Isospin Breaking and the precise Determination of V_{ud}

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Abstract

We examine a recent claim for a large first-order isospin breaking, due to $\rho - \omega$ mixing, in the superallowed O^+ to O^+ nuclear beta decays which provide the best determinations of V_{ud} . We find that no such effect exists, and present general arguments and explicit estimates to indicate that isospin breaking effects from the quark mass difference are not an obstacle to a precise measurement of V_{ud} .

1. If there exist only three generations, the weak mixing matrix [1,2] satisfies the unitarity relation

$$|V_{ud}|^2 + |V_{us}|^2 + |V_{ub}|^2 = 1. \quad (1)$$

A precise experimental test of (1) requires the measurement of the largest element V_{ud} to a level of about one part per thousand. Such an accuracy requires also that small and usually neglected effects, such as isospin breaking, must be under theoretical control when one calculates the matrix elements of the weak current. The most precise values of V_{ud} are obtained from a comparison of muon decay and $O^+ \rightarrow O^+$ superallowed Fermi nuclear beta decays. The electroweak and Coulombic radiative corrections have been studied in great detail, including several nuclear physics effects [3]. There is, in addition, another source of isospin violation, the mass difference between the up and down quarks which is more difficult to handle precisely. The $u-d$ mass difference has certain well-known aspects, such as its contribution to the proton-neutron mass difference. However, could it also have a "hidden" effect which would upset our ability to measure V_{ud} to the accuracy required?

In a recent paper, Lopez-Castro and Pestieau [4] presented a calculation of a previously neglected isospin breaking effect, $\rho-\omega$ mixing which arises if $m_u \neq m_d$. Within a vector dominance model [5] (in the following often abbreviated by VDM), they obtain a large correction to the Fermi form factors at zero momentum transfer:

$$F_+^{N-N'}(0) = \sqrt{2} \left(1 - \frac{\gamma_\rho}{\gamma_\omega} \frac{m_\rho^2}{m_\omega^2} \right), \quad (2)$$

where $\gamma_\rho, \gamma_\omega$ are the nucleon couplings of ρ (770) and ω (783); m_ρ is the ρ -mass, and m_ω describes $\rho-\omega$ mixing. Numerically,

$$\frac{\gamma_\rho}{\gamma_\omega} \frac{m_\rho^2}{m_\omega^2} \simeq (-2.1 \pm 0.2) 10^{-3}, \quad (3)$$

which is twice as large as the experimental error on V_{ud} and leads to a 2σ violation of the unitarity relation (1).

If this effect was correct, it would represent an important new development for precision tests of the standard model. In this note we reexamine this claim and show that $\rho-\omega$ mixing does not modify $F_+^{N-N'}(0)$ to first order in m_ρ^2 . (The incorrect result (2) is due to a violation of electromagnetic gauge invariance). We then argue that the $u-d$ mass difference does not affect $F_+(0)$ to first order in general. We conclude with an estimate of the second-order effects which show that these can safely be neglected.

2. We start by considering the specific case of $\rho-\omega$ mixing within the context of vector meson dominance. Including $\rho-\omega$ mixing, the general Lagrangian

has the form

$$L = L_K + L_{VDM} + L_{q^2} \quad (4)$$

consisting primarily of the kinetic energy terms

$$L_K = -\frac{1}{4}(\rho_{\mu\nu}\rho^{\mu\nu} + \omega_{\mu\nu}\omega^{\mu\nu} + F_{\mu\nu}F^{\mu\nu}) \quad (5)$$

and the vector couplings

$$\begin{aligned} L_{VDM} = & \bar{\psi}(i\partial^\mu\gamma_\mu - f_\rho\vec{\rho}\cdot\vec{\gamma} - f_\omega\frac{\omega}{2})\psi \\ & + \frac{1}{2}m_\rho^2\vec{\rho}^\mu\vec{\rho}_\mu + \frac{1}{2}m_\omega^2\omega_\mu\omega^\mu + m_\rho^2\omega_\mu\rho^\mu \\ & - \frac{em_\rho^2}{f_\rho'}\rho_\mu^3A^\mu - \frac{em_\omega^2}{f_\omega'}\omega_\mu A^\mu + \frac{1}{2}m_\gamma^2A_\mu A^\mu. \end{aligned} \quad (6)$$

Here, $\psi, \vec{\rho}_\mu, \omega_\mu, A^\mu$ are the nucleon, rho, omega and photon fields, respectively (ρ_μ^3 is the neutral component); $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$, etc. are the field strengths and f_ρ, f_ω are the ρ, ω couplings to the nucleons¹; the constants f_ρ' and f_ω' are related to f_ρ and f_ω in a way explained below².

Finally, L_{q^2} describes possible mixings in the field strength:

$$L_{q^2} = -\frac{1}{4}(\alpha F_{\mu\nu}\rho^{\mu\nu} + \beta F_{\mu\nu}\omega^{\mu\nu} + \gamma\rho_{\mu\nu}\omega^{\mu\nu}). \quad (7)$$

Because the matrix elements of L_{q^2} are all proportional to q^2 , (where q_μ is the momentum of the fields involved) L_{q^2} does not contribute to $F(0)$. [In principle, there could be direct couplings of the photon to the nuclei. In the spirit of VDM these should not be introduced. But even if we would include them, our conclusions remain unaltered (see footnote 4)].

The usual Lagrangian of VDM without ρ, ω mixing is obtained from (5) and (6) if we set $m_\rho^2 = 0$, and

$$f_\rho' = f_\rho, \quad f_\omega' = f_\omega, \quad m_\gamma^2 = e^2 \left(\frac{m_\rho^2}{f_\rho^2} + \frac{m_\omega^2}{f_\omega^2} \right). \quad (8)$$

This photon mass parameter is a well-known feature which insures the masslessness of the physical photons.

¹Our notation differs slightly from that used in [4]. We use f_ρ, f_ω instead of $\gamma_\rho, \gamma_\omega$; furthermore, our m_ρ^2 has a relative minus sign.

²We consider here only couplings to nucleons; if desired one can introduce other hadrons (pions, etc.) analogous to the nucleons.

Using L from Eqs. (5,6) we can now calculate the matrix elements of the charged and electromagnetic currents. From the diagrams of Fig.1 we get

$$\langle p' | J_\mu^+ | p \rangle = \sqrt{2} \bar{u}(p') \gamma_\mu \tau^+ u(p) \frac{m_\rho^2}{m_\rho^2 - q^2} \quad (9)$$

$$\begin{aligned} \langle p' | J_\mu^{em} | p \rangle = & \bar{u} \gamma_\mu \left\{ \frac{m_\rho^2}{m_\rho^2 - q^2} \frac{f_\rho \tau_3}{f_\rho' 2} + \frac{m_\omega^2}{m_\omega^2 - q^2} \frac{f_\omega 1}{f_\omega' 2} \right. \\ & \left. - \frac{m_\rho^2}{(m_\rho^2 - q^2)(m_\omega^2 - q^2)} \left(\frac{f_\rho \tau_3 m_\omega^2}{2 f_\omega'} + \frac{f_\omega m_\rho^2}{2 f_\rho'} \right) \right\} u. \end{aligned} \quad (10)$$

Higher orders in m_ρ^2 are neglected. The result (4) is obtained, if one sets $f_\rho' = f_\omega'$, $f_\omega' = f_\omega$ (the values of VDM without mixing). Then, at $q^2 = 0$, $q = p - p'$, (10) yields a correction of order m_ρ^2 to the electromagnetic form factor. We find

$$\langle p' | J_\mu^{em} | p \rangle = (\bar{u} \gamma_\mu (Q - m_\rho^2 (\frac{1}{2} \frac{f_\rho \tau_3}{f_\omega m_\rho^2} + \frac{1}{2} \frac{f_\omega}{f_\rho m_\rho^2})) \quad (11)$$

where $Q = \begin{pmatrix} 1 & 0 \\ 0 & 0 \end{pmatrix}$ is the charge matrix.

If we require the photon coupling to be proportional to Q , we must impose the condition

$$\frac{f_\rho}{f_\omega m_\rho^2} = - \frac{f_\omega}{f_\rho m_\omega^2} \quad (12)$$

to which we will return below³. Normalizing the (corrected) form factor of Eq. (11) to unity at $q^2 = 0$, the charged current form factor is altered with the result

$$\begin{aligned} \langle p' | J_\mu^+ | p \rangle &= \sqrt{2} \bar{u} \gamma_\mu \tau^+ u (1 + \frac{f_\rho}{f_\omega} \frac{m_\rho^2}{m_\omega^2} + O(q^2)) \\ \langle p' | J_\mu^{em} | p \rangle &= \bar{u} \gamma_\mu Q u (1 + O(q^2)) \end{aligned} \quad (13)$$

claimed in [4].

The relations (8) in usual VDM without mixing are obtained by imposing electromagnetic gauge invariance on the Lagrangian (4). In the presence of mixing, gauge invariance must be reexamined and we will show below that the relations (8) are indeed changed.

³If only couplings to the pion currents are considered, this condition need not be imposed since only the ρ couples to the pions and τ_3^2 is replaced by the charge matrix.

The gauge transformations

$$\begin{aligned} \psi(x) &\rightarrow e^{ieQ\Lambda(x)} \psi(x) \\ A_\mu &\rightarrow A_\mu - \partial_\mu \Lambda \end{aligned} \quad (14)$$

must be supplemented by

$$\begin{aligned} \rho_\mu^3 &\rightarrow \rho_\mu^3 - \frac{e}{f_\omega} \partial_\mu \Lambda \\ \omega_\mu &\rightarrow \omega_\mu - \frac{e}{f_\omega} \partial_\mu \Lambda, \end{aligned} \quad (15)$$

in order to maintain the gauge invariance of the couplings to the nucleons. The variations (14,15) induce the following changes in L :

$$\begin{aligned} \Delta L = & \rho_\mu \partial^\mu \Lambda \left(-\frac{em_\rho^2}{f_\rho} - \frac{em_\rho^2}{f_\omega} + \frac{em_\rho^2}{f_\rho'} \right) + \\ & \omega_\mu \partial^\mu \Lambda \left(-\frac{em_\omega^2}{f_\omega} - \frac{em_\rho^2}{f_\rho} + \frac{em_\omega^2}{f_\omega'} \right) + \\ & A_\mu \partial^\mu \Lambda \left(-m_\gamma^2 + \frac{e^2 m_\rho^2}{f_\rho f_\rho'} + \frac{e^2 m_\omega^2}{f_\omega f_\omega'} - \right. \\ & \left. \frac{1}{2} (\partial_\mu \Lambda)^2 \left(-\frac{e^2 m_\rho^2}{f_\rho^2} - \frac{e^2 m_\omega^2}{f_\omega^2} - \frac{2e^2 m_\rho^2}{f_\rho f_\omega} - m_\gamma^2 + \right. \right. \\ & \left. \left. \frac{2e^2 m_\rho^2}{f_\rho f_\rho'} + \frac{2e^2 m_\omega^2}{f_\omega f_\omega'} \right) \right). \end{aligned} \quad (16)$$

These vanish, if

$$\begin{aligned} \frac{m_\rho^2}{f_\rho'} &= + \frac{m_\rho^2}{f_\rho} \left(1 + \frac{f_\rho}{f_\omega} \frac{m_\rho^2}{m_\omega^2} \right) \\ \frac{m_\omega^2}{f_\omega'} &= \frac{m_\omega^2}{f_\omega} \left(1 + \frac{f_\omega}{f_\rho} \frac{m_\rho^2}{m_\omega^2} \right) \\ m_\gamma^2 &= e^2 \left(\frac{m_\rho^2}{f_\rho^2} + \frac{m_\omega^2}{f_\omega^2} + \frac{2m_\rho^2}{f_\rho f_\omega} \right) \end{aligned} \quad (17)$$

The conditions (17) replace (8); if $m_\rho^2 = 0$, they coincide.

If we now insert the value of f'_ρ of Eq.(17) into (10) and keep only terms linear in m_ρ^2 , the m_ρ^2 term is cancelled and we obtain

$$< p' | J_\mu^{em} | p > = \bar{u} \gamma_\mu Q_u (1 + O(q^2)) \quad (18)$$

as expected. Note that condition (12) need not be imposed as the terms proportional to m_ρ^2 vanish identically. Thus, ρ - ω mixing at first order only affects the q^2 -dependence and not the normalization at $q^2 = 0$. Eq. (16) shows that the effect claimed in [4] does not occur.

This result can be made obvious already at the Lagrangian level. Defining

$$\begin{aligned} \rho_\mu^3 &= \rho_\mu^3 - \frac{e}{f_\rho} A_\mu, & \rho_\mu^\pm &= \rho_\mu^\pm \\ \omega'_\mu &= \omega_\mu - \frac{e}{f_\omega} A_\mu \\ A'_\mu &= \kappa A_\mu, & \kappa &= \sqrt{1 + e^2 \left(\frac{1}{f_\rho^2} + \frac{1}{f_\omega^2} \right)} \simeq 1 + O(e^2) \end{aligned} \quad (19)$$

we get

$$\begin{aligned} L_K + L_{VDM} &= \\ &= \frac{1}{4} (F_{\mu\nu}^2 + \rho_{\mu\nu}^2 + \omega_{\mu\nu}^2) - \frac{e}{2\kappa f_\rho} F_{\mu\nu} \rho^{3\mu\nu} - \frac{e}{2\kappa f_\rho} F_{\mu\nu} \omega^{\mu\nu} \\ &+ \frac{1}{2} m_\rho^2 \rho_\mu^3 \rho^{\mu 3} + \frac{1}{2} m_\omega^2 \omega_\mu \omega^\mu + m_\rho^2 \rho_\mu^\pm \rho^{\mu \pm} \\ &+ \bar{\psi} (i \gamma^\mu \partial_\mu + f_\rho \gamma_\mu \vec{\rho}^\mu + \frac{\vec{\tau}}{2} - f_\omega \frac{\omega}{2} \gamma_\mu - \frac{e}{\kappa} Q \gamma_\mu A'_\mu) \psi. \end{aligned} \quad (20)$$

This equivalent form has an interesting structure. The physical photon, A'_μ , couples directly to the nucleons⁴ through a conserved current (the electric charge e is rescaled to e/κ). The mixing between ρ and γ , and ω and γ , are absorbed in the field strengths and thus explicitly vanish for $q^2 = 0$. The equivalence of (20) with the original form, first noted by Sakurai [5], offers an immediate proof that the form factors are not renormalized at zero momentum transfer.

3. The non-renormalization of $F_+(0)$ is in fact more general than our model calculation. The Ademollo-Gatto theorem [6] states that at $q^2 = 0$ vector form factors receive corrections only in **second order** in the symmetry breaking force.

⁴If we had introduced a direct photon-nucleon coupling in Eq (7) (see the discussion below Eq (7)), we would merely have to redefine the coupling constant in (20).

Since the proof of this theorem shows the (well-known) situations where the theorem does not apply, we sketch it here briefly in order to check the applicability of the theorem in the present case.

We start with the light-cone charge commutator

$$[T_+, T_-] = 2T_3, \quad (21)$$

where $T_\pm, T_3$ are the charges of vector currents corresponding to some $SU(2)$ subgroup of $SU(3)$ (not necessarily isospin); T_3 is diagonal.

When we take the matrix element of (21) between two equal states $|S\rangle$, the matrix element of T_3 is known absolutely since it the charge of a strictly conserved current. Thus we get

$$< S | [T_+, T_-] | S > = 2Q_s, \quad (22)$$

where Q_s is the charge of $|S\rangle$ with respect to T_3 . We next insert a complete set of states and separate the states $|r_\pm\rangle$ which lie in the same $SU(2)$ multiplet as $|S\rangle$:

$$\begin{aligned} (| < S | T_+ | r_+ > |^2 - | < S | T_- | r_+ > |^2) + \\ (\sum_{j+} | < S | T_+ | j_+ > |^2 - \sum_{j-} | < S | T_- | j_+ > |^2) = 2Q_s, \end{aligned} \quad (23)$$

where $|j_+ >, |j_+ >$ are states in other multiplets. Since the matrix elements $< S | T_\pm | j_\pm >$ are at least first order in an $SU(2)$ breaking force, the second piece on the l.h.s. is at least second order in the breaking. Note that $< f | T_+ | i > = F_+^{i \rightarrow f}(0)$.

If only one physical state $|r_+ >$ or $|r_- >$ exists, then its matrix element deviates from the symmetric value $\sqrt{2}Q_s$ only in second order in the symmetry breaking. This is the Ademollo-Gatto theorem. If, however, $|r_+ >$ is a mixture of physical states then only the sum of the squares of the respective matrix elements obeys the theorem. A well-known example of this effect are strangeness-changing weak decays, where the $T_\pm, T_3$ is generators of the U -spin- $SU(2)$ subgroup. If we take $|S\rangle = |K^+ >$, then $|r_+ > = \frac{1}{4} |\pi^0 > + \frac{\sqrt{3}}{2} |\eta_8 >$. Indeed one obtains

$$F_+^{K^+ \pi^0}(0) = 1 + \frac{3}{4} \left(\frac{m_d - m_u}{m_s - \bar{m}} \right), \quad (24)$$

but, in agreement with (23),

$$\frac{1}{4} |F_+^{K^+ \pi^0}(0)|^2 + \frac{3}{4} |F_+^{K^+ \eta}(0)|^2 = 1 + O\left(\frac{m_d - m_u}{m_s - \bar{m}}\right)^2 \quad (25)$$

In the present case, where T_+ is the usual isospin generator and $|S\rangle = |\text{proton}\rangle$, the only state which contributes in the first piece on the l.h.s. of

(23) is the neutron. Then, $\langle p|T_+|n \rangle = \sqrt{2}f_0(0)$ deviates from its symmetric value, $\sqrt{2}$, only in second order.

However, it could still happen that the corrections $\sum_j \langle s|T_+|j \rangle >^2$ are large, namely when $|j| > 1$ is (accidentally) degenerate with $|s \rangle$. To see this, we take the equation of motion of the charged current, i.e.,

$$i\partial_\mu \bar{u}\gamma^\mu d = (m_d - m_u)(\bar{u}d), \quad (26)$$

so that the matrix elements obey

$$(p_j + p_p)_\mu \langle j|\bar{u}\gamma_\mu d|p \rangle = \frac{m_d - m_u}{m_p^2 - m_j^2} \langle j|\bar{u}d|p \rangle, \quad (27)$$

and are enhanced if $(m_p - m_j)$ is of order $(m_d - m_u)$. Since there are no (multiparticle) hadronic states that near to the proton, the only dangerous states are $|p\gamma \rangle$, $|p\gamma\gamma \rangle$ with soft photons. However, they are taken into account explicitly by the electromagnetic corrections.

We expect therefore no exceptions to the Ademollo-Gatto theorem in $O^+ \rightarrow O^+$ superallowed beta decays.

4. Having convinced ourselves (and hopefully the reader) that isospin breaking from $m_u - m_d$ is absent to first order, we need to estimate the size of the second order effects. In the following we present estimates that these effects are small.

Consider first the deviations from unity of the nucleon vector form factor itself. A model independent estimate follows from a comparison with hyperon beta decay, where the deviation is proportional to $(m_u - m_d)^2$, $\hat{m} = \frac{1}{2}(m_u + m_d)$. The phenomenological fits [9] indicate an $SU(3)$ breaking of 4%.

Since most of it is almost certainly due to the axial current (for which the Ademollo-Gatto theorem does not apply), it is certainly very conservative to claim that the $SU(3)$ breaking in the vector current is less than 10%. Scaling this down to the level of isospin breaking for the proton-neutron form factor, one gets

$$F_1(0) \approx 1 \pm \frac{1}{10} \left(\frac{m_d - m_u}{m_p - \hat{m}} \right)^2 = 1 \pm 5 \cdot 10^{-5} \quad (28)$$

where the result [10]

$$\frac{m_d - m_u}{m_p - \hat{m}} = \frac{1}{43.5} \quad (29)$$

from chiral perturbation theory was used. This is far below the required accuracy.

Second, isospin breaking effects may occur in the nuclear wave function and matrix elements. The mass difference affects the matrix elements of the two-quark operator $H_m = m_u \bar{u}u + m_d \bar{d}d = \frac{1}{2}(m_u + m_d)(\bar{u}u + \bar{d}d) + \frac{1}{2}(m_u - m_d)(\bar{u}u - \bar{d}d)$

in the nuclear environment. Those involving quarks in the same nucleon (Fig. 2a) are included in the mass of the nucleon. Those with quarks in different nucleons (Fig. 2b) represent an isospin-breaking two-body force which can be taken into account by an infinite set of meson exchanges, as in Fig. 3, and lead to configuration mixing of states with different isospin.

Such a configuration mixing has been studied extensively for the coherent Coulomb effect [8]. The physics invoked is that the isotopic levels of nucleons consist of valence neutron and protons. There is an isospin violating interaction

$$V_C = Ze\alpha \left(\frac{\tau_3 + 1}{2} \right) \frac{1}{r} \quad (30)$$

between the nuclear core of charge Ze and the valence nucleons which mixes the energy levels of different isospins. We propose to treat the meson exchanges of Fig. 3 in an analogous fashion. For instance, taking the η -mesons as heavy (that is taking m_η larger than a typical nuclear momentum) results in a potential of the form

$$V_\eta = \frac{G_{\pi NN} G_{\eta NN}}{4\pi} \frac{m_\pi^2}{m_\eta^2} \frac{e^{-m_\pi r_{12}}}{r_{12}} \frac{(\vec{\sigma}\vec{p})_1}{2M} \frac{(\vec{\sigma}\vec{p})_2}{2M} (\tau_3)_1 \quad (31)$$

In order to compare this to the Coulomb interaction, V_C , we neglect the exponential and the spin factors (which would only decrease the effect). Then, the relative strength of V_η with respect to V_C is

$$\sigma = \frac{G_{\pi NN} G_{\eta NN}}{4\pi} \frac{m_\pi^2}{m_\eta^2} \frac{\vec{p}^2}{12M^2}. \quad (32)$$

There are then, in addition to the Coulombic correction δ_C , two further second order (in isospin-violation) corrections: the interference of V_η with V_C and a term coming from V_η alone. These corrections are then

$$\begin{aligned} \delta_{\text{int}} &\simeq \sigma \delta_C \\ \delta_\eta &\simeq \sigma^2 \delta_C. \end{aligned} \quad (33)$$

We can easily estimate σ . From chiral perturbation theory, we obtain

$$m_\eta^2 = \frac{1}{\sqrt{3}} \frac{m_d - m_u}{m_p - \hat{m}} (m_K^2 - m_\pi^2). \quad (34)$$

$SU(3)$ predicts the η -nucleon couplings to be

$$G_{\eta NN} = \left(3 - \frac{4D}{D + F} \right) G_{\pi NN}, \quad (35)$$

with $\frac{D}{D+F} = 0.63$ as obtained from fits to hyperon beta decays and the $SU(3)$ -Goldberger-Treiman relation which connects $\frac{D}{D+F}$ in weak decays to the strong couplings. Taking $< p^2 > \leq (250 MeV)^2$, we get with $\frac{C_{NN}^2}{4\pi} = 14$

$$\sigma = 6.10^{-4} . \quad (36)$$

Since δ_c is at most 3.10^{-3} , the corrections (33) may be safely neglected.

We can, in an analogous way, write an effective potential arising from $\rho - \omega$ mixing instead of $\eta - \pi^0$ mixing. Because of its short range, this should result in a small effect too. Our estimates confirm this expectation; $\rho - \omega$ mixing effects are at most of the order of the effects of $\eta - \pi^0$ mixing discussed here.

5. In summary, we have studied the influence of isospin breaking due to $m_d - m_u$ on the precise determination of V_{ud} . After showing explicitly that $\rho - \omega$ mixing (arising from $m_d \neq m_u$) does not give a first order effect within a vector dominance model, we argue that there is no reason to expect an exception from the Ademollo-Gatto theorem which requires generally that isospin breaking only affects vector form factors at second-order. We then estimate the second-order effects and conclude that they are completely negligible.

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Figure captions

Fig. 1. Vector dominance model for the weak and the electromagnetic form factor.

Fig. 2. Insertions of the isospin breaking two-quark operator H_m between two nucleons :
a) both quarks are in the same nucleon;
b) the two-quarks are in different nucleons.

Fig. 3. Meson-exchange model for the two-nucleon force induced by H_m as in Fig 2b.

Fig. 2

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Fig. 3