & \cdot & \cdot & d \\ \cdot & \cdot & \cdot & \cdot & \cdot & \cdot \end{bmatrix}, \tag{2.11}$$

where $s = \sin\Theta$, $c = \cos\Theta$, $t = \sin\varphi$, $d = \cos\varphi$, $r = \sin\psi$ and $l = \cos\psi$. The Cabibbo angle is introduced here through s and c, while t, d, r and l are model parameters.

Introducing the operator

$$\mathcal{O}_{ijkl}(A,B) = \{\bar{\psi}_i \gamma_\mu \gamma_5 A\psi_j, \bar{\psi}_k \gamma^\mu B\psi_l\}_{+} \tag{2.12}$$

for any matrices A and B, the pieces P^{SS} and P^{AA} of Equ. (2.10), transforming as a 405-plet or a 189-plet, respectively, can be written as

$$\begin{pmatrix} P^{SS} \\ P^{AA} \end{pmatrix} = (-)\left[\mathcal{O}(C_{+}, C_{-}) + \mathcal{O}(C_{-}, C_{+}) - \mathcal{O}(D_{+}, D_{-}) - \mathcal{O}(D_{-}, D_{+})\right] \begin{pmatrix} b^{SS} \\ b^{AA} \end{pmatrix}. \tag{2.13}$$

There are also two combinations of 280 and 280* representations

$$\begin{pmatrix} P^{AS} \\ P^{SA} \end{pmatrix} = (-)\,[\mathcal{O}(C_+, D_-) + \mathcal{O}(C_-, D_+) - \mathcal{O}(D_+, C_-) - \mathcal{O}(D_-, C_+)] \begin{pmatrix} b^{AS} \\ b^{SA} \end{pmatrix}. \tag{2.14}$$

Here b^{XY} incorporate sums of traceless three-dimensional SU (3) matrices t_a

$$\begin{aligned}
(b^{SS})_{ijkl} &= \frac{2}{3}\,\delta_{ij}\,\delta_{kl} + \frac{1}{4}(t_a)_{ij}(t_a)_{kl}, \\
(b^{AA})_{ijkl} &= \frac{1}{3}\,\delta_{ij}\,\delta_{kl} - \frac{1}{4}(t_a)_{ij}(t_a)_{kl}, \\
(b^{AS})_{ijkl} &= \frac{8}{9}\,\delta_{ij}\,\delta_{kl} - \frac{1}{6}(t_a)_{ij}(t_a)_{kl}, \\
(b^{SA})_{ijkl} &= \frac{1}{9}\,\delta_{ij}\,\delta_{kl} + \frac{1}{6}(t_a)_{ij}(t_a)_{kl}.
\end{aligned} \tag{2.15}$$

The product of neutral currents in this particular model is both strangeness- and parity-conserving, so it can be ignored. The enhancement (or hindrance) factors for the PV part are

$$H_{PV}^{\rm eff} = \frac{G}{2\sqrt{2}}(0.52\,P^{SS} + 0.72\,P^{AS} + 13.9\,P^{SA} + 3.73\,P^{AA}). \tag{2.16}$$

The Georgi-Glashow (GG) model. The main feature of this model[4)] is the absence of weak neutral currents. Using the quark spinor $\bar{\psi} = (\bar{p}, \bar{n}, \bar{\lambda}, \bar{p}', \bar{q}, \bar{q}', \bar{r}, \bar{r}')$, the charged-current product can be decomposed into

$$(\mathcal{J}_\mu^+ \mathcal{J}_-^\mu + \text{h.c.})_{PV} = \frac{G}{2\sqrt{2}}(P^{SS} + P^{SA} + P^{AS} + P^{AA} + P_1^*). \tag{2.17}$$

The operator P_1^* is defined in Appendix B, together with a complete set of operators which transform in the same way. The operators P^{XY} are defined by traceless model matrices

$$C_+ = (C_-)^\dagger = \begin{bmatrix}
\cdot & -c & -s & \cdot & \cdot & \cdot & \cdot & \cdot \\
\cdot & \cdot & \cdot & \cdot & \cdot & \cdot & c & -s \\
\cdot & \cdot & \cdot & \cdot & \cdot & \cdot & s & c \\
\cdot & s & -c & \cdot & \cdot & \cdot & \cdot & \cdot \\
\cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot \\
\cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot \\
\cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot \\
\cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot
\end{bmatrix},$$

$$D_+ = (D_-)^\dagger = \begin{bmatrix} \cdot & \cdot & \cdot & \cdot & -1 & \cdot & \cdot & \cdot \\ \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot \\ \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot \\ \cdot & \cdot & \cdot & \cdot & \cdot & -1 & \cdot & \cdot \\ \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & 1 & \cdot \\ \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & 1 \\ \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot \\ \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot \end{bmatrix},$$

$$G = \begin{bmatrix} 0 &&&&&&& \\ & 1 &&&&&& \\ && 1 &&&&& \\ &&& 0 &&&& \\ &&&& -1 &&& \\ &&&&& -1 && \\ &&&&&& 0 & \\ &&&&&&& 0 \end{bmatrix}, \tag{2.18}$$

$$F = \begin{bmatrix} \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot \\ \cdot & \cdot & \cdot & \cdot & c & -s & \cdot & \cdot \\ \cdot & \cdot & \cdot & \cdot & s & c & \cdot & \cdot \\ \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot \\ \cdot & c & s & \cdot & \cdot & \cdot & \cdot & \cdot \\ \cdot & -s & c & \cdot & \cdot & \cdot & \cdot & \cdot \\ \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot \\ \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot & \cdot \end{bmatrix}.$$

Introducing

$$P^{SS,AA} = K^{SS,AA} - \frac{\cos\beta}{\sin\beta} L^{SS,AA}, \tag{2.19}$$

we define

$$\begin{pmatrix} K^{SS} \\ K^{AA} \end{pmatrix} = (-)\,[\mathcal{O}(C_+, C_-) + \mathcal{O}(C_-, C_+) - \mathcal{O}(D_+, D_-) - \mathcal{O}(D_-, D_+) + \\ + \begin{pmatrix} -\frac{1}{5} & \\ & \frac{1}{3} \end{pmatrix} (\mathcal{O}(1, G) + \mathcal{O}(G, 1))] \begin{pmatrix} b^{SS} \\ b^{AA} \end{pmatrix}, \tag{2.20}$$

$$\begin{pmatrix} L^{SS} \\ L^{AA} \end{pmatrix} = [\mathcal{O}(C_+, D_-) + \mathcal{O}(D_-, C_+) + \mathcal{O}(C_-, D_+) + \mathcal{O}(D_+, C_-) + \\ + \begin{pmatrix} -\frac{1}{5} & \\ & \frac{1}{3} \end{pmatrix} (\mathcal{O}(1, F) + \mathcal{O}(F, 1))] \begin{pmatrix} b^{SS} \\ b^{AA} \end{pmatrix}.$$

Operators of mixed symmetry are

$$\begin{pmatrix} P^{SA} \\ P^{AS} \end{pmatrix} = \frac{1}{\sin\beta} [\mathcal{O}(D_+, C_-) - \mathcal{O}(C_-, D_+) + \mathcal{O}(D_-, C_+) + \mathcal{O}(C_+, D_-)] \begin{pmatrix} b^{SA} \\ b^{AS} \end{pmatrix}. \tag{2.21}$$

Analysis ends with the effective H_W

$$H_W^{\text{eff}} = \frac{G}{2\sqrt{2}} (0.44\, P^{SS} + 25.81\, P^{SA} + 0.66\, P^{AS} + 5.08\, P^{AA} + \sum_1^5 \beta_i P_i^*). \tag{2.22}$$

The constants β_i are given in Table 2 and the operators P_i^* are defined in Appendix B.

The Salam-Weinberg (SW) model. In this model the quark spinor is $\bar{\psi} = (\bar{p}, \bar{n}, \bar{\lambda}, \bar{p}')$. The charged-current model matrices correspond to the left corner 4×4 submatrix of C_+ in Equ. (2.11) with $d \equiv 1$ [3]. The products of currents can be decomposed into

$$(\mathcal{J}_\mu^+ \mathcal{J}_-^\mu + \text{h.c.}) = \frac{G}{2\sqrt{2}} (P^{SS} + P^{AA} + P_1^1),$$

$$\mathcal{J}_\mu^0 \mathcal{J}_0^\mu = \frac{G}{\sqrt{2}} (R^{SS} + R^{AA} + R_1^{15} + R_1^1). \tag{2.23}$$

The operators P_1^1, R_1^1 and R_1^{15} are defined in Appendix B. The operators transforming as 84— or 20-representations under SU(4), respectively, are

$$\begin{pmatrix} P^{SS} \\ P^{AA} \end{pmatrix} = (-) [\mathcal{O}(C_+, C_-) + \mathcal{O}(C_-, C_+) + \begin{pmatrix} -\frac{1}{5} \\ +\frac{1}{3} \end{pmatrix} \mathcal{O}(1,1)] \begin{pmatrix} b^{SS} \\ b^{AA} \end{pmatrix},$$

$$\begin{pmatrix} R^{SS} \\ R^{AA} \end{pmatrix} = (-1 + 2\sin^2\Theta_W) [\mathcal{O}(C_3, C_3) + \begin{pmatrix} -\frac{1}{20} \\ +\frac{1}{12} \end{pmatrix} \mathcal{O}(1,1)] \begin{pmatrix} b^{SS} \\ b^{AA} \end{pmatrix}. \tag{2.24}$$

The final result is

$$H_W^{\text{eff}} = \frac{G}{2\sqrt{2}} (0.58\, P^{SS} + 3.02\, P^{AA} + \sum_1^3 \gamma_i P_i^1) +$$

$$+ \frac{G}{\sqrt{2}} (0.58\, R^{SS} + 3.02\, R^{AA} + \sum_1^5 \beta_i R_i^{15} + \sum_1^3 \gamma_i R_i^1). \tag{2.25}$$

In order to have complete H_W^{eff}, one has to add the respective two-quark operator contributions to Equs. (2.16), (2.22) and (2.25). Using the model matrices defined above, one can calculate N and L matrices in the strict SU(n)-symmetry limit, as presented in Table 1. Details can be found in Appendix A.

3. *Nonleptonic hyperon-decay amplitudes and the PV $NN\pi$ amplitude*

In order to probe models of H_W, one has to connect the PV NNπ amplitude with strangeness-changing decay amplitudes. Unfortunately, this cannot be performed in an unambiguous way. In the literature there have been several attempts to deal with this problem[2,19-23]. We include only very brief comments, and for more details we refer the reader to the original literature. Our aim is to illustrate the uncertainty in theoretical estimates and to point out conclusions which are possible in spite of it, assuming, of course, that the necessary nuclear-physics calculations can be performed with accuracy.

For further reference it is useful to fix the normalization of H_W as follows

$$H_W = T^8\,(I = 1/2, \Delta\, S = 1) + \alpha T^8\,(I = 1) + \beta\, T^8\,(I = 0). \qquad (3.1)$$

Here all operators T transform as members of the same SU (3) octet. The parameters α and β contain model-dependent terms. Besides, if there are several nonequivalent tensors, say T, T', etc., α and β will depend on the corresponding reduced matrix-element ratios $< || T' || > / < || T || >$ etc.

In the early approaches, the necessary relations were established using current algebra and or SU (3) symmetry[2]. It was assumed that the differences between the initial and final baryon masses could be neglected. In this way the following sum rule (SR) was found

$$A\,(P_+^+) = \alpha \frac{2}{\sqrt{3}}\,[A\,(\Xi_-^-) - 2A\,(\Lambda_-^0)] \cong -\,\alpha\, 2.7 \cdot 10^{-7}. \qquad (3.2)$$

The methods employed to establish relations between PV amplitudes and SR Equ. (3.2) result in discrepancies when extended to p-wave amplitudes. The physical reason for this discrepancy can be traced to the SU(3)-symmetry breaking which is responsible for nonleptonic baryon decays in the first place. Besides, approaches based on current algebra and/or SU (3) symmetry work actually for off-mass-shell pions (i. e. $q^2 = 0$ instead of $q^2 = m_\pi^2$). Difficulties have been compensated by either introducing pole terms[24,25] or considering final-state interactions[26].

The vector-meson-pole model (VP)[24] leads to the expression[20]

$$A(P_+^+) = -\alpha \frac{2}{f_\pi}(F+D) - \beta \frac{C}{\sqrt{12}}\left(1+\frac{\delta}{\Phi}\right)(m_N - m_P)\frac{M_{K^*}^2}{M_\varrho^2} \cong$$
$$\cong -\alpha\, 1.06 \cdot 10^{-7} - \beta\, 0.64 \cdot 10^{-9}. \qquad (3.3)$$

The parameters appearing here are fixed by a fit to the s- and p-wave strangeness--changing amplitudes determined experimentally. The ϱ-pole term in Equ. (3.3) essentially gives the same contribution as would be obtained from the factorization-diagram (F) approximation[22,23].

The dominant contribution in the decuplet-pole model (DP)[25] is given by[20,25]

$$A(P_+^+) = -\alpha \frac{1}{f_\pi \sqrt{2}}(h_1 f_\omega + h_1 d_\omega) + \ldots = -\alpha\, 0.73 \cdot 10^{-7}. \qquad (3.4)$$

By taking into account final-state interactions (FI), SR (3.2) is modified into[19]

$$A(P_+^+) = \alpha \frac{2}{\sqrt{3}}\left[\frac{A(\Xi_-^-)}{0.983} - \frac{2A(\Lambda_-^0)}{1.22}\right] = -\alpha\, 1.22 \cdot 10^{-7}. \qquad (3.5)$$

The estimates Equs. (3.2—3.5) look natural as long as the operators T in Equ. (3.1) are considered as abstract entities endowed with the appropriate SU (3) transformation properties. Difficulties in understanding emerge as soon as one looks

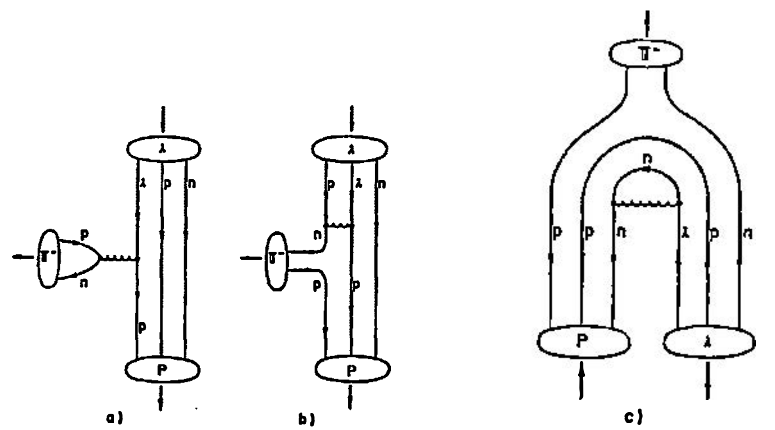

Fig. 2. The $\Lambda \to P\pi^-$ decay described by $H_W \sim (\bar{p}\lambda)(\bar{n}p)$. Hadron states are denoted by ellipses. Full lines correspond to quarks and heavy lines to IVB's.

at the quark content of weak-current products out of which the operators T are built. In free-quark models or in models in which baryon wave functions appear as products of quark fields[23,27,28] (PQF), contributions of certain operators in Equ. (3.1) do vanish identically[29]. The situation is illustrated in Figs. 2 and 3 for the Λ^0_- and P^+_+ amplitudes coming from the operators transforming as

$$T_\Lambda \sim (\bar{p}\,\lambda)\,(\bar{n}\,p) \tag{3.6}$$

and

$$T_P \sim (\bar{p}\,\lambda)\,(\bar{\lambda}\,p), \tag{3.7}$$

respectively. For local T_i operators, obtained from the Wilson expansion, the wavy IVB line is collapsed. The closed loops appearing in Figs. 2c and 3 exist no longer. The two λ lines in Fig. 3 remain unattached and the contribution vanishes

$$\langle N\pi^+ \mid T_P \mid P\rangle \equiv 0. \tag{3.8}$$

The Λ^0_- amplitude still receives contributions from the diagrams appearing in Fig. 2a, b, so this process is allowed.

This is altered when some radiative corrections, such as gluon exchanges among quarks, are considered. In an earlier Ref.[30], a simplified model with a real baryon loop was considered, as shown in Fig. 4. Baryon lines (BL) in this model emerge from the vertices containing weak currents composed of quarks. Each vertex is described by the standard SU (3) form-factor analysis of weak-current matrix elements

$$\langle B' \mid \bar{q}\,\Gamma_\mu\, q \mid B_i\rangle. \tag{3.9}$$

The pion vertex is determined by the standard baryon-pion coupling constant found empirically. The main shortcoming of this makeshift approximation is that one does not know form factors corresponding to currents containing charmed or unusually flavoured quarks c, for example

$$\langle B'\,(c) \mid \bar{c}\Gamma_\mu\, q \mid B_i\rangle. \tag{3.10}$$

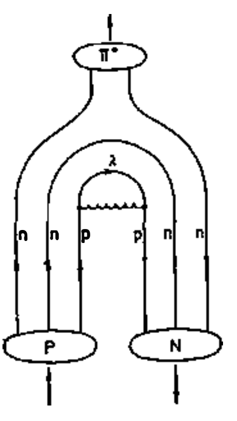

Fig. 3. The $P \to N\pi^+$ process described by $H_W \sim (\bar{p}\,\lambda)\,(\bar{\lambda}\,p)$.

The contributions from charmed and unusual baryons $B(c)$ were simply neglected in Ref[30]. This BL approximation can, from a certain point of view, be understood as going beyond the PQF-model approximation. The λ-quark loop of Fig. 3 exists even for local currents provided that gluon exchanges (GE), indicated by wavy lines in Fig. 5, are considered. However, such a diagram can be hindered very well by some effect similar to the Zweig-Iizuki (ZI) rule[31] by which the diagrams in Fig. 2a, b would be allowed. The situation is somewhat obscure, as the ZI rule is ordinarily connected with quark lines emerging from bound (baryon or meson) states. Some results were obtained when this rule was applied to quark lines emerging from currents[32]. It is difficult to guess whether a situation can emerge in which ZI hindered contributions can compete with ZI allowed ones, but which happen to be multiplied by very small numerical factors (i. e. sin Θ_C, $m_N - m_P$, etc.).

Two-quark operators lead simply to the insertion in one quark line and can easily be incorporated in diagrams similar to those in Figs. 2—4.

4. *Final results and conclusion*

Results in the form of simple sum rules have been obtained[2] for the conventional weak-interaction models[2,10], because in these models only one type of tensors T (3.1) appears. It has already been noted[30,33-34] and of*), that this is not so in the case of gauge-invariant models. The constants α and β can, therefore, be very complicated functions of the reduced matrix elements. Only some sort of dynamical calculation of these matrix elements can lead to definite theoretical predictions.

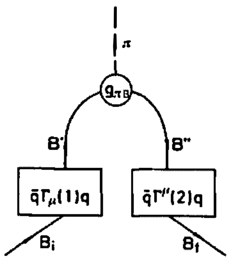

Fig. 4. The $B_i \to B_f \pi$ process. Full lines correspond to baryons and dashed lines to the pion. Boxes denote current form factors.

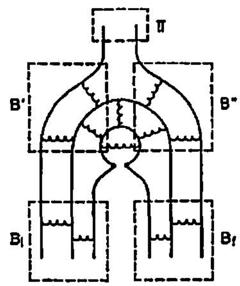

Fig. 5. The process described in Fig. 3 or Fig. 4. The IVB line is collapsed and gluon exchanges (wavy lines) are indicated. Dashed boxes represent bound states.

We first write the SU (3) tensorial decomposition of H_W^{eff} from Section 2. The SU (3) structure of the dominant terms in Equ. (2.16)

$$H_W^{\text{eff}}(DGG) = \frac{G}{2\sqrt{2}}\left\{+ 13.9\, c\, X^8_{1/2} + 3.73 \left[c^2 \left(\frac{10}{3\sqrt{5}} T^8_0 - \frac{4}{3\sqrt{2}} T^1_0 \right) + \right.\right.$$

$$+ s^2 \left(\frac{5}{\sqrt{15}} T^8_1(\lambda) - \frac{5}{3\sqrt{5}} T^8_0(\lambda) - \frac{4}{3\sqrt{2}} T^1_0(\lambda) \right) + sc \left(\frac{10}{\sqrt{30}} T^8_{1/2} + \right.$$

$$\left.\left.\left. + Y^8_{1/2} + \text{h. c.} \right) + U^8_1 + V^8_0 \right] + N_{23}\left(D^8_{1/2} - \text{h. c.}\right)\right\}. \tag{4.1}$$

Here the upper and lower tensorial indices refer to the SU (3) representation and the isospin content, respectively. The tensors T and D are composed of quarks

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*) V. I. Zakharov, private communication.

p, n and λ, while the tensors X, Y, U and V contain new quark flavours. The strangeness-conserving tensors $T(\lambda)$ contain a λ-quark pair; therefore, their contribution can be strongly suppressed by the ZI rule, as discussed in the preceding section. The same holds for tensors containing new quark flavours. The tensor D comes from the two-quark operator, see Table 1.

Table 1

Model	N		L	
	$\Delta S = 0$	$\Delta S = 1$	$\Delta S = 0$	$\Delta S = 1$
DGG	0	$N_{23} = -N_{32} = = s\,(tl - dr) - -c\,(dl - tr)$	0	0
GG	0	0	$2G$	0
SW	0	0	$1 + \left(\frac{1}{4} - \frac{1}{2}\sin^2\Theta_W - -\frac{1}{3}C_3\sin^2\Theta_W\right)$ (last three from neutral currents)	0

Table 2

Models	GG	SW
β_1	6.53	1.50
β_2	−39.36	0.11
β_3	30.03	−0.02
β_4	−44.60	−0.91
β_5	−12.51	−0.24

In the GG model, the SU (3) decomposition obtained from Equ. (2.22) is

$$H_W^{\mathrm{eff}}(\mathrm{GG}) = \frac{G}{2\sqrt{2}}\left\{25.81\,(X_1^8 + Y_0^8) + 5.08\left[c^2\left(\frac{10}{3\sqrt{5}}T_0^8 - \frac{4}{3\sqrt{2}}T_0^1\right) + \right.\right.$$

$$+ s^2\left(\frac{5}{\sqrt{15}}T_1^8(\lambda) - \frac{5}{3\sqrt{5}}T_0^8(\lambda) - \frac{4}{3\sqrt{2}}T_0^1(\lambda)\right) +$$

$$+ sc\left(\frac{10}{\sqrt{30}}T_{1/2}^8 + Z_{1/2}^8 + \mathrm{h.\,c.}\right) - \frac{1}{3}\left(\frac{10}{3\sqrt{5}}T_0^8 - \frac{8}{3\sqrt{8}}T_0^1 - \frac{5}{\sqrt{15}}T_1^8(\lambda) - \right.$$

$$\left.\left.\left. - \frac{5}{\sqrt{5}}T_0^8(\lambda) - \frac{4}{\sqrt{2}}T_0^1(\lambda)\right) + U_1^8 + V_0^8\right] + \hat{X}_1^8 + \hat{Y}_0^8 + W_0^1 + E_1^8 + E_0^8\right\}. \quad (4.2)$$

Here only the tensors T are composed of usually flavoured quarks. The tensor E is a two-quark operator defined by Equ. (2.7) and Table 1.

The decomposition of the effective Hamiltonian in the Salam-Weinberg model (2.25) is also very complicated

$$H_W^{\text{eff}}(\text{SW}) = \frac{G}{2\sqrt{2}}\Big\{3.02\Big[c^2\Big(\frac{10}{3\sqrt{5}}T_0^8 - \frac{4}{3\sqrt{2}}T_0^1\Big) + s^2\Big(\frac{5}{\sqrt{15}}T_1^8(\lambda) - $$
$$- \frac{5}{3\sqrt{5}}T_0^8(\lambda) - \frac{4}{3\sqrt{2}}T_0^1(\lambda)\Big) + sc\Big(\frac{10}{\sqrt{30}}T_{1/2}^8 + Y_{1/2}^8 + \text{h. c.}\Big) -$$
$$- \frac{1}{3}\Big(\frac{10}{3\sqrt{5}}T_0^8 - \frac{4}{3\sqrt{2}}T_0^1 - \frac{10}{3\sqrt{5}}T_0^8(\lambda) - \frac{8}{3\sqrt{2}}T_0^1(\lambda)\Big) + U_1^8\Big] + t_0^1\Big\} +$$
$$+ \frac{G}{\sqrt{2}}\Big\{3.02\,(2\sin^2\Theta_W - 1)\frac{1}{2}\Big[\frac{10}{9\sqrt{5}}T_0^8 - \frac{4}{9\sqrt{2}}T_0^1 + \frac{5}{\sqrt{15}}T_1^8(\lambda) +$$
$$+ \frac{5}{9\sqrt{5}}T_0^8(\lambda) + \frac{8}{9\sqrt{8}}T_0^1(\lambda) + Z_1^8 + V_0^8\Big] + \sin^2\Theta_W \frac{1}{3}(r_1^8 + s_0^8) +$$
$$+ \frac{1}{4}(1 - 2\sin^2\Theta_W)\,t_0^1\Big\} + \frac{G}{2\sqrt{2}}\Big[\frac{1}{3}\sin^2\Theta_W\, D_1^8 + \Big(1 + \frac{1}{4} - \frac{1}{2}\sin^2\Theta_W\Big)E_0^1\Big]. \tag{4.3}$$

The tensors r, s and t are products of the SU (4)-singlet part with the SU (3)-octet part. The two-quark operators D and E are defined by Equ (2.7) and Table 1.

One can immediately see from Equs. (4.1—4.3) that the parameters α and β appearing in Section 3 will in general be very complicated functions of the reduced matrix elements of tensor operators. For illustration, we write α and β for the DGG model

$$\alpha_{\text{DGG}} = \frac{1}{N_{23}} \frac{\langle N_f \| \Big(s^2 \frac{5}{\sqrt{15}}T_1^8 + U_1^8\Big) \| N_i\rangle}{\langle N_f \| \Big[D_1^8 + \frac{sc}{N_{23}}\Big(\frac{10}{\sqrt{30}}T_1^8 + Y_1^8\Big) + \frac{13\cdot 9c}{N_{23}}X_1^8\Big] \| N_i\rangle}, \tag{4.4}$$

$$\beta_{\text{DGG}} = \frac{1}{N_{23}} \frac{\langle N_f \| \Big[c^2 \frac{10}{3\sqrt{5}}\Big(1 - \frac{1}{2}\frac{s^2}{c^2}\Big)T_0^8 + V_0^8\Big] \| N_i\rangle}{\langle N_f \| \Big[D_0^8 + \frac{sc}{N_{23}}\Big(\frac{10}{\sqrt{30}}T_0^8 + Y_0^8\Big) + \frac{13\cdot 9c}{N_{23}}X_0^8\Big] \| N_i\rangle} + S. \tag{4.5}$$

Here S denotes the SU (3)-singlet isoscalar contribution coming from tensors such as T_0^1. If all tensors but T_l are unimportant, the DGG result is close to the standard Cabibbo model[2,20]. As can be seen from Equ. (3.3), the isoscalar contribution is most likely very small. The situation seems to be similar when all tensors containing p, n, and λ quarks are retained. Then

$$\alpha_{DGG} = \alpha_C \frac{1}{1 + \frac{N_{23}}{sc} \frac{\sqrt{30}}{10} \frac{\langle \| D_1^8 \| \rangle}{\langle \| T_1^8 \| \rangle}}; \quad \alpha_C = \frac{s}{c\sqrt{2}}. \tag{4.6}$$

According to the general concept of the model, $| \langle \| D_1 \| \rangle | \geq | \langle \| T_1 \| \rangle |$, while $N_{23} \sqrt{30}/10\ sc \cong 2.6$. The smallest value of the denominator is then $(-)1.6$, giving the maximal α which is slightly smaller than the value obtained in the Cabibbo model. The answer for the most general case can come only through dynamical calculations, in which a quark model, containing also unusally flavoured quarks, is employed. That still lies in the future.

In the framework of the ZI rule, only the contribution from the isoscalar term survives and therefore $\alpha = 0$ and the $A(P_+^+)$ amplitude is likely to be very small. The second factor in

$$\beta_{DGG} = \beta_C \frac{2c^2}{(2c^2 - s^2)\left(1 + \frac{N_{23}}{sc} \frac{\langle \| D_0^8 \| \rangle}{\langle \| T_0^8 \| \rangle}\right)}; \quad \beta_C = (2c^2 - s^2)/sc\sqrt{6} \tag{4.7}$$

is most probably of the order of unity or smaller, so $\beta \lesssim \beta_C \approx 3.7$. i.e.

$$| A(P_+^+) | \lesssim 2.3 \cdot 10^{-9}. \tag{4.8}$$

This is much smaller than the value in the Cabibbo model obtained from Equ. (3.3)

$$| A(P_+^+) | \cong 1.8 \cdot 10^{-8}. \tag{4.9}$$

General expressions for the GG model would be very lengthy and obscure, so we immediately apply the assumption that only the contributions from p, n and λ quarks are important. Then

$$\alpha_{GG} = \alpha_C \left(1 + \frac{1}{3s^2}\right) + e_1 \cong 8.6\, \alpha_C + e_1. \tag{4.10}$$

Unless there is a destructive interference between the first and second term(coming from the two-quark operator E) in Equ. (4.10), one should expect an enhancement over the Cabibbo-model value. The ratio β is again undetermined up to the two-quark contribution E, but it is likely to be comparable with the Cabibbo value. The ZI rule lowers the value of α

$$\alpha_{GG} \approx e_1, \tag{4.11}$$

while leaving β more or less unchanged.

For the SW model, the contribution from p, n and λ quarks indicates an enhancement over the result obtained in the Cabibbo model. Expression (4.3) is, however, more complicated than the simple expression used in deducing an optimistic prediction[29]. Keeping only the contributions from p, n and λ quarks, one obtains

$$a_{SW} = a_C \left[1 + \frac{1}{s^2} - \frac{2\sin^2\Theta_W}{s^2}(1+x)\right] = a_C(5 - 16\,x),$$
$$x = \frac{1}{6}\frac{\langle N_f \| (r_1^8 + D_1^8) \| N_i\rangle}{\langle \| T_1^8 \| \rangle}. \tag{4.12}$$

An enhancement due to the neutral-current contribution seems quite probable and dependent on the contribution of the two-quark operator D. Taking this to be negligible and estimating x on the basis of the diagrams presented in Fig. 4, we find[30]

$$2 \leq |\, a_{SW}/a_C \,| \leq 60. \tag{4.13}$$

As the approximations involved in computing (4.13) are rather drastic, this result can only serve as an indication. The ZI rule allows only neutral currents to contribute, giving

$$a_{SW} = \frac{\sqrt{30}}{5}\frac{\sin^2\Theta_W}{sc}x = a_C\frac{\sqrt{15}}{5}\frac{2\sin^2\Theta_W}{s^2}x = 12\,x\,a_C. \tag{4.14}$$

Unless something drastic happens with the isoscalar contribution, which does not seem likely, Equs. (3.2—3.5) would again indicate an enhancement somewhere in the region (4.13).

The possibility of enhancement for the models studied was always considered in comparison with a »standard« H_W of the form (3.1) with $a \equiv a_C$ and $\beta = \beta_C$.

This is somewhat misleading when the ZI rule is applied to the physical Cabibbo model. This model actually has

$$T_1^8 \equiv T_1^8(\lambda), \tag{4.15}$$

thus giving contributions only when the VP expression (3.3) is used.

Sum-rule expressions (3.2—3.5) depend for their validity on the assumptions concerning SU (3)-symmetry breaking. It seems, therefore, to be perfectly legitimate to strive to replace them by a dynamical calculation[21]. Besides, our previous discussion made it abundantly clear that the appearance of the multitude of nonequivalent tensors in the effective weak Hamiltonians makes sum rules inoperative unless supported by some dynamical calculation.

Purely dynamical calculation, based on the M. I. T. bag model[28], with two-quark operators omitted[21], gave for the SW H_W^{eff} (4.3)

$$A(P_+^+) = -\,1.9 \sin^2 \Theta_W A(\Lambda_-^0) \cong -2.6 \cdot 10^{-7}. \tag{4.16}$$

It is strange that this value falls inside the range of magnitudes obtained by combining Equs. (4.9) and (4.13)

$$4 \cdot 10^{-8} < |A(P_+^+)| < 10^{-6}. \tag{4.17}$$

The value in (4.16) should actually be compared with the Cabibbo model plus ZI-rule estimate, which, on the basis of (3.3) and (4.15), gives

$$A(P_+^+) \cong -0.24 \cdot 10^{-8}. \tag{4.18}$$

In respect to this, expression (4.16) is enhanced by two orders of magnitude. The result given in (4.18) corresponds to the factorization approximation[22] and to the earlier quark-model calculations[23].

Thus, in the case of the SW model, as compared with the Cabibbo model, there is every indication for the relative enhancement due to neutral currents. The absolute value of the $A(P_+^+)$ amplitude is, however, still questionable. Nevertheless, we may draw a tentative general conclusion that the experimental value

$$|A(P_+^+)| > 10^{-7} \tag{4.19}$$

will rule out both the DGG model and the »old« Cabibbo model. This conclusion assumes that:

— two-quark operators do not give contributions which would be much larger than four-quark operator contributions;

— two-quark operator contributions do not interfere destructively with four-quark operator contributions;

— contributions from unusually flavoured quarks are hindered.

This conclusion is much strengthened if the ZI rule applies to our case.

There are some indications for the existence of an $A(P_+^+)$ which is large[2, 35–38]. However, some of the calculated results might also be explained by the large isotensor ($\Delta I = 2$) component in the weak potential. It has been mentioned[39] that the appearance of the large isovector part in the ϱ-exchange PV potential for the SW model can also affect the analysis. More direct experimental information, coming from $\Delta I = 1$ processes, is, therefore, urgently needed.

Appendix A. The two-quark operator appearing in (2.1) is determined by the self-energy contribution (2.5). Divergences are to be removed by the wave function and mass renormalization[5,17], leaving as the leading finite contributions

$$\Sigma_{LR}(\hat{p}) = (-)\frac{i}{4\pi^2}\frac{fg}{M_W^2}[(\hat{p} - M_1)\mathcal{M}_2(1+\gamma_5)(\hat{p} - M_3)] \tag{A1}$$

for the left-right combination of currents, and

$$\Sigma_{LL}(\hat{p}) = \frac{i}{8\pi^2}\frac{f'g'}{M_W^2}[(\hat{p} - M_1)\hat{p}(1+\gamma_5)(\hat{p} - M_3)] \tag{A2}$$

for the left-left combination of currents. Here f, g, f' and g' are intermediate-vector-boson coupling constants, so $fg/M_W^2 \sim f'g'/M_W^2 \sim G_F/\sqrt{2}$. The masses M_1 and M_3 correspond ot the outgong (ingoing) particles, while $\mathcal{M}_2$ is the mass of the virtual quark. Expressions (A1) and (A2), corresponding to the operators with two derivatives (2.6) and three derivatives (2.7), respectively, are examples of the contributions one obtains by making all possible contractions in the product of the general charged currents $\mathcal{J}$ with left (C) and right (D) pieces

$$J_\pm^\mu = \bar{\Psi} C_\pm \gamma^\mu (1 - \gamma_5)\Psi + \bar{\Psi} D_\pm \gamma^\mu (1 + \gamma_5)\Psi. \tag{A3}$$

Intermediate states, corresponding to those contractions, introduce a diagonal matrix

$$\mathcal{M} + \mathcal{K}\hat{\gamma} \tag{A4}$$

in the space defined by Ψ, entering the PV part of the contribution as

$$\hat{\gamma}[C^-\mathcal{K}C^+ + C^+\mathcal{K}C^- - D^-\mathcal{K}D^+ - D^+\mathcal{K}D^-]\gamma_5 - [C^-\mathcal{M}D^+ - D^+\mathcal{M}C^- + \\ + C^+\mathcal{M}D^- - D^-\mathcal{M}C^+]\gamma_5 = \hat{\gamma}L\gamma_5 - N\gamma_5. \tag{A5}$$

When all quark masses are degenerate, the matrix is proportional to the unit matrix, thus leading to the L *matrix which conserves strangeness* ($\Delta S = 0$). The removal of quark-mass degeneracy also introduces $\Delta S = 1$ pieces in L to the order $\Delta m_q^2/M_W^2$. As L always contains C-C or D-D combinations, it is connected with the $\mathcal{O}_2^{(3)}$ (2.7) operator (see A2). On the other hand, the matrix N in which C and D mix is connected with $\mathcal{O}_2^{(2)}$ (2.6) operator (see A1). The matrix N always violates strangeness ($\Delta S = 1$) in its sector corresponding to standard p, n, λ quarks, irrespectively of the mass degeneracy. It is easy to see that its diagonal matrix elements vanish. As the matrix is diagonal

$$\mathcal{M}_{jk} = \mathcal{M}_j \delta_{jk},$$

one can write

$$N_{ii} = \sum_{j,k} \mathcal{M}_j \, [C^-_{ij} D^+_{ji} + C^+_{ij} D^-_{ji} - D^+_{ij} C^-_{ji} - D^-_{ij} C^+_{ji}]. \tag{A6}$$

As the matrices C and D are hermitic ($C^-_{ij} = C^+_{ji}$; $D^-_{ij} = D^+_{ji}$), the diagonal matrix element N_{ii} vanishes identically.

In an analogous way one can deal with the general neutral current

$$J^\mu_0 = \overline{\Psi} C_0 \gamma^\mu (1 - \gamma_5) \Psi + \overline{\Psi} D_0 \gamma^\mu (1 + \gamma_5) \Psi. \tag{A7}$$

Here C_0 and D_0 are diagonal matrices which commute with (A4). This gives

$$L = \mathcal{K} (C^2_0 - D^2_0), \qquad N = 0. \tag{A8}$$

Neutral currents induce only $\mathcal{O}^{(3)}_2$ strangeness-conserving operators.

Appendix B. In the Georgi-Glashow model there is one set of operators which mix.

$$\text{Set } P^\star_i \left(H = G + \frac{\cos \beta}{\sin \beta} F \right):$$

$$\begin{aligned}
P^\star_1 &= -\frac{1}{45} \{\mathcal{O}(1, H) + \mathcal{O}(H, 1) + 6\,\mathcal{O}(t, Ht) + 6\,\mathcal{O}(Ht, t)\},\\
P^\star_2 &= -\frac{1}{45} \{\mathcal{O}(1, H) + \mathcal{O}(H, 1)\},\\
P^\star_3 &= -\frac{1}{45} \{\mathcal{O}(1, H)\},\\
P^\star_4 &= -\frac{1}{45} \{\mathcal{O}(Ht, t)\},\\
P^\star_5 &= -\frac{1}{45} \{\mathcal{O}_1(H)\} - P^\star_4.
\end{aligned} \tag{B1}$$

The SW model has two sets of operators

Set R^{15}_i:

$$R^{15}_1 = \frac{1}{3} \sin^2 \Theta_W \{\mathcal{O}(C_3, 1)\},$$

$$R^{15}_2 = \frac{1}{3} \sin^2 \Theta_W \{\mathcal{O}(C_3\, t, t)\}$$

$$R_3^{15} = \frac{1}{3} \sin^2 \Theta_W \{\mathcal{O}(1, C_3)\}, \tag{B2}$$

$$R_4^{15} = \frac{1}{3} \sin^2 \Theta_W \{\mathcal{O}(t, C_3 t)\},$$

$$R_5^{15} = \frac{1}{3} \sin^2 \Theta_W \{\mathcal{O}_1(C_3)\} - R_2^{15}.$$

Set P_i^1:

$$P_1^1 = -\frac{1}{45} \{\mathcal{O}(1,1) + 6\,\mathcal{O}(t,t)\},$$

$$P_2^1 = -\frac{1}{45} \{\mathcal{O}(1,1)\}, \tag{B3}$$

$$P_3^1 = -\frac{1}{45} \{\mathcal{O}_1(1)\}.$$

Set R_i^1 is obtained by multiplication: $R_i^1 = \frac{1}{4} (1 - 2 \sin^2 \Theta_W) P_i^1$.

The constants β_i are given in Table 2, while

$$\gamma_1 = 2.43; \; \gamma_2 = -10.45 \text{ and } \gamma_3 = -7.62.$$

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NUKLEON PION AMPLITUDE KOJE NE ČUVAJU PARITET U BAŽDARNO INVARIJATNTNIM MODELIMA SLABIH, ELEKTROMAGNETSKIH I JAKIH MEĐUDJELOVANJA

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Sadržaj

U unificiranim baždarnim teorijama s uključenim jakim asimptotski slobodnim međudjelovanjima opisanim gluonima, izračunati su neleptonski efektivni Hamiltoniani. Detaljno su razmotreni dvo-kvarkovski operatori koji se javljaju u razvoju na malim udaljenostima.

Iz efektivnog Hamiltoniana moguće je dobiti dio nukleon-nukleon-pion amplitude koja narušava prostorni paritet. Opisani su problemi povezani s ovim izvođenjem, te uspoređene različite aproksimacije. Razmatrana je mogućnost postojanja analogona Zweig-Iizukinom pravilu. Postoje znakovi da neki modeli predviđaju da je nukleon-pion amplituda, koja ne čuva paritet, manja od 10^{-7} ($\hbar = c = 1$), što može poslužiti kao osnova za eksperimentalno razlučivanje.