QUARK MODEL CALCULATION OF THE PARITY VIOLATING NNπ COUPLING IN THE WEINBERG–SALAM MODEL

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We argue that there are no charged current contributions to the parity violating NNπ coupling except for small contributions from flavour symmetry breaking effects. From the neutral current product only the left–right chiral component contributes which is considerably enhanced due to gluon corrections and due to the lightness of current quark masses. The resulting parity violating NNπ coupling has a definite phase and is ~ 16 times stronger than the value used previously in nuclear physics calculations.

The successes of the Weinberg–Salam (WS) model [1] in explaining recent neutral current data in the lepton–lepton and lepton–hadron sector [2] has established the WS model as the prime candidate for neutral current phenomena.

In the hadron–hadron sector neutral currents are expected to substantially contribute to parity violating (p.v.), flavour conserving hadronic couplings which experimentally involves the whole spectrum of p.v. phenomena in nuclear interactions [3].

In this context the p.v. pion exchange force is particularly interesting because of its long range nature. Recently it has been argued that the neutral current contribution to the p.v. NNπ coupling $A_{\mathrm{NN}\pi}$ probably dominates over the charged current contribution [4–6]. Although we disagree with some of the arguments presented in refs. [5,6] to establish neutral current dominance we reach the same conclusion. In fact we shall argue that the charged current contribution to $A_{\mathrm{NN}\pi}$ is zero in the flavour symmetry limit.

Following ref. [8] (see also ref. [4]) we decompose the charged current product H^{W} and the neutral current product H^{Z} into 3 pieces O_+, O_- and O_0 transforming as $\underline{84}$, $\underline{20}$ and $\underline{15}$ representations of SU(4) [‡1].

$$\mathcal{H}_{\mathrm{eff}}^{\mathrm{W}} = -(G/2\sqrt{2}) \sin^2\theta_{\mathrm{C}} \sum_{\pm} \pm L_{\pm}^{\mathrm{W}} O_{\pm} \,, \tag{1}$$

$$\mathcal{H}_{\mathrm{eff}}^{\mathrm{Z}} = (G/2\sqrt{2}) \Big[(1-2\sin^2\theta_{\mathrm{W}}) \sum_{\pm} L_{\pm}^{\mathrm{Z}} O_{\pm} + \tfrac{1}{3} \sin^2\theta_{\mathrm{W}} L_0^{\mathrm{Z}} O_0 \Big] \,, \tag{2}$$

‡1 At first sight it is puzzling that only the symmetric 84, 20 and 15_{s} occur in the decomposition of the neutral current product since the neutral current product is *not* flavour symmetric at the SU(4) level. The neutral product is, however, flavour symmetric at the SU(3) level which explains the decomposition (2).

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where the relevant $\Delta I = 1$ components are given by

$$O_\pm = \tfrac{1}{2}(\bar{u}\gamma_\mu\gamma_5 u - \bar{d}\gamma_\mu\gamma_5 d)(\bar{s}\gamma_\mu s - \bar{c}\gamma_\mu c) + \tfrac{1}{2}(\bar{u}\gamma_\mu u - \bar{d}\gamma_\mu d)(\bar{s}\gamma_\mu\gamma_5 s - \bar{c}\gamma_\mu\gamma_5 c) \pm \text{Fierz transf.}, \quad (3)$$

$$O_0 = (\bar{u}\gamma_\mu\gamma_5 u - \bar{d}\gamma_\mu\gamma_5 d) \times (\bar{u}\gamma_\mu u + \bar{d}\gamma_\mu d + \bar{c}\gamma_\mu c + \bar{s}\gamma_\mu s) + (\bar{u}\gamma_\mu u + \bar{d}\gamma_\mu d + \bar{c}\gamma_\mu c + \bar{s}\gamma_\mu s) \times (\bar{u}\gamma_\mu\gamma_5 u - \bar{d}\gamma_\mu\gamma_5 d). \quad (4)$$

In the language of operator product expansions the operators $O_\pm$ in (1) and (2) are multiplicatively renormalizable four-quark operators of the lowest dimension that appear in the operator product expansion at short distances. The coefficients $L_\pm$ (we take $L_\pm^Z = L_\pm^W$) can be calculated in QCD to leading order [7]. In the free field limit $L_\pm$ and L_0^Z are equal to 1.

As will be discussed later on, there is in general a whole set of $\underline{15}$ operators of lowest dimension for which one has to find the appropriate linear combinations that are multiplicatively renormalized [8]. These operators can, however, always be rewritten in the effective form (2) for specific matrix elements. Note that the operators (3) and (4) are normally ordered and are therefore true four-quark operators which contribute only to connected quark diagrams.

From the last remark it is quite clear that there are *no* contributions from $O_\pm$ to the p.v. $NN\pi$ coupling in the valence quark approximation. The only contribution to $A_{NN\pi}$ comes from the operator O_0 which is entirely made from the product of *neutral currents.* Since the charged current contribution is zero in this approximation it also makes no sense to naively normalize the neutral current contribution to the charged current contribution as done in refs. [4–6].

Nonzero contributions to $A_{NN\pi}$ arise from the charged current product only when flavour symmetry breaking effects are included. First there is the factorizing contribution from the $\Delta I \neq 1$ piece $\approx \cos^2\theta_C$ $(\bar{u}d\bar{d}u)$ in the charged current product which gives a $\Delta I = 1$ contribution proportional to $(M_p - M_n)$ according to diagram Ia in fig. 1 [9,10]. One obtains $A_{pn\pi^+} \sim -0.1 \times 10^{-8}$. This contribution is so small that it will not be considered any further. Second there are the penguin type contributions from diagrams 2a and 2b in fig. 2 which do not cancel for $m_s \neq m_c$. Consideration of such mass breaking terms gives rise to additional contributions in the operator product expansion not included in (1) and (2) [11]. These new terms transform as $\underline{15}$ and have a left–right (LR) chiral structure. According to current lore these contributions are enhanced in fig. 2 since there is no helicity suppression in $\pi \to q_L\bar{q}_R$. Using estimates of refs. [11,12] one obtains

$$A_{pn\pi^+}^{\text{penguin}} \sim -(2\text{–}3) \times 10^{-8}. \quad (5)$$

Previously the charged current contribution to $A_{NN\pi}$ was obtained from the SU(3) sum rule [13]

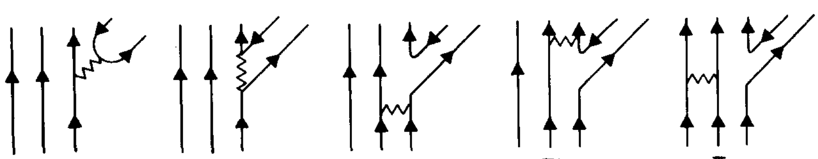
$$A_{pn\pi^+} = -(2/3)^{1/2} \operatorname{tg}\theta_C\,(2A_{\Lambda p\pi^-} + A_{\Xi^-\Lambda\pi^-}). \quad (6)$$

There are three types of contributions to the non-leptonic hyperon decay amplitudes appearing on the r.h.s. of eq. (6):

(i) contributions from connected quark diagrams or, equivalently, from the current algebra equal time commutator (ETC) term,

(ii) contributions from penguin type diagrams,

(iii) contributions from the factorizing diagrams I in fig. 1 which are proportional to $(M_{B_1} - M_{B_2})$.



Ia Ib IIa IIb III

Fig. 1. Valence quark diagrams for current × current contributions to $A_{pn\pi^+}$. Diagrams Ia and Ib are related by Fierz crossing and diagrams IIa and IIb by C-conjugation and crossing.

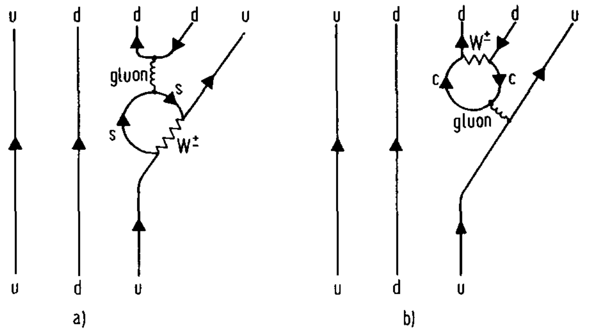


Fig. 2. Penguin type $W^{\pm}$ contribution to $A_{pn\pi^+}$.

Type (i) and (ii) [14] contributions satisfy sum rule (6), whereas type (iii) contributions do not since they are explicitly symmetry breaking. Type (i) contributions can be shown to give a zero r.h.s. in eq. (6) [15–17] leading to $A_{pn\pi^+} = 0$ in agreement with what was said before. This can also be phrased in the current algebra language where one finds $A_{pn\pi^+} \propto \langle n | H_{pc} | n\rangle = (F + D)$. Since $F/D = -1$ for connected quark diagrams [18,19] one finds again $A_{pn\pi^+} = 0$. The penguin type contributions (ii) *do* give a nonzero contribution to the r.h.s. of eq. (6). In fact, removing the factorizing type (iii) contributions from the experimental decay amplitudes (they destruct with the penguin contributions!) one calculates from the sum rule (6) $A_{pn\pi^+} = -6.2 \times 10^{-8}$ which is close to the estimate (5), and, quite naturally, close to the value $A_{pn\pi^+} \sim -5 \times 10^{-8}$ derived previously from eq. (6) [13]. The advantage of using the sum rule (6) instead of the direct estimate (5) is that the former evaluation is independent of the theoretical uncertainties inherent in estimating (5) [14]. Note that the phase of the penguin contribution to $A_{pn\pi^+}$ is definite.

O_0 to $A_{NN\pi}$ by considering the quark model diagrams in fig. 1. It turns out to be convenient to split O_0 into two pieces corresponding to the left–left (LL) and left–right (LR) chiral components of O_0. One has

$$O_0 = -\,[(\gamma_\mu)(-\gamma_\mu\gamma_5) + (-\gamma_\mu\gamma_5)(\gamma_\mu)]_{LL}$$
$$\times\,(\bar{u}u\bar{u}u - \bar{d}d\bar{d}d) + [(\gamma_\mu)(-\gamma_\mu\gamma_5) + (\gamma_\mu\gamma_5)(\gamma_\mu)]_{LR}$$
$$\times\,(\bar{u}u\bar{d}d - \bar{d}d\bar{u}u). \qquad (7)$$

The LL term in eq. (7) cannot contribute to $A_{NN\pi}$. This is clear for the factorizing diagram I just from its flavour content. The contributions to diagrams II and III vanish since the LL term in eq. (7) is flavour symmetric and therefore does not couple to the ground state baryons (Miura–Minakawa–Körner–Pati–Woo argument) [18,20].

Concerning the LR piece in eq. (7) we first discuss its contribution to the factorizing diagram I. Heuristically one expects this contribution to be enhanced since there is no helicity suppression in $\pi \to q_L\bar{q}_R$. The matrix element $\langle n\pi^+ | O_{0LR} | p\rangle$ is evaluated as in ref. [11]. The LR current product is Fierz transformed to a SP product which is rewritten into a VA current product using the equation of motion for the current quarks. One obtains

$$\langle n\pi^+ | O_{0LR} | p\rangle_I \qquad (8)$$
$$= -\bar{u}\,\tfrac{2}{3}\,\frac{M_p - M_n}{m_u - m_d}\,\frac{2m_\pi^2}{m_u + m_d} f_\pi \text{-} F_1(q^2 = m_\pi^2)u\,.$$

The exact numerical value of (8) depends on the choice of current quark masses for which there exist a number of estimates [21]. For definiteness we take $M_p - M_n = m_u - m_d$ and $m_u \sim m_d = 5$ MeV as in ref. [11], also we set $F(q^2 = m_\pi^2) = F(0) = 1$ which gives ~ -0.33 GeV^2 for (8). In the free quark limit $L_0^Z = 1$ this corresponds to (using $\sin^2\theta_W = \frac{1}{4}$)

$$A^{I}_{pn\pi^+}(L_0^Z = 1) = -11.5 \times 10^{-8}, \qquad (9)$$

which is already larger than the value $A_{pn\pi^+} \sim -5 \times 10^{-8}$ used in the preneutral current era [13].

The evaluation of the O_{0LR} contributions to diagrams II and III is more involved but may be done as described in detail for the LL charged current case in refs. [15,17]. One has

$$\langle n\pi^+ | O_0 | p\rangle_{II} = \bar{u}\,\{64 M_\pi (1$$
$$+\, O(M_\pi/M_p))H_2/L_-\}\,u\,,$$
$$\langle n\pi^+ | O_0 | p\rangle_{III} = \bar{u}\,\{-64\,M_\pi (1$$
$$+\, O(M_\pi/M_p))(M_\pi/M_p)H_3/L_-\}u\,, \qquad (10)$$

where H_2 and H_3 are wavefunction overlap determined in ref. [15] from a fit to non-leptonic hyperon decays [‡2]. Using $H_2 = 24 \times 10^{-3}$ GeV and $H_3 =$

‡2 A direct estimate of H_2 can be obtained from the decay $\Omega^- \to \Lambda K^-$ which has contributions from diagram II only. The above value of H_2 gives a partial rate $\Gamma_{\Omega^- \to \Lambda K^-}$ which is in good agreement with the recently measured partial rate $\sim 0.8 \times 10^{10}\ s^{-1}$ [17,22].

-19×10^{-3} GeV and $L_- = 3$ one has in the free field case ($L_0^Z = 1$)

$$A^{\mathrm{II}}_{\mathrm{pn}\pi^+}(L_0^Z = 1) = 2.5 \times 10^{-8} , \tag{11}$$

$$A^{\mathrm{III}}_{\mathrm{pn}\pi^+}(L_0^Z = 1) = -0.3 \times 10^{-8} . \tag{12}$$

The contribution from diagram III is quite small and will not be considered any further.

The matrix element $\langle \mathrm{n}\pi^+ | O_0 | \mathrm{p}\rangle$ can also be calculated using current algebra methods [23]. One obtains [‡3]

$$\langle \mathrm{n}\pi^+ | O_0 | \mathrm{p}\rangle_{\mathrm{ETC}} = f_\pi^{-1}\, 2\langle \mathrm{n} | O_0^{\mathrm{pc}} | \mathrm{n}\rangle , \tag{13}$$

where

$$O_0^{\mathrm{pc}} = 2(\bar{\mathrm{u}}\gamma_\mu \mathrm{u}\bar{\mathrm{u}}\gamma_\mu \mathrm{u} - \bar{\mathrm{d}}\gamma_\mu \mathrm{d}\bar{\mathrm{d}}\gamma_\mu \mathrm{d}) . \tag{14}$$

It should be noted that one *cannot* relate $\langle \mathrm{n} | O_0^{\mathrm{pc}} | \mathrm{n}\rangle$ by a SU(3) rotation to the corresponding one-particle matrix elements appearing in the same analysis of $\Delta S = 1$ nonleptonic hyperon decays, since the latter have a spatial (VV + AA) structure, whereas the former are pure VV. This fact has been overlooked in ref. [5]. In ref. [6] the ad hoc assumption $\langle \mathrm{B}' | \mathrm{VV} | \mathrm{B}\rangle = \langle \mathrm{B}' | \mathrm{AA} | \mathrm{B}\rangle$ was used for the relevant one particle matrix elements which is wrong at least in the quark model, where one finds $\langle \mathrm{B}' | (\mathrm{VV})_- | \mathrm{B}\rangle = \frac{1}{3}\langle \mathrm{B}' | (\mathrm{AA})_- | \mathrm{B}\rangle$ and $\langle \mathrm{B}' | (\mathrm{VV})_+ | \mathrm{B}\rangle = -\langle \mathrm{B}' | (\mathrm{AA})_+ | \mathrm{B}\rangle$ for the anti-symmetric and symmetric quark (or antiquark) flavour combinations. The latter relation is of course just the KKKPW argument [18,20] in disguise.

We shall in fact use the quark model to relate the unknown one-particle matrix element in (13) to the corresponding one appearing in the current algebra analysis of, e.g., $\Lambda \to \mathrm{p}\pi^-$:

$$\frac{\langle \mathrm{n} | 2(\bar{\mathrm{u}}\gamma_\mu \mathrm{u}\bar{\mathrm{u}}\gamma_\mu \mathrm{u} - \bar{\mathrm{d}}\gamma_\mu \mathrm{d}\bar{\mathrm{d}}\gamma_\mu \mathrm{d}) | \mathrm{n}\rangle}{\langle \mathrm{n} | [\bar{\mathrm{u}}\gamma_\mu(1-\gamma_5)\mathrm{s}\bar{\mathrm{d}}\gamma_\mu(1-\gamma_5)\mathrm{u}]_{\mathrm{pc}} | \Lambda\rangle} = \tfrac{2}{3}\sqrt{6} . \tag{15}$$

This then gives us

$$\langle \mathrm{n}\pi^+ | O_0 | \mathrm{p}\rangle_{\mathrm{ETC}} = \bar{\mathrm{u}}\, f_\pi^{-1}\, \tfrac{4}{3}\sqrt{6}(F/L_-)\mathrm{u} . \tag{16}$$

From the fit of Gronau [‡4] [19] one finds for the reduced one-particle matrix element $F \sim 30 \times 10^{-3}$ GeV. Using $L_- = 3$ one has in the free field case ($L_0^Z = 1$)

$$A^{\mathrm{ETC}}_{\mathrm{pn}\pi^+}(L_0^Z = 1) = 8.6 \times 10^{-8} . \tag{17}$$

The fact that the two values (11) and (17) have the same order of magnitude is not so surprising since there are good reasons to believe that the contributions of diagrams II are equivalent to the current algebra ETC term [15,17] in the symmetry limit. In fact one finds by explicit calculation

$$\frac{\langle \mathrm{n}\pi^+ | O_0 | \mathrm{p}\rangle_{\mathrm{ETC}}}{\langle \mathrm{n}\pi^+ | O_0 | \mathrm{p}\rangle_{\mathrm{II}}} = \frac{\langle \mathrm{p}\pi^- | O_-^{\Delta s=1} | \Lambda\rangle_{\mathrm{ETC}}}{\langle \mathrm{p}\pi^- | O_-^{\Delta s=1} | \Lambda\rangle_{\mathrm{II}}} . \tag{18}$$

The reason that the two estimates (11) and (17) do in fact differ by a factor ~3 comes from symmetry breaking effects and from differences in the details of the fits of refs. [15] and [19].

We shall now turn to the discussion of the renormalization constant L_0^Z in H_{eff}^Z. In ref. [8] the number of additional operators that mix with O_0 in the renormalization procedure was determined to be 4. The 5×5 renormalization mixing matrix has to be diagonalized to find the appropriate linear combinations that are multiplicatively renormalized. In ref. [25] this program was carried out numerically with the result

$$\mathcal{H}_{\mathrm{eff}}^Z(15) = (G/\sqrt{2})\tfrac{1}{3}\sin^2\theta_{\mathrm{W}}(1.50\,R_1^{15} + 0.11R_2^{15} - 0.02R_3^{15} - 0.91R_4^{15} - 0.24R_5^{15}) , \tag{19}$$

where

$$R_i^{15} = (O_i)^{\beta bj;\delta dl}_{\alpha ai;\gamma ck} : \bar{\psi}^{\alpha ai}\psi_{\beta bj}\bar{\psi}^{\gamma ck}\psi_{\delta dl} : , \tag{20}$$

$$O_1 = \{\gamma_\mu\gamma_5, C_3, 1; \gamma_\mu, 1, 1\}_+ ,$$

$$O_2 = \{\gamma_\mu\gamma_5, C_3, \lambda^m; \gamma_\mu, 1, \lambda^m\}_+ ,$$

$$O_3 = \{\gamma_\mu\gamma_5, 1, 1; \gamma_\mu, C_3, 1\}_+ ,$$

$$O_4 = \{\gamma_\mu\gamma_5, 1, \lambda^m; \gamma_\mu, C_3, \lambda^m\}_+ ,$$

$$O_5 = -O_2 + \mathrm{O}(C_3) . \tag{21}$$

[‡3] To be exact the current algebra commutation relation entering in the result (13) have to be evaluated separately for $O_\pm$ and the various pieces contributing to O_0 (see eq. (19)). The result is, however, the same as given by (13) and (14).

[‡4] The one-particle matrix element has recently been calculated using baryon wavefunctions derived from a harmonic oscillator potential [24]. The agreement with Gronau's fit value is satisfactory.

In (21) we have suppressed the Dirac indices α, β, γ and δ, the flavour indices a, b, c and d, and the colour indices i, j, k and l. The colour matrices λ^m are the usual SU(3) matrices where a summation over m is implied. C_3 is a flavour space matrix with only diagonal elements $1/2(u\bar{u})$; $-1/2(d\bar{d})$; $-1/2(s\bar{s})$; $1/2(c\bar{c})$. The contribution of the operator $O_5 = -O_2 + O(C_3)$ ($O(C_3)$ is defined in ref. [8]) vanishes on the mass shell [8] and will be dropped henceforth.

Because diagrams I and II (or the ETC term) have different colour configurations the effective L_0^Z resulting from (21) has to be separately evaluated for the two cases. Using $(\lambda^m)_i^j (\lambda^m)_l^k = 16/3\, \delta_i^k$ one obtains

$$(L_{0\,\mathrm{eff}}^Z)_{\mathrm{I}} = 7.0\,, \quad (L_{0\,\mathrm{eff}}^Z)_{\mathrm{II;ETC}} = -1.2\,. \tag{22}$$

Our final results for the neutral current contributions are then (for the charged current penguin type contribution see (5) and discussion thereafter)[‡5]

$$A_{\mathrm{pn}\pi^+}^{\mathrm{I}} = -80 \times 10^{-8}\,, \quad A_{\mathrm{pn}\pi^+}^{\mathrm{II}} = -3.0 \times 10^{-8}\,,$$

$$A_{\mathrm{pn}\pi^+}^{\mathrm{ETC}} = -10.6 \times 10^{-8}\,. \tag{23a,b,c}$$

Note that the phase of (23a) is fixed relative to strong interactions via the Goldberger–Treiman relation.

Whereas we are confident about the estimates (23b, c) concerning magnitudes we are not so sure about the phases relative to contribution (23a). The relative phases can in principle be calculated theoretically within a given model although in practise one determines the relative signs from fits to e.g. non-leptonic hyperon decays. The relative signs in (23b, c) have been determined by naively using the fits [15] and [19] which are, however, somewhat outdated since they do not take into account QCD corrections which have sizeable effects [11,12]. Using e.g. the results of the analysis of ref. [11] would change the sign of (23c) and reduce its magnitude considerably. Since, however, the factorizing part (23a) clearly dominates the contribution of diagram II (or the ETC term), the question of whether the contributions add or subtract is of no big practical importance at present given the uncertainties inherent in the estimate of the factorizing part. The question can, however, be settled with a new fit to the hyperon decays.

‡5 Neutral current penguins will not be considered here since their contributions are one order higher in the strong coupling constant α_s. See, however, ref. [26].

In conclusion we find that the dominant contribution to the $\Delta I = 1$ p.v. $\mathrm{NN}\pi$ coupling comes from the LR piece in the neutral current product of the Weinberg–Salam model. The matrix element of the LR operator is enhanced in a factorizing contribution due to the lightness of the u(d) current quark masses and due to the enhancement of the relevant operator O_0 from gluon corrections.

The exact strength of this contribution depends on the choice of current quark masses and on the details of calculating gluon corrections to the bare neutral current product. There are, however, no uncertainties in this number from unknown quark model matrix elements. Also the phase is calculable via the Goldberg–Treiman relation. The calculated p.v. $\mathrm{NN}\pi$ coupling is ~16 times stronger than the value used previously [13] in nuclear physics calculations [3]. Let us close by remarking that such a large p.v. long range force may be quite welcome for explaining some of the observed parity violations in nuclei (see e.g. ref. [4]).

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