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Flavor Changing Scalar Interactions

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Abstract

The smallness of fermion masses and mixing angles has recently been been attributed to approximate global $U(1)$ symmetries, one for each fermion type. The parameters associated with these symmetry breakings are estimated here directly from observed masses and mixing angles. It turns out that although flavor changing reaction rates may be acceptably small in electroweak theories with several scalar doublets without imposing any special symmetries on the scalars themselves, such theories generically yield too much CP violation in the neutral kaon mass matrix. Hence in these theories CP must also be a good approximate symmetry. Such models provide an alternative mechanism for CP violation and have various interesting phenomenological features.

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The inclusion of multiple scalar doublets at the weak scale in the standard $SU(2) \otimes U(1)$ electroweak theory entails the risk of flavor-changing neutral current processes with rates in excess of experimental bounds. To avoid this, most studies of such models have adopted the proposal[1] of a global symmetry that allows only one scalar doublet to couple to the right-handed quarks of each charge. Recently the need for such a symmetry has been challenged in an article[2] that attributes the various small ratios among quark mixing angles and quark and lepton masses to a set of approximate global $U(1)$ symmetries that act only on the quarks, but not on the scalars.[3] Specifically, it is assumed that every appearance of a fermion of the i'th generation in a Yukawa interaction of quarks or leptons with any scalar doublet ϕ_n is accompanied with a small dimensionless factor: ϵ_{Q_i} for left-handed quark doublets; ϵ_{U_i} or ϵ_{D_i} for right handed quarks; ϵ_{L_i} for left-handed lepton doublets; and ϵ_{E_i} for right-handed charged leptons. That is, writing the general Yukawa interaction in the form

$$\mathcal{L}_Y = -\lambda^U_{ijn}\bar{Q}_{Li}U_{Rj}\cdot\tilde{\phi}_n - \lambda^D_{ijn}\bar{Q}_{Li}D_{Rj}\cdot\phi_n - \lambda^E_{ijn}\bar{L}_{Li}E_{Rj}\cdot\phi_n + H.c., \qquad (1)$$

$$Q_{Lj} \equiv \begin{bmatrix} U_{Lj} \\ D_{Lj} \end{bmatrix} \quad \phi_n \equiv \begin{bmatrix} \phi_n^+ \\ \phi_n^0 \end{bmatrix} \quad \tilde{\phi}_n \equiv \begin{bmatrix} \phi_n^{0*} \\ -\phi_n^{+*} \end{bmatrix}$$

the Yukawa couplings are assumed to be of order

$$|\lambda^U_{ijn}| \approx \epsilon_{Q_i}\epsilon_{U_j} \qquad |\lambda^D_{ijn}| \approx \epsilon_{Q_i}\epsilon_{D_j} \qquad |\lambda^E_{ijn}| \approx \epsilon_{L_i}\epsilon_{E_j} \qquad (2)$$

for all n. (Here and below, "$\approx$" will be understood to indicate equality within a factor of order two or three.) Though there is no compelling theoretical justification for this assumption, it may be taken as representative of any theory of fermion-scalar couplings that attributes the small fermion masses and mixing angles to symmetries that act on the fermions rather than the scalars. With the aid of an additional somewhat ad hoc ansatz relating the ϵ's, it was shown in reference 2 that the rates of flavor-changing neutral current processes can be kept within experimental bounds without invoking any symmetry that restricts which scalars can interact with which quarks. We shall recover the same result here without using this ansatz. But as we shall see, there is a further problem with such multi-scalar models: they do not necessarily yield small violations of CP-conservation in the neutral kaon system.

To analyze this problem, the generations will be ordered so that, for $i < j$,

$$\epsilon_{Q_i} \leq \epsilon_{Q_j} \qquad \epsilon_{U_i} \leq \epsilon_{U_j} \qquad \epsilon_{D_i} \leq \epsilon_{D_j} \qquad \epsilon_{L_i} \leq \epsilon_{L_j} \qquad \epsilon_{E_i} \leq \epsilon_{E_j}. \tag{3}$$

The mass matrices arising from (1) may then be put into a real diagonal form by subjecting the fermions to transformations:

$$\begin{aligned} U_{Li} &\to V_{ij}^{U_L} U_{Lj} \qquad D_{Li} \to V_{ij}^{D_L} D_{Lj} \\ U_{Ri} &\to V_{ij}^{U_R} U_{Rj} \qquad D_{Ri} \to V_{ij}^{D_R} D_{Rj} \\ E_{Li} &\to V_{ij}^{E_L} E_{Lj} \qquad E_{Ri} \to V_{ij}^{E_R} E_{Rj}, \end{aligned} \tag{4}$$

with unitary matrices $V_{ij}^{U_L}$, etc., having elements

$$V_{ij}^{U_L} \approx \begin{cases} \epsilon_{Q_i}/\epsilon_{Q_j} & i \leq j \\ \epsilon_{Q_j}/\epsilon_{Q_i} & j \leq i \end{cases}, \tag{5}$$

and likewise for $V_{ij}^{D_L}$, $V_{ij}^{U_R}$, $V_{ij}^{D_R}$, $V_{ij}^{E_L}$, and $V_{ij}^{E_R}$. This transformation yields quark and lepton masses of order

$$m_{U_i} \approx \epsilon_{Q_i}\epsilon_{U_i}\langle\phi\rangle \qquad m_{D_i} \approx \epsilon_{Q_i}\epsilon_{D_i}\langle\phi\rangle \qquad m_{E_i} \approx \epsilon_{L_i}\epsilon_{E_i}\langle\phi\rangle \tag{6}$$

and a Cabibbo-Kobayashi-Maskawa (CKM) matrix of the form

$$V_{ij} \equiv [V^{U_L\dagger}V^{D_L}]_{ij} \approx \begin{cases} \epsilon_{Q_i}/\epsilon_{Q_j} & i \leq j \\ \epsilon_{Q_j}/\epsilon_{Q_i} & j \leq i \end{cases} , \tag{7}$$

where $\langle\phi\rangle$ is the common order of magnitude of the *conventionally* normalized complex neutral scalars, of order 247 GeV$/\sqrt{2} = 175$ GeV.

Now we will use experimental data to estimate the ϵ's. First, the ratios of the ϵ_{Q_i} are directly given by Eq. (5) in terms of the mixing angles. The ratio $\epsilon_{Q_1}/\epsilon_{Q_2}$ may be determined either from the Cabibbo angle

$$\epsilon_{Q_1}/\epsilon_{Q_2} \approx V_{us} = 0.218 \text{ to } 0.224$$

or less accurately from semi-leptonic B meson decays[4]

$$\epsilon_{Q_1}/\epsilon_{Q_2} \approx \frac{V_{ub}}{V_{cb}} \simeq 0.07 \ .$$

Given the theoretical uncertainties in extracting the ratio V_{ub}/V_{cb}, we regard these two estimates as being satisfactorily consistent, and we take $\epsilon_{Q_1}/\epsilon_{Q_2} = 0.2$. The second ratio of ϵ_{Q_i} is determined from

$$\epsilon_{Q_2}/\epsilon_{Q_3} \approx V_{cb} = 0.032 \text{ to } 0.054 \ .$$

Hence we take

$$\epsilon_{Q_1}/\epsilon_{Q_2} \approx .2 \qquad \epsilon_{Q_2}/\epsilon_{Q_3} \approx .04 \qquad \epsilon_{Q_1}/\epsilon_{Q_3} \approx .008 \ . \tag{8}$$

Using (8), (6), and the "experimental" values of the quark masses[5] , we have then also

$$\epsilon_{U_1} \approx .004/\epsilon_{Q_3} \qquad \epsilon_{U_2} \approx .2/\epsilon_{Q_3} \tag{9}$$

$$\epsilon_{D_1} \approx .006/\epsilon_{Q_3} \qquad \epsilon_{D_2} \approx .025/\epsilon_{Q_3} \qquad \epsilon_{D_3} \approx .03/\epsilon_{Q_3} \ . \tag{10}$$

The Yukawa couplings in Eq. (1) can then be estimated from Eq. (2), with the unknown ϵ_{Q_3} cancelling in all couplings.

Though it is not needed in estimating the Yukawa couplings, we can also estimate the factor ϵ_{Q_3} which is needed to determine the individual suppression factors. The top quark mass cannot be much less than $\langle\phi\rangle \simeq 175$ GeV, so if either of the quantities ϵ_{Q_3} and ϵ_{U_3} were much smaller than the other, then the larger would have to be much larger than unity, contrary to our assumption that the ϵ's are *suppression* factors. Thus Eq. (6) indicates that $\epsilon_{Q_3} \approx \epsilon_{U_3} \approx \sqrt{m_t/\langle\phi\rangle}$. But this actually applies to the Yukawa couplings defined at a renormalization scale of m_t, while we choose to quote the couplings defined at a renormalization scale of 1 GeV, which are larger by a factor $Z \approx 2$. We therefore estimate

$$\epsilon_{Q_3} \approx \sqrt{Z m_t/\langle\phi\rangle} \tag{11}$$

it being understood from now on that all ϵ's are defined at a renormalization scale of order 1 GeV.

With no measurable mixing angles for leptons, we cannot determine separate values for the leptonic suppression factors ϵ_{E_i} and ϵ_{L_i}. However the most stringent limits on scalar interactions were found in reference 2 to be set by the non-leptonic $K^0 \leftrightarrow \bar{K}^0$ and $B^0 \leftrightarrow \bar{B}^0$ transitions, to which we now turn. (The transitions $D^0 \leftrightarrow \bar{D}^0$ and $B_s^0 \leftrightarrow \bar{B}_s{}^0$ will be considered later.)

Exchange of neutral scalars produces two different kinds of parity-conserving $\Delta S = 2$ four-quark operators that can induce the transition $K^0 \leftrightarrow \bar{K}^0$:

$$\mathcal{L}_{\Delta S=2} = 2G(\bar{s}_R d_L)(\bar{s}_L d_R) + G'\left[(\bar{s}_L d_R)^2 + (\bar{s}_R d_L)^2\right] \tag{12}$$

with coupling constants

$$G = \sum_{nmN} \lambda^{D*}_{12n}\lambda^{D}_{21m} A_{nN} A^*_{mN}/m_N^2 \tag{13}$$

$$G' = \frac{1}{2}\sum_{nmN}[\lambda^{D}_{21n}\lambda^{D}_{21m} A_{nN} A_{mN} + \lambda^{D*}_{12n}\lambda^{D*}_{12m} A^*_{nN} A^*_{mN}]/m_N^2 \tag{14}$$

where

$$\langle 0|\phi^0_n(0)|N\rangle \equiv \frac{A_{nN}}{(2\pi)^{3/2}\sqrt{2\omega_N}} \tag{15}$$

and the sum over N runs over neutral Higgs scalar mass eigenstates. For an order-of-magnitude estimate of the $K_1^0 - K_2^0$ mass difference produced by this interaction, we will fall back on the vacuum insertion approximation (which is justified in quantum chromodynamics in the limit of a large number of colors):

$$\langle \bar{K}^0|\mathcal{O}_1\mathcal{O}_2|K^0\rangle \approx \langle \bar{K}^0|\mathcal{O}_1|0\rangle\langle 0|\mathcal{O}_2|K^0\rangle + \langle \bar{K}^0|\mathcal{O}_2|0\rangle\langle 0|\mathcal{O}_1|K^0\rangle \tag{16}$$

where each of $\mathcal{O}_1$ and $\mathcal{O}_2$ is either $(\bar{s}_L d_R)$ or $(\bar{s}_R d_L)$. The one-particle matrix elements of these operators can be calculated from the known matrix elements of the corresponding axial-vector current:

$$\langle 0|(\bar{s}_R d_L)|K^0\rangle = -\langle 0|(\bar{s}_L d_R)|K^0\rangle = \frac{m_K^2 F_K}{(2\pi)^{3/2}\sqrt{2m_K}2\sqrt{2}m_s} \tag{17}$$

where $F_K \simeq 230$ MeV is the kaon decay amplitude (as compared with $F_\pi \simeq$ 190 MeV.) This gives a $K_1^0 - K_2^0$ mass difference

$$\Delta M_K \approx \frac{(G-G')m_K^3 F_K^2}{4m_s^2} . \tag{18}$$

The flavor-changing suppression factors in G and G' turn out to be about the same

$$\epsilon_{Q_1}\epsilon_{D_2}\epsilon_{Q_2}\epsilon_{D_1} \approx \frac{1}{2}[\epsilon_{Q_2}^2\epsilon_{D_1}^2 + \epsilon_{Q_1}^2\epsilon_{D_2}^2] \approx 5 \times 10^{-8} . \tag{19}$$

The A_{nN} are of order unity, so barring unexpected cancellations, we expect that

$$G - G' \approx 5 \times 10^{-8}\, e^{i\delta}/m_H^2 \tag{20}$$

where δ is an unknown phase, and m_H is a weighted average of neutral scalar masses. Using this in (18) [with $m_s \simeq 180$ MeV] then yields

$$|\Delta M_K| \approx \frac{5\times 10^{-8} m_K^3 F_K^2}{4m_s^2 m_H^2} \approx \frac{3\times 10^{-5}\,\mathrm{eV}}{(m_H/300\,\mathrm{GeV})^2} . \tag{21}$$

The analysis we use to estimate ΔM_B parallels that used in Eqs. (12) to (21) for ΔM_K. The relevant coupling suppression factors are now

$$\begin{array}{ll} (\bar{b}_L d_R)(\bar{b}_R d_L) & \epsilon_{D_3}\epsilon_{Q_1}\epsilon_{D_1}\epsilon_{Q_3} \approx 10^{-6} \\ \frac{1}{2}[(\bar{b}_L d_R)^2 + (\bar{b}_R d_L)^2] & \frac{1}{2}[\epsilon_{D_3}^2\epsilon_{Q_1}^2 + \epsilon_{D_1}^2\epsilon_{Q_3}^2] \approx 2\times 10^{-5} \end{array} \tag{22}$$

(Note that the second suppression factor is an order of magnitude larger than the naive estimate $m_d m_b/\langle\phi\rangle^2$.) There have been many theoretical estimates of F_B, summarized by Buras and Harlander[6] . As a rough consensus value, we shall take $F_B \approx 230$ MeV. Following the same reasoning as for ΔM_K, we have then

$$|\Delta M_B| \approx \frac{2\times 10^{-5} m_B^3 F_B^2}{4 m_b^2 m_H^2} \approx \frac{10^{-2}\,\mathrm{eV}}{(m_H/300\,\mathrm{GeV})^2}\ . \tag{23}$$

There are also the more familiar box diagrams involving WW exchange. Assuming no accidental cancellations between these contributions, it seems reasonable to require that the scalar exchange contributions should not exceed twice the experimental values, $|\Delta M_K| = 3.5\times 10^{-6}$ eV and $|\Delta M_B| = (3.6\pm 0.7)\times 10^{-4}$ eV. This yields the conditions $m_H > 600$ GeV and $m_H > 1$ TeV, for K and B, respectively.

These Higgs masses are somewhat larger than seems plausible, but our analysis involves a number of dubious approximations, and it is entirely possible that we have overestimated the matrix elements for $K^0 \leftrightarrow \bar{K}^0$ and $B^0 \leftrightarrow \bar{B}^0$ transitions by factors of two or three. We conclude then that, as found in reference 2, the selection rule of reference 1 is not indispensable in keeping the scalar exchange contribution to the $K_1^0 - K_2^0$ and $B_1^0 - B_2^0$ mass differences at a reasonable size. But without this selection rule the Higgs scalars must be relatively heavy, and even so scalar exchange would be likely to make a large and perhaps dominant contribution to the $K_1^0 - K_2^0$ and $B_1^0 - B_2^0$ mas s differences.

Our conclusions shift when we consider the CP-violating part of the $K_1^0-K_2^0$ mass difference. This is usually expressed in terms of a parameter ϵ with $|\epsilon| \simeq 2.26 \times 10^{-3}$, which (for $|\epsilon'| \ll |\epsilon|$) is given by $\epsilon = \mathrm{Im}(\Delta M_K)/\sqrt{2}|\Delta M_K|$. If scalar exchange does indeed make a major contribution to the $K_1^0-K_2^0$ mass difference, then the phase δ in Eq. (20) would have to be quite small, of the order of a milliradian or less, in contradiction with the general expectation that all phases are of order unity. This leaves us with an interesting choice of alternatives:

- The CP-violating phases are indeed generically of order unity, but scalar exchange contributions to the $K_1^0 - K_2^0$ mass difference are much smaller than we have estimated, perhaps because of accidental cancellations in the calculation of the scalar-exchange contribution to the four-quark operator, or a gross failure of the vacuum insertion approximation used in calculating the $K^0 - \bar{K}^0$ matrix element, or both. This seems implausible unless the scalars are very heavy.

- The CP-violating phases are generically of order unity, but the scalar couplings *are* constrained by the selection rule of reference 1. This is of course automatic with just one scalar doublet, or in supersymmetric theories with just two scalar doublets. (However it is not at all automatic in supersymmetric theories with more than two scalar doublets. In particular, if several scalar doublets couple to the right-handed quarks of charge $-1/3$, and if as usually assumed these scalars have smaller vacuum expectation values than the doublets that couple to the right-handed quarks of charge 2/3, then the Yukawa couplings of these scalars would be correspondingly larger, leading to an even larger $K_1^0 - K_2^0$ mass difference.)

- All of the estimates in this paper are valid, but CP is a good approximate symmetry, with all the CP- violating phases like δ of order 10^{-3}.

The third alternative is admittedly a somewhat reactionary view of CP nonconservation. After the discovery of the process $K_2^0 \to \pi + \pi$ in 1964 it was widely assumed that this process is much slower than $K_1^0 \to \pi + \pi$ because CP is a good approximate symmetry for the weak interactions. Then following the discovery of a third generation of quarks and leptons in the 1970s, physicists became attracted to the idea that CP-violating phases are typically of order unity, and that CP only seems to be a good approximate symmetry because the third generation is weakly mixed with the first two. However, since we know that in any case we have to deal with quark masses and mixing angles that for mysterious reasons are very small, there is nothing absurd in supposing that CP-violating angles are also small. Indeed, apart from any consideration of scalar exchange effects, we may be driven to this assumption if theories with supersymmetry broken at the electroweak scale prove successful. Such theories have CP violating phases in the supersymmetry breaking interactions that generically lead to a neutron electric dipole moment three orders of magnitude larger than the present experimental limit[7] . This major problem of supersymmetrics models is avoided if we assume that CP-violating phases are generically of order 10^{-3}. In the balance of this paper we discuss the experimental consequences of this picture of CP violation, combined (where relevant) with our earlier assumptions regarding scalar couplings.

(1) Direct CP violating effects in the decays of K mesons will be unobservably small. The CKM contribution to $|\epsilon'/\epsilon|$ will be of order 10^{-6}, and the contribution from tree level exchange of scalar mesons will be even smaller. Hence these theories predict that the next round of experiments at CERN and Fermilab will *not* find a signal for $|\epsilon'/\epsilon|$ at the projected level of sensitivity of 10^{-4}. Such a null result would be extremely exciting since it would imply that the CKM matrix could not be the origin of the known CP violation (unless the top quark mass is found to take a value allowing a precise cancellation between two contributions to ϵ'/ϵ), thus implying an alternative source of CP violation, such as scalar exchange.

(2) All CP violating asymmetries which arise in particle decays must be of order 10^{-3} or less, since these asymmetries must be proportional to a CP violating phase. In particular CP violating effects in B meson decays will be too small to be observed in any experiment proposed to date. For example the angles α, β and γ of the unitarity triangle of the CKM matrix will be of order 10^{-3} and will be far too small to be observed at proposed B factories. Nevertheless such B factories could definitively exclude the CKM origin of

CP violation[8] .

(3) The most promising new positive signature of CP violation in our scheme is the neutron electric dipole moment. The electric dipole moment of the up quark arises from a one loop diagram with a virtual top quark and Higgs meson, and using the results of eqs. 8, 9 and 10 we estimate the resulting neutron electric dipole moment to be of order 10^{-26} e cm, close to the current experimental limit. In supersymmetric theories a comparable contribution would be expected from diagrams with internal superpartners. The electron electric dipole moment is expected to be of order 10^{-31} e cm.

(4) The predictions for the branching ratios for many rare K meson decays are not the same in our scheme as in the standard model. The most drastic change is for the $K_2^0 \to \pi\nu\bar{\nu}$ amplitude which is proportional to the CKM CP violating phase and therefore gets suppressed by two to three orders of magnitude. There is no tree level Higgs exchange contribution to this decay because the Higgs mesons do not couple to neutrinos.

(5) It is striking that for Higgs bosons with a typical mass of about 700 GeV and with couplings to quarks determined by Eqs. (8), (9) and (10), the tree level scalar exchange contribution to neutral K and B meson mass mixing turned out to be at about the level observed by experiment. Although this means that little can be learned about the CKM matrix from ΔM_K and ΔM_B, the case of D - $\bar{D}$ presents different opportunities. The analysis we use to estimate ΔM_D parallels that used in Eqs. (12) to (21) for ΔM_K and

ΔM_B. The relevant coupling suppression factors are now

$$
\begin{array}{ll}
(\bar{c}_L u_R)(\bar{c}_R u_L) & \epsilon_{U_1}\epsilon_{Q_2}\epsilon_{U_2}\epsilon_{Q_1} \approx 3 \times 10^{-7} \\
\frac{1}{2}[(\bar{c}_L u_R)^2 + (\bar{c}_R u_L)^2] & \frac{1}{2}[\epsilon^2_{U_1}\epsilon^2_{Q_2} + \epsilon^2_{U_2}\epsilon^2_{Q_1}] \approx 1 \times 10^{-6}
\end{array}
\tag{24}
$$

A theoretical estimate of F_D may be obtained from the previously quoted estimate $F_B \simeq 230$ MeV, using the relation (valid in the limit of large quark masses) $F_D/F_B \simeq \sqrt{m_b/m_c}$. This gives $F_D \approx 470$ MeV, so that

$$
|\Delta M_D| \approx \frac{10^{-6} m_D^3 F_D^2}{4 m_c^2 m_H^2} \approx \frac{2 \times 10^{-3}\ \mathrm{eV}}{(m_H/300\,\mathrm{GeV})^2}\ . \tag{25}
$$

If we take the typical Higgs mass as near 1 TeV to account for the observed values of $|\Delta M_K|$ and $|\Delta M_B|$, then the predicted value of $|\Delta M_D|$ is close to the current experimental limit, $|\Delta M_D| < 1.3 \times 10^{-4}$ eV. In the standard model ΔM_D is dominated by long distance contributions, which were originally estimated[9] to be in the range (0.3 to 0.01)$\times 10^{-4}$ eV, very much larger than the order 10^{-8}eV contribution from the short distance standard model box diagram. In this case, a positive observation of mass mixing at the level of 10^{-4} eV would not necessarily require new physics beyond the standard model. However a recent study[10] using heavy quark effective field theory and naive dimensional analysis suggests that the long distance standard model contribution to ΔM_D is in fact only modestly (about an order of magnitude) larger than the short distance contribution. Furthermore, a subsequent calculation[11] , which includes leading order QCD corrections, supports this low value of ΔM_D in the standard model. On this basis, we can conclude that a positive signal of neutral D meson mixing at the next round of searches at Fermilab, CESR and a tau/charm factory would provide evidence in favor of our scheme.

(6) For strange neutral beauty meson mixing $B_s^0 \leftrightarrow \bar{B}_s^0$ transitions, the relevant suppression factors are

$$\begin{array}{ll} (\bar{b}_L s_R)(\bar{b}_R s_L) & \epsilon_{D_3}\epsilon_{Q_2}\epsilon_{D_2}\epsilon_{Q_3} \approx 3\times 10^{-5} \\ \frac{1}{2}[(\bar{b}_L s_R)^2 + (\bar{b}_R s_L)^2] & \frac{1}{2}[\epsilon_{D_3}^2\epsilon_{Q_2}^2 + \epsilon_{D_2}^2\epsilon_{Q_3}^2] \approx 3\times 10^{-4}\ . \end{array} \tag{26}$$

Assuming that the experimental value of ΔM_B is dominated by scalar exchange, the scalar-mediated contribution to B_s mixing is predicted to be of order

$$(\Delta M_{B_s})_{scalar} \approx \left(\frac{\epsilon_{D_3}^2\epsilon_{Q_2}^2 + \epsilon_{D_2}^2\epsilon_{Q_3}^2}{\epsilon_{D_3}^2\epsilon_{Q_1}^2 + \epsilon_{D_1}^2\epsilon_{Q_3}^2}\right)\Delta M_B \approx 5\times 10^{-3}\,\mathrm{eV}\ . \tag{27}$$

(7) In theories with only one scalar doublet coupling to quarks of a given charge,[1] the positively charged scalars decay predominantly to $c\bar{s}$ and $\nu_\tau\bar{\tau}$, when the $t\bar{b}$ mode is kinematically forbidden. In the present class of theories the decay to $c\bar{b}$ completely dominates because the relevant products of ϵ_i are more than an order of magnitude larger for this mode than any other.

(8) Finally we consider exotic decay modes of the top quark. Our estimates indicate that in the class of theories we are considering Higgs particles would be too heavy to appear among the decay products of top quarks. But the phenomenology of such decays would be quite interesting, so it is worth considering the possibility that we have seriously overestimated neutral meson mass mixing, and that there are some Higgs scalars lighter than the top

quark. In most models with more than a single scalar doublet the exotic decays $t \to bh^+$ and $t \to ch^0$ will occur if they are kinematically allowed. (Here h^+ and h^0 are the lightest non-Goldstone mass eigenstates formed from linear combinations of the scalars destroyed by the fields ϕ_n^+ and ϕ_n^0 introduced in eq. 1.) As indicated above, the h^+ would decay predominantly through the channel $h^+ \to c\bar{b}$, and the h^0 decays predominantly via $h^0 \to b\bar{b}$, so that either of these exotic top quark decays yields $t \to b\bar{b}c$. However, as will be discussed below, the h^0 also has a large branching ratio to tau pairs.

The decays $t \to bh^+$ are induced by the Yukawa interaction $\lambda^U_{33n}\bar{Q}_{L3}U_{R3} \cdot \tilde{\phi}_n$, leading to a decay rate

$$\Gamma(t \to bh^+) \approx \frac{G_F m_t^3}{8\sqrt{2}\pi}\left(1 - \frac{m_{h^+}^2}{m_t^2}\right)^2 \tag{28}$$

If 46 GeV $\leq m_t \leq M_W$ then this exotic decay mode would dominate all others by a large factor, explaining how a top quark with mass less than m_W might not have been discovered. The charged Higgs h^+ decays predominantly to $c\bar{b}$. Using our values of the ϵ_i we compute the branching ratio to $\bar{\tau}\nu_\tau$ to be only $\approx 10^{-3}$. Hence, in this class of theories a successful search for the top quark at the Fermilab collider would require a technique to isolate candidate events with four b-type quarks and up to six jets. On the other hand, if $m_t > M_W$ we find

$$\frac{\Gamma(t \to bh^+)}{\Gamma(t \to bW^+)} \approx \left(\frac{1 - \frac{m_{h^+}^2}{m_t^2}}{1 - \frac{M_W^2}{m_t^2}}\right)^2 \frac{1}{1 + 2\frac{M_W^2}{m_t^2}} \tag{29}$$

which implies that a significant suppression of the conventional decay mode can occur. For example for a top quark mass of 100 GeV and a scalar mass of 50 GeV the conventional isolated lepton signature of the top quark will be suppressed by a factor of about 3. With sufficient statistics the top quark can still be discovered by the conventional mode, although a determination of its mass from the rate of these events could result in a considerable overestimate, about 25 GeV in the example given above. For $m_t \geq 150$ GeV the suppression of the conventional signal will be a factor of two or less.

Turning to the decay $t \rightarrow ch^0$, we note that this decay is of great interest since, unlike the decay to bh^+, this flavor-changing decay mode can only be large if the symmetry imposed in reference 1 is relaxed[12] . This decay is induced by the operator $\lambda^U_{32n}\bar{Q}_{L3}U_{R2} \cdot \tilde{\phi}_n$. The relevant coupling factor $\epsilon_{Q_3}\epsilon_{U_2} \approx 0.2$ is surprisingly large in this case[13] and such decays dominate (aside from the possible decay $t \rightarrow bh^+$) if the top quark is lighter than the W boson. The neutral Higgs h^0 decays predominantly to $\bar{b}b$. Using our values for the ϵ_i we find the branching ratio to tau pairs to be $\approx 10^{-1}$. Thus h^0 has much larger leptonic branching ratios than h^+. We expect the best signature at the Fermilab collider to occur when one neutral Higgs decays to b pairs and the other to tau pairs, with one tau giving an isolated electron and the other an isolated muon. For an integrated luminosity of $10 pb^{-1}$ and a top quark mass of 80 GeV, the Fermilab collider would produce ≈ 30 such events, with a signature $e + \mu +$ jets (from 2b and 2c quarks) + missing transverse energy.

A search for these events must take into account the softer p_T distribution of the isolated leptons compared to the distribution expected from conventional top quark decays.

For the case $m_t > M_W$, the exotic decay mode is no longer likely to dominate

$$\frac{\Gamma(t \to ch^0)}{\Gamma(t \to bW^+)} \approx \frac{\epsilon_{U_2}^2}{\epsilon_{U_3}^2} \left(\frac{1 - \frac{m_{h^0}^2}{m_t^2}}{1 - \frac{M_W^2}{m_t^2}} \right)^2 \frac{1}{1 + 2\frac{M_W^2}{m_t^2}} \ . \tag{30}$$

The decay $t \to ch^0$ does not significantly deplete the conventional decays, so the discovery of the top quark is not hindered by this process. However the discovery of such exotic, flavor-changing decays would not only reveal a Higgs boson but would strongly suggest a theory of several scalar doublets with approximate flavor and CP symmetries.

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