PERGAMON

Progress in Particle and Nuclear Physics 46 (2001) 413–457

Progress in Particle and Nuclear Physics

http://www.elsevier.nl/locate/npe

Beta Decay Beyond the Standard Model

P. HERCZEG

Theoretical Division, Los Alamos National Laboratory, P.O. Box 1663, MS B283, Los Alamos, NM 87545

Abstract

We review and discuss beta decay interactions in extensions of the Standard Model, and the role of beta decay experiments in obtaining information on them.

PII: S0146-6410(01)00149-1

Contents

1 Introduction

Nuclear and neutron beta decay played a prominent role in the developments that led to the Standard Model (SM) [1].[1] The new possibilities for experimental studies of beta decay that opened up after the discovery of parity violation resulted soon in the recognition of the V-A structure of the weak interactions. This period culminated in the formulation of the universal V-A current x current theory of the weak interactions. While the V-A theory remained consistent with all data, there was strong motivation to search for deviations from the V-A structure in beta decay (and other processes). One of the reasons was CP-violation, discovered in 1964. The V-A theory could not account for this effect. While it was recognized the the observed CP-violation could be due to a new force, it was equally possible that with appropriate modifications the weak interactions could account for it. Another motivation to search for new interactions was the non-renormalizability of the V-A theory.

Today the motivation to search for new interactions is not weaker. Despite the remarkable success of the SM, for many theoretical reasons, and especially because of the large number of undetermined parameters of the model, the existence of new physics is expected. In fact, we have already the first strong experimental evidence, in the form of neutrino oscillations, that some extension of the SM is required. The origin of CP-violation is still an open question, although one of the possibilities is that the SM can account for it. Regardless, among the new interactions there may be new sources of CP-violation. It is interesting that without some new source of CP-violation the baryon asymmetry of the universe cannot be generated.

In this review article we shall discuss beta decay interactions in extensions of the SM [3]. We shall review the existing bounds on new interactions provided by beta decay experiments, and consider the constraints on them from other sources. The purpose is to assess what sensitivities would be required in beta decay experiments, to obtain new information. In the next section we consider the general form of possible new $d \to ue^-\bar{\nu}_e$ interactions, and the resulting effective Hamiltonian for nucleon beta decay. In Sections 3, 4, and 5 we focus on the time-reversal (T) invariant components of $d \to ue^-\bar{\nu}_e$ interactions containing vector and axial-vector quark currents (V,A-type interactions), scalar quark currents (S-type interactions), and tensor quark currents (T-type interactions), respectively. In Section 6 we discuss the T-violating components of the $d \to ue^-\bar{\nu}_e$ interactions. In Section 7 we summarize our conclusions.

[1] In the following we shall understand the electroweak component of the Standard Model to be the $SU(2)_L \times U(1)$ gauge theory [2], containing three fermion families and one Higgs doublet, and only left-handed neutrinos.

2 General Considerations

In the SM the $d \to ue^-\bar{\nu}_e$ (and $u \to de^+\nu_e$) transition underlying beta decay arises from W-exchange, and has the $V - A$ form [2]

$$H = (GV_{ud}/\sqrt{2})\bar{e}\gamma_\lambda(1-\gamma_5)\nu_e\bar{u}\gamma^\lambda(1-\gamma_5)d + \text{H.c.} \quad , \tag{1}$$

where $G/\sqrt{2} = g^2/8M_W^2$, and V_{ud} is the ud-element of the Kobayashi-Maskawa matrix. The field $\frac{1}{2}(1-\gamma_5)\nu_e$ in the interaction (1) represents a massless two-component neutrino, which is the $T_z = +1/2$ state of the $\mathrm{SU}(2)_L$ doublet involving the electron.

In many extensions of the SM there are new contributions to $d \to ue^-\bar{\nu}_e$. We shall consider here only such contributions from new physics that can be represented by nonderivative local four-fermion couplings. [3] The most general $d \to ue^-\bar{\nu}_e$ four-fermion interaction involving the neutrino states $\nu_e^{(L)}$ and $\nu_e^{(R)}$, where $\nu_e^{(L)}$ is the neutrino state in the $W^+ \to e^+\nu_e^{(L)}$ amplitude and $\nu_e^{(R)}$ is a right-handed singlet state, [4,5] can be written as

$$H_\beta = H_{V,A} + H_{S,P} + H_T \quad , \tag{2}$$

where

$$\begin{aligned} H_{V,A} = &\bar{e}\gamma^\lambda(1-\gamma_5)\nu_e^{(L)}[a_{LL}\bar{u}\gamma_\lambda(1-\gamma_5)d + a_{LR}\bar{u}\gamma_\lambda(1+\gamma_5)d] \\ &+\bar{e}\gamma^\lambda(1+\gamma_5)\nu_e^{(R)}[a_{RR}\bar{u}\gamma_\lambda(1+\gamma_5)d + a_{RL}\bar{u}\gamma_\lambda(1-\gamma_5)d] \\ &+ \text{H.c.} \quad , \end{aligned} \tag{3}$$

[2] Our metric, γ matrices and $\sigma_{\lambda\mu}$ are the same as in Ref. [4].

[3] Accordingly, we shall not consider beta decay interactions involving second-class currents [5]. Interactions involving second class currents cannot be introduced without spoiling the renormalizability of the theory or without having to face severe theoretical and phenomenological difficulties [6]. For reviews of the present experimental limits on second-class currents see Ref. [7].

[4] There is no experimental evidence at present for the existence of right-handed neutrinos; the singlet neutrino which may be required by neutrino oscillation data [8] can have either helicity. The bound on the number of light neutrinos from standard big-bang nucleosynthesis allows still the existence of an additional (singlet) neutrino, even if it had interactions as strong as the weak interaction [9]. In the presence of a large $\nu_e - \bar{\nu}_e$ asymmetry the number of such neutrinos allowed by big-bang nucleosynthesis is much larger [10].

[5] Couplings involving neutrino states other than $\nu_e^{(L)}$ and $\nu_e^{(R)}$ are possible, but for these in most cases additional constraints apply.

$$H_{S,P} = \bar{e}(1-\gamma_5)\nu_e^{(L)}[A_{LL}\bar{u}(1-\gamma_5)d + A_{LR}\bar{u}(1+\gamma_5)d] \\ +\bar{e}(1+\gamma_5)\nu_e^{(R)}[A_{RR}\bar{u}(1+\gamma_5)d + A_{RL}\bar{u}(1-\gamma_5)d] \\ + \text{H.c.} \quad , \tag{4}$$

$$H_T = \alpha_{LL}\bar{e}\,\frac{\sigma_{\lambda\mu}}{\sqrt{2}}(1-\gamma_5)\nu_e^{(L)}\bar{u}\,\frac{\sigma_{\lambda\mu}}{\sqrt{2}}(1-\gamma_5)d \\ +\alpha_{RR}\bar{e}\,\frac{\sigma_{\lambda\mu}}{\sqrt{2}}(1+\gamma_5)\nu_e^{(R)}\bar{u}\,\frac{\sigma_{\lambda\mu}}{\sqrt{2}}(1+\gamma_5)d + \text{ H.c.} \tag{5}$$

The Hamiltonian (3) has vector (V) and axial-vector (A) structure, the one in Eq. (4) has scalar (S)-type (proportional to $\bar{u}d$) and pseudoscalar (P)-type (proportional to $\bar{u}\gamma_5 d$) terms, and the Hamiltonian (5) contains tensor (T)-type interactions.

In the Hamiltonians (3) - (5) the first and second subscript on the coupling constants gives the chirality of the neutrino and of the d-quark. Note that there are no tensor couplings of the α_{LR}- and α_{RL}-type, due to the relation $\sigma_{\lambda\nu}\gamma_5 = \frac{1}{2}\, i\epsilon_{\lambda\mu\alpha\beta}\sigma^{\alpha\beta}$. The interactions (3) - (5) are time reversal invariant if the coupling constants are real.

The fields e, u, and d in Eqs. (3) - (5) are the mass eigenstates. The neutrino states $\nu_e^{(L)}$ and $\nu_e^{(R)}$ are in general linear combinations of the left-handed and the right-handed components of the neutrino mass-eigenstates ν_i:

$$\nu_e^{(L)} = \sum_i U_{ei}\nu_{iL} \quad , \tag{6}$$

$$\nu_e^{(R)} = \sum_i V_{ei}\nu_{iR} \quad , \tag{7}$$

where $\nu_{iL} = \frac{1}{2}(1-\gamma_5)\nu_i$, $\nu_{iR} = \frac{1}{2}(1+\gamma_5)\nu_i$; U_{ei} and V_{ei} are (in a basis where the charged leptons are diagonal) elements of the neutrino mixing matrix.

The constant a_{LL} in Eq. (3) contains the SM contribution, and can therefore be written as $a_{LL} = (a_{LL})_{SM} + a'_{LL}$, where $(a_{LL})_{SM} = g^2V_{ud}/8m_W^2$, and a'_{LL} represents new $V-A$ interactions. Let us consider the decay of the nucleon due to the interaction (2). Neglecting the induced form factors[6], the effective interaction describing $n \to pe^-\bar{\nu}_e$ is given by

$$H_\beta^{(N)} \simeq H_{V,A}^{(N)} + H_S^{(N)} + H_T^{(N)} \quad , \tag{8}$$

where

[6] For a review, see N. C. Mukhopadhyay, Ref. [7].

$$H^{(N)}_{V,A} = \bar{e}\gamma_\lambda(C_V + C'_V\gamma_5)\nu_e\bar{p}\gamma^\lambda n + \bar{e}\gamma_\lambda\gamma_5(C_A + C'_A\gamma_5)\nu_e\bar{p}\gamma^\lambda\gamma_5 n + \text{ H.c.} \quad , \tag{9}$$

$$H^{(N)}_S = \bar{e}(C_S + C'_S\gamma_5)\nu_e\bar{p}n + \text{H.c.} \tag{10}$$

$$H^{(N)}_T = \bar{e}\frac{\sigma_{\lambda\mu}}{\sqrt{2}}(C_T + C'_T\gamma_5)\nu_e\bar{p}\frac{\sigma_{\lambda\mu}}{\sqrt{2}}n + \text{H.c.} \tag{11}$$

Here we have to remember that the $C_i - C'_i$ components of the Hamiltonians (9) - (11) involve $\nu_e^{(L)}$, and the $C_i + C'_i$ components $\nu_e^{(R)}$. In Eqs. (9) - (11)

$$C_V = g_V(a_{LL} + a_{LR} + a_{RR} + a_{RL}) \quad , \tag{12}$$

$$C'_V = g_V(-a_{LL} - a_{LR} + a_{RR} + a_{RL}) \quad , \tag{13}$$

$$C_A = g_A(a_{LL} - a_{LR} + a_{RR} - a_{RL}) \quad , \tag{14}$$

$$C'_A = g_A(-a_{LL} + a_{LR} + a_{RR} - a_{RL}) \quad , \tag{15}$$

$$C_S = g_S(A_{LL} + A_{LR} + A_{RR} + A_{RL}) \quad , \tag{16}$$

$$C'_S = g_S(-A_{LL} - A_{LR} + A_{RR} + A_{RL}) \quad , \tag{17}$$

$$C_T = 2g_T(\alpha_{LL} + \alpha_{RR}) \quad , \tag{18}$$

$$C'_T = 2g_T(-\alpha_{LL} + \alpha_{RR}) \quad , \tag{19}$$

where the constants $g_V \equiv g_V(0)$, $g_S \equiv g_S(0)$ and $g_T \equiv g_T(0)$ are defined by

$$\langle p|\bar{u}\gamma_\lambda d|n\rangle = g_V(q^2)\bar{u}_p\gamma_\lambda u_n \quad , \tag{20}$$

$$\langle p|\bar{u}\gamma_\lambda\gamma_5 d|n\rangle = g_A(q^2)\bar{u}_p\gamma_\lambda\gamma_5 u_n \quad , \tag{21}$$

$$\langle p|\bar{u}d|n\rangle = g_S(q^2)\bar{u}_p u_n \quad , \tag{22}$$

$$\langle p|\bar{u}\sigma_{\lambda\mu}d|n\rangle = g_T(q^2)\bar{u}_p\sigma_{\lambda\mu}u_n \quad . \tag{23}$$

CVC predicts $g_V = 1$, and in the absence of new interactions the experimental value of g_A is $g_A = -1.2670 \pm 0.0035$ [11]. The constants g_S and g_T were calculated in Ref. [12] in connection with a study of neutral current interactions of a general Lorentz structure. Employing a quark model with spherically symmetric wave functions, g_S and g_T are given by $g_S = -\frac{1}{2} + \frac{9}{10}g_A \simeq 0.6$, $g_T = \frac{5}{3}(\frac{1}{2} + \frac{3}{10}g_A) \simeq 1.46$. The uncertainty in these predictions has been estimated to be about 30% to 60% [12]. Including an uncertainty of this size, one has

$$0.25 \lesssim g_S \lesssim 1 \ , \tag{24}$$

$$0.6 \lesssim g_T \lesssim 2.3 \ . \tag{25}$$

In the Hamiltonian (8) we neglected the contribution from H_P (the part of the Hamiltonian (4) involving the pseudoscalar quark current $\bar{u}\gamma_5 d$), since this vanishes in the nonrelativistic approximation for the nucleons. The interaction (8) is identical with the general beta decay interaction considered in Ref. [13]. The general formulas for observables in allowed beta decays can be found in this reference.

With the neutrinos (6) and (7) the observed beta decay probability is the sum of the probabilities of decays into the energetically allowed neutrino mass-eigenstates. In the following we shall assume that the neutrinos that can be produced in beta-decay are light enough that the effects of their masses on the decay probability can be neglected.[7] In particular, we shall neglect the terms arising from the interference between amplitudes involving neutrinos of different chirality. As it is easily seen, under the above assumption the effect of neutrino mixing can be taken into account by multiplying in observables the coupling constants a_{Lk}, A_{Lk} $(k = L, R)$ and α_{LL} by $\sqrt{u_e}$, and a_{Rk}, A_{Rk} $(k = L, R)$ and α_{RR} by $\sqrt{v_e}$, where

$$u_e = \sum_i{}' |U_{ei}|^2 \ , \tag{26}$$

$$v_e = \sum_i{}' |V_{ei}|^2 \ . \tag{27}$$

The prime on the summation in Eqs. (26) and (27) indicates that the sum extends only over the neutrinos that are light enough to be produced in beta decay.

The terms in the Hamiltonians (3) - (5) involving the right-handed neutrino state $\nu_e^{(R)}$ can manifest themselves in beta decay only if either the right-handed

[7] An exception could be the end of the $e^{\pm}$ spectrum , where neutrino masses as small as of the order of electronvolts can already be important. The neutrino mass effect at the end of the $e^{\pm}$ spectrum depends also on the structure of the interaction involved (see Ref. [14]). Admixtures of heavier neutrinos in $\nu_e^{(L)}$ can be probed through searches for kinks in the Kurie plot [15].

neutrinos are sufficiently light, or (for Majorana neutrinos) if there is mixing between the heavy right-handed neutrinos and the light ones. In the latter scenario the effects of the $\nu_e^{(R)}$-terms are expected to be suppressed by the light-heavy neutrino mixing angles, which should be small. Note that if all the neutrinos are light, we have $u_e = v_e = 1$, as a consequence of the unitarity of the neutrino mixing matrix.

When discussing constraints on beta decay interactions, we shall need to consider also muon decay, and the decays $\pi \to e\nu_e$ and $\pi \to \mu\nu_\mu$. As for beta decay, we shall assume also for these processes that the neutrinos that can be produced in them are light enough that in observables their masses can be neglected. Neutrino mixing in these processes can be then taken into account, as in beta decay, by multiplying in observables the coupling constants by square roots of sums (in muon decay also square roots of products of sums) analogous to (26) and (27). We note that under the above assumption the sums u_e and v_e in muon decay and in $\pi \to e\nu_e$ are equal to u_e and v_e in beta decay, and so are the sums u_μ and v_μ (defined as u_e and v_e, except for the replacements $U_{ei} \to U_{\mu i}$ and $V_{ei} \to V_{\mu i}$) in muon decay and in $\pi \to \mu\nu_\mu$.

3 New V,A Interactions: T-Invariant Contributions

3.1 Model Independent Considerations

The most general form of the Hamiltonian for $d \to ue^-\nu_e^{(L,R)}$ constructed from vector and axial-vector currents is given by Eq. (3). For given neutrino states $\nu_e^{(L)}$ and $\nu_e^{(R)}$ the Hamiltonian (3) contains 8 real parameters (four complex coupling constants). One of these is an overall phase, which does not enter the observables. We can choose therefore a_{LL} to be real and positive. Defining $\bar{a}_{ik} = a_{ik}/a_{LL}$ $(ik = LR, RR, RL)$, a set of the remaining six parameters is, for example, $|\bar{a}_{LR}|, |\bar{a}_{RR}|, |\bar{a}_{RL}|$, the phases $e^{i\varphi_L} = \bar{a}_{LR}/|\bar{a}_{LR}|$, $e^{i\varphi_R} = \bar{a}_{RL}\bar{a}^*_{RR}/|\bar{a}_{RR}||\bar{a}_{RL}|$ and $e^{i\varphi_{RL}} = \bar{a}^*_{RL}/|\bar{a}_{RL}|$. We shall not consider further the phase φ_{RL} since in observables it is involved only in terms proportional to neutrino mass. The reason is that such terms arise from the interference of amplitudes involving neutrinos of different chiralities.

The Hamiltonian (9) can be written in the form

$$H_\beta^{(N)} = (a_{LL})_{SM}(1+\bar{a}'_{LL})g_V(1+\bar{a}_{LR})[\bar{e}\gamma_\mu(1-\gamma_5)\nu_e^{(L)}\bar{p}\gamma^\mu(1-\lambda\gamma_5)n$$
$$+\bar{e}\gamma_\mu(1+\gamma_5)\nu_e^{(R)}\bar{p}\gamma^\mu(x+\gamma_5\lambda y)n] + \text{H.c.} \quad , \tag{28}$$

where $\bar{a}'_{LL} = a'_{LL}/(a_{LL})_{SM}$, and[8]

$$\lambda = \left(\frac{g_A}{g_V}\right)\frac{1-\bar{a}_{LR}}{1+\bar{a}_{LR}} \quad , \tag{29}$$

$$x = \frac{\bar{a}_{RR}+\bar{a}_{RL}}{1+\bar{a}_{LR}} \quad , \tag{30}$$

$$\lambda y = \left(\frac{g_A}{g_V}\right)\frac{\bar{a}_{RR}-\bar{a}_{RL}}{1+\bar{a}_{LR}} \quad , \tag{31}$$

As follows from Eq. (28), normalized observables (asymmetries and polarizations) can involve 5 parameters: $|\lambda|$, $|x|$, $|y|$, the phase of λ, and the relative phase of x and λy. The rate depends also on a_{LL}. As seen from Eq. (28), as long as the induced form factors are neglected (as we do here), the number of parameters at the nucleon level remains the same as at the quark level, since the only change is that $1+\bar{a}_{LR}$ is replaced at the nucleon level by $g_V(1+\bar{a}_{LR})$, and $(1-\bar{a}_{LR})$ and $(\bar{a}_{RR}-\bar{a}_{RL})$ get multiplied by g_A.

In the following we shall keep in λ, x, y and λy only the lowest order terms in the $\bar{a}_{ik}$'s. In this approximation we have

$$Re\lambda \simeq (g_A/g_V)(1-2Re\bar{a}_{LR}) \quad , \tag{32}$$

$$Im\lambda \simeq -2(g_A/g_V)Im\bar{a}_{LR} \simeq -2(Re\lambda)Im\bar{a}_{LR} \quad , \tag{33}$$

$$x \simeq \bar{a}_{RR}+\bar{a}_{RL} \quad , \tag{34}$$

$$y \simeq \bar{a}_{RR}-\bar{a}_{RL} \quad , \tag{35}$$

$$\lambda y \simeq (Re\lambda)(\bar{a}_{RR}-\bar{a}_{RL}) \quad , \tag{36}$$

$$Re x^*\lambda y \simeq (Re\lambda)(|\bar{a}_{RR}|^2-|\bar{a}_{RL}|^2) \quad , \tag{37}$$

$$Im x^*\lambda y \simeq -(Re\lambda)Im\bar{a}^*_{RR}\bar{a}_{RL} \quad . \tag{38}$$

Recall that if we allow $u_e \neq 1$, $v_e \neq 1$ (see Eqs. (26) and (27)), a_{LL} becomes

[8] Note that the parameters x and y, defined in Eqs. (30) and (31), are identical to the parameters x and y introduced in Ref. [16] only for $\bar{a}_{RL} = \bar{a}_{LR}$. Note also that in Ref. [16] the coupling constants (and therefore x and y) are assumed to be real.

in observables

$$a_{LL}^{(e)} \equiv a_{LL}\sqrt{u_e} \quad ; \tag{39}$$

the replacements for the $\bar{a}_{ik}$'s are

$$\begin{aligned} &\bar{a}_{LR} \to \bar{a}_{LR} \quad , \\ &\bar{a}_{RR} \to \bar{a}_{RR}^{(e)} \equiv \bar{a}_{RR}\sqrt{\tilde{v}_e} \quad , \\ &\bar{a}_{RL} \to \bar{a}_{RL}^{(e)} \equiv \bar{a}_{RL}\sqrt{\tilde{v}_e} \quad , \end{aligned} \tag{40}$$

where $\tilde{v}_e = v_e/u_e$.

As seen from Eq. (28), in beta decay information on $\bar{a}'_{LL}$ can be obtained only from the decay rate, through its effect on V_{ud}. If analogous LL-type V,A interactions exist for all the three families, and their coupling constants are the same in the weak eigenstate basis, a limit on $Re\bar{a}'_{LL}$ can be obtained using the unitarity of the 3-family Kobayashi-Maskawa matrix, assuming that there is no new contribution to muon decay, that $u_e = 1$ for all the processes involved and $u_\mu = 1$ in muon decay.[9] In the presence of the new $V - A$ interactions the measured ui-element $\bar{V}_{ui}$ $(i = d, s, b)$ of the Kobayashi-Maskawa matrix is related then to the true matrix elements V_{ui} as[10]

$$|\bar{V}_{ui}|^2 \simeq |V_{ui}|^2(1 + 2Re\bar{a}'_{LL}) \qquad (i = d, s, b) \quad . \tag{41}$$

Using the experimental values of $\bar{V}_{ud}$, $\bar{V}_{us}$, and $\bar{V}_{ub}$ recommended in Ref. [11], we find

$$\sum_i |\bar{V}_{ui}|^2 = 0.9959 \pm 0.0026 \quad . \tag{42}$$

It follows, using the unitarity relation for V_{ui} $(i = d, s, b)$, and the relation (41), that

$$-4 \times 10^{-3} < Re\bar{a}'_{LL} < 8 \times 10^{-5} \qquad (90\% \text{ c.l.}) \quad . \tag{43}$$

It should be noted that the value of $\bar{V}_{ud}$, and therefore the bound (43), has unknown uncertainties, mainly from charge symmetry breaking effects in the

[9] For a recent application, see Ref. [17].

[10] In the general case the right-hand side of Eq. (41) has to be multiplied by $(u_e)_i/u_e u_\mu$, where $(u_e)_i$ is the quantity (26) in the process which provides $\bar{V}_{ui}$.

nucleus [18]. A value of $\bar{V}_{ud}$ free of nuclear structure uncertainties is provided by the neutron lifetime τ_n and asymmetry parameter A_n, but at present this is not as accurate as the $\bar{V}_{ud}$ from the ft-values. This gives one of the motivations for improving the accuracy in the measurements of τ_n and A_n.[11] A further source of V_{ud} not affected by nuclear structure uncertainties is the beta decay of the pion $\pi^+ \to \pi^0 e^+ \nu_e$ [20]. A precision measurement of the rate of this decay is in progress at PSI [21].

An a'_{LL} interaction contributes also to the ratio $R_\pi \equiv \Gamma(\pi \to e\nu_e)/\Gamma(\pi \to \mu\nu_\mu)$. Since $|\bar{a}'_{LL}| < 1$, R_π is then given by (see Ref. [22])

$$R_\pi \simeq (R_\pi)_{SM} \frac{u_e}{u_\mu}(1 + 2Re\bar{a}'_{LL}) \quad , \tag{44}$$

where $(R_\pi)_{SM}$ is the SM value of R_π.

A lower bound on the quantity u_e/u_μ follows from a comparison of the predicted and the experimental mass of the W [22]. Assuming that there is no appreciable new contribution to muon decay, the predicted value $(m_W)_p$ of is given by

$$(m_W)_p = \left(\frac{\pi\alpha}{\sqrt{2}G_F}\right)^{1/2} \left[\sin^2\theta_W(1 - \Delta r)\right]^{-1/2} [u_e u_\mu]^{1/4} \quad , \tag{45}$$

where Δr represents radiative corrections [23]. Using in Eq. (45) $\sin^2\theta_W = 0.2230 \pm 0.0004$ and $\Delta r = 0.0354 \pm 0.0012$ [11], and identifying $(m_W)_p$ with the experimental value $(m_W)_{expt} = 80.419 \pm 0.056$ GeV [11], we find $[u_e u_\mu]^{1/4} > 0.998$ (90% c.l.). Since $u_e \leq 1$, $u_\mu \leq 1$, we obtain

$$0.992 < u_e/u_\mu < 1.008 \qquad \text{(90\% c.l.)}. \tag{46}$$

Using the experimental value $(R_\pi)_{expt} = 0.12303 \pm 0.0036$ [24], and $(R_\pi)_{SM} = (1.2352 \pm 0.0005)$ [25], we find

$$-9 \times 10^{-3} < Re\bar{a}'_{LL} < 5 \times 10^{-3} \quad \text{(90\% c.l.)} \; . \tag{47}$$

For $u_e = u_\mu = 1$ the bounds are

$$-5 \times 10^{-3} < Re\bar{a}'_{LL} < 4 \times 10^{-4} \; \text{(90\% c. l.)} \; . \tag{48}$$

[11] Note that the ft-values of $0^+ \to 0^+$ transitions are needed even then if one wants to search for right-handed currents (see Ref. [19]).

It should be noted that the bounds (48) would disappear, and the bounds (47) would become insignificant, if there is also a contribution from a new LL-type V, A interaction to $\pi \to \mu\nu_\mu$ with a coupling constant equal to a'_{LL}.

We note further that beta decay cannot provide precise information on $Re\bar{a}_{LR}$, because of the theoretical uncertainties in the calculations of g_A. However $Im\bar{a}_{LR}$ can be accessed through the T-odd D-correlation (see Section 6). If the right- and left-handed quark mixing matrices are equal, charged current universality implies, under analogous assumptions as the ones required for the limits on $Re\bar{a}'_{LL}$, limits on $Re\bar{a}_{LR}$. These are the same as the limits (43) for $Re\bar{a}'_{LL}$. Bounds on $Re\bar{a}_{LR}$ are provided also by R_π. These are the same as the bounds for $(-Re\bar{a}'_{LL})$ (see Eqs. (47) and (48)); the bounds in Eq. (48) disappear, and those in Eq. (47) become insignificant, if there is a contribution to $\pi \to \mu\nu_\mu$ from a muonic interaction of the same type and with a coupling constant equal to $Re\bar{a}_{LR}$.

It is easy to show that in pure Fermi and pure Gamow-Teller transitions normalized observables depend only on $|x|^2$ and $|y|^2$, respectively. In mixed transitions normalized observables can depend also on the other parameters. In observables g_V and g_A are multiplied by the Fermi matrix element M_F and Gamow-Teller matrix element M_{GT}, respectively.

Let us consider now $\bar{a}^{(e)}_{RR}$ and $\bar{a}^{(e)}_{RL}$. The best present limit on $\bar{a}^{(e)}_{RR}$ with $\bar{a}^{(e)}_{RL} = 0$, and on $\bar{a}^{(e)}_{RL}$ with $\bar{a}^{(e)}_{RR} = 0$ from experiments investigating nuclear and neutron beta decays is [12]

$$|\bar{a}^{(e)}_{RR}| < 6.3 \times 10^{-2} \qquad (90\% \text{ c.l.}) \ , \tag{49}$$

and

$$|\bar{a}^{(e)}_{RL}| < 3.7 \times 10^{-2} \qquad (90\% \text{ c.l.}) \ , \tag{50}$$

respectively. An experiment in preparation at CERN-ISOLDE [27], measuring the longitudinal polarizations of positrons emitted by a polarized nucleus in opposite directions with respect to the nuclear polarization, aims to improve the limits (49) and (50) to 3×10^{-2} [26]. Improved limits are expected also from the planned new generation neutron and nuclear beta decay experiments [28].

A bound on $\bar{a}^{(e)}_{Rk}$ ($k = R$ or L) follows from charged current universality.

---

[12] See Figs. 1 and 2 in Ref. [26] (where the limits on $\bar{a}^{(e)}_{RR}$ and $\bar{a}^{(e)}_{RL}$ are those given on δ and ζ). When $\bar{a}_{RR}$ and $\bar{a}_{RL}$ are relatively real, Figs. 1 and 2 give limits on $\bar{a}^{(e)}_{RR}$ and $\bar{a}^{(e)}_{RL}$ also when both are nonzero.

Assuming $u_\mu = 1$ and the absence of new muon decay interactions, and that $u_e = 1$ for all the pertinent processes, the relations between $\bar{V}_{ui}$ and V_{ui} ($i = d, s, b$) in the presence of an a_{Rk}-interaction are

$$
\begin{aligned}
|\bar{V}_{ud}|^2 &= |V_{ud}|^2(1 + |\bar{a}^{(e)}_{Rk}|^2) \qquad (k = R \text{ or } L) \ , \\
|\bar{V}_{uj}|^2 &= |V_{uj}|^2(1 + |\bar{c}^{(e)}_{Rk}|^2) \qquad (k = R \text{ or } L;\ j = s, b) \ ,
\end{aligned}
\tag{51}
$$

where $c^{(e)}_{Rk}$ depends on the coupling constants of Rk-type interactions involving the second and the third quark family, and on the mixing matrix of the right-handed quarks. Independently of the values of $c^{(e)}_{Rk}$, we find

$$
|\bar{a}^{(e)}_{Rk}| < 5.1 \times 10^{-2} \qquad (90\% \text{ c.l.}) \quad (k = R \text{ or } L) \ . \tag{52}
$$

The ratio R_π is given by (see Ref. [22])

$$
R_\pi = (R_\pi)_{SM}\frac{u_e}{u_\mu}(1 + |\bar{a}^{(e)}_{Rk}|^2) \qquad (k = R \text{ or } L) \ , \tag{53}
$$

yielding the bound

$$
|\bar{a}^{(e)}_{Rk}| < 8.9 \times 10^{-2} \qquad (90\% \text{ c.l.}) \quad (k = R \text{ or } L) \ . \tag{54}
$$

For $u_e = u_\mu = 1$ the bound is $|\bar{a}^{(e)}_{Rk}| < 5.4 \times 10^{-2}$ (90% c.l.).

R_π is not sensitive to an a_{Rk} interaction if there is an analogous contribution to $\pi \to \mu\nu_\mu$ with the same coupling constant.

If both $\bar{a}_{RR}$ and $\bar{a}_{RL}$ are nonzero, $|\bar{a}^{(e)}_{Rk}|$ in Eqs. (51-52) and (53-54) are replaced by $|\bar{a}^{(e)}_{RR} + \bar{a}^{(e)}_{RL}|$ and $|\bar{a}^{(e)}_{RR} - \bar{a}^{(e)}_{RL}|$, respectively. An observable, which vanishes when only one of $\bar{a}_{RR}$ and $\bar{a}_{RL}$ is nonzero, is the ratio P^F_L/P^{GT}_L of $e^\pm$-longitudinal polarizations in Fermi and Gamow-Teller decays. For V,A-interactions P^F_L/P^{GT}_L is given by

$$
\frac{P^F_L}{P^{GT}_L} \simeq 1 - 8Re\bar{a}^{(e)*}_{RR}\bar{a}^{(e)}_{RL} \ . \tag{55}
$$

Measurements of e^+-longitudinal polarizations yielded $P^F_L/P^{GT}_L = 1.0010 \pm 0.0027$ [29], which implies

$$
-6.8 \times 10^{-4} < Re\bar{a}^{(e)*}_{RR}\bar{a}^{(e)}_{RL} < 4.3 \times 10^{-4} \qquad (90\% \text{ c.l.}) \ . \tag{56}
$$

A stringent constraint on the interactions of light ($\lesssim$ 10 MeV) right-handed neutrinos comes from the observed neutrino pulse from supernova 1987A [30,31]. The observed $\bar{\nu}_e$-luminosity is consistent with the standard supernova model, and this implies severe constraints on possible new cooling mechanisms of the supernova core. The requirement that the process $e^- p \to \nu_e^{(R)} n$ does not carry away most of the energy that can be radiated by the supernova, leads for V,A interactions to the conclusion [31] that the coupling constants have to satisfy either the upper bound

$$\left(\frac{1}{6} g_V^2 |\bar{a}_{RR}^{(e)} + \bar{a}_{RL}^{(e)}|^2 + \frac{1}{2} g_A^2 |\bar{a}_{RR}^{(e)} - \bar{a}_{RL}^{(e)}|^2\right)^{1/2} \lesssim 1.2 \times 10^{-5} \ , \tag{57}$$

or the lower bound

$$\left(\frac{1}{6} g_V^2 |\bar{a}_{RR}^{(e)} + \bar{a}_{RL}^{(e)}|^2 + \frac{1}{2} g_A^2 |\bar{a}_{RR}^{(e)} - \bar{a}_{RL}^{(e)}|^2\right)^{1/2} \gtrsim 2 \times 10^{-2} \ . \tag{58}$$

For other types of $d \to ue^- \bar{\nu}_e^{(R)}$ interactions the constraints are probably similar. The bounds (57) and (58) could be evaded if $\nu_e^{(R)}$ has some additional interaction, which can trap it. A special interaction of this kind has been suggested in Ref. [32]. In the following while we shall bear in mind the bounds (57) and (58), we shall not invoke them in our discussions, since they do not diminish the importance of terrestrial experiments.

In conclusion, the experimental limit (50) is the best present model independent limit on $\bar{a}_{RL}^{(e)}$ (assuming $\bar{a}_{RR}^{(e)} = 0$). For $\bar{a}_{RR}^{(e)}$ the best limit is also the direct limit from beta decay (Eq. (49)), since (52) is not as rigorous.

New V,A-type $d \to ue^- \bar{\nu}_e$ interactions involving right-handed currents can arise at the tree-level from the exchange of new charged gauge bosons (as, for example, in left-right symmetric models), in models with new fermions which have right-handed couplings to the W and which mix with the known fermions, and from the exchange of leptoquarks. In all the above cases the resulting $d \to ue^- \bar{\nu}_e$ interactions can be represented by contact nonderivative four-fermion interactions. Such contact $d \to ue^- \bar{\nu}_e$ interactions can arise also in composite models, from the exchange of constituents.

New $V - A$ $d \to ue^- \bar{\nu}_e$ interactions are present, for example, in models with leptoquarks, and among contact interactions.

3.2 Left-Right Symmetric Models

Left-right symmetric models (L-R models) [33] are attractive extensions of the standard electroweak model, which provide a framework for the understanding of the origin of parity violation in the weak interaction. The simplest models are based on the gauge group $SU(2)_L \times SU(2)_R \times U(1)$. In the following we shall refer to $SU(2)_L \times SU(2)_R \times U(1)$ models as "L-R models".

The fermions in $SU(2)_L \times SU(2)_R \times U(1)$ models are assigned to representations of the group in a left-right symmetric manner. The left- (right-) handed fermions are in doublets of $SU(2)_L$ $[SU(2)_R]$ and singlets of $SU(2)_R$ $[SU(2)_L]$:

$$
\begin{aligned}
\begin{pmatrix} u' \\ d' \end{pmatrix}_L &: \qquad (T_L, T_R, Y) = \left(\frac{1}{2}, 0, \frac{1}{3}\right) \\
\begin{pmatrix} u' \\ d' \end{pmatrix}_R &: \qquad (T_L, T_R, Y) = \left(0, \frac{1}{2}, \frac{1}{3}\right) \\
\begin{pmatrix} \nu_e' \\ e' \end{pmatrix}_L &: \qquad (T_L, T_R, Y) = \left(\frac{1}{2}, 0, -1\right) \\
\begin{pmatrix} \nu_e' \\ e' \end{pmatrix}_R &: \qquad (T_L, T_R, Y) = \left(0, \frac{1}{2}, -1\right)
\end{aligned}
\tag{59}
$$

and the same assignments for the second and third family. $\vec{T}_L$, $\vec{T}_R$, and Y are the generators of $SU(2)_L$, $SU(2)_R$, and $U(1)$, respectively. The electric charges is given by $Q = T_{3L} + T_{3R} + \frac{1}{2}Y$.

In addition to the observed gauge bosons W and Z (called in L-R models W_1 and Z_1), the model contains a second charged gauge boson, W_2 (see Eq. (62)), and a second massive neutral gauge boson, Z_2. The model requires at least one Higgs field of the type $\phi(\frac{1}{2}, \frac{1}{2}, 0)$. In addition, at least one Higgs field of a different type must be introduced to break the gauge symmetry down to electromagnetic gauge invariance.

In Eq. (59) the primed fields are the gauge-group eigenstates. They are linear combinations of the mass-eigenstates. In terms of the mass eigenstates

$$Q^{(u)} = \begin{pmatrix} u \\ c \\ t \end{pmatrix}, \qquad Q^{(d)} = \begin{pmatrix} d \\ s \\ b \end{pmatrix}$$
$$E = \begin{pmatrix} e \\ \mu \\ \tau \end{pmatrix}, \qquad n = \begin{pmatrix} \nu_1 \\ \nu_2 \\ \vdots \end{pmatrix} \tag{60}$$

the coupling of the gauge group eigenstate charged gauge bosons W_L and W_R to the quarks and the leptons is given by

$$\mathcal{L} = (g_L/\sqrt{2})(\bar{Q}_L^{(u)}\gamma_\lambda V_L Q_L^{(d)} + \bar{n}_L\gamma_\lambda U^\dagger E_L)W_L \tag{61}$$
$$+ (g_R/\sqrt{2})(\bar{Q}_R^{(u)}\gamma_\lambda V_R Q_R^{(d)} + \bar{n}_R\gamma_\lambda V^\dagger E_R)W_R + \text{ H.c.} \ ,$$

where g_L and g_R are the $SU(2)_L$ and the $SU(2)_R$ gauge coupling constants, respectively; $\psi_L = \frac{1}{2}(1-\gamma_5)\psi$, $\psi_R = \frac{1}{2}(1+\gamma_5)\psi$ ($\psi = Q^{(u)}, Q^{(d)}, \ldots$). The number of neutrino mass-eigenstates is three in the case of Dirac neutrinos, and six if the neutrinos are Majorana fermions. The matrices U_L, U_R and U, V are the quark and leptonic mixing matrices, respectively. The fields W_L and W_R are linear combinations of the mass-eigenstates W_1 and W_2:

$$W_L = \cos\zeta W_1 + \sin\zeta W_2$$
$$W_R = e^{i\omega}(-\sin\zeta W_1 + \cos\zeta W_2) \ , \tag{62}$$

where ζ is a mixing angle and ω is a CP-violating phase.

If the right-handed mixing angles are equal to the left-handed ones and $g_R = g_L$, parity violation in the interaction (61) is due only to the difference in the masses of the W_2 and the W_1.

The Hamiltonian responsible for nuclear beta decay resulting from (61) is of the form (3) with [34,35]

$$a_{LL} \simeq (g_L^2 \cos\theta_1^L / 8m_1^2) \ , \tag{63}$$

$$\bar{a}_{RR} \simeq e^{i\alpha}(\cos\theta_1^R / \cos\theta_1^L)(g_R^2 m_1^2 / g_L^2 m_2^2) \ , \tag{64}$$

$$\bar{a}_{LR} \simeq -e^{i(\alpha+\omega)}(\cos\theta_1^R/\cos\theta_1^L)(g_R\zeta/g_L) \quad , \tag{65}$$

$$\bar{a}_{RL} \simeq -e^{i\omega}(g_R\zeta/g_L) \quad , \tag{66}$$

where m_1, m_2 are the masses of W_1, W_2; $\cos\theta_1^L = (U_L)_{ud}$ and $e^{i\alpha}\cos\theta_1^R = (U_R)_{ud}$. Note that for the phases φ_L and φ_R (see Section 3.1) one has the relation $\varphi_R = -\varphi_L(= -\alpha - \omega)$.

Let us consider $\bar{a}_{RR}$. Stringent limits on $g_R^2 m_1^2/g_L^2 m_2^2$ come from the $K_L - K_S$ mass difference Δm_K if one requires that each individual contribution to Δm_K, corresponding to box diagrams with a given pair of internal quarks, is smaller than the experimental value of Δm_K, and if one assumes some reasonable restrictions on fine-tuned cancellations [36]. For models with manifest left-right symmetry (where $g_R = g_L$, $\theta_i^R = \theta_i^L$ $(i = 1,2,3)$), and there are no CP-violating phases beyond the Kobayashi-Maskawa phases in the left-handed and right-handed sectors) [37,38] [13] and pseudomanifest left-right symmetry (where $g_R = g_L$, $\theta_i^R = \theta_i^L$, but the new CP-violating phases are allowed) [38] this limit is $\bar{a}_{RR} \simeq m_1^2/m_2^2 \lesssim 3.6 \times 10^{-3}$ [40,36,34,41]. For general L-R models (models with nonmanifest left-right symmetry) [38] the limit depends on the form of V_R. The conclusion of the analysis in Ref. [36] is that the weakest bound is obtained when V_R is a unit matrix. Then

$$\bar{a}_{RR} \simeq \frac{g_R^2 m_1^2}{g_L^2 m_2^2} \lesssim 7.5 \times 10^{-2} \quad . \tag{67}$$

Further constraints on $g_R^2 m_1^2/g_L^2 m_2^2$ can be deduced from the results of searches for new charged vector bosons in high-energy $\bar{p}p$ collisions at the Fermilab Tevatron [42]. These experiments set upper limits on the product $\sigma B_{e\nu}$ of the $p\bar{p} \to W_R X$ cross-section and the $W_R \to e\nu_e$ branching ratio for given values of m_2. The ratio of $\sigma B_{e\nu}$ and $(\sigma B_{e\nu})_{SM}$ ($\sigma B_{e\nu}$ evaluated with SM coupling constants) is given by [14]

$$\frac{\sigma B_{e\nu}}{(\sigma B_{e\nu})_{SM}} = \frac{g_R^2 \cos^2\theta_1^R}{g_L^2 \cos^2\theta_1^L} \quad , \tag{68}$$

which for $V_R = I$ is $\sigma B_{e\nu}/(\sigma B_{e\nu})_{SM} = g_R^2/g_L^2$. Thus from the experimental limits on $\sigma B_{e\nu}$ one can deduce limits on g_R/g_L. For g_R/g_L it can be shown [34] that the internal consistency of the model equires $g_R/g_L \geq 0.55$. Inspection of the experimental results shows that this holds from about $m_2 = 600$ GeV, and

[13] Investigations of beta decay in manifestly left-right symmetric models include Refs. [37], [16], and [39].

[14] The dependence of W_R-searches on the parameters of L-R models was considered in Ref. [43].

that $|\bar{a}^{(e)}_{RR}|$ ($\lesssim |\bar{a}_{RR}|$) can be as large as the present experimental limit (49) in the range [45]

$$800 \text{ GeV } \gtrsim m_2 < 1.8 \text{ TeV} \quad . \tag{69}$$

For $m_2 \geq 1.8$ TeV one would need to have $g_R \geq \sqrt{4\pi}$, for which the perturbative treatment of g_R could not hold. [15] The branching ratio would be $B_{e\nu}$ smaller, and therefore the constraints on $g_R^2 m_1^2/g_L^2 m_2^2$ weaker, if the W_R decays also to some new particles [43].

Stringent experimental bounds exist also on the mass of the Z_2. However for L-R models with an arbitrary Higgs sector the masses of the W_2 and the Z_2 are not related through known parameters. [16]

For $\bar{a}_{RL}$ it can be shown that $|\bar{a}_{RL}| \leq C|\bar{a}_{RR}|$, where C is a constant of the order of unity, except for Higgs representations with unreasonably high T_R [50, 36]. This does not mean, of course that $\bar{a}_{RL}$ can be neglected relative to $\bar{a}_{RR}$. A limit

$$|\bar{a}^{(e)}_{RL}| < 0.067 \qquad (90\% \text{ c.l.}) \tag{70}$$

on $\bar{a}^{(e)}_{RL}$ comes from the experimental value of the ρ-parameter in muon decay (Ref. [35]; see also Ref. [51]).

3.3 Exotic Fermions

The $d \to ue^-\bar{\nu}_e$ interaction can contain terms with right-handed currents even from W-exchange, if there are quarks and/or leptons whose right-handed components are in non-singlet representations of the $SU(2)$ component of the SM gauge group, and which mix with the usual quarks and leptons. Fermions with noncanonical $SU(2) \times U(1)$ assignments are referred to as "exotic". Such fermions occur in many extensions of the SM [52]. The new fermions, except the neutrinos, must be heavy (new charged and heavy neutral leptons heavier than about 90 GeV and 45 GeV, respectively, and new quarks heavier than about 200 GeV), as dictated by limits on direct production [11].

[15] A model where $g_R \neq g_L$ at the W_R scale has been constructed in Ref. [46]. In this model $g_R < g_L$. However, models with $g_R > g_L$ are also possible [47].
[16] See Ref. [48]. For L-R models with some specific choices for the Higgs bosons, data on neutral current interactions imply for $g_R = g_L$ lower bounds on both m_{Z_2} and m_2 of the order of (1-2) TeV [49]. The best lower bound on m_{Z_2} from direct searches for new neutral gauge bosons is 630 MeV (90% c.l.) [11], assuming that the Z_2 has no decays into final states involving some new particles.

We shall assume that the electric charge and color assignments of the new fermions are the standard ones, in which case the non-singlet fermions can only be in doublets [52].

In the presence of ordinary-exotic fermion mixing the coupling of the W to charged currents involving the usual quarks and charged leptons and the light neutrinos is given by [52]

$$\begin{aligned}\mathcal{L} = &(g\sqrt{2})(\bar{Q}_L^{(u)}\gamma_\lambda(A_L^{u\dagger}A_L^d)Q_L^d + \bar{n}_{\ell L}\gamma_\lambda(A_L^{\nu\dagger}A_L^e)E_L\\ &+ \bar{Q}_R^u\gamma_\lambda(F_R^{u\dagger}F_R^d)Q_R^{(d)}\\ &+ \bar{n}_{\ell R}^c\gamma_\lambda(F_R^{\nu\dagger}F_R^e)E_R)W^\lambda + \text{H.c.} \quad ,\end{aligned} \tag{71}$$

where $\bar{Q}^{(u)} \equiv (\bar{u}, \bar{c}, \bar{t})$, $\bar{Q}^{(d)} \equiv (\bar{d}, \bar{s}, \bar{b})$, $\bar{E} \equiv (e, \mu, \tau)$, $\bar{n}_{\ell L} \equiv (\nu_{1L}, \nu_{2L}, \ldots)$ and [17] $\bar{n}_{\ell R}^c \equiv (\bar{\nu}_{1R}^c, \bar{\nu}_{2R}^c, \ldots) = C(\bar{n}_{\ell L})^T$. In Eq. (71) the matrices A_L^k and F_R^k ($k = u, d, e, \nu$) relate, respectively, the ordinary and the exotic fermion weak eigenstates to the light fermion mass-eigenstates.

Ordinary-exotic fermion mixing induces generally flavor-changing neutral currents (FCNC) between the ordinary fermions. For FCNC transitions among the usual charged fermions there are stringent constraints on the strength of the corresponding interactions from limits on processes such as $\mu \to 3e$ or $K_L \to \mu\mu$. The mixing of ordinary fermions with exotic ones leads also to deviations from the SM predictions in flavor-conserving neutral current processes and in charged current processes. If one is interested in the latter effects, one can work in the limit where FCNC transitions are absent [52]. The matrices A_L^k and F_R^k ($k = u, d, e$) have then greatly simplified forms [52]. In such a framework the beta decay interaction resulting from the Lagrangian (71) is of the form (3), with [52,53,51]

$$a_{LL} \simeq (g^2 V_{ud}/8m_W^2) \quad , \tag{72}$$

$$\bar{a}_{LR} \simeq s_R^u s_R^d (\hat{V}_R)_{ud} \quad , \tag{73}$$

$$\bar{a}_{RL} \simeq e^{i\varphi_e} s_R^e \quad , \tag{74}$$

$$\bar{a}_{RR} \simeq \bar{a}_{RL}\bar{a}_{LR} \quad , \tag{75}$$

where $s_R^i \equiv \sin\theta_R^i$ ($i = u, d, e$), θ_R^i are light-heavy mixing angles, and $\hat{V}_R$

[17] We follow here the notation of Ref. [52], where all the right-handed neutrinos are denoted by n_R^c.

is a matrix defined in Ref. [52]. The quantities u_e and v_e are given here by $u_e = \sum_i |(A_L^\nu)_{ei}|^2$ and $v_e = \sum_i |(F_R^\nu)_{ei}|^2$ (denoted in Ref. [52] by $(c_L^{\nu_e})^2$ and $(s_R^{\nu_e})^2$, respectively). The CP-violating phase φ_e (which is the phase φ_{LR} here) has no detectable effect, as discussed in Sec. 3.1. A further CP-violating phase can reside in $(\hat{V}_R)_{ud}$.

Since $\bar{a}_{LR}$ is small ($|\bar{a}_{LR}| \lesssim 10^{-2}$; see Ref. [51] for a review of the bounds), it follows from the relation (75) that $\bar{a}_{RR}$ can be neglected relative to $\bar{a}_{RL}$.

Global analyses [18] of the constraints on ordinary-exotic fermion mixings yielded

$$|\bar{a}_{RL}^{(e)}| < 4.2 \times 10^{-2} \qquad (90\% \text{ c.l.}) \ . \tag{76}$$

The limit (76) originates mainly from muon decay data. Constraints on $\bar{a}_{RL}^{(e)}$ come also from high-energy neutrino-electron scattering, which constrains s_R^e [52]. Using the experimental and SM value of coupling constants for the $\nu_\mu - e$ scattering given in Ref. [11], we find (note that $\tilde{v}_e \lesssim 1$) $|a_{RL}^{(e)}| < 6 \times 10^{-2}$ (90% c.l.). Thus, the best present limit on $\bar{a}_{RL}^{(e)}$ in exotic fermion models is the beta decay limit (50) [45].

3.4 Leptoquark Exchange

Leptoquarks (LQs) are bosons which couple to lepton-quark pairs [56]. They appear in many theories that go beyond the SM, for example in grand unified theories [57], superstring inspired models [58], supersymmetric models with R-parity violation [59], and composite models [60]. LQs which do not induce proton decay could be light enough to cause observable effects in some low energy processes [61]. As we shall see, among the possible LQs coupled to the first fermion family some can give rise to new beta decay interactions.

Assuming that the LQ-fermion couplings are dimensionless, the spin of the LQs can be only zero or one.

The most general $SU(2)_L \times U(1) \times SU(3)_c$ invariant lepton number conserving (for Dirac neutrinos) and baryon number conserving Lagrangian for the couplings of spin-zero LQs to a SM family contains 9 LQ states, characterized by definite SM quantum numbers and a definite fermion number [62,63]. If a right-handed neutrino is added, as we do here, there is an additional LQ state, and two of the LQs can have an additional coupling [63]. Similarly, there are 9 spin-one LQ states [62] (10 if a right-handed neutrino is included [22]).

[18] See Refs. [52] and [54]. For a review, see Ref. [55].

The $d \to ue^-\bar{\nu}_e$ transition can be mediated by LQs of electric charge $Q = \frac{2}{3}$ (giving rise to $\nu_e^{(L,R)} \to u$ transitions) and $Q = \frac{1}{3}$ (inducing $\nu_e^{(L,R)} \to d^c$ transitions) [53].

Inspection shows that from the possible spin-one LQ states the ones that contribute to $d \to ue^-\bar{\nu}_e$ are (in the notation of Ref. [62]) are the $Q = \frac{2}{3}$ states U_1 and $(U_3)_0$, and the $Q = \frac{1}{3}$ states $(V_2)_-$ and $(\tilde{V})_2)_+$ [22,51], where the second subscript $(0, \pm)$ represents the value of T_z ($T_z = 0, \pm\frac{1}{2}$). The couplings of these LQs to the first SM family, extended by a right handed neutrino, are

$$\mathcal{L}_{U_1} = \left\{ \frac{1}{2} h_{1L} [\bar{u}\gamma_\mu(1-\gamma_5)\nu_e^{(L)} + \bar{d}\gamma_\mu(1-\gamma_5)e] \right. \\ \left. + \frac{1}{2} h_{1R} [\bar{d}\gamma_\mu(1+\gamma_5)e + \frac{1}{2} h_{1R}^{(\nu)} \bar{u}\gamma_\mu(1-\gamma_5)\nu_e^{(R)} \right\} U_1^\mu + \text{ H.c.} \ , \quad (77)$$

$$\mathcal{L}_{(U_3)_0} = \frac{1}{2} h_{3L} [\bar{u}\gamma_\mu(1-\gamma_5)\nu_e^{(L)} - \bar{d}\gamma_\mu(1-\gamma_5)e](U_3)_0^\mu + \text{ H.c.} \ , \quad (78)$$

$$\mathcal{L}_{(V_2)_-} = \left[\frac{1}{2} g_{2L} \bar{d}^c \gamma_\mu(1-\gamma_5)\nu_e^{(L)} + \frac{1}{2} g_{2R} \bar{u}^c \gamma_\mu(1+\gamma_5)e \right] (V_2)_-^\mu + \text{ H.c.} \ , \quad (79)$$

$$\mathcal{L}_{(\tilde{V}_2)_+} = - \left[\frac{1}{2} \tilde{g}_{2L} \bar{u}^c \gamma_\mu(1-\gamma_5)e + \frac{1}{2} \tilde{g}_{2R} \bar{d}^c \gamma_\mu(1+\gamma_5)\nu_e^{(R)} \right] (\tilde{V}_2)_+^\mu + \text{ H.c.} \quad (80)$$

From the possible spin-zero LQ states the ones which contribute to $d \to ue^-\bar{\nu}_e$ are the $Q = \frac{1}{3}$ states S_1 and $(S_3)_0$, and the $Q = \frac{2}{3}$ states $(R_2)_-$ and $(\tilde{R}_2)_+$ [22,51]. Their couplings are given by

$$\mathcal{L}_{S_1} = \left[\frac{1}{2} g_{1L} (\bar{u}^c(1-\gamma_5)e - \bar{d}^{c'}(1-\gamma_5)\nu_e^{(L)}) \right. \\ \left. + \frac{1}{2} g_{1R} \bar{u}^c(1+\gamma_5)e + \frac{1}{2} g_{1R}^{(\nu)} \bar{d}^c(1+\gamma_5)\nu_e^{(R)} \right] S_1 + \text{ H.c.} \ , \quad (81)$$

$$\mathcal{L}_{(S_3)_0} = -\frac{1}{2} g_{3L} [\bar{u}^c(1-\gamma_5)e + \bar{d}^c(1-\gamma_5)\nu_e^{(L)}](S_3)_0 + \text{ H.c.} \ , \quad (82)$$

$$\mathcal{L}_{(R_2)_-} = \left[\frac{1}{2} h_{2L} \bar{u}(1-\gamma_5)\nu_e^{(L)} - \frac{1}{2} h_{2R} \bar{d}(1+\gamma_5)e \right] (R_2)_- + \text{ H.c.} \ , \quad (83)$$

$$\mathcal{L}_{(\tilde{R}_2)_+} = - \left[\frac{1}{2} \tilde{h}_{2L} \bar{d}(1-\gamma_5)e + \frac{1}{2} \tilde{h}_{2R} \bar{u}(1+\gamma_5)\nu_e^{(R)} \right] (\tilde{R}_2)_+ + \text{ H.c.} \quad (84)$$

The fields in Eqs. (77-84) are the gauge group eigenstates (to simplify writing, we have omitted the primes on them).

In this section we can restrict attention to chiral LQs (LQs that couple to left-handed or to right-handed quarks, but not to both), since only such LQs (or the chiral parts of nonchiral LQs) can generate V,A-type beta decay interactions. The U_1 and the S_1 can be replaced by two-independent chiral LQs, the U_{1L}, U_{1R}, and S_{1L}, S_{1R}, where the second subscript indicates the chirality of the quark in the couplings. The R_2 ($\tilde{R}_2$) and V_2 ($\tilde{V}_2$) can be replaced by R_{2L}, R_{2R} ($\tilde{R}_{2L}, \tilde{R}_{2R}$) and V_{2L}, V_{2R} ($\tilde{V}_{2L}, \tilde{V}_{2R}$). The $(U_3)_0$ and $(S_3)_0$ are chiral.

The beta decay interactions from U_{1L}, S_{1L}, $(U_3)_0$ and $(S_3)_0$ are of V-A form. U_{1R} and S_{1R} gives rise to a_{RR} interactions. a_{LR}- and a_{RL}-interactions can be induced only through LQ mixing [51] [19] To avoid further constraints (see Ref. [64]), we shall assume that the right-handed quarks are to a good approximation mass eigenstates. The exchange of U_{1R} yields then (see Eq. (77))

$$\bar{a}_{RR} = \frac{h_{1R}^{*} h_{1R}^{(\nu)}}{4M_{1R}^2} \left(\frac{\sqrt{2}}{G_F V_{ud}}\right) \quad , \tag{85}$$

where M_{1R} is the mass of U_{1R}. We shall write (85) in the form

$$\bar{a}_{RR} = \gamma_V \frac{h_{1R}^{(\nu)}}{h_{1R}} \quad , \tag{86}$$

where $\gamma_V = \frac{1}{4}(|h_{1R}|^2/M_{1R}^2)(\sqrt{2}/G_F V_{ud})$. A stringent limit on γ_V comes from a recent measurement of parity violation in ^{133}Cs atoms [65]. The h_{1R}-coupling of the U_{1R} (see Eq. (77)) gives a contribution

$$(Q_W)_U \simeq 411.5\gamma_V \tag{87}$$

to the weak charge Q_W of ^{133}Cs. Identifying this contribution with the difference [20] $\Delta Q_W = (Q_W)_{expt} - (Q_W)_{SM} = 0.44 \pm 0.44$ between the experimental and the SM values of Q_W, we find [45]

$$\gamma_V < 2.6 \times 10^{-3} \qquad\qquad (90\%\ \text{c. l.}) \ . \tag{88}$$

Further constraints come from collider experiments: from searches for LQ production at the Fermilab Tevatron [68], HERA [69,70], at LEP [71], and (for high-mass LQs) from searches for neutral-current [72] and charged current contact interactions [73,74].

[19] We shall consider a_{LR}-interactions induced by LQs in Section 6, in connection with the T-odd D-coefficient. A discussion of other effects of LQ mixing will be included in Ref. [45].

[20] The experimental value of the weak charge of $^{133}C_s$ is $(Q_W)_{expt} = -72.65(28)_{expt}(34)_{theor}$ [66]. This result was verified by W. R. Johnson, and independently by V. A. Dzuba [67]. The SM value is $(Q_W)_{SM} = -73.09 \pm 0.03$ [11].

The Tevatron results set a lower bound of 150 GeV on the mass of U_{1R} (assuming the presence of the $h_{1R}^{(\nu)}$ coupling). Analysis [45] of the available data shows that $\bar{a}_{RR}$ from U_{1R}-exchange can be as large as the present experimental limit (Eq. (49)) for M_{1R} in the range

$$210 \text{ GeV} < M_{1R} < 510 \text{ GeV} \quad ; \tag{89}$$

$\bar{a}_{RR} = 10^{-2}$ (for example) is possible for 150 GeV $< M_{1R} <$ 3.2 TeV. These conclusions are set by the Tevatron results [68] and the results of H1 at HERA [69] (the other constraints from collider results are either weaker or not relevant), and by the condition $|h_{1R}^{(\nu)}| < \sqrt{4\pi}$, which we impose, so that it would be possible for the perturbative treatment of $h_{1R}^{(\nu)}$ to be adequate. Note that a large $\bar{a}_{RR}$ requires a very large ratio $h_{1R}^{(\nu)}/h_{1R}$.

A similar analysis [45] shows that $\bar{a}_{RR}$ due to S_{1R}-exchange can be at the present experimental limit for S_{1R} masses in the range

$$79 \text{ GeV} < M_{1R} < 260 \text{ GeV} \quad ; \tag{90}$$

$\bar{a}_{RR} = 10^{-2}$ is possible for 79 GeV $< M_{1R} <$ 1.6 TeV.

3.5 Contact Interactions

Physics characterized by a mass-scale $\Lambda > G_F^{-1/2}$, where $G_F^{-1/2}(\simeq 300 \text{ GeV})$ is the Fermi scale, can be described up to energies of order Λ by nonrenormalizable contact interactions invariant under the SM gauge group [75,76]. The lowest dimension contact terms relevant to beta decay are dimension 6 four-fermion interactions. Such contact interactions arise in composite models (where the quarks and leptons are bound states of more fundamental particles), induced by the exchange of constituents [75]. They can provide also a description of processes mediated by the exchange of heavy bosons.

We shall consider here, and in Sections 4.2 and 5.2, the contribution of contact interactions to beta decay [45,77]. The most general $SU(2)_L \times U(1) \times SU(3)_c$ invariant contact four-fermion interaction relevant to charged-current processes in the first family is given by [78]

$$\begin{aligned} \mathcal{L}_L = {} & \alpha_L \bar{L}\gamma_\mu \tau^a L \bar{Q}\gamma^\mu \tau^a Q + \beta_L^{(d)} \bar{L} e_R \bar{d}_R Q \\ & + \beta_L^{(u)} \bar{L} e_R \bar{Q} u_R + \gamma_L \bar{L}\sigma_{\mu\nu} e_R \bar{Q}\sigma_{\mu\nu} u_R + \text{H. c.} \quad , \end{aligned} \tag{91}$$

As in Eqs. (77-84), the fields in Eq. (91) are the gauge group eigenstates. $L = (\nu_{eL}, e_L)$, $Q = (u_L, d_L)$ are the left-handed lepton and quark doublets, and the $\tau^{a\prime}$s are the Pauli matrices. With the subscript L on the coupling constants we indicate the helicity of the neutrino involved.

In composite models the coupling constants $\alpha_L, \ldots \beta_L^{(d)}, \ldots$ are written customarily as $\epsilon g^2/\Lambda_i^{(\epsilon)}$, where $\epsilon = \pm 1, i = \alpha_L, \ldots \beta_L^{(d)}, \ldots$; g is a strong coupling constant taken to be $\sqrt{4\pi}$, and $\Lambda_i^{(\epsilon)}$ is the compositness scale associated with the corresponding operator. For contact terms originating from heavy boson (B) exchange the coupling constants are proportional to g_B^2/m_B^2.

The first term in (91) yields a V-A $d \to ue^-\bar{\nu}_e$ interaction; the other contain scalar-, pseudoscalar-, and tensor-type $d \to ue^-\bar{\nu}_e$ couplings, which we shall consider in Sections 4.2 and 5.2.

The coupling constant of the V-A interaction is given by (assuming that the same contact interaction is present for all the three families)

$$\bar{a}'_{LL} = \frac{\alpha_L}{\sqrt{2}G_F} \,. \tag{92}$$

In addition to the bounds (43) and (47) from charged current universality and R_π, the $\bar{a}'_{LL}$ in Eq. (92) is constrained by experimental results on the $e^\pm p \to \overset{(-)}{\nu}_e X$ reaction, obtained by the H1 and ZEUS collaborations at HERA [73,74]. An analysis [79,80] of these data yielded $|\Lambda^\pm| \gtrsim (2-5)$ TeV for $\Lambda^\pm$ associated with the α_L- term. This implies

$$|\bar{a}'_{LL}| \lesssim 0.2 \text{ to } 0.03 \,, \tag{93}$$

which is weaker than the bounds (43) and (47) from charged current universality and R_π.

The α_L-term in the Lagrangian (91) contains also a neutral current electron-quark interaction. Inspection shows that the limit on $\bar{a}'_{LL}$ from atomic parity violation is weaker than the bounds (43) and (47) [21], and so is the limit from other neutral current data [72].

If we include among the fermions a right-handed neutrino, four additional

[21] We note that for the α_L-term Q_W is suppressed due to the isovector character of the interaction.

contact terms are possible [45]:

$$\mathcal{L}_R = \alpha_R \bar{e}_R \gamma_\mu \nu_{eR} \bar{u}_R \gamma^\mu d_R + \beta_R^{(d)} \bar{L} \nu_{eR} \bar{Q} d_R$$
$$+ \beta_R^{(u)} \bar{L} \nu_{eR} \bar{u}_R Q + \gamma_R \bar{L} \sigma_{\mu\nu} \nu_{e_R} \bar{Q} \sigma^{\mu\nu} d_R + \text{H. c.} \tag{94}$$

As in the Lagrangian (91), the fields in Eq. (94) are the gauge group eigenstates.

The first term in (94) is an a_{RR}- interaction with

$$\bar{a}_{RR} = \frac{\alpha_R}{4} \left(\frac{\sqrt{2}}{G_F V_{ud}} \right) \tag{95}$$

(assuming that mixing in the right-handed sector can be neglected); the remaining terms are, again, of scalar-, pseudoscalar-, and tensor structure.

Note that a_{LR} and a_{RL} type interactions cannot appear in (91) and (94), since these violate $SU(2)_L \times U(1)$ invariance.

The data on the $e^\pm p \to \overset{(-)}{\nu}_e X$ reaction constrain also the a_{RR}-interaction. This term was not included in the analysis in Refs. [79] and [80]. We expect on $\bar{a}_{RR}$ a limit which is weaker than the limit (93) since, neglecting terms proportional to the neutrino mass, there is no interference in the $e^\pm p \to \overset{(-)}{\nu}_e X$ cross-section between the a_{RR}-term and the SM contribution.

4 Scalar Interactions: T-Invariant Contributions

4.1 Model Independent Considerations

The most general scalar and pseudoscalar $d \to u e^- \bar{\nu}_e$ interaction is given in Eq. (4). The S-type and P-type components of (4) can be written as

$$H_S = \left[a_{LS} \bar{e}(1-\gamma_5)\nu_e^{(L)} + a_{RS} \bar{e}(1+\gamma_5)\nu_e^{(R)} \right] \bar{u} d + \text{H. c.} \ , \tag{96}$$

$$H_P = \left[a_{LP} \bar{e}(1-\gamma_5)\nu_e^{(L)} + a_{RP} \bar{e}(1+\gamma_5)\nu_e^{(R)} \right] \bar{u} \gamma_5 d + \text{H. c.} \ , \tag{97}$$

where

$$a_{LS} = A_{LL} + A_{LR} \ , \tag{98}$$

$$a_{RS} = A_{RR} + A_{RL} \quad , \tag{99}$$

$$a_{LP} = -A_{LL} + A_{LR} \quad , \tag{100}$$

$$a_{RP} = A_{RR} - A_{RL} \quad . \tag{101}$$

As mentioned in Section 2, for beta decay only H_S is important. For $n \to pe^-\bar{\nu}_e$ it yields the interaction (10) with

$$C_S - C'_S = 2g_S a_{LS} \quad , \tag{102}$$

$$C_S + C'_S = 2g_S a_{RS} \quad , \tag{103}$$

where g_S is defined in Eq. (22).

Scalar-type beta decay interactions would show up in the allowed approximation only in Fermi (or mixed) transitions. T-invariant contributions to observables can depend on $g_S Re\bar{a}_{LS}$ (through interference with the SM contribution, $g_S^2(|\bar{a}_{LS}|^2 + |\bar{a}_{RS}|^2)$ $[= \frac{1}{2}(|C_S|^2 + |C'_S|^2)\sqrt{2}/G_F V_{ud}]$, and $g_S^2(|\bar{a}_{LS}|^2 - |\bar{a}_{RS}|^2)$ $[= -ReC_S C'^*_S)\sqrt{2}/G_F V_{ud}]$, where $\bar{a}_{kS} = a_{kS}\sqrt{2}/G_F V_{ud}$ $(k = L, R)$. The present experimental limit on $g_S Re\bar{a}_{LS}$ is

$$|g_S Re\bar{a}_{LS}| < 10^{-2} \qquad (90\%\ \text{c.l.}) \quad , \tag{104}$$

obtained from the ratio P_L^F/P_L^{GT} of positron polarizations in a Fermi and a Gamow-Teller transition [82].[22] For $g_S\bar{a}_{RS}$ the experimental limit (deduced from a measurement of $e - \nu$ correlation in ^{32}Ar β decay) is at the 0.1 level [83]. New experiments sensitive to scalar interactions are under preparation [84].

Stringent constraints on H_P, which we shall use later on, come from the ratio $R_\pi \equiv \Gamma(\pi \to e\nu_e)/\Gamma(\pi \to \mu\nu_\mu)$. The Hamiltonian H_P gives a contribution to R_π given by [22]

$$R_\pi = (R_\pi)_{SM} \frac{u_e}{u_\mu}(|1 + \omega_e \bar{a}_{LP}|^2 + \omega_e^2 |\bar{a}_{RP}^{(e)}|^2) \quad , \tag{105}$$

where[23] $\omega_e = m_\pi^2/m_e(m_u + m_d) \simeq 2.65 \times 10^3$, and $\bar{a}_{RP}^{(e)} = \bar{a}_{RP}\sqrt{\bar{v}_e}$. We have not included in (105) the contribution $\omega_e^2|Im\bar{a}_{LP}|^2$ from $Im\bar{a}_{LP}$, which is in-

[22] A limit on $g_S Re\bar{a}_{LS}$ comparable to (104) follows from ft-values of Fermi beta decays, which are modified in the presence of $Re\bar{a}_{LS}$ by the Fierz interference term [81]. However, because of nuclear structure uncertainties, this limit is not as reliable as (104).

[23] We used $m_u = 5.1$ MeV and $m_d = 9.3$ MeV [85].

dependently known to be small (see Section 6). Taking $(R_\pi)_{expt}/(R_\pi)_{SM}$ at 90% c. l., and using for u_e/u_μ the range (46), the bounds on $Re\bar{a}_{LP}$ and $a_{RP}^{(e)}$ are [22]

$$Re\bar{a}_{LP} \simeq -7.5 \times 10^{-4}$$

or (106)

$$|Re\bar{a}_{LP}| < 3.2 \times 10^{-6}$$

if Rea_{LP} alone contributes,

$$|\bar{a}_{RP}^{(e)}| < 3.4 \times 10^{-5} \tag{107}$$

if only $a_{RP}^{(e)}$ contributes, and

$$|Re\bar{a}_{LP}| < 7.5 \times 10^{-4} \quad , \qquad |\bar{a}_{RP}^{(e)}| < 4.0 \times 10^{-4} \tag{108}$$

if Rea_{LP} and $a_{RP}^{(e)}$ contribute simultaneously.

For $u_e/u_\mu = 1$ the bound in the second equation in (106) and the bound in Eq. (107) are more stringent: $|Re\bar{a}_{LP}| < 1.7 \times 10^{-6}$ and $|\bar{a}_{RP}^{(e)}| < 2.1 \times 10^{-5}$, respectively.

R_π is not sensitive to $\bar{a}_{LP}$ and $\bar{a}_{RP}^{(e)}$ if there are analogous contributions $\bar{b}_{LP}$ and $\bar{b}_{RP}^{(\mu)} = \bar{b}_{RP}\sqrt{\tilde{v}_\mu}$ ($\tilde{v}_\mu = v_\mu/u_\mu$) to $\pi \to \mu\nu_\mu$ for which $\bar{a}_{LP}/\bar{b}_{LP} = m_e/m_\mu$ and $\bar{a}_{RP}^{(e)}/\bar{b}_{RP}^{(\mu)} = m_e/m_\mu$ [86]. For other cases the constraints from $(R_\pi)_{expt}$ could still be weakened by accidental cancellations between the electronic and the muonic terms. In the following, while keeping this possibility in mind, we shall ignore it in our considerations.

Scalar interactions can arise at the tree level from the exchange of Higgs bosons, spin-zero or spin-one leptoquarks, and in supersymmetric models with R-parity violation from the exchange of sleptons. They can appear also in composite models, in the form of contact interactions.

4.2 Contact Interactions

S-type $d \to ue^-\bar{\nu}_e$ contact interactions are contained in the $\beta_k^{(i)}$ terms ($i = d, u$; $k = L, R$) in Eqs. (91) and (94). This interaction is of the same form as the general Hamiltonian (96). The coupling constants a_{LS} and a_{RS} are given by [77]

$$a_{LS} = \frac{1}{4}\left(\beta_L^{(d)} - \beta_L^{(u)}\right)^* \quad , \tag{109}$$

$$a_{RS} = \frac{1}{4}\left(\beta_R^{(d)} + \beta_R^{(u)}\right) \quad . \tag{110}$$

The coupling constants of the P-type terms (contained in the same contact terms) are

$$a_{LP} = \frac{1}{4}\left(\beta_L^{(d)} + \beta_L^{(u)}\right)^* \quad , \tag{111}$$

and

$$a_{RP} = \frac{1}{4}\left(\beta_R^{(d)} - \beta_R^{(u)}\right) \quad . \tag{112}$$

The constraints from $(R_\pi)_{expt}$ require that the values of $\beta_L^{(d)}$ and $\left(-\beta_L^{(u)}\right)$, and of $\beta_R^{(d)}$ and $\beta_R^{(u)}$ be extremely close.

In addition to limits from beta decay, bounds on a_{LS} and $a_{RS}^{(e)}$ follow also from the analysis [79] of data on the $e^\pm p \to \overset{(-)}{\nu}_e X$ reaction. For $\beta_L^{(d)}$ and $\beta_L^{(u)}$ the lower bounds on the corresponding Λ's are at the 1 TeV level. The bounds on $\beta_R^{(d)}$ and $\beta_R^{(u)}$ (which were not included in the analysis) should be the same. It follows that the limits from $e^\pm p \to \overset{(-)}{\nu}_e X$ are $|\bar{a}_{LS}| \lesssim 0.4$ and $|\bar{a}_{RS}^{(e)}| \lesssim 0.4$. These are weaker than the limits from beta decay experiments.

4.3 Leptoquark Exchange

A possible source of S-type beta decay interactions is the exchange of spin-zero or spin-one non-chiral leptoquarks. As follows from Eqs. (77-84), all the LQ states that can contribute to beta decay, except $(S_3)_0$ and $(U_3)_0$ (which are chiral), can give rise to S-type interactions. In all cases the S-type interaction is accompanied by a P-type one with $|a_{kP}| = |a_{kS}|$ ($k = L, R$) [53]. Thus the scalar interactions from LQ exchange are constrained by the bounds (106-108): Eq. (106) for the interaction from $(R_2)_-$ and $(V_2)_-$, Eq. (107) for the one from $(\tilde{R}_2)_+$ and $(\tilde{V}_2)_+$, and the bounds in Eq. (108) for the interaction from S_1 and U_1 [22] [24]. The bounds from additional constraints on the LQ parameters are

[24] S_1- and U_1-exchange gives rise also to an a'_{LL}-interaction, which changes in the first term in Eq. (105) the quantity $|1 + \omega_e \bar{a}_{LP}|$ to $|1 + \bar{a}'_{LL} + \omega_e \bar{a}_{LP}|$. However,

weaker.

4.4 Supersymmetric Models with R-Parity Violation

In supersymmetric models where the superpotential can contain gauge invariant renormalizable supersymmetric interactions that violate the conservation of R-parity ($R = (-1)^{3B+L+2s}$, where B and L are the baryon and lepton number, respectively, and s is the spin of the particle; thus $R = +1$ for the particles of the SM, and $R = -1$ for their superpartners), there are contributions to processes with the usual particles from the exchange of single sfermions. In this section we shall consider $d \to ue^-\bar{\nu}_e$ interactions of this kind in the framework of the minimal supersymmetric standard model (MSSM) [87,88].

In the MSSM, unlike in the SM, the conservation of lepton number (L) and of baryon number (B) is not automatic: the superpotential can contain renormalizable and gauge invariant L- and B-violating terms. The general forms of these are

$$W_{\not{L}} = \frac{1}{2}\lambda_{ijk}L_iL_jE_k^c + \lambda'_{ijk}L_iQ_jD_k^c + \mu_iL_iH_u \tag{113}$$

$$W_{\not{B}} = \frac{1}{2}\lambda''_{ijk}U_i^cD_j^cD_k^c \quad , \tag{114}$$

where $i, j, k = 1, 2, 3$ are family indices, and summations over i, j, k are implied; L_i, Q_i (E_i^c, U_i^c, D_i^c) are the $SU(2)$-doublet (singlet) lepton and quark superfields. The constants λ_{ijk} are antisymmetric under the interchange $i \leftrightarrow j$, and λ''_{ijk} is antisymmetric under $j \leftrightarrow k$.

The couplings in $W_{\not{L}}$ and $W_{\not{B}}$ violate invariance under R-parity. If both the λ'_{ijk} and the λ''_{ijk} terms are present, some of the products $\lambda'_{ijk}\,\lambda''_{\ell mn}$ would have to be extremely small to prevent too rapid proton decay. One way to deal with this problem is to postulate R-parity invariance. This would eliminate both $W_{\not{L}}$ and $W_{\not{B}}$. Another possibility is that B is conserved, but the L-violating terms are present. This scenario is obtained by demanding invariance under "baryon parity" (under baryon parity Q_i, U_i^c, and D_i^c are odd, and the remaining superfields are even). The model we shall consider in the following is the R-parity violating MSSM ($\not{R}$MSSM), defined as the MSSM with $W_{\not{L}}$ included in the superpotential [88].

The first two terms in Eq. (113) give rise to new contributions to the $d \to ue^-\bar{\nu}_e$ transition. There are two types of contributions. The first, which originate

since $|\bar{a}'_{LL}| < 1$, it has a negligible effect on the first bound in Eq. (108). The second bound in Eq. (108) is not affected at all.

from the λ'-couplings alone, are $V - A$ interactions, proportional to $|\lambda'_{11k}|^2$ ($k = 1,2,3$) [89]. The second type depend on both λ_{ijk} and λ'_{ijk}. These are given by [22]

$$H_j = \frac{\lambda^*_{1j1}\lambda'_{j11}}{4m^2_{\tilde{e}_{jL}}}\bar{e}(1-\gamma_5)\nu_e\bar{u}(1+\gamma_5)d + \text{H. c.} \tag{115}$$

($j = 2,3$). The Hamiltonians in Eq. (115) have both an a_{LS}-type and an a_{LP}-type component, with $|a_{LS}| = |a_{LP}|$. It follows that the a_{LS}-interactions are constrained by R_π. From Eq. (106) we have (using $|g_S| \le 1$)

$$g_S Re\bar{a}_{LS} \simeq -7.5 \times 10^{-4}$$

or (116)

$$|g_S Re\bar{a}_{LS}| < 3.2 \times 10^{-6} \ .$$

4.5 Higgs Exchange

A possible source of scalar-type $d \to ue^-\bar{\nu}_e$ interactions is the exchange of charged Higgs bosons[25]. Charged Higgs bosons appear in many extensions of the SM. The simplest case is the standard $SU(2) \times U(1)$ model with a Higgs sector containing two Higgs doublets. We shall consider the version of the two-Higgs doublet model, where flavor-changing neutral currents are eliminated at the tree level by demanding that only one Higgs doublet couples to the same right-handed field [91]. The couplings of the charged Higgs boson H^+ to the fermions in such a model are proportional to the mass of the right-handed fermion involved in the couplings. In the minimal supersymmetric standard model the presence of two Higgs doublets coupled this way is required [87].

The exchange of H^+ gives rise to Hamiltonians of the forms (96) and (97) with

$$\bar{a}_{LS} = -\bar{a}_{LP} \simeq \frac{m_d m_e}{m_H^2}\tan^2\beta \tag{117}$$

($\bar{a}_{RS} = 0$, since the model does not contain right-handed neutrinos). In (117) we neglected the contibution from the $\bar{u}_R d_L$ coupling since for large $\tan\beta$ it is suppressed by the factor $\cot\beta/\tan\beta$ [91]. In Eq. (117) m_H is the mass of the H^+, and $\tan\beta = v_u/v_d$, where v_u and v_d are the vacuum expectation values of the Higgs fields coupled to u_R and d_R, respectively.

[25] Effects of charged Higgs bosons in beta decay have been studied previously in the papers in Ref. [90].

The H^+ couples also to the second and the third family, with couplings proportional to m_μ and m_τ, respectively. This implies that $\bar{a}_{LP}$ (and therefore $\bar{a}_{LS}$) is not constrained by $(R_\pi)_{expt}$ (see Section 4.1). But $\bar{a}_{LS}$ is small. Even for $\tan\beta \simeq 65$ (which, roughly, is the upper limit on $\tan\beta$ [87]) and $m_{H^+} > 69$ GeV (the experimental limit from H^+-searches [11]) one has,

$$|\bar{a}_{LS}| \lesssim 4 \times 10^{-6} \ . \tag{118}$$

In models with more complicated Higgs sectors the pattern and size of the charged Higgs couplings could be different. In the most general conceivable scenario the beta decay interaction from charged Higgs exchange has the same form as the one from the $\beta_i^{(k)}$-terms ($k = u, d$; $i = L, R$) in the contact interactions in Eqs. (91) and (94). Using the same notation for the Higgs interaction as for the contact terms, the corresponding $\bar{a}_{LS}$ and $\bar{a}_{RS}$ are given by Eqs. (109) and (110). The constants $\bar{a}_{LP}$ and $\bar{a}_{RP}$, given by Eqs. (111) and (112), are constrained by the bounds (106-108) from $(R_\pi)_{expt}$.

5 Tensor Interactions

5.1 Model Independent Considerations

The most general tensor-type $d \to u e^- \bar{\nu}_e$ interaction is given in Eq. (5). We can rewrite this in the form

$$\begin{aligned} H = \Big[& a_{LT} \bar{e} \frac{1}{\sqrt{2}} \sigma_{\lambda\mu} (1-\gamma_5) \nu_e \\ & + a_{RT} \bar{e} \frac{1}{\sqrt{2}} \sigma_{\lambda\mu} (1+\gamma_5) \nu_e \Big] \, \bar{u} \frac{1}{\sqrt{2}} \sigma^{\lambda\mu} d + \text{H. c.} \ , \end{aligned} \tag{119}$$

where

$$a_{LT} = 2\alpha_{LL} \ , \tag{120}$$

$$a_{RT} = 2\alpha_{RR} \ . \tag{121}$$

The $n \to p e^- \bar{\nu}_e$ interaction induced by the Hamiltonian (119) is of the form (11), with

$$C_T - C_T' = 2 g_T a_{LT} \tag{122}$$

$$C_T + C_T' = 2 g_T a_{RT} \tag{123}$$

Tensor-type interactions would manifest themselves in the allowed approximation only in Gamow-Teller (or mixed) beta decays. T-invariant contributions to observables can depend on $g_T Re\bar{a}_{LT}$ (through interference with the SM contribution), $g_T^2(|\bar{a}_{LT}|^2 + |\bar{a}_{RT}|^2)[= \frac{1}{2}(|C_T|^2 + |C_T'|^2)\sqrt{2}/G_F V_{ud}]$ and $g_T^2(|\bar{a}_{LT}|^2 - |\bar{a}_{RT}|^2)[= -ReC_T C_T'^* \sqrt{2}/G_F V_{ud}]$, where $\bar{a}_{kT} = a_{kT}\sqrt{2}/G_F V_{ud}$ ($k = L, R$).

The best limits on $Re\eta_{LT}$ and η_{RT} from beta decay experiments are

$$|g_T Re\bar{a}_{LT}| < 1.3 \times 10^{-2} \qquad (90\% \text{ c. l.}) \tag{124}$$

(from a measurement of P_L^F/P_L^{GT} [82]), and

$$|g_T \bar{a}_{RT}| < 8 \times 10^{-2} \qquad (68\% \text{ c. l.}) \tag{125}$$

(implied by a limit on $g_T^2(|\bar{a}_{LT}|^2 + |\bar{a}_{RT}|^2)$, obtained from ^{6}He β decay [92].

Constraints on a_{LT} and a_{RT} of any origin come from the ratio R_π. Although a tensor interaction does not contribute directly to R_π, electromagnetic radiative corrections to the operators in Eq. (119) induce P-type interactions of strength

$$a_{LP} = \frac{1}{4} k_0 a_{LT} \tag{126}$$

and

$$a_{RP} = \frac{1}{4} k_0 a_{RT} \tag{127}$$

where $k_0 \simeq -2.8 \times 10^{-2}$ [93]. Using $|g_T| \lesssim 2.3$, we have from Eqs. (106-108), taking into account the bound (124),

$$|g_T Re\bar{a}_{LT}| < 1.1 \times 10^{-3} \tag{128}$$

if only a_{LT} contributes,

$$|g_T \bar{a}_{RT}^{(e)}| < 1.1 \times 10^{-2} \tag{129}$$

if a_{LT} is absent, and

$$\begin{gathered} |g_T Re\bar{a}_{LT}| < 0.25 \ , \\ |g_T \bar{a}_{RT}^{(e)}| < 0.13 \end{gathered} \tag{130}$$

if a_{LT} and a_{RT} contribute simultaneously.

For $u_e/u_\mu = 1$ the bound in Eq. (128) and the bound in Eq. (129) become $|g_T Re\bar{a}_{LT}| < 6 \times 10^{-4}$ and $|g_T \bar{a}^{(e)}_{RT}| < 7 \times 10^{-3}$, respectively.

Tensor-type $d \to ue^-\bar{\nu}_e$ interactions can arise from the exchange of spin-zero leptoquarks, and as contact interactions in composite models.

5.2 Contact Interactions

The γ_L and γ_R terms in (91) and (94) contain tensor-type charged-current interactions with [77]

$$a_{LT} = -\gamma_L^* \ , \tag{131}$$

$$a_{RT} = \gamma_R \ . \tag{132}$$

The limit on $\bar{a}_{LT}$ from $e^\pm p \to \overset{(-)}{\nu_e} X$ [79] is $|\bar{a}_{LT}| \lesssim 0.4$; for $\bar{a}_{RT}$ the bound should be the same.

5.3 Leptoquark Exchange

The $S-$ and P-type $d \to ue^-\bar{\nu}_e$ interactions generated by the exchange of the spin- zero LQs $(R_2)_-$, $(\tilde{R}_2)_+$, and S_1 are accompanied by tensor-type $d \to ue^-\bar{\nu}_e$ interactions. For these $|a_{kT}| = |a_{kP}|$ $(= |a_{kS}|)$ $(k = L, R)$ [53]. Consequently, for a_{LT} from $(R_2)_-$-exchange, a_{RT} from $(\tilde{R}_2)_+$-exchange and for a_{LT} and a_{RT} from S_1-exchange, the bounds (106), (107) and (108) apply, respectively [77].

6 Time Reversal Violation

At present there is no unambigous direct evidence for T-violation in the fundamental interactions [94]. But T-violation is intimately connected with CP-violation by the CPT theorem. Strong evidence for the validity of CPT invariance comes from the properties of $K^0 - \bar{K}^0$ mixing [94]. In the following we shall assume its validity and use the terms "T-violation" and "CP-violation" interchangably.

CP-violation has been seen in the mixing of the neutral kaons, and recently also in the $K^0 \to 2\pi$ decay amplitudes (through a nonzero value of the parameter ϵ'/ϵ) [94]. The latter result implies the existence of a non-superweak CP-violating interaction. One of the major questions in the field of CP-violation is the origin of the observed effects. The most economical possibility is that both ϵ and ϵ'/ϵ are due to the Kobayashi-Maskawa phase δ_{KM} in the SM. Future experimental investigations of CP-violation in B-decays and in some rare kaon decays will give further information on this possibility. Another important question is whether there are sources of CP-violation other than δ_{KM}, independently of their relevance or lack of it for the observed CP-violation.[26] As we shall discuss later on, searches for T-violation in beta decay aim to contribute to the second question.

T-violating interactions can be probed in beta decay through searches for T-odd correlations in the beta decay probability [13]. Sensitive experimental information is available on the coefficients D and R of the correlation $< \vec{J} > \cdot\vec{p}_e \times \vec{p}_\nu / JE_eE_\nu$ and $< \vec{\sigma}_e > \cdot < \vec{J} > \times\vec{p}_e/JE_e$ ($\vec{\sigma}_e$ and $\vec{J}$ are the electron and the nuclear spins), respectively.

Contributions to T-odd correlations arise not only from T-violating interactions, but also from (T-invariant) electromagnetic final state interactions. We shall write $D = D_t + D_f, R = R_t + R_f$, where D_t, R_t, and D_f, R_f are the T-violating and T-invariant contributions, respectively.

The D-correlation is sensitive to V,A-type T-violating interactions; the R-correlation probes scalar- and tensor-type T-violating interactions. In first order in the new beta decay interactions D_t and R_t resulting from the Hamiltonian (2) are given by [13].

$$D_t \simeq a Im\bar{a}_{LR} \quad , \tag{133}$$

$$R_t \simeq \frac{1}{2g_A}(a \mp b) g_T Im\bar{a}_{LT} - \frac{a}{2g_V} g_S Im\bar{a}_{LS} \quad , \tag{134}$$

where a and b are constants containing the nuclear matrix elements. The upper (lower) sign in the first term in Eq. (134) is for decays with $e^-(e^+)$ in the final state.

The best limit on D_t/a comes at present from ^{19}Ne decay. For ^{19}Ne decay $a \simeq -1.03$. The experimental value $(D)_{Ne} = (0.1 \pm 0.6) \times 10^{-3}$ [96] yields

$$|Im\bar{a}_{LR}| < 1.1 \times 10^{-3} \qquad (90\% \text{ c.l.}) \quad . \tag{135}$$

D_f has been estimated for this case to be $\sim 2 \times 10^{-4} p_e/(p_e)_{max}$ [97].

[26] It is interesting to mention in this connection that the Kobayashi-Mashawa phase in the SM is not sufficient to generate the baryon asymmetry of the universe [95].

A new experiment is under way at NIST to measure D in neutron decay with an expected sensitivity of 3×10^{-4}. The initial run yielded $(D)_n = [-0.6 \pm 1.2(stat) \pm 0.5(syst)] \times 10^{-3}$ [98]. D_f is smaller in neutron decay than in ^{19}Ne decay by an order of magnitude [96].

For R a measurement in ^{19}Ne decay gave [99]

$$(R)_{Ne} = 0.079 \pm 0.053 \tag{136}$$

This implies the limits $|g_S Im\bar{a}_{LS}| < 0.3$ (90% c.l.) and $|g_T Im\bar{a}_{LS}| < 0.5$ (90% c.l.) on a scalar and a tensor interaction, respectively.

A recent experiment measuring R in $^8Li \to {}^8Be + e^- + \bar{\nu}_e$ decay yielded [100]

$$(R))_{Li} = (-0.2 \pm 4.0) \times 10^{-3} \quad . \tag{137}$$

The results (136) and (137) are complementary, since R in ^{19}Ne decay can have contributions from both scalar and tensor interactions, while in ^8Li decay only from tensor interactions. Subtracting the contribution from R_f, the experimental result (137) gives $(R_t)_{Li} = (-0.9 \pm 4.0) \times 10^{-3}$ [100]. This value, and Eq. (134) with $(a)_{Li} \simeq 0$ and $(b)_{Li} = 4/3$ gives

$$|g_T Im\bar{a}_{LT}| < 1.4 \times 10^{-2} \qquad \text{(90\% c.l.)} \quad . \tag{138}$$

The best limit on $|g_S Im\bar{a}_{LS}|$ from beta decay experiments is about $|g_S Im\bar{a}_{LS}| \lesssim 0.1$, implied by the limit on $g_S^2(|\bar{a}_{LS}|^2 + |\bar{a}_{RS}|^2)$ obtained in the experiment of Ref. [83].

An experiment to measure R in neutron decay to an accuracy of 5×10^{-3} is being developed at PSI [101]. In neutron decay $R_f \simeq 10^{-3}$; the constants a and b in Eq. (134) are $a \simeq 0.87$ and $b \simeq 2.2$.

For $Im\bar{a}_{LS}$ and $Im\bar{a}_{LT}$ (and also for $Im\bar{a}_{LP}$) of any origin the bounds

$$|Im\bar{a}_{LS}|, |Im\bar{a}_{LT}| \lesssim 10^{-4} \tag{139}$$

have been deduced [102] from experimental limits on P,T-violating electron-nucleon interactions, provided by atomic and molecular electric dipole moment searches. The Ima_{LS}- and Ima_{LT}-interactions contribute to the $e - N$ interactions through diagrams involving W-exchange in addition to the scalar or tensor interaction. [27]. The bounds (139) are considerably more stringent than

[27] For an a_{LR} interaction no significant limit follows this way, since the corresponding diagrams are suppressed by $m_e m_u$ or $m_e m_d$ [102].

the present direct limits and the expected limit from R in neutron decay. However, the theoretical uncertainties associated with them could be large, and therefore improved direct limits on $Im\bar{a}_{LS}$ and $Im\bar{a}_{LT}$ would be still useful. A discussion of R_t in the extensions of the SM considered in Section 4 will be included in Ref. [103].

In the SM D_t and R_t are extremely small: the Kobayashi-Maskawa phase contributes only in second order in the weak interaction, and the contribution from the θ-term in the QCD Lagrangian is constrained by the stringent bound $|\theta| \lesssim 4\times 10^{-10}$ from the experimental limit on the electric dipole moment of the neutron. As a consequence, in the SM $|D_t/a|, |R_t/a| \lesssim 10^{-12}$ [104]. Thus D_t and R_t of observable size can come only from sources of CP-violation beyond those present in the SM.

We shall consider now the D-coefficient in extensions of the SM.

As we have seen in Section 3, an a_{LR}-interaction can arise in left-right symmetric models and in models with exotic fermions. It can arise also from leptoquark exchange, if (as expected) LQs of the same electric charge but different SM quantum numbers mix. In the contact interaction Lagrangian (91) an a_{LR}-interaction is forbidden by the requirement of invariance under the SM gauge group.

In L-R models we have from Eq. (65) [34]

$$D_t/a = -\frac{g_R}{g_L}\zeta\frac{\cos\theta_1^R}{\cos\theta_1^L}\sin(\alpha+\omega) \tag{140}$$

The phase $(\alpha+\omega)$ generates also a quark-quark interaction, which contributes to the electric dipole moment (EDM) of the neutron and to the EDMs of atoms and molecules. From the experimental limit on the EDM of the ^{199}Hg atom $d(^{199}Hg)$ [105] one can deduce [106]

$$|D_t/a| \lesssim 10^{-4}/k \quad , \tag{141}$$

where the constant k is expected to be of the order of 10. Calculations find values of ranging from k of the order of unity to k of the order of 100. The theoretical uncertainty in k is difficult to asses.

In models with exotic fermions the constant $\bar{a}_{LR}$ is given by Eq. (73). The elements of $\hat{V}_R$ are complex in general. D_t is given by [53]

$$D_t/a \simeq -s_R^u s_R^d (\hat{V}_R)_{ud}\sin\phi \quad , \tag{142}$$

where we have written $(\hat{V}_R)_{ud} = e^{i\phi}(\hat{V}_R)'_{ud}$, with $(\hat{V}_R)'_{ud}$ real. The limit on D_t/a

from $d(^{199}Hg)$ is the same here as in L-R models.

In models with leptoquarks CP-violation in the effective LQ-fermion couplings can originate from LQ mixing and from fermion mixing, and for spin-zero LQs also from complex LQ-fermion couplings.[28] Inspection shows that from spin-zero LQs D_t can arise from R_2 and $\tilde{R}_2$, if $(R_2)_-$ and $(\tilde{R}_2)_+$ mix; similarly, from spin-one LQs D_t can be generated by V_2 and $\tilde{V}_2$, if there is mixing between $(V_2)_-$ and $(\tilde{V}_2)_+$ [108,106].

D_t from the $(R_2)_- - (\tilde{R}_2)_+$ system comes from the h_{2L} and $\tilde{h}_{2L}$ couplings in Eqs. (83) and (84). The mass eigenstates B_1 and B_2 are related to $(R_2)_-$ and $(\tilde{R}_2)_+$ as $(R_2)_- = B_1 \cos\alpha + B_2 \sin\alpha$, $(\tilde{R}_2)_+ = (-B_1 \sin\alpha + B_2 \cos\alpha)e^{i\psi}$. Ignoring for simplicity fermion mixing, we obtain for maximal mixing from Eqs. (83) and (84)

$$D_t/a \simeq \frac{1}{16}\left(\frac{1}{m_1^2} - \frac{1}{m_2^2}\right)\left[h'_{2L}\tilde{h}'_{2L}\sin\left(\varphi_L - \tilde{\varphi}_L - \psi\right)\right]\left(G_V V_{ud}/\sqrt{2}\right)^{-1} \quad ,(143)$$

where we have written $h_{2L} = h'_{2L}e^{i\varphi_L}$ and $\tilde{h}_{2L} = \tilde{h}'_{2L}e^{i\tilde{\varphi}_L}$, with h_{2L} and h'_{2L} real.

The couplings h_{2L} and $\tilde{h}_{2L}$ induce also a P,T-violating quark-quark interaction, an EDM for the electron, and quark electric and chromoelectric dipole moments (which contribute to the neutron EDM). The quark-quark interaction (which is generated at one-loop level from diagrams involving W-exchange and containing a LQ propagator in one of the vertices) does not lead to a significant constraint on D_t/a, since it is suppressed by m_u^2 or m_d^2. The electron EDM and the quark electric and chromoelectric dipole moment do not arise at the one-loop level. Based on dimensional estimates of the dipole moments, the conclusion is that they allow D_t/a to be as large as the present experimental limit on D_t/a. The discussion and conclusions for D_t from the $(V_2)_- - (\tilde{V}_2)_+$ system is analogous.

7 Conclusions

In this article we reviewed and discussed new beta decay interactions in extensions of the Standard Model. Our aim was to consider the sources and the structure of such interactions, and the constraints on them from outside of beta decay, to assess what sensitivities would be required in beta decay experiments to obtain new information on them.

[28] Leptoquark interactions as a possible origin of the observed CP-violation was considered in the papers in Ref. [107].

We shall consider first T-invariant contributions.

From new V, A $d \to ue^-\bar{\nu}_e$ interactions T-even asymmetries and polarizations probe a_{RR}- and a_{RL}-type interactions (see Eq. (3)). New V, A interactions can arise in left-right symmetric models, in models with exotic fermions, in models involving leptoquarks, and also in composite models. We find that for all these models improvements of the beta decay limits on a_{RR} and/or a_{RL} would provide new information.

There are two types of scalar-pseudoscalar $d \to ue^-\bar{\nu}_e$ interactions for neutrinos of each helicity: one involving d_R and the other containing u_R (see Eq. (4)). If only one type is present, one has $|a_{kS}| = |a_{kP}|$ $(k = L, R)$. The ratio $R_\pi \equiv \Gamma(\pi \to e\nu_e)/\Gamma(\pi \to \mu\nu_\mu)$ sets in such a case stringent limits (the weakest being $\sim 8 \times 10^{-4}$) on a_{kP} (see Eqs. (106-108)), and therefore also on a_{LS}. This is the situation in the R-parity violating minimal supersymmetric standard model, and in leptoquark models. An exception is charged Higgs exchange in models where the ratio of the electronic and muonic Higgs-fermion coupling constant is equal to m_e/m_μ (R_π is not sensitive to such contributions). But then the Higgs contribution to beta decay is extremely small. If both types of scalar-pseudoscalar couplings are present, a_{LS} and a_{RS} are not constrained by R_π, but the absolute values of the coupling constant of the two types of couplings must be extremely close. This is the case for scalar-type interactions from Higgs exchange in general Higgs models, and from contact terms.

There is only one type of tensor $d \to ue^-\bar{\nu}_e$ interaction for neutrinos of each helicity. Such interactions can come from contact terms or from spin-zero leptoquark exchange. Tensor interactions contribute to R_π only through electromagnetic radiative corrections, and therefore the constraints from R_π are weaker (see Eqs. (128-130). As a consequence, improvements of the beta decay limits on a_{LT} and a_{RT} would give new information on contact interactions. For tensor interactions from leptoquark exchange this is not so, because the tensor interactions from leptoquark exchange are accompanied by pseudoscalar-type interactions of equal strength. Consequently, the same bounds (Eqs. (106-108)) apply as for the accompanying scalar interactions.

Turning to T-violation, the D-coefficient is sensitive to V,A type T-violating $d \to ue^-\bar{\nu}_e$ interactions. It can receive tree level contributions in left-right symmetric models, in models with exotic fermions, and from leptoquark exchange. In left-right symmetric models and in models with exotic fermions the limit on D_t/a from the electric dipole moment of the ^{199}Hg atom is better than the limit from beta decay by 1-2 orders of magnitude, but it is not as reliable as the direct limit. For D_t from leptoquark exchange the conclusion based on dimensional estimates of the relevant two-loop contributions to the electron electric dipole moment, and the electric and chromoelectric quark dipole moments is that D_t/a can be as large as the present experimental limit on D_t/a.

The R coefficient is sensitive to scalar- and tensor-type T-violating $d \to ue^-\bar{\nu}_e$ interactions. For these the limits obtained from P, T-violating electron-nucleon interactions are considerably more stringent than from R. However, the uncertainties in the indirect limits could be large, and therefore improved direct limits, even if weaker than from the $e - N$ interactions, are still useful.

Note Added

After this article was completed, a new experimental limit, $|d(^{199}Hg)| < 2.1 \times 10^{-28}$ ecm (95% c.l.), was reported on the electric dipole moment of the ^{199}Hg atom [109]. This improves the previous limit [105], and therefore the limit on D_t/α from $d(^{199}Hg)$ in left-right symmetric models and in models with exotic fermions (Eq. (141)), by a factor of 4.

Acknowledgments

I would like to thank Prof. A. Faessler for inviting me to write this article. I am grateful to R. N. Mohapatra for many useful conversations. I am indebted to J. A. Behr, T. J. Bowles, J. Deutsch, V. Egorov, G. Greene, C. Gould, A. Hime, W. R. Johnson, S. K. Lamoreaux, D. Počanić, N. Severijns, G. Steigman, P. A. Vetter, D. J. Vieira, W. S. Wilburn, and A. R. Young for informative conversations and communications. I would like to thank also the Referee for careful reading of the manuscript and useful comments. This work was supported by the Department of Energy, under contract W-7405-ENG-36.

References

[1] For a review, see R. E. Marshak, Riazuddin, and C. P. Ryan, *Theory of Weak Interactions in Particle Physics* (John Wiley & Sons, Inc., 1969); C. S. Wu and S. A. Moszkowski, *Beta Decay* (Interscience Publishers, John Wiley & Sons, New York, 1966).

[2] S. Weinberg, Phys. Rev. Lett. **19**, 1264 (1967); A. Salam, in *Elementary Particle Theory: Relativistic Groups and Analycity (Nobel Symposium No. 8)*, edited by N. Svartholm (Almqvist and Wiksell, Stockholm, 1969), p. 367; S. L. Glashow, Nucl. Phys. **22**, 579 (1961); S. L. Glashow, J. Iliopoulos, and L. Maiani, Phys. Rev. D 2, 1285 (1970).

[3] Previous reviews of the field include J. Deutsch and P. Quin, in *Precision Tests of the Standard Electroweak Model*, edited by P. Langacker, Advanced Series on Directions in High Energy Physics - Vol 14 (World Scientific, Singapore, 1995), p. 706; P. Herczeg, ibid., p. 786.

[4] J. D. Bjorken and S. D. Drell, *Relativistic Quantum Mechanics* (McGraw-Hill, New York, 1963).

[5] S. Weinberg, Phys. Rev. **112**, 1375 (1958).

[6] B. R. Holstein and S. B. Treiman, Phys. Rev. D **13**, 3059 (1976); P. Langacker, Phys. Rev. D **14**, 2340 (1976) and Phys. Rev. D **15**, 2386 (1976).

[7] L. Grenacs, Ann. Rev. Nucl. Part. Sci. **35**, 455 (1985); N. C. Mukhopadhyay, in *Physics Beyond the Standard Model*, proceedings of the Fifth International WEIN Symposium (WEIN'98), Santa Fe, New Mexico, USA, June 14-19, 1998, edited by P. Herczeg, C. M. Hoffman, and H. V. Klapdor-Kleingrothaus (World Scientific, Singapore, 1999), p. 222.

[8] For a review, see A. Yu. Smirnov, talk given at the XIXth International Conference on Neutrino Physics and Astrophysics, Neutrino 2000, June 16-21, 2000, Sudbury, Canada, hep-ph/0010097.

[9] For a review, see K. A. Olive, G. Steigman, and T. P. Walker, astro-ph/9905320.

[10] H-S. Kang and G. Steigman, Nucl. Phys. B **372**, 494 (1992); J. Kneller, R. Scherrer, G. Steigman, and T. P. Walker, to be submitted (G. Steigman, private communiation); see also S. Esposito, G. Mangano, G. Miele, and O. Pisanti, astro-ph/0005571.

[11] Particle Data Group: Review of Particle Physics, Eur. Phys. J. C **15**, 1 (2000).

[12] S. L. Adler *et al.*, Phys. Rev. D **11**, 3309 (1975).

[13] J. D. Jackson, S. B. Treiman, and H. W. Wyld, Jr., Phys. Rev. **106**, 517 (1957). Our constants C_K, C'_K, are the same as in this reference, except for the opposite sign of C'_V, C_A, C'_S, C'_P and C'_T.

[14] G. J. Stephenson, Jr., and T. Goldman, Phys. Lett. B **440**, 89 (1998); G. J. Stephenson, Jr., T. Goldman, and B. H. J. McKellar, Phys. Rev. D **62**, 093013 (2000), and references quoted therein.

[15] R. E. Shrock, Phys. Lett. B **96**, 159 (1980); J. Deutsch, M. Lebrun, and R. Prieels, Nucl. Phys. A **518**, 149 (1990).

[16] B. R. Holstein and S. B. Treiman, Phys. Rev. D **16**, 2369 (1977).

[17] G. Altarelli, G. F. Giudice, and M. L. Mangano, Nucl. Phys. B **506**, 29 (1997).

[18] For a review, see I. S. Towner and J. C. Hardy, in *Physics Beyond the Standard Model*, proceedings of the Fifth International WEIN Symposium (WEIN'98), Santa Fe, New Mexico, USA, June 14-19, 1998, edited by P. Herczeg, C. M. Hoffman, and H. V. Klapdor-Kleingrothaus (World Scientific, Singapore, 1999), p. 338.

[19] A. S. Carnoy, J. Deutsch, R. Prieels, N. Severijns, and P. A. Quin, J. Phys. G **18**, 823 (1992).

[20] W. J. Marciano and A. Sirlin, Phys. Rev. Lett. **56**, 22 (1986); Phys. Rev. D **35**, 1672 (1987).

[21] D. Počanić *et al.*, Experimental Proposal PSI R-89.01 (Paul Scherrer Institute, Villigen, 1988).

[22] P. Herczeg, Phys. Rev. D **52**, 3949 (1995).

[23] A. Sirlin, Phys. Rev. D **22**, 970 (1980).

[24] D. I. Britton *et al.*, Phys. Rev. D **49**, 28 (1994); Phys. Rev. Lett. **68**, 3000 (1992); G. Czapek et al., Phys. Rev. Lett **70**, 17 (1993). See also D. A. Bryman, Comments Nucl. Part. Phys. **21**, 101 (1993).

[25] W. J. Marciano and A. Sirlin, Phys. Rev. Lett. **71**, 3629 (1993).

[26] J. Deutsch, in *Physics Beyond the Standard Model*, proceedings of the Fifth International WEIN Symposium (WEIN'98), Santa Fe, New Mexico, USA, June 14-19, 1998, edited by P. Herczeg, C. M. Hoffman, and H. V. Klapdor-Kleingrothaus (World Scientific, Singapore, 1999), p. 322.

[27] N. Severijns *et al.*, Proposal CERN/ISC 96-11, ISC/P80 (Experiment IS356).

[28] For a review of neutron decay studies, see G. L. Greene, C. R. Gould, W. M. Snow, and F. Plasil, Summary of the Workshop on Fundamental Physics with Pulsed Neutron Beams, Triangle Park, NC, June 2000, to appear in the Proceedings (World Scientific, Singapore). Proposals for nuclear beta decay studies include A. Hime, Fundamental Symmetries with Trapped Atoms, LDRD proposal, Los Alamos National Laboratory, June 1998; S. G. Crane *et al.*, "Parity Violation Observed in Beta Decay of Magnetically Trapped ^{82}Rb Atoms", submitted to Phys. Rev. Lett.

[29] A. S. Carnoy, J. Deutsch, T. A. Girard, and R. Prieels, Phys. Rev. C **43**, 2825 (1991).

[30] G. Raffelt and D. Seckel, Phys. Rev. Lett. **60**, 1793 (1988).

[31] R. Barbieri and R. N. Mohapatra, Phys. Rev. D **39**, 1229 (1989).

[32] K. S. Babu, R. N. Mohapatra and I. Z. Rotstein, Phys. Rev. D **45**, R3312 (1992).

[33] J. C. Pati, and A. Salam, Phys. Rev. D **10**, 275 (1974); R. N. Mohapatra and J. C. Pati, Phys. Rev. D **11**, 566 (1975); **11**, 2558 (1975); G. Senjanović and R. N. Mohapatra, Phys. Rev. D **12** 1502 (1975). For a review, see R. N. Mohapatra, *Unification and Supersymmetry* (Springer-Verlag, New York, 1986).

[34] P. Herczeg, Phys. Rev. D **28**, 200 (1983).

[35] P. Herczeg, in *Weak and Electromagnetic Interactions in Nuclei*, proceedings of the International Symposium, Heidelberg, Germany, July 1-5, 1986, edited by H. V. Klapdor (Springer-Verlag, Berlin, 1986), p. 528.

[36] P. Langacker and S. Uma Sankar, Phys. Rev. D **40**, 1569 (1989).

[37] M. A. B. Bég, R. V. Budny, R. N. Mohapatra, and A. Sirlin, Phys. Rev. Lett. **38**, 1252 (1977).

[38] R. N. Mohapatra, F. E. Paige, and D. P. Sidhu, Phys. Rev. D **17**, 2642 (1987); See also R. N. Mohapatra, in *New Frontiers in High Energy Physics*, edited by A. Perlmutter and L. F. Scott (Plenum, New York, 1978), p. 337.

[39] A. S. Carnoy, J. Deutsch, and B. R. Holstein, Phys. Rev. D **38**, 1636 (1988).

[40] G. Beall, M. Bander, and A. Soni, Phys. Rev. Lett. **48**, 848 (1982). See also Ref. [36].

[41] H. Harari and M. Leurer, Nucl. Phys. B **233**, 221 (1984).

[42] F. Abe *et al.*, Phys. Rev. Lett. **74**, 2900 (1995); S. Abachi *et al.*, Phys. Rev. Lett. **76**, 3271 (1996).

[43] T. G. Rizzo, Phys. Rev. D **50**, 325 (1994).

[44] M. Cvetič, P. Langacker, and B. Kayser, Phys. Rev. Lett.**68**, 2871 (1992).

[45] P. Herczeg, "New V,A Interactions in Beta Decay" (in preparation).

[46] D. Chang, R. N. Mohapatra, and M. K. Parida, Phys. Rev. Lett. **52**, 1072 (1984); Phys. Rev. D **30**, 1052 (1984).

[47] R. N. Mohapatra, private communication.

[48] M. Cvetič and P. Langacker, Phys. Rev. D **42**, 1797 (1990); M. Cvetič *et al.*, Ref. [44].

[49] J. Erler and P. Langacker, Phys. Lett. B **456**, 68 (1999).

[50] E. Massó, Phys. Rev. Lett. **52**, 1956 (1984).

[51] P. Herczeg, Ref. [3].

[52] P. Langacker and D. London, Phys. Rev. D **38**, 886 (1988).

[53] P. Herczeg, in *Fundamental Symmetries in Nuclei and Particles*, edited by H. Henrikson and P. Vogel (World Scientific, Singapore, 1990), p. 46.

[54] E. Nardi, E. Roulet, and D. Tommasini, Nucl. Phys. B **386**, 239 (1992).

[55] D. London, in *Precision Tests of the Standard Electroweak Model*, edited by P. Langacker, Advanced Series on Directions in High Energy Physics - Vol 14 (World Scientific, Singapore, 1995), p. 951.

[56] For a review, see S. Davidson, D. Bailey, and B. A. Campbell, Z. Phys. C **61**, 613 (1994).

[57] The relevant papers include J. C. Pati and A. Salam, Phys. Rev. D **10**, 275 (1974); G. Senjanović and A. Šokorec, Z. Phys. C **20**, 255 (1983); R. N. Mohapatra, G. Segrè, and L. Wolfenstein, Phys. Lett. B **145**, 433 (1984); W. Buchmüller and D. Wyler, Phys. Lett. B **177**, 377 (1986); P. H. Frampton, Mod. Phys. Lett. A **7**, 559 (1992); P. H. Frampton and T. W. Kephart, Phys. Rev. D **42**, 3892 (1990); P. H. Frampton and B-H. Lee, Phys. Rev. Lett. **64**, 619 (1990).

[58] For a review, see J. L. Hewett and T. G. Rizzo, Phys. Rep. **183**, 193 (1989).

[59] R. N. Mohapatra, Phys. Rev. D **34**, 3457 (1986).

[60] See, for example, B. Schremmp and F. Schremmp, Phys. Lett. B **153**, 101 (1985).

[61] A few of the pertinent papers are E. Baver and M. Leurer, Phys. Rev. D **51**, 260 (1995); M. Leurer, Phys. Rev. Lett. **71**, 1324 (1993); Phys. Rev. D **49**, 333 (1994); W. Buchmüller, R. Rückl, and D. Wyler, Phys. Lett. B **191**, 442 (1987); P. Langacker, M. Luo, and A. K. Mann, Rev, Mod. Phys. 64, 87 (1992); P. Herczeg, Ref. [22]; Y. Kizikuri, Phys. Lett. B **185**, 183 (1987); O. Shanker, Nucl. Phys. B **206**, 253 (1982); Nucl. Phys. B **204**, 375 (1982). See also Refs. [57-60].

[62] W. Buchmüller, R. Rückl, and D. Wyler, Ref. [61].

[63] S. M. Barr, Phys. Rev. D **34**, 1567 (1986).

[64] M. Leurer, Ref. [61].

[65] S. C. Bennett and C. E. Wieman, Phys. Rev. Lett. **82**, 2484 (1999).

[66] A. Derevianko, Phys. Rev. Lett **85**, 1618 (2000).

[67] W. R. Johnson, private communication.

[68] See M. Strovink, in *Physics Beyond the Standard Model*, proceedings of the Fifth International WEIN Symposium (WEIN'98), Santa Fe, New Mexico, USA June 14-19, 1998, edited by P. Herczeg, C. M. Hoffman, and H. V. Klapdor-Kleingrothaus (World Scientific, Singapore, 1999), p. 442.

[69] C. Adloff *et al.* (H1 Collaboration), Eur. Phys. J. C **14**, 553 (2000).

[70] J. Breitweg *et al.* (ZEUS Collaboration), Eur. Phys. J. C **16**, 253 (2000).

[71] P. Abreu *et al.* (DELPHI Collaboration), Phys. Lett. B **446**, 62 (1999); G. Abbiendi *et al.* (OPAL Collaboration), Eur. Phys. J. C **6**, 1 (1999); M. Acciarri, *et al.* (L3 Collaboration), Phys. Lett. B **433**, 163 (1998).

[72] V. Barger and K. Cheung, Phys. Lett. B **480**, 149 (2000).

[73] H1 Collaboration, Phys. Lett B **379**, 319 (1996); Z. Phys. C **74**, 191 (1997); ZEUS Collaboration, Z. Phys. C **72**, 47 (1996).

[74] H1 Collaboration, talk at the XXIX International Coference on High Energy Physics, Vancouver, July 1998, Abstract 751 (quoted in Ref. [64] below); ZEUS Collaboration, talk at the XXIX International Coference on High Energy Physics, Vancouver, July 1998, Abstract 553 (quoted in Ref. [64] below); V. Chekelian, talk at the XXXIV Rencontres de Moriond, March, 1999.

[75] E. J. Eichten, K. D. Lane, and M. E. Peskin, Phys. Rev. Lett. **50**, 811 (1983).

[76] C. J. C. Burges and H. J. Schnitzer, Nucl. Phys. B **228**, 464 (1983); W. Buchmüller and D. Wyler, Nucl. Phys. B **268**, 621 (1986).

[77] P. Herczeg, "Scalar and Tensor Interactions in Beta Decay" (in preparation).

[78] F. Cornet and J. Rico, Phys. Let. B **412**, 343 (1997). For a review, see D. Zeppenfeld and K. Cheung, in *Physics Beyond the Standard Model*, proceedings of the Fifth International WEIN Symposium (WEIN'98), Santa Fe, New Mexico, USA, June 14-19, 1998, edited by P. Herczeg, C. M. Hoffman, and H. V. Klapdor-Kleingrothaus (World Scientific, Singapore, 1999), p. 462.

[79] F. Cornet and Rico, Ref. [78].

[80] F. Cornet, M. Relaño, and J. Rico, Phys. Rev. D **61**, 037701 (2000).

[81] I. S. Towner and J. C. Hardy, Ref. [18]; W. E. Ormand, B. A. Brown, and B. R. Holstein, Phys. Rev. C **40**, 2914 (1989).

[82] A. S. Carnoy, J. Deutsch, T. A. Girard, and R. Prieels, Phys. Rev. C **43**, 2825 (1991).

[83] E. G. Adelberger *et al.*, Phys. Rev. Lett. **83**, 1299 (1999); E. G. Adelberger, Phys. Rev. Lett. **70**, 2856 (1993), (E) ibid., **71**, 469 (1993).

[84] See A. Gorelov *et al.*, to appear in Hyperfine Interactions (2000); V. Egorov *et al.*, Nucl. Phys. A **621**, 745 (1997), and Letter of Intent E341a, SIRa facility, Ganil, Caen, France; P. A. Vetter *et al.*, to appear in proceedings of Spin 2000, Osaka, Japan, October 2000 (AIP Conference Proceedings, 2001).

[85] H. Leutwyler, Phys. Lett. B **378**, 313 (1996).

[86] B. McWilliams and L.-F. Li, Nucl. Phys. B **179**, 62 (1981); R. E. Schrock, Phys. Rev. D **24**, 1232 (1981); K. Mursula, M. Roos, and F. Scheck, Nucl. Phys. B **219**, 321 (1983).

[87] For a recent review, see S. P. Martin, hep-ph/9709356, v3:1999 [extended version of a contribution to the book *Perspectives in Supersymmetry*, edited by G. L. Kane (World Scientific, Singapore, 1998)].

[88] Recent reviews of R-parity violating supersymmetric models include G. Bhattacharyya, in *Physics Beyond the Standard Model*, proceedings of the Fifth International WEIN Symposium (WEIN'98), Santa Fe, New Mexico, USA, June 14-19, 1998, edited by P. Herczeg, C. M. Hoffman, and H. V. Klapdor-Kleingrothaus (World Scientific, Singapore, 1999), p. 60; H. Driner, hep-ph/9707435, in *Perspectives in Supersymmetry*, edited by G. L. Kane (World Scientific, Singapore, 1998); S. P. Martin, Ref. [87].

[89] R. N. Mohapatra, Phys. Rev. D **34**, 3457 (1986); V. Barger, G. F. Giudice, and T. Han, Phys. Rev. **40**, 2987 (1989).

[90] B. McWilliams and L.-F. Li, Ref. [85]; H. E. Haber, G. L. Kane, and T. Sterling, Nucl. Phys. B **161**, 493 (1979); W. E. Ormand, B. A. Brown, and B. R. Holstein, Ref. [81].

[91] For a brief review, see F. M. Borzumati and C. Greub, Phys. Rev. D **58**, 074004 (1998).

[92] F. Gluck, Nucl. Phys. A **628**, 493 (1998); C. H. Johnson, F. Pleasonton, and T. A. Carlson, Phys. Rev. **132**, 1149 (1963).

[93] M. B. Voloshin, Phys. Lett. B **283**, 120 (1992).

[94] For a recent review, see L. Wolfenstein, "CP-Violation", hep-ph/0011400, to be published in *Lecture Notes of TASI-2000*, edited by J. L. Rosner (World Scientific, Singapore, 2000).

[95] For a review see, for example, A. Riotto and M. Trodden, Ann. Rev. Nucl. Part. Sci. **49**, 35 (1999).

[96] A. L. Hallin, F. P. Calaprice, D. MacArthur, L. E. Pillonen, M. B. Schneider, and D. F. Schreiber, Phys. Rev. Lett **52**, 337 (1984).

[97] C. G. Callan, Jr. and S. B. Treiman, Phys. Rev. **162**, 1494 (1967).

[98] L. J. Lising *et al.*, (emiT collaboration), Phys. Rev. C **62**, 055501 (2000).

[99] M. B. Schneider, F. P. Calaprice, A. T. Hallin, D. W. MacArthur, and D. F. Schreiber, Phys. Rev. Lett. **51**, 1239 (1983).

[100] J. Sromicki, *et al.*, Phys. Rev. C **53**, 932 (1996).

[101] I. C. Barnett *et al.*, spokesperson: J. Sromicki.

[102] I. B. Khriplovich, Nucl. Phys. B **352**, 385 (1991); I. B. Khriplovich and S. K. Lamoreaux, *CP-Violation Without Strangeness* (Springer-Verlag, 1997).

[103] P. Herczeg, "T-Violation in Beta Decay" (in preparation).

[104] P. Herczeg and I. B. Khriplovich, Phys. Rev. D **56**, 80 (1997).

[105] J. P. Jacobs, W. M. Klipstein, S. K. Lamoreaux, B. R. Heckel, and E. N. Forston, Phys. Rev. A **52**, 3521 (1995).

[106] P. Herczeg, in *Symmetries and Fundamental Interactions in Nuclei*, edited by W. C. Haxton and E. M. Henley (World Scientific, Singapore, 1995), p. 89.

[107] S. M. Barr, Ref. [63]; S. M. Barr and E. M. Freire, Phys. Rev. D **41**, 2129 (1990); L. J. Hall and L. J. Randall, Nucl. Phys. B **274**, 157 (1986).

[108] P. Herczeg, in *Progress in Nuclear Physics*, edited by W.-Y Pauchy Hwang, S.-C. Lee, C.-E Lee, and D. J. Ernst (North-Holland, New York, 1991).

[109] M. V. Romalis, W. C. Griffith, and E. N. Fortson, hep-ex/0012001.