Limit on second class polar vector couplings in semileptonic weak interactions

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Although a great deal of experimental work has succeeded in limiting a possible second class axial coupling to about ten percent of the size of weak magnetism, published limits on the size of a possible second class polar vector current are one order of magnitude less sensitive. We summarize the present situation and demonstrate that muon capture provides the strictest present limit.

Over the past decade, a great deal of experimental effort has been expended in verifying the validity of the conserved vector current (CVC) hypothesis[1] and the lack of so-called second class currents[2] in the semileptonic $\Delta S=0$ weak interaction. Thus, the former has been studied via (i) shape factor measurements in allowed Gamow-Teller beta decay in $A=12$ and 20 systems,[3] (ii) rate and recoil polarization measurements in muon capture on ^{12}C (Ref. 4), (iii) ft value measurements in 0^+-0^+ Fermi transitions,[5] and good agreement with CVC has been observed in each case.[6] Correspondingly, a series of experiments involving (i) delayed particle-beta correlations in allowed Gamow-Teller decay of $A=8$ and 20 systems[7] and (ii) polarization-beta correlations in the superallowed $A=12$ and 19 beta decay[8] has shown the absence of second class *axial* currents—at least at the level of 10% of weak magnetism.[9]

However, corresponding attention has not been focused upon the possible presence of second class *polar* vector current effects. In this paper I would like to assess the current limits which can be placed upon such terms. The reason why it is difficult to detect the existence of second class vector effects may be seen by considering the most general matrix element of a polar vector current between nucleons

$$\langle p_{p_2}|V_\mu|n_{p_1}\rangle=\bar{u}(p_2)\left[g_V(q^2)\gamma_\mu+g_S(q^2)\frac{q_\mu}{2m_N}-i\sigma_{\mu\nu}q^\nu\frac{1}{2m_N}g_M(q^2)\right]u(p_1)\,, \tag{1}$$

where $q=p_1-p_2$. Here, $g_V(q^2),g_M(q^2)$ are the usual charge and weak magnetism form factors whose values at $q^2=0$ are given in terms of static electromagnetic moments of neutron and proton states

$$\begin{aligned}g_V(q^2=0)&=Q(\mathrm{p})-Q(\mathrm{n})=1.00\,,\\ g_M(q^2=0)&=\mu_a(\mathrm{p})-\mu_a(\mathrm{n})=3.70\,.\end{aligned} \tag{2}$$

It is these properties (or more precisely their nuclear analogs) which have been verified via the above-mentioned CVC tests. If, however, the weak vector current can be written as

$$V_\mu=V_\mu^{\mathrm{I}}+V_\mu^{\mathrm{II}}\,, \tag{3}$$

where V_μ^{I} is the conventional first class current which is related to the electromagnetic current by CVC,

$$[I_\pm,J_\mu^{\mathrm{EM}}]=\mp V_\mu^\pm\,, \tag{4}$$

while V_μ^{II} is a second class polar vector current, then the CVC relations given in Eq. (1) are unchanged, but the additional CVC prediction[10]

$$g_S(q^2)=0 \tag{5}$$

is violated. More generally, if we consider a nuclear matrix element between parent (daughter) states having spins J (J') and projection M (M') along some axis of quantization, we can write[11]

$$\begin{aligned}l^\mu\langle\beta_{p_2}|V_\mu|\alpha_{p_1}\rangle=\delta_{JJ'}\delta_{MM'}\left[a(q^2)\frac{P\cdot l}{2M}+e(q^2)\frac{q\cdot l}{2M}\right]+i\frac{1}{2M}b(q^2)C_{J'1;J}^{M'k;M}(\vec{q}\times\vec{l})_k\\ +C_{J'2;J}^{M'k;M}\left[\frac{1}{2M}f(q^2)C_{11;2}^{nn';k}l_nq_{n'}+\frac{1}{(2M)^3}g(q^2)P\cdot l\left[\frac{4\pi}{5}\right]^{1/2}Y_z^k(\hat{q})\vec{q}^{\,2}\right]\,,\end{aligned} \tag{6}$$

where

$$\begin{aligned}P_\mu&=(p_1+p_2)_\mu\,,\\ q_\mu&=(p_1-p_2)_\mu\,,\end{aligned} \tag{7}$$

and

$$M=\tfrac{1}{2}(M_1+M_2) \tag{8}$$

is the mean nuclear mass. Here, l^μ is the lepton current and is given by

$$l^\mu=\bar{u}(p)\gamma_\mu(1+\gamma_5)v(k)\,. \tag{9}$$

If α,β are members of the same isotopic multiplet, then CVC predicts $a(0)$, $b(0)$, and $g(0)$ in terms of static electromagnetic $E0$, $M1$, and $E2$ moments of initial and final states, respectively, while $V_\mu^{\rm II}$ contributes only to $e(q^2),f(q^2)$ which must vanish if no second class vector currents are present.[12] If α,β are not members of a common isotopic system then e,f need not vanish and thus we shall concentrate our discussion on the case of superallowed transitions. Also, in impulse approximation only $e(q^2)$ receives contributions from $g_S(q^2)$—$f(q^2)$ remains vanishing—and so we shall consider only $e(q^2)$ in the following.[11] [Note also that since $f(q^2)$ carries $\Delta J=2$, it will only contribute to transitions wherein a large number of additional structure functions complicate the analysis.]

It is now straightforward to see why the absence of second class vector currents is so difficult to verify since

$$q^\mu \bar{u}(p)\gamma_\mu(1+\gamma_5)v(k)=m_e\bar{u}(p)(1+\gamma_5)v(k)\ , \tag{10}$$

any effects in $e(q^2)$ must be[11]

$$\mathscr{O}\left[\frac{m_e^2}{m_N E}\right] \lesssim 0.1\%\ . \tag{11}$$

Of course, in the presence of SU(2) symmetry breaking, it is no longer required that $e(q^2)$ is strictly zero. We show in the Appendix, however, that such symmetry breaking effects are quite small,

$$e(q^2=0)<0.05\ , \tag{12}$$

in comparison with current experimental limits and can be neglected for the purposes of the present discussion.

The most straightforward approach to seeking $e(q^2)\neq 0$ effects is the analysis of careful spectral shape factor measurements. The simplest of these to interpret are experiments involving 0^+-0^+ Fermi transitions, for which $e\neq 0$ implies that if the shape factor is parametrized as

$$S(E)=1+\lambda_1 E+\lambda_2 E^2+\lambda_3\frac{1}{E}\ , \tag{13}$$

we have[13]

$$\lambda_3 \cong \tfrac{2}{15}R^2 m_e^2 \Delta + \frac{m_e^2}{M}\frac{e}{a}\ , \tag{14}$$

where $\Delta=M_1-M_2$ is the nuclear mass difference, and R is the nuclear radius. There are two such measurements which have been published. One by Thies, Appel, and Behrens involved the transition $^{33}K^m \to {}^{38}Ar$ and yielded[14]

$$\lambda_3^{\rm exp}=(-0.03\pm 0.08)m_e$$

or (15)

$$\frac{1}{A}e=-77\pm 205\ .$$

A second, by Kruger, Appel, Thies, and Behrens, utilized $^{34}Cl \to {}^{34}S$ and yielded[15]

$$\frac{dS}{dE}=\left[\lambda_1+2\bar{E}\lambda_2-\lambda_3\frac{1}{\bar{E}_2}\right]^{\rm exp}=(-0.03\pm 0.08)/m_e\ , \tag{16}$$

where $\bar{E}\sim 2.2$ MeV is the mean beta energy for the transition. Comparison with the theoretical expectation[11]

$$\left[\frac{dS}{dE}\right]^{\rm th}=\frac{R^2}{15}\left[4\Delta-2\frac{m_e^2}{\bar{E}^2}\Delta-8\bar{E}\right]-\frac{e}{a}\frac{m_e^2}{M\bar{E}^2} \tag{17}$$

gives

$$\frac{1}{A}e=45\pm 120\ . \tag{18}$$

An alternate approach to bounding second class vector effects is measurement of the nuclear-spin—β-momentum correlation in the superallowed decay[8]

$$^{19}Ne \to {}^{19}F+e^+ +\nu_e\ .$$

The theoretical expression for the correlation coefficient

$$A(m_e)=2\frac{ac\frac{1}{\sqrt{3}}+c^2+\cdots}{a^2+c^2+\cdots} \tag{19}$$

is *very* sensitive to the Gamow-Teller matrix element c (since $c^{\rm exp}\approx -1.6$) which in turn is determined from the ft value as a function of the scalar form factor e. The detailed analysis has been given elsewhere,[16] with the result

$$\frac{1}{A}e=15\pm 26\ . \tag{20}$$

Thus, from spectral analysis of nuclear beta decay the available limits on the absence of second class vector effects are surprisingly crude.

An alternative approach involves examination of ft values of 0^+-0^+ Fermi transitions, wherein we expect[17]

$$ft^{0^+\text{-}0^+}=ft^{0^+\text{-}0^+}({\rm CVC})\Big/\left[1+\frac{e}{a}\frac{m_e}{m_N}\left\langle\frac{m_e}{E}\right\rangle\right]\ . \tag{21}$$

Since $\langle m_e/E\rangle \propto R/\alpha Z$, the existence of a second class vector term e would be revealed by a systematic deviation from constancy of the experimental ft values. Thus, a two parameter fit by Towner and Hardy gave[18]

$$\frac{1}{A}e=2\pm 11\ , \tag{22}$$

while a more recent analysis by Szybisz and Silbergleit yielded[19]

$$\frac{1}{A}e=-3.2\pm 8\ . \tag{23}$$

Although it is possible to improve this analysis by normalizing to the standard ^{14}O ft value, present anomalies in the measurement of the corresponding Q value need to be cleared up before this can be performed definitively.[20] Even then the limits are not particularly stringent[19]

$$\frac{1}{A}e=-4.4\pm 5.2\ . \tag{24}$$

Thus, the most precise ($\sim 0.1\%$) nuclear beta decay measurements are unable to limit the size of the second class vector form factor e to less than one or two times the size of the weak magnetism term. In view of this, it is perhaps somewhat surprising that with one additional assumption, one can use the $\sim 10\%$ measurements on muon capture to provide a limit which is nearly a factor of 3 better than those quoted above.

The additional assumption—besides CVC and the absence of second class axial currents—is the validity of the partial conservation of axial-vector current (PCAC) hypothesis for the determination of the induced pseudoscalar form factor.[11] Actually, this hypothesis has been nicely verified in recent sets of experiments involving (i) analysis of muon capture and beta decay[21]

$$^{16}N(0^-, 120\ keV) \leftrightarrow {}^{16}O\ ,$$

and (ii) rate and recoil polarization measurements involving[22]

$$\mu^- + {}^{12}C \rightarrow {}^{12}B + \nu_\mu\ .$$

If we now consider muon capture between states of the same isotopic multiplet only two cases are available kinematically,

$$\mu^- + p \rightarrow n + \nu_\mu\ ,$$

$$\mu^- + {}^3He \rightarrow {}^3H + \nu_\mu\ .$$

Allowing for the presence of a scalar coupling we find then[23]

$$\frac{\Gamma_{tot}(g_S \neq 0, g_T = 0)}{\Gamma_{tot}(g_S = 0, g_T \neq 0)} = \frac{g_V^2 + 3g_A^2 + 2\frac{m_\mu}{2M}\left\{ g_V g_S + g_A \left[g_A + 2(g_V + g_M) - g_P \frac{m_\mu}{2M} \right] \right\}}{g_V^2 + 3g_A^2 + 2\frac{m_\mu}{2M} g_A \left[g_A + 2(g_V + g_M) - g_T - g_P \frac{m_\mu}{2M} \right]}\ , \tag{25}$$

where the notation is standard. Thus, experiments which have previously *assumed* $g_S = 0$ in order to place limits on g_T can correspondingly be utilized—using the result $g_T = 0$ verified in nuclear beta decay—to limit the size of a possible second class vector coupling g_S. In this way Kim and Primakoff have used the muon capture lifetime measured for ^{3}He

$$\Gamma(\mu^- + {}^3He \rightarrow {}^3H + \nu_\mu)_{exp} = \begin{cases} 1505 \pm 46\ sec^{-1}\ (\text{Ref. 24}) \\ 1465 \pm 67\ sec^{-1} (\text{Ref. 25}) \underset{av}{\rightarrow} 1489 \pm 40\ sec^{-1} \end{cases} \tag{26}$$

together with electron scattering data on ^{3}H and ^{3}He to determine[26]

$$g_T(q^2 = -m_\mu^2) = 0.4 \pm 2.0\ , \tag{27}$$

which becomes, using $g_T = 0$ and Eq. (25),

$$g_S(q^2 = -m_\mu^2) = -0.5 \pm 2.4\ . \tag{28}$$

Also, a CERN-Saclay-Bologna collaboration has averaged new capture rates measured in liquid hydrogen together with older values obtained with gaseous targets

$$\Gamma(\mu^- + p \rightarrow n + V_\mu) = \begin{cases} 515 \pm 85\ sec^{-1}\ (\text{Ref. 27}) \\ 464 \pm 42\ sec^{-1}\ (\text{Ref. 28})\ \text{liquid}\ H_2 \\ 460 \pm 20\ sec^{-1}\ (\text{Ref. 29}) \\ 686 \pm 88\ sec^{-1}\ (\text{Ref. 30}) \\ 651 \pm 57\ sec^{-1}\ (\text{Ref. 31})\ \text{gaseous}\ H_2 \end{cases} \tag{29}$$

to obtain[32]

$$g_T(q^2 = 0) = 0.3 \pm 1.9\ , \tag{30}$$

which becomes, using $g_T = 0$ and Eq. (25),

$$g_S(q^2 = 0) = -0.4 \pm 2.3\ . \tag{31}$$

These analyses, $|g_S| \lesssim 2$, limit the size of a possible scalar coupling to about half the value of weak magnetism and are the best that can be accomplished with present techniques. The improved 0^+-0^+ measurements suggested by Koslowsky *et al.* could be very helpful in decreasing the range of these limits.[33] It should be noted that we are yet some distance from the value of g_S suggested by Blin-Stoyle on dimensional grounds[34] in a CVC violating theory.

As a final comment, we note that the limits derived here on the size of the scalar coupling are relevant also in limiting the value of possible charged Higgs couplings,[35] for which there exists a scalar interaction of the form

$$\frac{G}{\sqrt{2}} \cos\theta_c \bar{u} d \bar{u}_e (1 + \gamma_5) v_{\nu_e} \lambda m_e\ . \tag{32}$$

Our bounds given previously imply

$$|\lambda| < \frac{1}{m_N} \sim 10^{-3}\ MeV^{-1}\ . \tag{33}$$

In a simple single Higgs doublet model we expect[35]

$$\lambda \sim \frac{m_u + m_d}{m_N^2} < 4\times 10^{-7} \tag{34}$$

for $m_N > 5$ GeV as given by Cornell data.[36] However, in models wherein one has two or more such doublets the couplings have no such constraint, Eq. (34), and the limits on e may be of some use.

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APPENDIX: SCALAR COUPLING INDUCED BY SYMMETRY BREAKING

In order to deduce the size of the scalar coupling induced by symmetry violation, we note that in the presence of quark mass differences and electromagnetism the naive CVC condition

$$i\partial^\mu V_\mu^\pm = 0 , \tag{A1}$$

is replaced by

$$i(\partial^\mu - ieA^\mu)V_\mu^\pm = \pm(m_d - m_u)S^\pm , \tag{A2}$$

where

$$S^+ = \bar{u}d = (S^-)^\dagger \tag{A3}$$

is a scalar density. The Ademollo-Gatto theorem,[38] of course, guarantees that $a(q^2=0)$ will be unchanged from its unbroken value up to terms of second order in SU(2) breaking. That is not the case for $e(q^2)$ however, which is nonzero to first order in the breaking.

First, omit electromagnetism and consider just the effect of the u,d quark mass difference. We then have for a matrix element of Eq. (A2) between two states which are isotopic analogs

$$a(q^2)\Delta + e(q^2)\frac{q^2}{2M} = S(q^2)(m_d - m_u) , \tag{A4}$$

where we have defined

$$\langle \beta | S | \alpha \rangle = \delta_{JJ'}\delta_{MM'}S(q^2) . \tag{A5}$$

Assuming $a(q^2)$ to be dominated by the ρ and $S(q^2)$ by δ (980 MeV), we have[39]

$$a(0)\Delta = S(0)(m_d - m_u) ,$$
$$e(0) = S(0)(m_d - m_u)\left[\frac{1}{m_\delta^2} - \frac{1}{m_\rho^2}\right]2M . \tag{A6}$$

Since electromagnetism has been hypothetically turned off we have

$$m_d - m_u = m_n - m_p = \Delta , \tag{A7}$$

and hence $S(0) = a(0)$. Then,

$$e(0) = a(0)(m_d - m_u)2M\left[\frac{1}{m_\delta^2} - \frac{1}{m_\rho^2}\right]$$
$$\cong Aa(0) = -0.01 . \tag{A8}$$

A similar result can be derived from the MIT bag model[40] where, writing the quark wave function as

$$\Psi(\vec{r}) = \begin{bmatrix} iu(r)\chi \\ v(r)\vec{\sigma}\cdot\hat{r}\chi \end{bmatrix} , \tag{A9}$$

it is straightforward to demonstrate that[41]

$$e = -a\tfrac{1}{3}2M \int d^3r[u_u(r)v_d(r) - u_d(r)v_u(r)]$$
$$\approx a\frac{2M}{3}(m_d - m_u)\int d^3r u_u(r)v_u(r)R\frac{3\omega - 2}{2\omega(\omega-1)}$$
$$\approx -0.02aA . \tag{A10}$$

Thus, by either approach, the value of the scalar coupling induced by quark mass difference effects is small and negative.

Of course, electromagnetism also breaks SU(2) invariance and therefore can generate an effective scalar coupling. However, any such effects are included in the standard radiative correction[42] and therefore need not be considered further.

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