Strongly Coupled Charged Scalar in B and T Decays

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Limits on charged-scalar Yukawa couplings from τ and B decays are discussed. They (and other existing limits) are consistent with "strong" couplings (~ 1 for the third generation) even if the lightest scalar mass is in the range 20 GeV $\lesssim M_\phi \lesssim 100$ GeV but saturated in this case. $B(B\to\tau\nu X)$ and (for $M_\phi > m_t$) $B(T\to\tau\nu X)$ may be then as large as 30% and 70%, respectively. Both for $M_\phi < m_t$ and $M_\phi > m_t$ the potentially possible top-quark signature in $p\bar{p}$ collisions is $\tau+2$ jets final state. The upper limit for $\tau\to\nu\eta\pi$ is $\simeq 0.003\%$.

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In most "beyond the standard" models several Higgs doublets are present and consequently the weak forces are mediated, in addition to the intermediate vector bosons, by charged scalar particles. Models with two and three scalar doublets have been explicitly studied and some constraints on the quark Yukawa couplings have been derived from K^0-$\bar{K}^0$, D^0-$\bar{D}^0$, and B^0-$\bar{B}^0$ mixing[1,2] and from CP nonconservation.[3] One appealing possibility is that in three- (or more-) doublet models the hierarchy of the vacuum expectation values is such that the Yukawa couplings are of the same order for the members of a given heavy generation and ~ 1 for the third generation (we shall call such couplings "strong").

On the phenomenological side a recent thorough analysis[4] of the high-precision data on muon decay, inverse muon decay, π_{l2} decays, and nuclear Gamow-Teller transitions shows that in those reactions the effective scalar coupling might be of the order of 10% of G_F, provided it is proportional to the lepton mass.

In this paper we analyze the limits on the charged-scalar couplings from the τ and B decays. In three- (or more-) doublet models with "natural" absence at the tree level of flavor-changing neutral currents there are two independent sets of Yukawa couplings, driving the up- and down-quark masses, respectively. Our bounds are consistent with both being ~ 1 for the third generation and they are in fact saturated by such strong couplings if the lightest scalar mass is in the range 20 GeV $\lesssim M_\phi \lesssim 100$ GeV. In this case the branching ratios $B(B\to\tau\nu X)$ and (for $M_\phi > m_t$) $B(T\to\tau\nu X)$ may be as large as 30% and 70%, respectively. Thus, those channels are very restrictive for potential scalar exchange.

τ decays.—Our notation (for effective couplings) is shown in Fig. 1. We also define the effective scalar couplings G,

$$G^L_{ij,l} = Y^L_{ij} Y_l / M_\phi^2, \quad G_{l,l} = Y_l Y_l / M_\phi^2, \text{ etc.,} \tag{1}$$

and the ratios H,

$$G^L_{ij,l}/G_F \equiv H^L_{ij,l}, \quad G_{l,l}/G_F \equiv H_{l,l}. \tag{2}$$

As has already been mentioned, from the e-μ sector one gets the limit[4]

$$H_{e,\mu} < 0.22 \tag{3}$$

provided that

$$Y_e = m_e/v_l, \quad Y_\mu = m_\mu/v_l. \tag{4}$$

With scalar exchange included, the rate for the decay

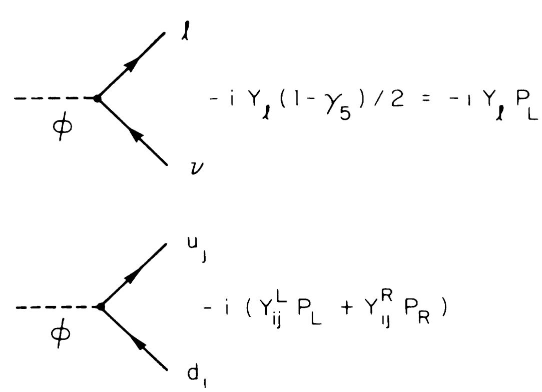

FIG 1. Our notation for effective couplings.

$\tau \to \mu \nu \bar{\nu}$ reads

$$\Gamma_{\tau\to\mu\nu\bar{\nu}} = \frac{G_F^2 m_\tau^5}{192\pi^3}(1+\tfrac{1}{32}H_{\tau,\mu}^2)\left[1 - \frac{m_\mu}{m_\tau}\frac{H_{\tau,\mu}}{\sqrt{2}(1+\frac{1}{32}H_{\tau,\mu}^2)} - 8\frac{m_\mu^2}{m_\tau^2} + O\left(\frac{m_\mu^4}{m_\tau^4}\right)\right]. \tag{5}$$

An analogous formula holds for $\tau \to e\nu\bar{\nu}$ but because of Eq. (4) the scalar-exchange contribution is negligible as compared to the contribution to $\tau \to \mu\nu\bar{\nu}$. Now, the ratio[5]

$$R \equiv \frac{\Gamma(\tau\to\mu\nu\bar{\nu}) - \Gamma(\tau\to e\nu\bar{\nu})}{\Gamma(\tau\to e\nu\bar{\nu})} = 0.016 \pm 0.037 \tag{6}$$

can be used to place a limit on $H_{\tau,\mu}$. Using Eq. (5) we have

$$R = -\frac{H_{\tau,\mu}}{\sqrt{2}}\frac{m_\mu}{m_\tau} + \frac{H_{\tau,\mu}^2}{32} - 8\frac{m_\mu^2}{m_\tau^2}\left(1+\frac{H_{\tau,\mu}^2}{32}\right) \tag{7}$$

and, from (6),

$$|H_{\tau,\mu}| \lesssim 2.5. \tag{8}$$

Let us now study the scalar-exchange contribution to the hadronic τ-decay modes. The rate for $\tau \to \pi^- \nu$ is measured with good accuracy. The full amplitude reads

$$M = \frac{G_F}{\sqrt{2}}\{f_\pi k^\mu \cos\theta_C \bar{\nu}_\tau\gamma_\mu(1-\gamma_5)\tau\phi_\pi^* + i\frac{Y_\tau}{2\sqrt{2}G_F M_\phi^2}(Y_{ud}^R - Y_{ud}^L)\langle\pi^-|\bar{d}\gamma^5 u|0\rangle\bar{\nu}_\tau(1+\gamma_5)\tau\}, \tag{9}$$

where k^μ is the pion four-momentum, f_π and ϕ_π are its decay constant and wave function, respectively, and θ_C is the Cabibbo angle. One gets

$$\Gamma(\tau\to\pi\nu) = \frac{G_F^2 f_\pi^2\cos^2\theta_C}{16\pi}m_\tau^3\left(1-\frac{m_\pi^2}{m_\tau^2}\right)^2\left|1 + i\frac{(H_{ud,\tau}^R - H_{ud,\tau}^L)\langle\pi^-|\bar{d}\gamma^5 u|0\rangle\phi_\pi}{2\sqrt{2}m_\tau f_\pi\cos\theta_C}\right|^2. \tag{10}$$

In spite of the apparent correction the prediction in Eq. (10) cannot be distinguished from the one in the standard model provided $Y_\tau/Y_\mu = m_\tau/m_\mu$. Indeed, the pion decay constant f_π is measured in the decay $\pi \to \mu\bar{\nu}$ whose decay rate is in this case subject to the same (numerically) correction. Thus, presence of the scalar contribution leads to a rescaling of f_π but does not affect the standard-model prediction for $\tau \to \pi\nu$ and no bound can be obtained from this process. An analogous result holds for the decay $\tau \to K\nu$ and the kaon decay constant f_K.

We are nevertheless able to get a certain constraint from the fact that the effective $\tilde{f}_\pi$ and $\tilde{f}_k$ satisfy $\tilde{f}_\pi/\tilde{f}_K \approx 1/1.27$. Thus

$$1.27 f_\pi\left|1 + i\frac{(H_{ud,\tau}^R - H_{ud,\tau}^L)\langle\pi|\bar{d}\gamma^5 u|0\rangle\phi_\pi}{2\sqrt{2}m_\tau f_\pi\cos\theta_C}\right| = f_K\left|1 + i\frac{(H_{us,\tau}^R - H_{us,\tau}^L)\langle K|\bar{s}\gamma^5 u|0\rangle\phi_k}{2\sqrt{2}m_\tau f_K\sin\theta_C}\right|^2, \tag{11}$$

and the pseudoscalar current matrix elements can be evaluated from partial conservation of axial-vector current[6]:

$$|\langle\pi|\bar{d}\gamma^5 u|0\rangle| = \sqrt{2}f_\pi m_\pi^2/(m_u+m_d), \tag{12}$$

and analogously for the kaon matrix element. Since we expect $f_\pi \approx f_K$, under more specific assumptions about the scalar sector Eqs. (11) and (12) can provide an approximate constraint on the coupling $H_{su,\tau}$ and $H_{du,\tau}$. This we explore in the context of the three-doublet model.

B decays.— In the standard model we have

$$\Gamma(B\to e\nu X) = (G_F^2 m_b^5/192\pi^3)[C_{ce}|V_{bc}|^2 + C_{ue}|V_{bu}|^2], \tag{13}$$

where the numerical factors C_{ce} and C_{ue} contain phase-space effects and QCD corrections for the transitions $b \to ce\nu$ and $b \to ue\nu$, respectively. Equation (13), rewritten in the form

$$\Gamma(B\to e\nu X) = \frac{G_F^2 m_b^5}{192\pi^3}C_{ce}|V_{bc}|^2[1+R], \tag{14}$$

where

$$R = C_{ue}|V_{bu}|^2/C_{ce}|V_{bc}|^2, \tag{15}$$

is used to determine $|V_{bc}|^2$ given the experimental data for $B(B\to e\nu X)$, the B-meson lifetime, the ratio R (from the momentum spectrum of leptons), and the "theoretical" value of C_{ce}, $C_{ce} \approx 0.4$ (with 5% error).[7] Since we focus on scalar couplings proportional to the lepton mass the result (13) does not change even with the scalar exchange present. Let us now consider the channel $b \to c\tau\nu$ (and similarly $b \to c\bar{c}s$ but, as we shall see in the next section, in multidoublet models the scalar contribution to the latter is strongly suppressed):

$$\Gamma(B\to X\tau\nu)=(G_F^2m_b^5/192\pi^3)[C_{c\tau}|V_{bc}|^2+C_{u\tau}|V_{bu}|^2+\Delta_\tau|V_{bc}|^2], \tag{16}$$

where

$$\Delta_\tau=\tfrac{1}{32}(|H^L_{bc,\tau}|^2+H^R_{bc,\tau}|^2)C_{c\tau}/|V_{bc}|^2,\quad C_{c\tau}\cong 0.06. \tag{17}$$

(In the first approximation we neglect the interference between $V-A$ and scalar amplitudes. This is justified *a posteriori* by the value of the upper bound on Δ as compared to $C_{c\tau}$.)

To our knowledge there is no experimental upper limit for $B(B\to X\tau\nu)$ but we can place a limit on Δ from its corresponding reflection on the total width and on the predicted $B(B\to e\nu X)$. The total width is

$$\Gamma(B\to X)=(G_F^2m_b^5/192\pi^3)[C_c|V_{bc}|^2+C_u|V_{bu}|^2+(\Delta_\tau+\Delta_c)|V_{bc}|^2], \tag{18}$$

where Δ_c describes the scalar contribution to $b\to c\bar{c}s$. The coefficients C_c and C_u are subject to various uncertainties mainly due to the choice of the quark mass values and to the strong interaction effects.[8] A fair choice[9] $C_c=2.75$ together with the experimental branching ratio $B(B\to e\nu X)=(12.3\pm0.9)\%$ leaves some room for scalar contribution. Using Eqs. (13) and (18) and $C_{ce}\approx 0.4$, and setting $R=C_{ue}|V_{bu}|^2/C_{ce}|V_{bc}|^2=0$ (experimentally $R<0.03$), we get

$$\Delta_\tau+\Delta_c\lesssim 0.7-0.8. \tag{19}$$

Multidoublet models.—As is well known, in such models the flavor-changing neutral currents are "naturally" (irrespectively of the values of the Yukawa couplings) eliminated at the tree level if only one scalar doublet couples to the up quarks and one other scalar is coupled to the down quarks.[10] In addition, for $Y_l=m_l/v$, only one scalar must couple to leptons (one of the two or one other).

In the mass-eigenstate basis the charged-scalar couplings to fermions read

$$L_Y^+=2^{3/4}G_F^{1/2}\bar{U}\Big[M_UK\textstyle\sum_i Y_i\Phi_i^+P_L-KM_D\sum_i X_i\Phi_i^+P_R\Big]D-\bar{\nu}M_LP_RL\sum_i X_i\Phi_i^++\text{H.c.}, \tag{20}$$

where K is the Kobayashi-Maskawa matrix and M_U, M_D, and M_L are mass matrices for the up and down quarks and the leptons, respectively. The explicit expressions for the constants X_i and Y_i in terms of the vacuum expectation values are well known in the two- and three-doublet models.[11] In the two-doublet model $X=1/Y$. In the three-doublet model there is no such constraint and both couplings may approach the perturbativity limit

$$2^{3/4}G_F^{1/2}|Y_i|m_t\sim 2^{3/4}G_F^{1/2}|X_i|m_b=O((4\pi)^{1/2}). \tag{21}$$

These limits are consistent with the constraints derived from the renormalization-group approach, Eq. (18) in Bagger, Dimopoulos, and Masso,[12] which give $|Y_i|\lesssim 5$, $|X_i|\lesssim 50$.

We have assumed leptons and down quarks to get masses from the same scalar doublet since, as we shall see, in both two- and three-doublet models this gives the strongest coupling of leptons to scalars.

In several previous papers the limits on Y_i's have been derived from the K^0-$\bar{K}^0$ and B^0-$\bar{B}^0$ mixing[1,2] and from the CP nonconservation in the kaon system.[13] Roughly speaking, for light scalars ($M_\phi\approx 20$–100 GeV), those limits coincide with the perturbativity limit (21) [one gets[2] $|Y|\lesssim 2(M_\Phi/m_t)^{1/2}$]. As long as $X_i\lesssim Y_i$ the effective scalar couplings are then some orders of magnitude below the experimental limits (3), (8), and (19), even for $M_\phi\approx 20$ GeV. At the tree level, only t decays are sensitive to the Y coupling. However, it is perfectly possible that $X_i\gg Y_i$ (e.g., both Yukawa couplings are equally important for the third generation), and then X_i couplings dominate in all but t decay sectors. The following discussion, including Eq. (24) and (25) holds for both two- and three-doublet models.

Neglecting in the first approximation all but the lightest Higgs-boson-exchange contribution to the effective coupling we get from (20)

$$H^R_{su,\tau}/H^R_{du,\tau}=\tan\theta_C m_s/m_d\approx(m_s/m_d)^{1/2}. \tag{22}$$

Equation (11) can now be used to place limits on H^R's if we make a plausible assumption that $1<1.27f_\pi/f_K<1.27$. One gets

$$H^R_{du,\tau}\lesssim 0.3,\quad H^R_{su,\tau}\lesssim 1.3. \tag{23}$$

When expressed in terms of X the limits (8) and (23) are similar and give

$$|X|\lesssim 2[M_\phi/(1\text{ GeV})], \tag{24}$$

whereas the limit (3) is somewhat weaker.

It turns out that the strongest limit for X is provided by B decays, Eq. (19). We get

$$|X|\lesssim 0.9[M_\phi/(1\text{ GeV})]. \tag{25}$$

The difference between the bounds (25) and (24) is important since cross sections are proportional to X^4. For instance the bound (25) makes the scalar contribution to $\tau\to\mu\nu\nu$ smaller than 0.4%. On the other hand, if (24) were saturated the $B\to\tau\nu X$ decay would saturate the total B width, in contradiction to experimental evidence for $B\to e\nu X$ and $B\to\mu\nu X$ and simple quark counting. Since Eq. (25) is the strongest limit presently available we can use Eqs. (6), (18), and (19) to estimate the upper limit for $B(B\to\tau\nu X)$:

$$B\lesssim 30\%, \tag{26}$$

to be compared to the standard-model result $\simeq 2.5\%$.

For the top-quark decay $t \to b\tau\nu$ we have

$$\Gamma(t \to b\tau\nu) = \frac{G_F^2 m_t^2}{192\pi^3}\left\{1 + \frac{(m_t m_\tau)^2}{4}\left[(YX)^2 + \left(\frac{m_b}{m_t}\right)^2 X^4\right]\frac{1}{M_\phi^4}\right\}. \tag{27}$$

Taking the upper limit $(X/M_\phi)^4 = 0.7\ \mathrm{GeV}^{-4}$, $Y = \frac{1}{2}\times(m_b/m_t)X$ [for $M_\phi \approx m_t$ this agrees with[2] $|Y| < 2(M_\phi/m_t)^{1/2}$], and $m_t \approx 50$ GeV, we get the bound (if $M_\phi > m_t$)

$$B(T \to \tau\nu X) \lesssim 70\% \tag{28}$$

and also the top-hadron lifetime about 2 times shorter than in the standard model. The branching ratios for the semileptonic decays into electrons and muons would be correspondingly smaller than in the standard model. If $M_\phi < m_t$ the top-quark total decay width is within very good approximation given by $t \to b\phi$ and the branching ratios for other decay modes are very small.

Given the bound (25) we can estimate the contribution of scalar exchange to channels strongly suppressed in the standard model, like, e.g., the $G = -1$, $P = +1$ decay mode $\tau \to \nu\pi\eta$. This is of interest in the context of the "missing mode" puzzle in τ decays.[14] The order-of-magnitude estimate is given by the inclusive rate $\tau \to \phi\nu \to q\bar{q}\nu$. We get 0.003% as the upper limit for $\tau \to \nu du$ and 0.06% for $\tau \to \nu su$.

Finally, let us comment on the K^0-$\bar{K}^0$, B^0-$\bar{B}^0$, and D^0-$\bar{D}^0$ mixing. It is not difficult to see that the first two mainly constrain the Y couplings and the last one is sensitive to the X couplings.[1,2] Our calculation with both terms present [Lagrangean (20)] shows that the K^0-$\bar{K}^0$ and B^0-$\bar{B}^0$ mixing is consistent with the derived bounds. In addition with the charged-scalar exchange included there is no longer a lower limit on the top-quark mass from the B^0-$\bar{B}^0$ mixing.[15] For D^0-$\bar{D}^0$ mixing we get $\Delta m_D \approx 10^{-16}$ GeV [for $M_\phi = 50$ GeV and using (25)].

Conclusions.— Strongly coupled scalars, with couplings ~ 1 for the third generation, have several theoretically attractive features, as advocated for instance in Ref. 3. The bounds on the Yukawa couplings derived in this paper from τ and B decays are consistent with strong couplings. (The specific set of parameters proposed in Ref. 3 in a model for the CP violation seems, however, to be ruled out by our bounds.) If the mass of the lightest scalar is in the range 20 GeV $\lesssim M_\phi \lesssim 100$ GeV those bounds are saturated by such strong couplings. Thus, strongly coupled scalars in that mass range, if they exist, should be soon observed directly and indirectly.

The branching ratios for $B \to \tau\nu X$ and (if $M_\phi > m_t$) for $T \to \tau\nu X$ are most promising and most restrictive for potential scalar exchange. With the present limits saturated they may be as large as 30% and 70%, respectively. For $M_\phi > m_t$ the top-quark decay signature would then be $t \to \tau$+jet and for $m_t > M_\phi$ the decay $t \to b\phi^+ \to \tau$+jet could be the dominant one. In both cases the top-quark signature in $p\bar{p}$ collisions could potentially be the final state τ+2 jets (the second jet is expected from $W \to tb$ decay) and it could have been missed so far, even if light.

A large contribution of $B \to \tau\nu X$ with the subsequent decay $\tau \to e\nu\nu$ would reveal itself in the shape of the momentum spectrum of electrons in $B \to eX$. Such an analysis is perhaps possible even with the presently available data.

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