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BETA-DECAY BEYOND THE STANDARD MODEL

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ABSTRACT

We discuss possible sources of new interactions in beta-decay and the role of beta-decay experiments in obtaining information on them.

1. INTRODUCTION

It is a great pleasure and honor for me to have the opportunity to contribute to this Symposium celebrating Felix. The subject I have chosen for my talk has been one of the fields of his interest for many years. In the fifties Felix and his collaborators made important contributions¹⁾ concerning the properties of the beta-decay interaction to the sequence of experimental results that paved the way to the $V - A$ theory.

In 1968, when I came to Caltech as a postdoctoral fellow, associated half-time with the group led by Felix, the $V - A$ theory was, as ten years earlier, consistent with all pertinent data. Nevertheless, 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 that the observed CP-violation could be due to a new force, it was equally possible that with appropriate modifications the weak interaction could account for it. Another motivation to search for new interactions was the non-renormalizability of the $V - A$ theory. Felix was keenly interested in the implications of the new

theoretical models, and devoted his time to experiments that probed further the nature of the weak interactions and the origin of CP-violation.

Today the motivation to search for new interactions is not weaker. Although the minimal standard model is consistent with all observations, for many theoretical reasons, and especially because of the large number of undetermined parameters in the model, the presence of new physics is expected at some level.

In this talk we shall discuss possible sources of new interactions in beta-decay, and the role of beta-decay experiments in obtaining information on them.

In the next section we discuss beta-decay in left-right symmetric models, and then in models with new fermions. Sections 3 and 4 deal with scalar and tensor interactions, respectively. In Section 5 we analyze the effects of leptoquark exchange. The last section contains our conclusions.

In considering new contributions to beta-decay we shall assume that the quarks and leptons are elementary, and that their interactions are described by a renormalizable gauge theory. Then the $d \rightarrow ue^- \bar{\nu}_e$ (and $u \rightarrow de^+ \nu_e$) transition underlying beta-decay will be described by an interaction, which is included in the most general local (assuming that the masses of the exchanged particles are much larger than the momentum transfer) nonderivative four-fermion interaction. Neglecting induced form factors and the q^2 - dependence of the remaining form factors, the general beta-decay interaction has then the familiar form²⁾

$$\begin{aligned}
H_\beta = & \bar{e} \gamma_\lambda (C_V + C'_V \gamma_5) \nu_e \bar{p} \gamma^\lambda n \\
& + \bar{e} \gamma_\lambda \gamma_5 (C_A + C'_A \gamma_5) \nu_e \bar{p} \gamma^\lambda \gamma_5 n \\
& + \bar{e} (C_S + C'_S \gamma_5) \nu_e \bar{p} m + \bar{e} \gamma_5 (C_P + C'_P \gamma_5) \nu_e \bar{p} \gamma_5 n \\
& + \bar{e} \frac{\sigma^{\lambda\mu}}{\sqrt{2}} (C_T + C'_T \gamma_5) \nu_e \bar{p} \frac{\sigma^{\lambda\mu}}{\sqrt{2}} n + H.c.
\end{aligned} \tag{1}$$

The constants C_K, C'_K are real if time-reversal (T)-invariance holds. The terms proportional to C_P and C'_P do not contribute in the nonrelativistic approximation for the nucleons.

In writing the Hamiltonian (1) we have assumed that only one neutrino state is involved in beta-decay, but we shall consider where important the more

general case. We shall always assume that the kinematic effects of the masses of the neutrinos that can be produced in beta-decay can be neglected.³⁾

In the minimal standard model the effective interaction describing the transition $d \rightarrow ue^- \bar{\nu}_e$ is the V-A interaction

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

where $G/\sqrt{2} = g^2/8M_W^2$, and $U_{ud} = \cos\theta_1$ is the ud-element of the Kobayashi-Maskawa matrix. The neutrino is a massless two-component particle. The beta-decay Hamiltonian corresponding to (2) is of the form (1) with

$$C_V/g_V = -C_V'/g_V = C_A/g_A = -C_A'/g_A = -(G/\sqrt{2})U_{ud} \quad (3)$$

and all the other constants equal to zero. In Eq. (3) $g_V = g_V(0)$, $g_A = g_A(0)$, where $g_V(q^2)$ and $g_A(q^2)$ are the usual vector and axial-vector form factors for the $n \rightarrow p$ transition ($g_V = 1$ in the CVC approximation, and $g_A \simeq +1.25$).

Although the minimal standard model interaction dominates beta-decay, the data still allow sizable deviations from the V-A structure.⁴⁾ Recent experimental searches for deviations from the predictions of the minimal standard model focused among T-even observables on the longitudinal polarization of beta-rays from unpolarized nuclei. For the Hamiltonian (1), assuming that all the neutrino mass-eigenstates can be emitted in beta decay, the $e^\mp$ longitudinal polarization P_L in the direction parallel to $\vec{p}_e$ in allowed transitions is⁵⁾

$$P_L = G \frac{p_e}{E_e} / (1 + b \frac{m_e}{E_e}) , \quad (4)$$

where

$$\begin{aligned} G\xi = & 2|M_F|^2 \left[\mp \operatorname{Re}(C_S C_S'^* - C_V C_V'^*) - \frac{\alpha Z m_e}{p_e} \operatorname{Im}(C_S C_V'^* + C_S' C_V^*) \right] \\ & + 2|M_{GT}|^2 \left[\mp \operatorname{Re}(C_T C_T'^* - C_A C_A'^*) \right. \\ & \left. + \frac{\alpha Z m_e}{p_e} \operatorname{Im}(C_T C_A'^* + C_T' C_A^*) \right] , \end{aligned} \quad (5)$$

$$\begin{aligned} \xi = & |M_F|^2 (|C_S|^2 + |C_S'|^2 + |C_V|^2 + |C_V'|^2) \\ & + |M_{GT}|^2 (|C_A|^2 + |C_A'|^2 + |C_T|^2 + |C_T'|^2) \end{aligned} \quad (6)$$

and b is the Fierz interference term, given by

$$b\xi = \pm 2(1 - \alpha^2 Z^2)^{1/2} \text{Re} [|M_F|^2 (C_S C_V^* + C_S' C_V'^*) - |M_{GT}|^2 (C_T C_A^* + C_T' C_A'^*)] \quad (7)$$

From T-odd observables experimental information is available on the coefficients D and R of the correlations $\langle \vec{J} \rangle \cdot \vec{p}_e \times \vec{p}_\nu / J E_e E_\nu$ and $\vec{\sigma} \cdot \langle \vec{J} \rangle \times \vec{p}_e / J E_e$, respectively ($\vec{\sigma}$ = electron spin, $\vec{J}$ = nuclear spin).⁶⁾ D and R can be written as $D = D_t + D_f$, $R = R_t + R_f$, where D_t, R_t represent the T-violating contributions and D_f, R_f are the T-invariant contributions due to electromagnetic final state interactions. D_t and R_t are given by⁵⁾

$$D_t \xi = 2\delta_{J'J} M_F M_{GT} \left(\frac{J}{J+1} \right)^{1/2} \text{Im}(C_S C_T^* + C_S' C_T'^* + C_V C_A^* + C_V' C_A'^*) \quad (8)$$

$$R_t \xi = \pm 2\lambda_{J'J} |M_{GT}|^2 \text{Im}(C_T C_A^* + C_T' C_A'^*) + 2\delta_{J'J} M_F M_{GT} \left(\frac{J}{J+1} \right)^{1/2} \text{Im}(C_S C_A^* + C_S' C_A'^* + C_V C_T^* + C_V' C_T'^*) \quad (9)$$

where $\lambda_{J'J}$ is an angular momentum factor, defined in Ref. 5.

In observables we shall keep only the lowest order terms in new beta-decay couplings. Note that then

$$\begin{aligned} C_S C_A^* + C_S' C_A'^* &\simeq - (C_S C_A^* + C_S' C_A'^*) \simeq g_A (C_S C_V^* + C_S' C_V'^*) \\ &\simeq - g_A (C_S C_V^* + C_S' C_V'^*) \\ &- [(g^2/8M_W^2) U_{ud}]^2 g_A (\bar{C}_S - \bar{C}_S') \quad , \end{aligned} \quad (10)$$

$$\begin{aligned} C_T C_A^* + C_T' C_A'^* &\simeq - (C_T C_A^* + C_T' C_A'^*) \simeq g_A (C_T C_V^* + C_T' C_V'^*) \\ &\simeq - g_A (C_T C_V^* + C_T' C_V'^*) \\ &- [(g^2/8M_W^2) U_{ud}]^2 g_A (\bar{C}_T - \bar{C}_T') \quad , \end{aligned} \quad (11)$$

where we have introduced the notation $\bar{C}_K \equiv C_K / (-g^2/8M_W^2) U_{ud}$, $\bar{C}_K' \equiv C_K' / (-g^2/8M_W^2) U_{ud}$.

2. NEW V,A INTERACTIONS

2.1 Model Independent Considerations

The most general interaction for $d \rightarrow ue^{-}\bar{\nu}$ constructed from vector and axial-vector currents can be written in the form

$$\begin{aligned} H^{V,A} = & \bar{e}\gamma^\lambda(1-\gamma_5)(\Sigma_i U_{ei}\nu_i) [a_{LL}\bar{u}\gamma_\lambda(1-\gamma_5)d \\ & + a_{LR}\bar{u}\gamma_\lambda(1+\gamma_5)d] \\ & + \bar{e}\gamma^\lambda(1+\gamma_5)(\Sigma_i V_{ei}\nu_i)[a_{RL}\bar{u}\gamma_\lambda(1-\gamma_5)d \\ & + a_{RR}\bar{u}\gamma_\lambda(1+\gamma_5)d] + H.c. \end{aligned} \quad (12)$$

where u, d are mass-eigenstates, U_{ei} and V_{ei} are, respectively, elements of the left-handed and the right-handed leptonic mixing matrices, and ν_i are the light neutrino mass-eigenstates (i.e. the neutrino mass-eigenstates that can be produced in beta-decay). Apart from U_{ei} and V_{ei} , the Hamiltonian (12) contains seven independent parameters (four complex numbers minus an overall phase). The observables are quadratic functions of the a_{ik} 's; a set of independent parameters $|a_{ik}|^2 (i, k = L, R)$ and the phases of $a_{LL}a_{LR}^*$, $a_{RR}a_{RL}^*$, and $a_{LL}a_{RR}^*$. The last of these appears only in observables proportional to the neutrino masses, which we shall not consider. For beta-decay into ν_i the coupling constants C_V, C_V', C_A and C_A' can be read off from Eq. 12. Thus $C_A = g_A[(a_{LL} - a_{LR})U_{ei} - (a_{RL} - a_{RR})V_{ei}]$, etc. Normalized observables (such as polarizations and asymmetries) depend only on the ratios $\eta_{ik} = a_{ik}/a_{LL} (ik = LR, RL, RR)$.

As the neutrinos are not observed, the beta-decay probability is the sum of the probabilities of decays into the energetically allowed neutrino mass-eigenstates. The effect of this is taken into account (assuming, as we do, that the kinematic effects of the neutrino masses can be neglected) by replacing in observables a_{LL}, η_{RL} and η_{RR} by $a_{LL}^{(*)} = a_{LL}\sqrt{u_e}$, $\eta_{RL}^{(*)} \equiv \eta_{RL}\sqrt{\tilde{v}_e}$, and $\eta_{RR}^{(*)} \equiv \eta_{RR}\sqrt{\tilde{v}_e}$, respectively, where $u_e = \Sigma_i |U_{ei}|^2$, $v_e = \Sigma_i |V_{ei}|^2$, $\tilde{v}_e = v_e/u_e$ (the summation is over the matrix elements associated with the light neutrinos).

In pure transitions normalized observables are either independent of the η_{ik} 's or proportional to (see Ref. 7)

$$1 - 2|\eta_{RR}^{(\epsilon)} + \eta_{RL}^{(\epsilon)}|^2 \quad (\text{Fermi transitions}) \quad ,$$

and

$$1 - 2|\eta_{RR}^{(\epsilon)} - \eta_{RL}^{(\epsilon)}|^2 \quad (\text{Gamow - Teller transitions}) \quad .$$

The rates involve $|a_{LL}|^2$, and in Gamow-Teller decays and in mixed transitions also g_A ; normalized T-even observables in mixed transitions depend on $(g_A M_{GT}/g_V M_F) \text{Re}[(1 - \eta_{LR})/(1 + \eta_{LR})]$, and some also on $\text{Re}[(\eta_{RR}^{(\epsilon)} + \eta_{RL}^{(\epsilon)})(\eta_{RR}^{(\epsilon)} - \eta_{RL}^{(\epsilon)})^*]$, in addition to $|\eta_{RR}^{(\epsilon)} + \eta_{RL}^{(\epsilon)}|^2$ and $|\eta_{RR}^{(\epsilon)} - \eta_{RL}^{(\epsilon)}|^2$. The parameter η_{LR} enters in lowest order always in the above combination, so that even in cases where the nuclear matrix elements are known, uncertainties in the theoretical value of g_A prevent the possibility of setting stringent limits on $\text{Re}\eta_{LR}$ and $|\eta_{LR}|$.

The T-violating component of the D-coefficient is given by (see Eq. 8)

$$D_t = a_D \text{Im}(\eta_{LR} + \eta_{RL}^{(\epsilon)} \eta_{RR}^{(\epsilon)*}) \quad . \quad (13)$$

The constant a_D contains the nuclear matrix elements for ^{19}Ne - and for n-decay $a_D \simeq -1.03$ and $a_D \simeq 0.87$, respectively.

What information is available on the parameters η_{ik} ?

From experimental results on pure Gamow-Teller transitions⁸⁾ and on the e^+ -longitudinal polarization in a Fermi transition⁹⁾ one can deduce¹⁰⁾

$$|\eta_{RR}^{(\epsilon)}| < 0.13 \quad (14)$$

and

$$|\eta_{RL}^{(\epsilon)}| < 0.13 \quad (15)$$

Comparisons of e^+ -polarization in Fermi and Gamow-Teller decays^{11,12)} yield $P_L^F/P_L^{GT} = 1.00 \pm 0.003$ (see Ref. 12), implying $(P_L^F/P_L^{GT} - 1) \simeq -8 \text{Re}(\eta_{RR}^{(\epsilon)} \eta_{RL}^{(\epsilon)*})$

$$-5.2 \times 10^{-4} < (-\text{Re}\eta_{RR}^{(\epsilon)} \eta_{RL}^{(\epsilon)*}) < 7.4 \times 10^{-4} \quad (90\% \text{ c.l.}) \quad (16)$$

The best limit on D_t/a_D comes from ^{19}Ne -decay. Barring a cancellation between the two terms in D (see Eq. 13), the experimental value $D = (0.1 \pm 0.6) \times 10^{-3}$ [Ref. 13] yields

$$-0.86 \times 10^{-3} < \text{Im } \eta_{LR} < 1.05 \times 10^{-3} \quad (90\% \text{ c.l.}), \quad (17)$$

and

$$-0.86 \times 10^{-3} < \text{Im } \eta_{RL}^{(*)} \eta_{RR}^{(*)*} < 1.05 \times 10^{-3} \quad (90\% \text{ c.l.}), \quad (18)$$

From Eqs. (16) and (18) one obtains

$$|\eta_{RR}^{(*)}| |\eta_{RL}^{(*)}| < 1.3 \times 10^{-3} \quad . \quad (19)$$

There are further constraints on $\eta_{RL}^{(*)}$ and $\eta_{RR}^{(*)}$ from beta-ray asymmetries in polarized ^{19}Ne -and n-decays. In the case when $\eta_{RL}^{(*)}, \eta_{RR}^{(*)}$ are real and $\eta_{RR}^{(*)}$ positive, the combined beta-decay constraints yield $(-0.07) \lesssim \eta_{RL}^{(*)} \lesssim 0.05$ and $\eta_{RR}^{(*)} \lesssim 0.13$ (Ref. 14). For the general case ($\eta_{RL}^{(*)}, \eta_{RR}^{(*)}$ complex) the constraints have not been investigated yet.

A stringent limit on $|\eta_{RL}^{(*)}|$ and $|\eta_{RR}^{(*)}|$ has been derived from the energetics of Supernova 1989A. For right-handed neutrinos lighter than about 10 MeV the limit obtained (from the $e^-_R p \rightarrow \nu^{(R)}_e n$ reaction) is (Ref. 15; see also Ref. 16)

$$|\eta_{RL}^{(*)}| < 3 \times 10^{-5} \quad (20)$$

and

$$|\eta_{RR}^{(*)}| < 3 \times 10^{-5} \quad , \quad (21)$$

unless $|\eta_{RL}^{(*)}| > 2 \times 10^{-2}$ or $|\eta_{RR}^{(*)}| > 0.054$. The large values of $|\eta_{RL}^{(*)}|$ or $|\eta_{RR}^{(*)}|$, required to avoid the constraints (20) and (21) are not ruled out yet. A new class of beta-decay experiments involving measurements of the longitudinal positron polarization from polarized nuclei¹⁷⁾, with a new technique for e^+ -polarization measurements¹⁸⁾ is expected to improve the limits on $|\eta_{RL}^{(*)}|$ and $|\eta_{RR}^{(*)}|$ (Ref. 19).

For η_{LR} the best limit is

$$|\eta_{LR}| < 0.1 \quad (22)$$

provided by data on inclusive neutrino and antineutrino scattering.²⁰⁾

Right-handed interactions in beta-decay can arise if there are new charged gauge bosons which have right-handed couplings to the usual leptons and/or quarks, or if there are new quarks and leptons which mix with the usual ones and which have right-handed couplings to the W . Below we consider such possibilities. Still another possible source of right-handed interactions is leptoquark-exchange, which we shall consider in Section 5.

2.2 Left-Right Symmetric Models

Left-right symmetric models are attractive extensions of the standard electroweak model, which provide an understanding of the origin of parity-violation in the weak interactions.²¹⁾ The simplest models are based on the gauge group $SU(2)_L \times SU(2)_R \times U(1)$ (Ref. 22). These contain two distinct charged gauge boson fields W_L and W_R , which couple to left-handed and right-handed fermion currents, respectively. The gauge boson mass-eigenstates W_1 and W_2 are related to W_L and W_R as

$$\begin{aligned} W_L &= \cos \zeta W_1 + \sin \zeta W_2 \\ W_R &= e^{i\omega} (-\sin \zeta W_1 + \cos \zeta W_2) \end{aligned} \quad (23)$$

where ζ is a mixing angle and ω is a CP-violating phase. The exchange of W_1 and of W_2 generate an effective beta-decay interaction of the form 12 (Ref. 23).

Let us consider first the simplest version of $SU(2)_L \times SU(2)_R \times U(1)$ models, where the $SU(2)_L$ and $SU(2)_R$ gauge coupling constants g_L and g_R are equal ($g_L = g_R \equiv g$), the left-handed and the right-handed quark mixing matrices U_L and U_R are the same, $\tilde{v}_e, \tilde{v}_\mu, \dots = 1$, and CP-invariance holds (manifestly left-right symmetric models). In these the parameters $a_{LL}^{(\epsilon)}$, η_{LR} , $\eta_{RL}^{(\epsilon)}$ and $\eta_{RR}^{(\epsilon)}$ are given by

$$\begin{aligned} a_{LL}^{(\epsilon)} &= a_{LL} \simeq (g^2 \cos^2 \zeta / 8m_1^2) \cos \theta_1^L \\ \eta_{LR} &= \eta_{RL}^{(\epsilon)} \simeq -\zeta \\ \eta_{RR}^{(\epsilon)} &= \eta_{RR} \simeq m_1^2 / m_2^2 \equiv \delta, \end{aligned} \quad (24)$$

where m_1 and m_2 are the masses of W_1 and W_2 , and $\cos \theta_1^L$ is the ud -element of U_L .

The parameters ζ and δ are subject to the constraints described in Section 2.1. In addition, one has $|\zeta| < 0.033, \delta < 0.04$ from muon decay data, and the more stringent (but less reliable) bounds $|\zeta| \lesssim 4 \times 10^{-3}, \delta \lesssim 3 \times 10^{-3}$ from nonleptonic transitions.²⁴⁾ There is also a bound $|\zeta| \lesssim 10^{-2} - 10^{-3}$ implied by charged current universality.^{25,26)}

The above bounds are sufficient for the supernova limits (20) and (21) to become applicable, so that $|\zeta|, \delta \lesssim 3 \times 10^{-3}$. As a consequence, in manifestly left-right symmetric models the right-handed current effects in beta-decay are expected to be too small to be observable.

In the most general version of $SU(2)_L \times SU(2)_R \times U(1)$ models (models with non-manifest left-right symmetry and $\tilde{v}_e, \tilde{v}_\mu, \dots \neq 1$) muon decay gives a limit $|\eta_{RL}^{(e)}| < 0.047$, but no constraint on $\eta_{RR}^{(e)}$ (Ref. 7). The $K^0 \rightarrow \bar{K}^0$ transition sets a weak upper bound on $|\eta_{RR}^{(e)}|$, not better than 7.5×10^{-2} (Ref. 26). The best limit on $|\eta_{RR}^{(e)}| |\eta_{RL}^{(e)}|$ is from beta-decay data (Eq. 19). The upper limits on $|\eta_{RL}^{(e)}|$ and $|\eta_{RR}^{(e)}|$ are already near the values, where the supernova limits (20) and (21) become applicable.

For η_{LR} one has the constraint $|\eta_{LR}| \lesssim 4 \times 10^{-3}$ from nonleptonic K -decays and the D -coefficient.⁷⁾ There is no constraint on η_{LR} for light right-handed neutrinos from charged-current universality (see Ref. 16).

The T-violating component of the D -coefficient receives a contribution in first-order in $W_{1,2}$ -exchange²⁷⁾, given by

$$D_t \simeq -a_D \zeta \frac{g_R}{g_L} (\cos \theta_1^R / \cos \theta_1^L) \sin(\alpha + \omega) \quad , \quad (25)$$

where $\cos \theta_1^R e^{i\alpha} = (U_R)_{ud}$. In neglecting the second term in Eq. (13) we have assumed that $g_R^2 m_1^2 / g_L^2 m_2^2$ can be neglected relative to one, as indicated by the analysis in Ref. 16. The phase $(\alpha + \omega)$ contributes also to the electric dipole moment of the neutron D_n and to ϵ'/ϵ (Ref. 27). The limits on $|D_t a_D|$ from these observables are $\simeq 3 \times 10^{-5}$ (Ref. 28), but they are not as reliable as the limit (17).

2.3 Exotic Fermions

The interaction responsible for $d \rightarrow ue^{-}\bar{\nu}$ can contain terms involving right-handed currents even in models where the gauge group remains the standard $SU(2) \times U(1)$ group. This happens if new quarks and leptons exist which mix with the usual ones, and whose right-handed components are in non-singlet representations of $SU(2)$ (Ref. 29). If the electric charge and color assignments of the new fermions are the standard ones, the only possibility for the non-singlet right-handed fermions is to be in doublet representations (mirror-fermions or vector doublets).³⁰⁾ The new fermions have to be heavy (heavier than about 20 GeV) in view of limits from direct production at accelerators (see Ref. 30).

Requiring that the light-heavy fermion mixing does not lead to flavor-changing neutral currents involving the usual quarks and charged leptons (which are severely constrained by experiment), there is a considerable simplification in the form of the mixing matrices.³⁰⁾ In this framework the parameters $a_{LL}^{(e)}$, η_{LR} , $\eta_{RL}^{(e)}$ and $\eta_{RR}^{(e)}$ are in lowest order in light-heavy mixings given by

$$a_{LL}^{(e)} \simeq (-g^2/8M_W^2) \sqrt{u_e} \quad , \quad (26)$$

$$\eta_{LR} \simeq s_R^u s_R^d (\hat{V}_R)_{ud} \quad , \quad (27)$$

$$\eta_{RL}^{(e)} \simeq s_R^e \sqrt{\tilde{v}_e} \quad , \quad (28)$$

$$\eta_{RR}^{(e)} \simeq \eta_{LR} \eta_{RL}^{(e)} \quad . \quad (29)$$

Here $s_R^u \equiv \sin \theta_R^u$, etc.; θ_R^e , θ_R^u are the light-heavy mixing angles associated with the electron, u -quark, and the d -quark; the quantities u_e and $\tilde{v}_e$ were already defined in Sectin 2.1 (Ref. 31). The matrix $\hat{V}_R$ is unitary for $n_R^u = n_R^d = m_R^u$, where n_R^u, n_R^d and m_R^u is the number of charge $-\frac{2}{3}$ singlet states, the number of charge $-(\frac{1}{3})$ singlet states and the number of charge $-\frac{2}{3}$ doublet states, respectively.³⁰⁾

A comprehensive analysis of constraints on mixings between the usual and possible new fermions was made in Ref. 30. The results for the parameters

involved in (26) - (29) are $|s_R^e \sqrt{\tilde{v}_e}| < 0.042$ and $|s_R^u s_R^d \text{Re}(\hat{V}_R)_{ud}| < 6 \times 10^{-3}$, so that

$$|\eta_{RL}^{(e)}| < 0.042 \quad (30)$$

$$|\text{Re} \eta_{LR}| < 6 \times 10^{-3} \quad (31)$$

The limits (30), (31), and Eq. 29 imply

$$|\text{Re} \eta_{RR}^{(e)}| < 2.5 \times 10^{-4} \quad (32)$$

The elements of $\hat{V}_R$ are in general complex. The phase in $(\hat{V}_R)_{ud}$ contributes to D_t (Ref. 32). Writing $(\hat{V}_R)_{ud} = e^{i\phi} (\hat{V}_R')_{ud}$, where $(\hat{V}_R')_{ud}$ is real, we have (see Eqs. 27 - 29)

$$\begin{aligned} D_t &\simeq a_D [1 - (s_R^e \sqrt{\tilde{v}_e})^2] \text{Im} \eta_{LR} \simeq a_D \text{Im} \eta_{LR} \\ &\simeq a_D s_R^u s_R^d (\hat{V}_R')_{ud} \sin \phi \quad (33) \end{aligned}$$

The bounds (17) imply

$$|s_R^u s_R^d (\hat{V}_R')_{ud} \sin \phi| < 1 \times 10^{-3} \quad (34)$$

An upper limit of $\sim 3 \times 10^{-5}$ on $|s_R^u s_R^d (\hat{V}_R')_{ud} \sin \phi|$ follows again from D_n and ϵ'/ϵ .

Eqs. (31) and (34) yield the limit

$$|\eta_{LR}| < 6.1 \times 10^{-3} \quad (35)$$

From Eqs. (29), (30) and (34) we have

$$|\text{Im} \eta_{RR}^{(e)}| < 4.4 \times 10^{-5} \quad (36)$$

which, when combined with the bound (32), yields

$$|\eta_{RR}^{(e)}| < 2.0 \times 10^{-4} \quad (37)$$

From (35) and (37) one obtains $|\text{Re} \eta_{RR}^* \eta_{RL}| \leq |\eta_{RR}^{(e)}| |\eta_{RL}^{(e)}| < 1.1 \times 10^{-5}$.

3. SCALAR INTERACTIONS

The most stringent bound on the T-invariant component of scalar-type beta-decay interaction comes from the value³³⁾

$$b_F = (1.2 \pm 5) \times 10^{-3} \quad (38)$$

of the Fierz interference term (Eq. 7), derived from ft-values of superallowed $0^+ \rightarrow 0^+$ decays. This implies (see Eqs. 7 and 10)

$$|\text{Re}(\overline{C}_S - \overline{C}'_S)| = |b_F| < 9.4 \times 10^{-3} \quad (90\% \text{ c.l.}) \quad (39)$$

The best limit on T-violating scalar couplings is set by a measurement of the R-coefficient in ^{19}Ne -decay.³⁴⁾ The experimental value $R = -0.079 \pm 0.053$ yields (see Eqs. 9 and 10)

$$|\text{Im}(\overline{C}_S - \overline{C}'_S)| < 0.64 \quad (90\% \text{ c.l.}) \quad (40)$$

The limits on $|C_S|$ and $|C'_S|$ are an order of magnitude weaker than (39) (Ref. 4).

Scalar-type couplings can arise from charged Higgs-boson exchange. Charged Higgs bosons are present already in the standard model if the Higgs sector contains more than one Higgs doublet. The couplings of the charged Higgs boson to fermions in such models are not determined. The possibility that the corresponding scalar beta-decay couplings have values near the limits (39) and (40) cannot be ruled out, although there is no known natural model with such features (see Ref. 35).

Another possible source of scalar-type beta-decay interactions is lepto-quark exchange, which we shall discuss in Section 5.

Denoting the coupling of a charged Higgs boson (H) to $\bar{u}d$ and for example to $\bar{\nu}_e e$ by f'' and f' , C_S is given by $(f' f'' / m_H^2) g_S$, where $g_S \equiv g_S(0)$ is defined by $\langle p | \bar{u}d | n \rangle := g_S(q^2) \bar{u}_p u_n$. g_S was estimated in Ref. 36 and in Ref. 37, obtaining $g_S \simeq 1$ and $g_S \simeq 0.6$, respectively.

4. TENSOR INTERACTIONS

The experimental bounds on tensor couplings in beta-decay are

$$|\text{Re}(\overline{C}_T - \overline{C}'_T)| < (1.3 \times 10^{-2}) \quad (95\% \text{ c.l.}), \quad (41)$$

mainly from the Fierz interference term in ^{22}Na (Ref. 4), and

$$|\text{Im}(\overline{C}_T - \overline{C}'_T)| < (9.1 \times 10^{-2}) \quad (95\% \text{ c.l.}) \quad (42)$$

from the result $\text{Im}(C_T C_A'^* + C_T' C_A^*) = -0.020 \pm 0.047$ (Ref. 38), derived from data on $e^\mp$ longitudinal polarizations. The limits on $|C_T|$ and $|C_T'|$ are an order of magnitude weaker than (41) (Ref. 4).

A mechanism that could generate tensor couplings would be the exchange of spin-two bosons, coupled to lepton and tensor currents. However, there is no known renormalizable gauge theory involving tensor bosons. The only other possible mechanism for elementary fermions which gives rise to tensor couplings in beta-decay at the tree-level is the exchange of spin-zero leptoquarks (see Section 5).

The constants C_T, C_T' are given in terms of the coupling constants a_{TT}, a_{TT}' of the four-fermion couplings $\bar{e}(\sigma^{\mu\nu}/\sqrt{2})\nu_e \bar{u}(\sigma_{\mu\nu}/\sqrt{2})d$ and $\bar{e}(\sigma^{\mu\nu}/\sqrt{2})\gamma_5 \nu_e \bar{u}(\sigma_{\mu\nu}/\sqrt{2})d$ as $C_T = a_{TT}g_T, C_T' = a_{TT}'g_T$, where $g_T \equiv g_T(0)$ is defined by $\langle p|\bar{u}(\sigma_{\mu\nu}/\sqrt{2})d|n \rangle \simeq g_T(q^2)\bar{u}_p(\sigma_{\mu\nu}/\sqrt{2})u_n$. A calculation³⁷⁾ yields $a_T \simeq 1.2g_A$.

5. LEPTOQUARK EXCHANGE

Leptoquarks are bosons which induce quark-lepton transitions. They appear in many extensions of the minimal standard model. In models where they do not induce proton decay, leptoquarks could be light enough to cause observable effects in some low-energy processes.³⁸⁾ Both spin-one and spin-zero leptoquarks occur. The $d \rightarrow ue^- \nu$ transition can be mediated by either Q ($\equiv$ electric charge) $= \frac{2}{3}$ or $Q = -\frac{1}{3}$ leptoquarks. We shall denote spin-one and spin-zero leptoquarks of electric charge Q by $X^{(Q)}$ and $Y^{(Q)}$, respectively. The

four-fermion interactions $H_{X(Q)}$ for $d \rightarrow ue^-\nu_e$ generated by $X^{(Q)}$ -leptoquark exchange can be written in the form

$$H_{X(2/3)} = \sum_{i,j=V,A} f_{ij} \bar{e} \Gamma_i d \bar{u} \Gamma_j \nu_e + H.c., \quad (43)$$

$$H_{X(-1/3)} = \sum_{i,j=V,A} h_{ij} \bar{e} \Gamma_i u^c \bar{d}^c \Gamma_j \nu_e + H.c., \quad (44)$$

where $\Gamma_V = \gamma_\mu$, $\Gamma_A = \gamma_\mu \gamma_5$, and u^c, d^c are the charge-conjugates of u, d . The f_{ij} 's and h_{ij} 's are proportional to the product of the fermion-leptoquark couplings involved, and to m_{LQ}^{-2} , where m_{LQ} is the mass of the leptoquark. We shall assume that the left-handed neutrino state involved in the leptoquark interaction is the same as the one participating in the minimal standard model interaction. In Eqs. (43) and (44) we have taken for simplicity the left-handed and the right-handed neutrino states $\nu_e^{(L)}$ and $\nu_e^{(R)}$ to be the left-handed and right-handed components of the same neutrino field ν_e , and shall assume that ν_e is a mass-eigenstate. For general neutrino states the beta-decay coupling constants associated with interaction terms involving $\nu_e^{(L)}$ ($\nu_e^{(R)}$) have to be multiplied by $\sqrt{u_e}$ ($\sqrt{v_e}$). The four-fermion interactions induced by $Y^{(2/3)}$ - and $Y^{(-1/3)}$ -leptoquark exchange ($H_{Y(2/3)}$ and $H_{Y(-1/3)}$) have the same form as (43) and (44), respectively, except for $V, A \rightarrow S, P$ ($\Gamma_S = 1, \Gamma_P = \gamma_5$).

We shall consider now the effects of the various leptoquark-exchange Hamiltonians on beta-decay.⁴⁰⁾ Leptoquarks do not contribute in lowest order to leptonic and nonleptonic processes. Leptoquark contributions to semileptonic processes are therefore not constrained significantly by them.

$X^{(2/3)}$ -leptoquarks. After a Fierz transformation the Hamiltonian (43) yields a beta-decay interaction of the form (1) with

$$C_V/g_V = C_A/g_A = \frac{1}{2} (f_{VV} + f_{AA}) , \quad (45)$$

$$C'_V/g_V = C'_A/g_A = \frac{1}{2} (f_{VA} + f_{AV}) , \quad (46)$$

$$C_S/g_S = -C_P/g_P = -f_{VV} + f_{AA} , \quad (47)$$

$$C'_S/g_S = -C'_P/g_P = -f_{VA} + f_{AV} , \quad (48)$$

$$C_T = C'_T = 0 \quad , \quad (49)$$

where g_V, g_A and g_S have been defined earlier, and $g_P \equiv g_P(0)$ is defined by $\langle p | \bar{u} \gamma_5 d | n \rangle = g_P(q^2) \bar{u}_p \gamma_5 u_n$.

A stringent constraint on the pseudoscalar-type couplings comes from the experimental result⁴¹⁾

$$R_\pi \equiv \frac{\Gamma(\pi^+ \rightarrow e^+ \nu_e)}{\Gamma(\pi^+ \rightarrow \mu^+ \nu_\mu)} = (1.218 \pm 0.014) \times 10^{-4} \quad (50)$$

on $\pi \rightarrow \ell \nu$ decays. R_π is given by

$$R_\pi = (1.233 \times 10^{-4}) (Q^{(e)}/Q^{(\mu)}) \quad , \quad (51)$$

where $Q^{(e)}$ and $Q^{(\mu)}$ are sensitive to new semileptonic interactions involving axial-vector and pseudoscalar quark currents⁴²⁾ ($Q^{(e)}/Q^{(\mu)} = 1$ in the minimal standard model). The experimental result (50) implies

$$0.9663 < \frac{Q^{(e)}}{Q^{(\mu)}} < 1.0095 \quad (90\% \text{ c.l.}) \quad (52)$$

In the limits (52) we have included the theoretical uncertainty of 0.3% in the calculation of the coefficient of $Q^{(e)}/Q^{(\mu)}$ in Eq. 51. Neglecting the leptoquark contribution to a_{LL} , and assuming that there is no cancellation between the contribution of the pseudoscalar couplings involving the electron and the muon⁴³⁾, Eq. 52 yields⁴⁴⁾

$$| \text{Re} (\bar{C}_P - \bar{C}'_P) / g_P | < 1.2 \times 10^{-3} \quad (53)$$

$$| \text{Im} (\bar{C}_P - \bar{C}'_P) / g_P | < 0.2 \times 10^{-4} \quad . \quad (54)$$

The bounds (53) and (54), and the relations (47) and (48) imply (with $g_S \sim 1$)

$$| \text{Re} (\bar{C}_S - \bar{C}'_S) | \lesssim 1.2 \times 10^{-3} \quad , \quad (55)$$

$$| \text{Im} (\bar{C}_S - \bar{C}'_S) | \lesssim 0.2 \times 10^{-4} \quad . \quad (56)$$

For the Fierz-interference term b_F (Eq. 7) and the T-violating component of R in ^{10}Ne -decay (55) and (56) set the limits

$$|b_F| \lesssim 1.2 \times 10^{-3} \quad (57)$$

and

$$|(R_t)_{Ne}| \lesssim 1.6 \times 10^{-4} \quad , \quad (58)$$

respectively.

The V, A part of the $d \rightarrow ue^- \nu_e$ interaction (minimal standard model contribution plus the V, A part of the leptoquark contribution) is of the form (12) with

$$a_{LL} = a_{LL}^{(o)} + \frac{1}{4} (f_{VV} + f_{AA} - f_{VA} - f_{AV}) \quad , \quad (59)$$

$$a_{RR} = \frac{1}{4} (f_{VV} + f_{AA} + f_{VA} + f_{AV}) \quad , \quad (60)$$

$$a_{LR} = 0 \quad , \quad (61)$$

$$a_{RL} = 0 \quad , \quad (62)$$

where $a_{LL}^{(o)}$ is the minimal standard model contribution [$a_{LL}^{(o)} = -(g^2/8M_W^2)U_{ud}$]. As a consequence, the T-violating component of the D-coefficient vanishes. Neglecting the scalar contribution to P_L^F , we have $P_L^F = P_L^{GT} \approx 1 - 2|\eta_{RR}|^2$ (and therefore $P_L^F/P_L^{GT} = 1$).

The only constraints on the V, A -part of the interaction (as well as on the V, A -part of the other leptoquark-exchange interactions to be discussed below) are those described in Section 2.1.

$\Lambda^{(-1/3)}$ -leptoquarks. The beta-decay interaction resulting from the Hamiltonian (44) is of the form (1) with

$$C_V/g_V = -C_A/g_A = -\frac{1}{2} (h_{VV} + h_{AA}) \quad , \quad (63)$$

$$C'_V/g_V = -C'_A/g_A = -\frac{1}{2} (h_{VA} + h_{AV}) \quad , \quad (64)$$

$$C_S/g_S = C_P/g_P = h_{VV} + h_{AA} \quad , \quad (65)$$

$$C'_S/g_S = -C'_P/g_P = -h_{VA} + h_{AV} , \quad (66)$$

$$C_T = C'_T = 0 . \quad (67)$$

In view of the relations (65) and (66), the bounds (53) - (58) apply also in this case.

The constants a_{ik} of the V, A -part of the interaction, with the minimal standard model contribution included, are

$$a_{LL} = a_{LL}^{(o)} , \quad (68)$$

$$a_{LR} = -\frac{1}{4}(h_{VV} + h_{AA} - h_{VA} - h_{AV}) , \quad (69)$$

$$a_{RR} = 0 , \quad (70)$$

$$a_{RL} = -\frac{1}{4}(h_{VV} + h_{AA} + h_{VA} + h_{AV}) , \quad (71)$$

The T-violating component of the D-coefficient is now generally non-zero. We have from Eq. (13)

$$D_t = (a_D/(g^2/8M_W^2) U_{ud}) \frac{1}{4} \text{Im} (h_{VV} + h_{AA} - h_{VA} - h_{AV}) \quad (72)$$

Leptoquark contributions to D_t are constrained only by the experimental limit on D .

For P_L^F and P_L^{GT} we have in this case $P_L^F = P_L^{GT} \simeq 1 - 2|\eta_{RL}|^2$ (so that again $P_L^F/P_L^{GT} = 1$).

$\chi^{(2/3)}$ -leptoquarks. A novel feature of the beta-decay interaction generated by spin-zero leptoquark-exchange is the presence of tensor terms, in addition to V, A, S, P -couplings.

The V, A -part of the beta-decay interaction corresponding to $H_{Y(2/3)}$ is the same as for $H_{X(-1/3)}$, except for $h_{VV} + h_{AA} \rightarrow -\frac{1}{2}(-f_{SS} + f_{PP})$ and $h_{VA} + h_{AV} \rightarrow -\frac{1}{2}(-f_{SP} + f_{PS})$. Thus D_t is generally non-zero (given by Eq. 72 with the above substitutions), and also for P_L^F, P_L^{GT} the same conclusions hold as for $H_{X(-1/3)}$.

The S, P, T -couplings are related as

$$C_S/g_S = C_P/g_P = C_T/g_T , \quad (73)$$

$$C'_S/g_S = C'_P/g_P = C'_T/g_T \quad , \quad (74)$$

where $C_S = -\frac{1}{4}(f_{SS} + f_{PP})$, $C'_S = -\frac{1}{4}(f_{SP} + f_{PS})$. Consequently, the limits (53) and (54) imply

$$|b_{GT}| = |Re(\bar{C}_T - \bar{C}'_T)/g_A| \lesssim 10^{-3}g_T \quad (75)$$

$$|Im(\bar{C}_T - \bar{C}'_T)| < (6.2 \times 10^{-4})g_T \quad , \quad (76)$$

and also the bounds (55) - (57). For R_t in ^{10}Ne -decay we obtain

$$|(R_t)_{N_e}| \lesssim |1.9 - 1.6(1 - \frac{g_S}{g_T})| \times 10^{-4}|g_T| \quad . \quad (77)$$

$\chi^{(-1/3)}$ -leptoquarks. The V, A -part of the interaction resulting from $H_{Y(-1/3)}$ is the same as for $H_{X(2/3)}$, except for $f_{VV} + f_{AA} \rightarrow -\frac{1}{2}(-h_{SS} + h_{PP})$, $f_{VA} + f_{AV} \rightarrow -\frac{1}{2}(-h_{SP} + h_{PS})$.

The S, P, T -couplings are related as

$$C_S/g_S = C_P/g_P = -C_T/g_T \quad (78)$$

$$C'_S/g_S = C'_P/g_P = -C'_T/g_T \quad (79)$$

and therefore the bounds (55) - (57), (75) and (76) hold. For $(R_t)_{N_e}$ we find

$$|(R_t)_{N_e}| \lesssim |1.9 - 1.6(1 + \frac{g_S}{g_T})| \times 10^{-4}|g_T| \quad . \quad (80)$$

6. CONCLUSIONS

In this talk we discussed possible sources of new interactions in beta-decay⁴⁵⁾, their effects on beta-decay observables, and the constraints on the beta-decay coupling constants. We can summarize the main points as follows:

- Right-handed interactions can arise in beta-decay at the tree level in left-right symmetric models (or other models involving new gauge bosons with right-handed couplings), in models with exotic fermions, and in models involving leptoquarks.

An improvement of the existing limit on $|\eta_{RR}^{(\epsilon)}|$ through beta-decay measurements would further constrain models with $X^{(2/3)}$ - and $Y^{(-1/3)}$ -type leptoquarks, as well as general left-right symmetric models. In the above models $|\eta_{RR}^{(\epsilon)}|$ is not constrained by muon-decay data. The $K^0 \rightarrow \bar{K}^0$ transition does not impose a significant limit on $|\eta_{RR}^{(\epsilon)}|$ generated by leptoquark exchange, and constrains $|\eta_{RR}^{(\epsilon)}|$ in general left-right symmetric models only weakly. In models with exotic fermions various data combined set already a stringent upper limit ($\simeq 3 \times 10^{-4}$) on $|\eta_{RR}^{(\epsilon)}|$.

Improvements of the limit on $|\eta_{RL}^{(\epsilon)}|$ through beta-decay measurements would set more stringent constraints on models with $X^{(-1/3)}$ - and $Y^{(2/3)}$ -type leptoquarks. For general left-right symmetric models the best present limit on $|\eta_{RL}^{(\epsilon)}|$ comes from muon-decay, and in models with exotic fermions from a combination of data (including muon-decay).

The beta-decay limits for $\text{Re}(\eta_{RR}^{(\epsilon)} \eta_{RL}^{(\epsilon)*})$ and $\text{Im}(\eta_{RR}^{(\epsilon)} \eta_{RL}^{(\epsilon)})$ combined provide a complementary constraint for $|\eta_{RR}^{(\epsilon)}|$ and $|\eta_{RL}^{(\epsilon)}|$.

It should be noted that the present limits on $|\eta_{RR}^{(\epsilon)}|$ and $|\eta_{RL}^{(\epsilon)}|$ are close to the values where the limits derived from the energetics of Supernova 1989A become applicable.

- More sensitive searches for a nonzero D -coefficient would be of importance.

The T-violating component of the D -coefficient receives a tree-level contribution both in left-right symmetric models and in models with exotic fermions, and also in models with $X^{(-1/3)}$ - and $Y^{(2/3)}$ -type leptoquarks. For D_t in left-right symmetric models and in models with exotic fermions more stringent limits than the present experimental limit on D follow from the experimental results on D_n and ϵ'/ϵ . However these bounds are not rigorous, in view of the uncertainties in the calculations. The contribution to D_t from leptoquark exchange can be as large as the present experimental upper limit for D .

- For scalar type beta-decay interactions generated by charged Higgs-boson exchange values of the T-invariant and T-violating components of $C_S - C'_S$ as large as the present limits cannot be ruled out. Although there

is no compelling model where the size of the couplings can be so large, improvements of the present limits would be worthwhile.

Scalar interactions are generated also by leptoquark exchange. However, limits on pseudoscalar couplings derived from the $\pi \rightarrow e\nu/\pi \rightarrow \mu\nu$ ratio (whose validity require the absence of some cancellations) set stringent limits on such interactions (cf. Eqs. 55 and 56).

- The only possible source of tensor interactions in beta-decay at the tree level appears to be the exchange of spin-zero leptoquarks. For tensor interactions generated this way the stringent limits from $\pi \rightarrow e\nu/\pi \rightarrow \mu\nu$ apply (see Eqs. 75 and 76). In spite of these points, improved limits on tensor couplings, even above these limits, would be desirable.
- In models where the quarks and leptons are composite, which we did not discuss here, new beta-decay interactions could arise as a result of constituent exchange⁴⁶⁾. The strength of such interactions is of the order of $4\pi/\Lambda^2$, where Λ is the compositeness scale. With $\Lambda \gtrsim (1-3)$ TeV (the current limits from other processes), the corresponding beta-decay constants could have values near the present upper bounds.

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43. If $(\bar{C}_P - C'_P)/(C_P^\mu - C_P^{\mu'}) = m_e/m_\mu$, where $C_P^\mu, C_P^{\mu'}$ are coupling constants analogous to C_P, C'_P , but associated with couplings involving $\bar{\mu}\nu_\mu$ and $\bar{\mu}\gamma_5\nu_\mu$, one would have $Q^{(*)}/Q^{(\mu)} = 1$ like in the minimal standard model. This happens for Higgs couplings in some models, but here such a cancellation would have to be regarded as accidental.
44. In our paper (Ref. 42) and in O. Shanker (Ref. 39) the assumption was made that the pseudoscalar contributions to $\pi \rightarrow \ell\nu$ are small relative to the minimal standard model contribution. The upper limit in Eq. 53 then becomes 10^{-5} . However, there is no justification for such an assumption. The upper limit in Eq. 49 of Ref. 42 should be 2.9×10^{-5} , instead of 5.7×10^{-5} .
45. The same mechanisms can, of course, generate new charged current quark-lepton four-fermion interactions involving the other leptons. The corresponding coupling constants do not respect in general $e - \mu - \tau$ universality. $e - \mu - \tau$ universality holds only for left-right symmetric models with $u_\ell = v_\ell = 1 (\ell = e, \mu, \tau)$ and for the contributions of η_{LR} to normalized observables in general left-right symmetric models.

46. For a review see Eichten, E., et al., Rev. Mod. Phys. 56, 579 (1984); see also Jodidio, A., et al., Phys. Rev. D34, 1967 (1986) and references quoted therein.