Light scalar as the messenger of electroweak and flavor symmetry breaking

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We propose a new framework for understanding the hierarchies of fermion masses and mixings. The masses and mixings of all standard model (SM) charged fermions other than top arise from higher dimensional operators involving a messenger scalar S and flavon scalars F_i. The flavons spontaneously break SM flavor symmetries at around the TeV scale. The SM singlet scalar S couples directly to the Higgs H and spontaneously breaks another $U(1)$ at the electroweak scale. At the TeV scale, SM quarks and charged leptons have renormalizable couplings to S, but not to H or F_i. These couplings involve new heavy vectorlike fermions. Integrating out these fermions produces a pattern of higher dimensional operators that reproduce the observed hierarchies of the SM masses and mixings in terms of powers of the "little hierarchy": the ratio of the electroweak scale to the flavor-breaking scale. The framework has important phenomenological implications. Flavor-changing neutral currents are within experimental limits but $D^0 - \bar{D}^0$ mixing and $B_s \to \mu^+ \mu^-$ could be close to current sensitivities. The neutral scalar s of the messenger field mixes with the light Higgs of the SM, which can have strong effects on Higgs decay branching fractions. The s mass eigenstate may be lighter than the Higgs, and could be detected at the Tevatron or the LHC.

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I. INTRODUCTION

Explaining the fermion mass hierarchy and mixing pattern is an outstanding challenge of particle physics [1–3]. The fermion masses are parametrized by the standard model (SM) Yukawa interactions of chiral fermions with a single Higgs doublet. It is technically natural for the dimensionless Yukawa couplings to take small values, since global chiral flavor symmetries are restored (at tree level) in the limit that these couplings vanish, but it is a total mystery why these values are spread over more than 5 orders of magnitude in a suggestive pattern of intergenerational and intragenerational hierarchies.

Although the gauge sector of the SM is well established, little is yet known about the Higgs sector. Higgs physics may be much richer than the minimal SM formulation, presenting new dynamics at the TeV scale that will be accessible to experiments at the LHC. Most work on extended Higgs sectors has been motivated by frameworks for understanding the naturalness and hierarchy problem of the SM Higgs boson, but not by the hierarchy problems of the SM flavor sector. One reason is that models that attempt to generate the flavor-breaking patterns of the SM Yukawas from new TeV scale dynamics are strongly constrained by experimental searches for flavor-changing neutral currents (FCNCs) and charged lepton flavor violation (CLFV).

The top quark Yukawa coupling has a value close to 1, suggesting that a SM Yukawa coupling is the correct

explanation for the top mass. The smallness of the other Yukawas suggests that some or all of the other quarks and the charged leptons do not couple directly to the electroweak symmetry breaking order parameter, which in the SM is represented by the vacuum expectation value (vev) of the Higgs scalar. Thus a good starting point to construct theories of flavor is to specify a field or mechanism to act as the messenger of electroweak symmetry breaking to the other quarks and leptons.

One simple choice for a messenger is a TeV mass scalar leptoquark, postulated to have a renormalizable coupling between the top quark and the SM leptons [4,5]. Radiative corrections can then generate a natural hierarchy of fermion masses related to powers of a loop factor.

An even simpler choice for a messenger is an electroweak mass scalar that transforms as a SM singlet and extends the Higgs sector of the SM. In this work, we explore this idea of an extended Higgs sector related to the generation of the fermion mass hierarchy. We present a simple framework where the Higgs doublet H couples directly to a complex scalar S that is a SM singlet and is charged under a new local $U(1)_S$ symmetry carried by a vector boson Z'. All of the SM fermions are singlets under this new $U(1)_S$ (apart from small effects from $Z - Z'$ mixing), which is broken spontaneously at the electroweak scale by the vacuum expectation value of S.

In our framework the singlet scalar S is the messenger to SM fermions of both flavor breaking and electroweak symmetry breaking. All SM fermions apart from the third generation quark doublet q_{3L} and right-handed top u_{3R} are assumed to carry a nonzero charge under a gauged chiral flavor symmetry forbidding all SM dimension 4 Yukawa

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couplings except that of the top quark. We assume that the flavor symmetry is spontaneously broken at a scale $\gtrsim 1$ TeV by the vacuum expectation of one or more complex scalar "flavon" fields F_i. The flavor charges of the SM fermions forbid any dimension 4 couplings to either F_i or to the Higgs field H.

We introduce new fermions that are vectorlike under both the SM gauge symmetries and $U(1)_S$; these fermions naturally acquire masses $\gtrsim$ TeV that we will generically denote as M, and have dimension 4 couplings to both F_i and to H. Integrating out these heavy fermions gives higher dimension effective couplings of the SM fermions to H that replace the role of Yukawa couplings in the SM. These couplings contain explicit flavor breaking in the form of $\langle F_i \rangle / M$, which we take to be of order 1, as well as being suppressed by powers of $S^\dagger S / M^2$, whose vev we take to be of order $1/50$.

In our framework all of the observed SM fermion mass hierarchies are generated from powers of $\langle S \rangle / M \sim 1/7$, which is essentially the ratio of the electroweak scale to the TeV scale, often called the "little hierarchy." We can be agnostic about the source of the little hierarchy itself, since many possibilities have been proposed. The additional challenge of our framework is to achieve simultaneously the appropriate flavon physics at the TeV scale.

Models in our framework have, in addition to the SM particle content, a light singlet scalar s that mixes with the Higgs boson h. Exchanges of s between SM fermions are a new source of FCNC. There is an extra Z' at the electroweak (EW) scale, but apart from small $Z - Z'$ mixing effects it does not couple to SM fermions. There may be other Z's and one or more flavon scalars at the TeV scale. We predict a host of new heavy fermions around the TeV scale; these are also a source of new FCNC and CLFV effects. We show that flavon charge patterns that reproduce the observed SM fermion masses and mixings also supply enough extra suppression of FCNC and CLFV effects to satisfy current experimental bounds.

In addition to explaining the hierarchy of fermion masses and mixings, models in our framework have many interesting phenomenological implications. Mixing of the singlet s with the Higgs boson h can cause large deviations from the SM predictions for the Higgs decay branching fractions, potentially observable at the Tevatron or LHC. The s mass eigenstate itself will also be produced at the LHC, and could be confused with h if it turns out to be the lightest mass eigenstate. While new FCNC effects are suppressed, we predict contributions to $D^0 - \bar{D}^0$ mixing, $B_s \to \mu^+ \mu^-$, that are close to the current value or limit. The exotic top quark decays $t \to ch$ and $t \to cs$ can have branching fractions on the order of 10^{-3}.

Our paper is organized as follows. In Sec. II, we present the basic outline of our framework. In Sec. III, we discuss the constraints on the model parameters from the low energy phenomenology. Section IV contains the phenome-

nological implications and predictions of the model, especially for the new top decays and Higgs signals at the Tevatron and LHC. In Sec. V, we outline a possible ultraviolet completion realizing our proposal. Section VI contains our conclusions and further discussion.

II. MODEL AND FORMALISM

We extend the gauge symmetry of the SM by a $U(1)_S$ local symmetry and an additional local flavon symmetry which in the simplest case would be a $U(1)_F$. All of the SM fermions are neutral with respect to $U(1)_S$, while all of the SM fermions apart from the third generation quark doublet q_{3L} and right-handed top u_{3R} are charged under the chiral $U(1)_F$. We introduce a complex scalar field S which has charge 1 under $U(1)_S$, is neutral under the flavon symmetry, and is a SM singlet. We also introduce one or more complex scalar fields F_i, the flavon scalars. In the simplest case there would be a single flavon scalar F that has charge 1 under $U(1)_F$, is neutral under $U(1)_S$, and is a SM singlet. The Higgs field H is taken as neutral under $U(1)_S \times U(1)_F$. We assume that the flavon charges of the SM fermions are such that only the top quark has an allowed dimension 4 Yukawa interaction.

The S field is assumed to develop a vev that spontaneously breaks the $U(1)_S$ symmetry. In frameworks where the little hierarchy between the electroweak scale and the TeV scale is generated, this could occur naturally by extending the Higgs sector to include S with a mixed potential. The pseudoscalar component of S is then "eaten" to give mass to the $U(1)_S$ Z' gauge boson. Notice that the vev of S does not in itself break any of the global flavor symmetries of the Yukawa-less SM; S is only a messenger of flavor breaking, just as it is also a messenger of electroweak breaking. This is the fundamental distinction that allows S to exist at the electroweak scale without inducing unacceptably large flavor violating effects.

The flavon scalars F_i are assumed to develop vevs that spontaneously break the local flavon symmetry at the TeV scale, with the pseudoscalar components of the F_i eaten to give the flavon gauge bosons mass. To preserve the little hierarchy, we assume that the direct mixing between the F_i and the extended Higgs sector is negligible.

In this framework the Yukawa interactions of the lighter quarks and leptons are replaced by higher dimension operators that couple these fermions to H, S, and the F_i. As we will show later in an explicit example, these can be generated as effective couplings by integrating out new heavy fermions at the TeV scale. These effective couplings should respect all of the SM gauge symmetries, as well as $U(1)_S$ and the flavon symmetries. In particular, the $U(1)_S$ charged field S can only appear as powers of $S^\dagger S / M^2$, where M denotes a generic TeV scale parameter. Powers of F_i / M and $F_i^\dagger / M$ can also appear, but the exact form depends on the flavon charge assignments of the SM fermions. Since we will assume that vevs of the F_i are of

order M, we can absorb the F_i/M dependence into the dimensionless complex couplings h_{ij}, where i, j are generation labels; all these couplings we will then take to be of order 1.

$$
\begin{aligned}
\mathcal{L}^{\text{Yuk}} = &\, h_{33}^u \bar{q}_{3L} u_{3R} \bar{H} + \left(\frac{S^\dagger S}{M^2}\right)(h_{33}^d \bar{q}_{3L} d_{3R} H + h_{22}^u \bar{q}_{2L} u_{2R} \bar{H} + h_{23}^u \bar{q}_{2L} u_{3R} \bar{H} + h_{32}^u \bar{q}_{3L} u_{2R} \bar{H}) \\
&+ \left(\frac{S^\dagger S}{M^2}\right)^2 (h_{22}^d \bar{q}_{2L} d_{2R} H + h_{23}^d \bar{q}_{2L} d_{3R} H + h_{32}^d \bar{q}_{3L} d_{2R} H + h_{12}^u \bar{q}_{1L} u_{2R} \bar{H} + h_{21}^u \bar{q}_{2L} u_{1R} \bar{H} + h_{13}^u \bar{q}_{1L} u_{3R} \bar{H} \\
&+ h_{31}^u \bar{q}_{3L} u_{1R} \bar{H}) + \left(\frac{S^\dagger S}{M^2}\right)^3 (h_{11}^u \bar{q}_{1L} u_{1R} \bar{H} + h_{11}^d \bar{q}_{1L} d_{1R} H + h_{12}^d \bar{q}_{1L} d_{2R} H + h_{21}^d \bar{q}_{2L} d_{1R} H + h_{13}^d \bar{q}_{1L} d_{3R} H \\
&+ h_{31}^d \bar{q}_{3L} d_{1R} H) + \text{H.c.} \tag{1}
\end{aligned}
$$

Note that the above interactions are very similar to those proposed in Ref. [6], except our interactions involve suppression by powers of $(\frac{S^\dagger S}{M^2})$, instead of $(\frac{H^\dagger H}{M^2})$. We will refer to this as the Babu-Nandi texture. The hierarchies among the fermion masses and mixings are obtained from a single small dimensionless parameter,

$$
\epsilon \equiv \frac{v_s}{M}, \tag{2}
$$

where v_s is the vev of S. As was shown in [6], a good fit to the observed fermion masses and mixings is obtained with $\epsilon \sim 0.15$. The couplings h_{ij} are all of order 1; the largest coupling needed is $h_{23}^u = 1.4$, while the smallest coupling needed is $h_{22}^u = 0.14$.

The Babu-Nandi texture is not unique, and it does not predict any precise fermion mass relations, since there are slightly more unspecified order 1 parameters than there are Yukawa parameters in the SM.

A. Fermion masses and Cabibbo-Kobayashi-Maskawa mixing

The gauge symmetry of our model is the usual $SU(3)_c \times SU(2)_L \times U(1)_Y$ of the SM, plus two additional local symmetries: $U(1)_S$ and the flavon symmetry. The SM symmetry is broken spontaneously by the usual Higgs doublet H at the electroweak scale. We assume that the extra $U(1)_S$ symmetry is also broken spontaneously at the electroweak scale by a SM singlet complex scalar field S. The flavon symmetry, $U(1)_F$ in the simplest case, is broken spontaneously above a TeV by a SM singlet scalar flavon field F. The pseudoscalar part of the complex scalar field S is absorbed by the Z' gauge boson $U(1)_S$ to get its mass. Thus after symmetry breaking the remaining scalars at the electroweak scale are neutral bosons h and s. Parameterizing the Higgs doublet and singlet in the unitary gauge as

$$
H = \begin{pmatrix} 0 \\ \frac{h}{\sqrt{2}} + v \end{pmatrix} \qquad S = \left(\frac{s}{\sqrt{2}} + v_s\right), \tag{3}
$$

The observed SM fermion mass hierarchy is generated from the following low energy effective interactions:

with $v \simeq 174$ GeV, and defining an additional small parameter

$$
\beta \equiv \frac{v}{M}, \tag{4}
$$

we obtain, from Eqs. (1)–(4) the following mass matrices for the up and down quark sector:

$$
\begin{aligned}
M_u &= \begin{pmatrix} h_{11}^u \epsilon^6 & h_{12}^u \epsilon^4 & h_{13}^u \epsilon^4 \\ h_{21}^u \epsilon^4 & h_{22}^u \epsilon^2 & h_{23}^u \epsilon^2 \\ h_{31}^u \epsilon^4 & h_{32}^u \epsilon^2 & h_{33}^u \end{pmatrix} v, \\[6pt]
M_d &= \begin{pmatrix} h_{11}^d \epsilon^6 & h_{12}^d \epsilon^6 & h_{13}^d \epsilon^6 \\ h_{21}^d \epsilon^6 & h_{22}^d \epsilon^4 & h_{23}^d \epsilon^4 \\ h_{31}^d \epsilon^6 & h_{32}^d \epsilon^4 & h_{33}^d \epsilon^2 \end{pmatrix} v.
\end{aligned} \tag{5}
$$

The charged lepton mass matrix is obtained from M_d by replacing the couplings h_{ij} appropriately. Note that these mass matrices are the same as in [6], and as was shown there, good fits to the quark and charged lepton masses, as well as the Cabibbo-Kobayashi-Maskawa mixing angles are obtained by choosing $\epsilon \sim 0.15$, and all the couplings h_{ij} of order 1. To leading order in ϵ, the fermion masses are given by

$$
\begin{aligned}
(m_t, m_c, m_u) &\simeq (|h_{33}^u|, |h_{22}^u|\epsilon^2, |h_{11}^u - h_{12}^u h_{21}^u/h_{22}^u|\epsilon^6)v, \\
(m_b, m_s, m_d) &\simeq (|h_{33}^d|\epsilon^2, |h_{22}^d|\epsilon^4, |h_{11}^d|\epsilon^6)v, \\
(m_\tau, m_\mu, m_e) &\simeq (|h_{33}^\ell|\epsilon^2, |h_{22}^\ell|\epsilon^4, |h_{11}^\ell|\epsilon^6)v,
\end{aligned} \tag{6}
$$

while the quark mixing angles are

$$
\begin{aligned}
|V_{us}| &\simeq \left| \frac{h_{12}^d}{h_{22}^d} - \frac{h_{12}^u}{h_{22}^u} \right| \epsilon^2, \\[6pt]
|V_{cb}| &\simeq \left| \frac{h_{23}^d}{h_{33}^d} - \frac{h_{23}^u}{h_{33}^u} \right| \epsilon^2, \\[6pt]
|V_{ub}| &\simeq \left| \frac{h_{13}^d}{h_{33}^d} - \frac{h_{12}^u h_{23}^d}{h_{22}^u h_{33}^d} - \frac{h_{13}^u}{h_{33}^u} \right| \epsilon^4.
\end{aligned} \tag{7}
$$

Generically, all of the h_{ij} can be nonvanishing, but in a particular ultraviolet completion flavon charge conserva-

tion may push some of them to higher order in ϵ or to vanish altogether. However from (6) and (7) we see that the Babu-Nandi texture is rather robust: the only flavor off-diagonal couplings needed to reproduce the observed mixings are one or more of h_{12}^d, h_{12}^u, one or more of h_{23}^d, h_{23}^u, and one or more of h_{13}^d, h_{13}^u; the rest can either vanish or appear at higher order in ϵ.

B. Yukawa interactions and FCNC

Our model has flavor-changing neutral current interactions in the Yukawa sector. Using Eqs. (1)–(4), the Yukawa interaction matrices Y_u^h, Y_d^h, Y_u^s, Y_d^s for the up and down sector, for h^0 and s^0 fields are obtained to be

$$\sqrt{2}Y_u^h = \begin{pmatrix} h_{11}^u \epsilon^6 & h_{12}^u \epsilon^4 & h_{13}^u \epsilon^4 \\ h_{21}^u \epsilon^4 & h_{22}^u \epsilon^2 & h_{23}^u \epsilon^2 \\ h_{31}^u \epsilon^4 & h_{32}^u \epsilon^2 & h_{33}^u \end{pmatrix},$$

$$\sqrt{2}Y_d^h = \begin{pmatrix} h_{11}^d \epsilon^6 & h_{12}^d \epsilon^6 & h_{13}^d \epsilon^6 \\ h_{21}^d \epsilon^6 & h_{22}^d \epsilon^4 & h_{23}^d \epsilon^4 \\ h_{31}^d \epsilon^6 & h_{32}^d \epsilon^4 & h_{33}^d \epsilon^2 \end{pmatrix},$$

(8)

with the charged lepton Yukawa coupling matrix Y_ℓ obtained from Y_d by replacing $h_{ij}^d \rightarrow h_{ij}^\ell$.

$$\sqrt{2}Y_u^s = \begin{pmatrix} 6h_{11}^u \epsilon^5 \beta & 4h_{12}^u \epsilon^3 \beta & 4h_{13}^u \epsilon^3 \beta \\ 4h_{21}^u \epsilon^3 \beta & 2h_{22}^u \epsilon \beta & 2h_{23}^u \epsilon \beta \\ 4h_{31}^u \epsilon^3 \beta & 2h_{32}^u \epsilon \beta & 0 \end{pmatrix},$$

$$\sqrt{2}Y_d^s = \begin{pmatrix} 6h_{11}^d \epsilon^5 \beta & 6h_{12}^d \epsilon^5 \beta & 6h_{13}^d \epsilon^5 \beta \\ 6h_{21}^d \epsilon^5 \beta & 4h_{22}^d \epsilon^3 \beta & 4h_{23}^d \epsilon^3 \beta \\ 6h_{31}^d \epsilon^5 \beta & 4h_{32}^d \epsilon^3 \beta & 2h_{33}^d \epsilon \beta \end{pmatrix},$$

(9)

with the charged lepton Yukawa coupling matrix Y_ℓ obtained from Y_d by replacing $h_{ij}^d \rightarrow h_{ij}^\ell$.

There are several important features that distinguish our model from the proposals in [6–8]:

(i) Note, from Eqs. (5) and (8), in our model, the Yukawa couplings of h to the SM fermions are exactly the same as in the SM. This is because the fermion mass hierarchy in our model is arising from $(\frac{S^\dagger S}{M^2})$. This is a distinguishing feature of our model from that proposed in [6,7] where the Yukawa couplings of h are flavor-dependent, because the hierarchy there arises from $(\frac{H^\dagger H}{M^2})$.

(ii) In our model, we have an additional singlet Higgs boson whose couplings to the SM fermions are flavor-dependent as given in Eq. (9). Again, this is because the hierarchy in our model arises from $(\frac{S^\dagger S}{M^2})$. In particular, s^0 does not couple to the top quark, and its dominant fermionic coupling is to the bottom quark. This will have interesting phenomenological implications for the Higgs searches at the LHC.

(iii) We note from Eqs. (5)–(8) that the mass matrices and the corresponding Yukawa coupling matrices for h are proportional as in the SM. Thus there are

no flavor-changing Yukawa interactions mediated by h. However, this is not true for the Yukawa interactions of the singlet Higgs as can be seen from Eqs. (5) and (9). Thus the s exchange will lead to flavor violation in the neutral Higgs interactions.

C. Higgs sector and the Z'

The Higgs potential of our model, consistent with the SM and the extra $U(1)_S$ symmetry, can be written as

$$V(H, S) = -\mu_H^2(H^\dagger H) - \mu_S^2(S^\dagger S) + \lambda_H(H^\dagger H)^2 + \lambda_S(S^\dagger S)^2 + \lambda_{HS}(H^\dagger H)(S^\dagger S). \quad (10)$$

Note that after absorbing the three components of H in $W^\pm$ and Z, and the pseudoscalar component of S in Z', we are left with only two scalar Higgs, h^0 and s^0. The squared mass matrix in the (h^0, s^0) basis is given by

$$\mathcal{M}^2 = 2v^2 \begin{pmatrix} 2\lambda_H & \lambda_{HS}\alpha \\ \lambda_{HS}\alpha & 2\lambda_S\alpha^2 \end{pmatrix}, \quad (11)$$

where $\alpha = v_s/v$.

The mass eigenstates h and s can be written as

$$h^0 = h\cos\theta + s\sin\theta, \qquad s^0 = -h\sin\theta + s\cos\theta, \quad (12)$$

where θ is the mixing angle in the Higgs sector.

In the Yukawa interactions discussed above, as well as in the gauge interactions involving the Higgs fields, the fields appearing are h^0 and s^0, and these can be expressed in terms of h and s using Eq. (12).

The mass of the Z' gauge boson is given by

$$m_{Z'}^2 = 2g_E^2 v_s^2. \quad (13)$$

Note that the Z' does not couple to any SM particles directly. The Z' coupling to the SM particles will be only via dimension 6 or higher operators. Such couplings will be generated by the vectorlike fermions in the model to be discussed in Sec. V.

III. PHENOMENOLOGICAL IMPLICATIONS: CONSTRAINTS FROM EXISTING DATA

In this section, we discuss the constraints on our model from the existing experimental results. As can be seen from Eq. (9), the exchange of s gives rise to tree level FCNC processes. This will cause $K^0 - \bar{K}^0$ mass splitting, $D^0 - \bar{D}^0$ mixing, $K_L \rightarrow \mu^+\mu^-$, $B_s^0 \rightarrow \mu^+\mu^-$, as well as contributions to the electric dipole moment of neutron and electron, and other rare processes that we discuss below.

A. $K^0 - \bar{K}^0$ mixing

In our model, this arises from the tree level s exchange between $d\bar{s}$ and $\bar{s}d$, and is proportional to $\beta^2\epsilon^{10}$. Taking $\beta \sim \epsilon \sim 0.15$, and the values of the couplings h_{12}^d and h_{21}^d

to be of order 1, the contribution to $\Delta m_K^{\text{Higgs}} \simeq 10^{-16}$ to 10^{-17} GeV, for an s mass of 100 GeV. The experimental value of Δm_K is 3.5×10^{-15} GeV [9]. Thus, since the contribution goes like m_s^{-4}, s can be lighter than 100 GeV. Note that $\epsilon = \upsilon_s/M$ is fixed to be ~ 0.15 to explain fermion mass hierarchy and the Cabibbo-Kobayashi-Maskawa mixing. However, $\beta = \upsilon/M$ is a parameter in our model. Although the Δm_K constraint allows a somewhat larger value of β, we shall see that $D^0 - \bar{D}^0$ mixing constrains $\beta \sim \epsilon$.

B. $D^0 - \bar{D}^0$ mixing

This contribution is again due to the tree level s exchange between $u\bar{c}$ and $\bar{u}c$, and is proportional to $\beta^2 \epsilon^6$, and hence is enhanced compared to Δm_K. Again, taking the couplings h_{12}^u and h_{21}^u to be of order 1 and $\beta \sim \epsilon$, we get $\Delta m_D \sim 10^{-14}$ GeV for $m_s = 100$ GeV. This is to be compared with the current experimental value of 1.6×10^{-14} GeV [9,10]. Thus Δm_D gives a much stronger restriction on the model parameters. β can not be much larger than ϵ, and s can not be much lighter than 100 GeV. If our proposal is correct, an electroweak singlet scalar should be observed at the LHC.

C. Other rare processes

In our model, the tree level s exchange between $d\bar{s}$ and $\mu^+\mu^-$ will contribute to $K_L \to \mu^+\mu^-$. This contribution is proportional to $\beta^2\epsilon^{10}$, and leads to a contribution to this branching ratio $\sim 10^{-14}$ for $\beta \sim \epsilon$ and $m_s \sim 100$ GeV. This is very small compared to the current experimental value of $\sim 6.9 \times 10^{-9}$ [9]. Similarly, the contribution to the other rare processes such as $K_L \to \mu e$, $K \to \pi \bar{\nu}\nu$, $\mu \to e\gamma$, $\mu \to 3e$, $B_d - \bar{B}_d$ mixing, etc. are several orders of magnitude below the corresponding experimental limits.

D. Constraint on the mass of s

Experiments at LEP2 have set a lower limit of 114.4 GeV for the mass of the SM Higgs boson, from nonobservation of the associated production $e^+e^- \to Zh$. In our model, since the singlet Higgs can mix with the doublet h, there will be a limit for m_s depending on the value of the mixing angle, θ. For $\sin^2\theta \gtrsim 0.25$, the bound of 114.4 applies also for m_s [11]. However, s can be lighter if the mixing is small.

E. Constraint on the mass of the Z'

We have assumed that the extra $U(1)$ symmetry in our model is spontaneously broken at the EW scale. But the corresponding gauge coupling, g_E, is arbitrary and hence the mass of Z' is not determined in our model. However, very accurately measured Z properties at LEP1 put a constraint on the $Z - Z'$ mixing to be $\sim 10^{-3}$ or smaller [9,12]. In our model, the Z' does not couple to any SM particle

directly. $Z - Z'$ mixing can take place at the one loop level with the new vectorlike fermions in the loop. The mixing angle is

$$\theta_{ZZ'} \sim \frac{g_Z g_E}{16\pi^2}\left(\frac{m_Z}{M}\right)^2, \qquad (14)$$

where M is the mass of the vectorlike fermions with masses in the TeV scale. Even with $g_E \sim 1$, we get $\theta_{ZZ'} \sim 10^{-4}$ or less. Thus there is no significant bound for the mass of this Z' from the LEP1. This Z' can couple to the SM particles via dimension 6 operators with the interaction of the form

$$L = \frac{\bar{\psi}_L \sigma^{\mu\nu} \psi_R H Z'^{\mu\nu}}{M^2}. \qquad (15)$$

As was shown in [13], no significant bound on $m_{Z'}$ emerges from these interactions.

IV. PHENOMENOLOGICAL IMPLICATIONS: NEW PHYSICS SIGNALS

Motivated to explain the observed mass hierarchy in the fermion sector, we have constructed a model which has a complex singlet Higgs (in addition to the usual doublet), a new $U(1)_S$ gauge symmetry at the EW scale, and a new set of vectorlike fermions at the TeV scale. Thus our model has new particles such as a scalar Higgs and a new Z' boson at the EW scale, and heavy vectorlike quarks and leptons. The model has many phenomenological implications for the production and decays of the Higgs bosons, top quark physics, a new scenario for Z' physics, and the production and decays of the vectorlike fermions.

A. Higgs signals

1. Higgs coupling to the SM fermions

As can be seen from (8), the couplings of the doublet Higgs h to the SM fermions are identical to that in the SM, whereas the couplings of the singlet Higgs have a different flavor dependence. In particular, the singlet Higgs s does not couple to the top quark, whereas its couplings to $(b, \tau; c, s, \mu; u, d, e)$ involve the flavor-dependent factors $(2,2;2,4,4;6,6,6)$, respectively, in the limit of zero mixing between h and s. Including the mixing, these factors will be modified. Thus our model will be distinguished from the SM by the fact that the Higgs has nonstandard couplings to fermions predicted in terms of two model parameters: the ratio of vevs α and the mixing angle θ.

2. Higgs decays

The couplings of the Higgs bosons h and s to the fermions and the gauge bosons can be obtained from Eqs. (8) and (9), and are given in Table I.

TABLE I.   Yukawa and gauge couplings of h and s.

Interaction	Coupling	Interaction	Coupling
$s \to u\bar{u}$	$\frac{m_u}{v\sqrt{2}}(\sin\theta + \frac{6\cos\theta}{\alpha})$	$h \to u\bar{u}$	$\frac{m_u}{v\sqrt{2}}(\cos\theta - \frac{6\sin\theta}{\alpha})$
$s \to d\bar{d}$	$\frac{m_d}{v\sqrt{2}}(\sin\theta + \frac{6\cos\theta}{\alpha})$	$h \to d\bar{d}$	$\frac{m_d}{v\sqrt{2}}(\cos\theta - \frac{6\sin\theta}{\alpha})$
$s \to \mu^+\mu^-$	$\frac{m_\mu}{v\sqrt{2}}(\sin\theta + \frac{4\cos\theta}{\alpha})$	$h \to \mu^+\mu^-$	$\frac{m_\mu}{v\sqrt{2}}(\cos\theta - \frac{4\sin\theta}{\alpha})$
$s \to s\bar{s}$	$\frac{m_s}{v\sqrt{2}}(\sin\theta + \frac{4\cos\theta}{\alpha})$	$h \to s\bar{s}$	$\frac{m_s}{v\sqrt{2}}(\cos\theta - \frac{4\sin\theta}{\alpha})$
$s \to \tau^+\tau^-$	$\frac{m_\tau}{v\sqrt{2}}(\sin\theta + \frac{2\cos\theta}{\alpha})$	$h \to \tau^+\tau^-$	$\frac{m_\tau}{v\sqrt{2}}(\cos\theta - \frac{2\sin\theta}{\alpha})$
$s \to c\bar{c}$	$\frac{m_c}{v\sqrt{2}}(\sin\theta + \frac{2\cos\theta}{\alpha})$	$h \to c\bar{c}$	$\frac{m_c}{v\sqrt{2}}(\cos\theta - \frac{2\sin\theta}{\alpha})$
$s \to b\bar{b}$	$\frac{m_b}{v\sqrt{2}}(\sin\theta + \frac{2\cos\theta}{\alpha})$	$h \to b\bar{b}$	$\frac{m_b}{v\sqrt{2}}(\cos\theta - \frac{2\sin\theta}{\alpha})$
$s \to t\bar{t}$	$\frac{m_t}{v\sqrt{2}}\sin\theta$	$h \to t\bar{t}$	$\frac{m_t}{v\sqrt{2}}\cos\theta$
$s \to ZZ$	$\frac{2m_Z^2}{v\sqrt{2}}\sin\theta$	$h \to ZZ$	$\frac{2m_Z^2}{v\sqrt{2}}\cos\theta$
$s \to Z'Z'$	$\frac{m_{Z'}^2}{v\alpha\sqrt{2}}\cos\theta$	$h \to Z'Z'$	$\frac{m_{Z'}^2}{v\alpha\sqrt{2}}\sin\theta$
$s \to W^+W^-$	$\frac{2m_W^2}{v\sqrt{2}}\sin\theta$	$h \to W^+W^-$	$\frac{2m_W^2}{v\sqrt{2}}\cos\theta$
		$h \to ss$	λ_{hss}

The coupling of h to s given by

$$\lambda_{hss} = \frac{m_h^2}{4v}\{(1-\mu)\sin2\theta[\cos^3\theta - \alpha\sin^3\theta \\ + \sin2\theta(\alpha\cos\theta - \sin\theta)] \\ + 3\sin2\theta[\sin\theta(1+\mu - (1-\mu)\cos2\theta) \\ - \cos\theta(1+\mu - (1-\mu)\cos2\theta)/\alpha]\},$$

where $\mu = m_s^2/m_h^2$.

Because of the flavor dependence of the couplings, the branching ratios (BR) for h to various final states are altered substantially from those in the SM. These branching ratios for h to various final states are shown in Figs. 1–4 for values of the mixing angle $\theta = 0$, 20°, 26°, and 40°, respectively.

For $\theta = 0$, i.e. no mixing, these BR's are the same as for the SM Higgs. Note that for both $\theta = 20°$ and 26°, the gg and the $\gamma\gamma$ BR's are enhanced substantially compared to

FIG. 2 (color online).   Branching ratio of $h \to 2x$, for $\theta = 20°$ and $\alpha = 1$.

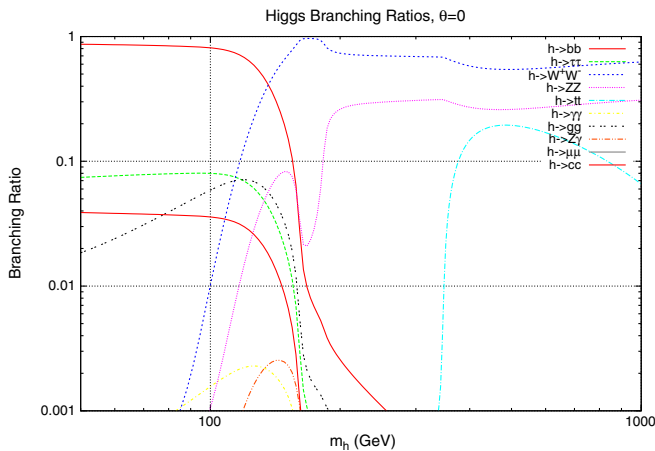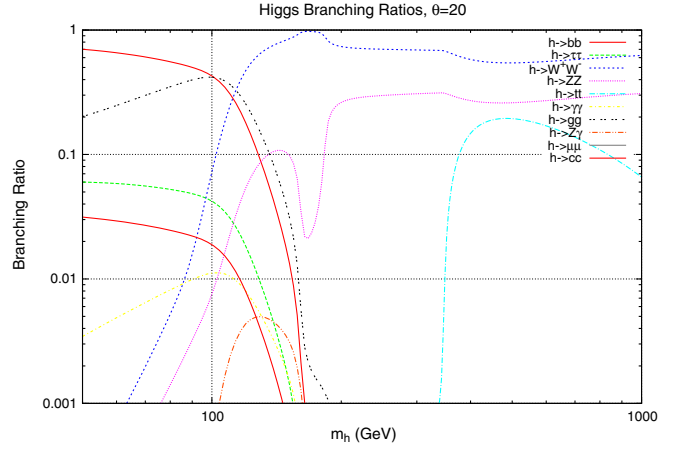FIG. 3 (color online).   Branching ratio of $h \to 2x$, for $\theta = 26°$ and $\alpha = 1$.

FIG. 1 (color online).   Branching ratio of $h \to 2x$, for $\theta = 0$ and $\alpha = 1$ [18].

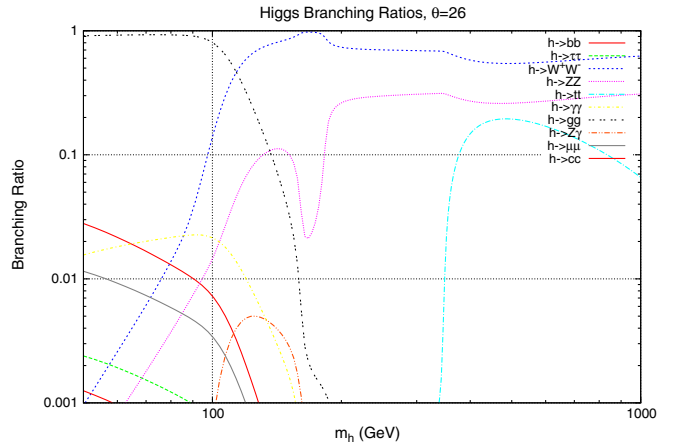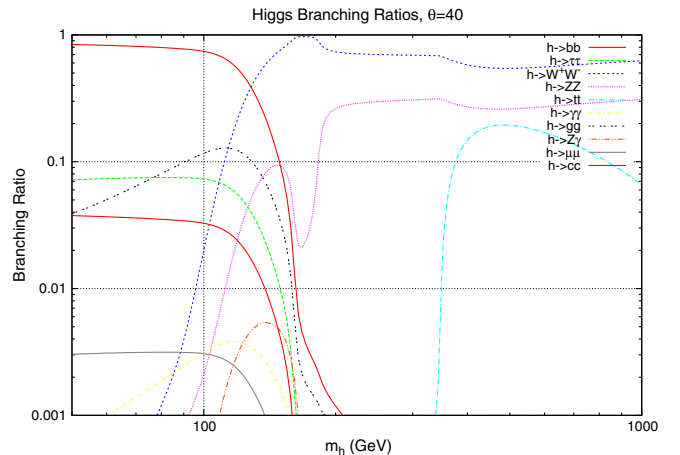FIG. 4 (color online).   Branching ratio of $h \to 2x$, for $\theta = 40°$ and $\alpha = 1$.

the SM. This is due to drastic reduction for the $b\bar{b}$ mode from an approximate cancellation in the corresponding coupling as can be seen from Table I. In particular, for $\theta = 26°$, the effect is quite dramatic. For a light Higgs (m_h around 115 GeV), the usually dominant $b\bar{b}$ mode is highly suppressed and the $\gamma\gamma$ mode is enhanced by a factor of almost 10 compared to the SM. This is to be contrasted with the proposal of Refs. [6,7] in which the $h \to \gamma\gamma$ mode is reduced by about a factor of 10. Thus the Higgs signal in this mode for a Higgs mass of ~114–140 GeV gets a big enhancement, making its potential discovery via this mode much more favorable at the LHC. Such a signal may be observable at the Tevatron for a Higgs mass ~114 as the luminosity accumulates, but would require about 10 fb^{-1} or more of data [14].

Another interesting effect is the Higgs signal via the WW^* for the light Higgs. In the SM, this mode becomes important for the Tevatron search for Higgs masses greater than about 135 GeV, where the BR to WW^* is approximately equal to that of $b\bar{b}$. Currently Tevatron experiments have excluded a SM Higgs with mass around 170 GeV (where the BR to WW^* is around 100%) for this mode [15]. In our framework, for $\theta = 20°$ for example, the crossover between the WW^* mode and the $b\bar{b}$ mode takes place sooner than 135 GeV. Thus the Tevatron experiments will be more sensitive to the lower mass range than for a SM Higgs, and should be able to exclude masses much smaller than 160 GeV.

For a heavy Higgs, $m_h > 200$ GeV, the Higgs will be accessible via the golden mode $h \to ZZ$. However, in this case, both h and s decay via this mode with comparable BR's (see Figs. 2 and 4 for $\theta = 20°$ and $40°$). So initially it will be hard to tell whether we are seeing h or s, a case of Higgs look-alikes. An accurate measurement of this cross section times the BR, and the mass of the observed Higgs, we will be able to distinguish a heavy h from a heavy s, since the production cross sections depend on the mixing angle.

B. Top quark physics

In the SM, the $t \to ch$ mode is severely suppressed with a BR $\sim 10^{-14}$ [16]. In our model, as can be seen from Eqs. (8) and (9), although $t \to ch$ is zero at tree level, we have a large coupling for $t \to cs \sim 2\epsilon\beta$ (note s here denotes the singlet Higgs, not the strange quark). This gives rise to a significant BR for the $t \to cs$ mode for a Higgs mass of up to about 150 GeV. If the mixing between the h and s is substantial, both decay modes, $t \to cs$ and $t \to ch$ will have BR $\sim 10^{-3}$. With a very large $t\bar{t}$ cross section, $\sigma_{t\bar{t}} \sim 10^3$ pb at the LHC, this could be an observable production mode for Higgs bosons at the LHC.

C. Z' physics

Our model has a Z' boson near the EW scale from the spontaneous breaking of the extra $U(1)$ symmetry. As
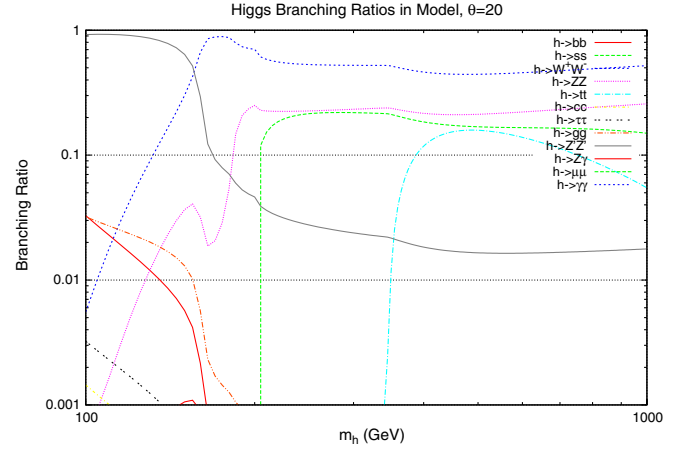

FIG. 5 (color online). Branching ratio of $h \to 2x$ including $h \to ss$ and $h \to Z'Z'$ where $m_{Z'} = 40$ GeV and $m_s = 100$ GeV. Here $\theta = 20°$ and $\alpha = 1$.
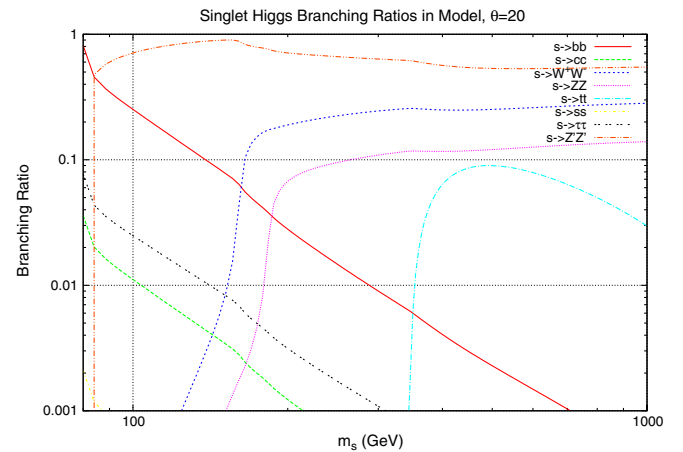

FIG. 6 (color online). Branching ratio of $s \to 2x$ including $s \to Z'Z'$ where $m_{Z'} = 40$ GeV. Here $\theta = 20°$ and $\alpha = 1$.

discussed before, since the $Z - Z'$ mixing is very small, $\sim 10^{-4}$ or less, its mass is not constrained by the very accurately measured Z properties at LEP. Its mass can be as low as a few GeV from the existing constraints. This Z' does not couple to SM particles with dimension 4 operators. It does couple to s at tree level via the $sZ'Z'$ interaction. Thus it can be produced via the decay of s (or h if there is a substantial mixing between h and s). This gives an interesting signal for the Higgs decays: $s \to Z'Z'$, $h \to Z'Z'$ if allowed kinematically. In Figs. 5 and 6, we give the BR's for h and s decays for a Z' mass of 40 GeV. The Z' will decay to the SM particles via the $Z - Z'$ mixing with the same branching ratio as the Z. Thus the clear final state signal will be $l^+l^-l^+l^-$ pairs ($l = e, \mu$) with each pair having the invariant mass of the Z'. Such a signal will be easily detectable at the LHC. If the Z' happens to be very light, (say a few GeV), and the mixing angle is extremely tiny, there is a possibility that the Z''s may produce dis-

placed vertices at the detector. Both of these will be very unconventional signals for Higgs bosons at the LHC.

D. $B_s^0 \to \mu^+ \mu^-$

In our framework this decay gets a contribution from an FCNC interaction mediated by the s exchange. The amplitude for this decay is $A \sim 4h_{22}^d h_{22}^\ell \epsilon^6 \beta^2$. Taking $\beta \sim \epsilon$, $A \sim 4h_{22}^d h_{22}^\ell \epsilon^8$, and with the couplings h_{22}^d, $h_{22}^\ell \sim 1$, we obtain the branching ratio, $BR(B_s^0 \to \mu^+ \mu^-) \sim 10^{-9}$. The current experimental limit for this BR is 4.7×10^{-8} [9], and thus there is a possibility that this decay could be observed at the Tevatron.

E. Production and decay of heavy fermions

Our framework requires vectorlike quarks and leptons, both $SU(2)$ doublets and weak singlets, with masses around the TeV scale. The heavy quarks can be pair produced at high energy hadron colliders via the strong interaction. For example, for a 1 TeV vectorlike quark, the production cross section at the LHC is ~ 60 fb [17]. We need several such vectorlike quarks for our model. So the total production cross section could be as large as a few hundred fb. These will decay to the light quarks of the same electric charge and Higgs bosons (h or s): Thus the signal will be two high p_T jets together with the final states arising from the Higgs decays. For a heavy Higgs, in the golden mode ($h \to ZZ$, $s \to ZZ$), this will give rise to two high p_T jets plus four Z bosons. In the case of a light Z', the final state signal will be two high p_T jets plus up to eight charged leptons in the final state (with each lepton pair having the invariant mass of the Z').

V. ULTRAVIOLET COMPLETION

We present two concrete examples of models from which an effective action like Eq. (1) can be derived. The first example only reproduces the second and third generation quark couplings, but its simplicity serves to introduce the basic issues and mechanisms. The second example is a complete three generation TeV scale model of quark flavor. The correct lepton couplings can be obtained from a copy of the same structure used for the down-type quarks. We assume that neutrino masses benefit from some additional seesaw mechanism, although it is not obvious that we cannot obtain them by refining the TeV scale flavon model.

A. Two generation model

For this pedagogical example we will employ two important simplifications:

(i) We only reproduce the second and third generation quark couplings. In the next subsection we extend this to include the first generation.

(ii) We will choose charge assignments such that the couplings h_{32}^u, h_{32}^d, and h_{23}^d are higher order in ϵ. As already mentioned nonzero values of these cou-

TABLE II. Charge assignments in the two generation model for the scalar fields H, S, F, and the SM quark fields q_{3L}, q_{2L}, u_{3R}, u_{2R}, d_{3R}, and d_{2R}. Also listed are the color triplet weak doublet heavy quark pairs Q_{iL}, Q_{iR} and the color triplet weak singlet heavy quark pairs U_{iL}, U_{iR}, D_{iL}, D_{iR}.

Field	$U(1)_Y$	$U(1)_S$	$U(1)_F$	Field	$U(1)_Y$	$U(1)_S$	$U(1)_F$
H	1/2	0	0	U_{1L}	2/3	1	0
S	0	1	0	U_{1R}	2/3	1	1
F	0	0	1	U_{2L}	2/3	-1	3
q_{3L}	1/6	0	0	U_{2R}	2/3	-1	3
q_{2L}	1/6	0	2	D_{1L}	$-1/3$	-1	-1
u_{3R}	2/3	0	0	D_{1R}	$-1/3$	-1	-1
u_{2R}	2/3	0	3	D_{2L}	$-1/3$	2	3
d_{3R}	$-1/3$	0	-1	D_{2R}	$-1/3$	2	2
d_{2R}	$-1/3$	0	3	D_{3L}	$-1/3$	1	3
Q_{1L}	1/6	-1	-1	D_{3R}	$-1/3$	1	3
Q_{1R}	1/6	-1	0				
Q_{2L}	1/6	1	1				
Q_{2R}	1/6	1	2				
Q_{3L}	1/6	-1	3				
Q_{3R}	1/6	-1	2				
Q_{4L}	1/6	2	2				
Q_{4R}	1/6	2	1				

plings are not needed to reproduce the observed SM quark masses and mixings.

With these simplifications we postulate a TeV scale model with the field content shown in Table II, where the hypercharges are listed along with the charge assignments under $U(1)_S$ and $U(1)_F$. The Higgs doublet H is the only scalar that carries the hypercharge, while the SM singlet S is the only scalar carrying the $U(1)_S$ charge. The SM singlet flavon F is the only scalar carrying the $U(1)_F$ charge. The SM quarks are neutral under $U(1)_S$. The third generation up-type quark fields also carry no $U(1)_F$ charge, while the other quark fields have flavor-dependent nonzero $U(1)_F$ charges.

We introduce four pairs of new color triplet weak doublet fermion fields Q_{iL}, Q_{iR}, two pairs of color triplet up-type weak singlets U_{iL}, U_{iR}, and three pairs of color triplet down-type weak singlets D_{iL}, D_{iR}. Each pair is vectorlike with respect to the SM gauge group and $U(1)_S$, thus no anomalies are introduced with respect to these gauge groups, and each vectorlike pair naturally acquires a Dirac mass of order M [when they have the same $U(1)_F$ charge] or of order the vev of F (when their $U(1)_F$ charges differ by one). We assume that both the vev of F and M are greater than, but of order of, a TeV. Any residual anomaly in $U(1)_F$ can be handled either by introducing more heavy fermions or using the Green-Schwarz mechanism above the TeV scale.

With these charge assignments the only dimension 4 couplings involving the second and third generation SM quarks are

$$f_1 \bar{q}_{3L} u_{3R} \bar{H} + f_2 \bar{q}_{3L} Q_{1R} S + f_3 \bar{D}_{1L} d_{3R} S^\dagger + f_4 \bar{q}_{2L} Q_{2R} S^\dagger$$

$$+ f_5 \bar{U}_{1L} u_{3R} S + f_6 \bar{q}_{2L} Q_{3R} S + f_7 \bar{U}_{2L} u_{2R} S^\dagger$$

$$+ f_8 \bar{D}_{3L} d_{2R} S + \text{H.c.}, \tag{16}$$

where the f_i are dimensionless coupling constants. Thus the top quark receives the correct mass from electroweak symmetry breaking for $|f_1| \simeq 1$. The other couplings involve the S scalar, but not the Higgs H or the flavon F. Both electroweak symmetry breaking and flavor symmetry breaking are communicated to the rest of the SM quark sector via a Froggart-Nielsen type mechanism, integrating out the heavy TeV scale fermions from tree level diagrams that connect SM quark left doublets to SM quark right singlets and to H or $\bar{H}$.

The renormalizable couplings involving just the heavy fermions are

$$f_9 \bar{Q}_{1R} Q_{1L} F + f_{10} \bar{Q}_{1L} D_{1R} H + M \bar{D}_{1R} D_{1L} + f_{11} \bar{Q}_{2R} Q_{2L} F$$

$$+ f_{12} \bar{Q}_{2L} U_{1R} \bar{H} + f_{13} \bar{U}_{1R} U_{1L} F + f_{14} \bar{Q}_{3R} Q_{3L} F^\dagger$$

$$+ f_{15} \bar{Q}_{3L} U_{2R} \bar{H} + M \bar{U}_{2R} U_{2L} + f_{16} \bar{Q}_{2L} Q_{4R} S^\dagger$$

$$+ f_{17} \bar{Q}_{4L} Q_{2R} S + f_{18} \bar{Q}_{4R} Q_{4L} F^\dagger + f_{19} \bar{Q}_{4L} D_{2R} H$$

$$+ f_{20} \bar{D}_{2R} D_{2L} F^\dagger + f_{21} \bar{D}_{2L} D_{3R} S + M \bar{D}_{3L} D_{3R} + \text{H.c.} \tag{17}$$

Thus, integrating out the heavy fermions in the tree level diagram composed from the couplings

$$f_2 \bar{q}_{3L} Q_{1R} S + f_9 \bar{Q}_{1R} Q_{1L} F + f_{10} \bar{Q}_{1L} D_{1R} H + M \bar{D}_{1R} D_{1L}$$

$$+ f_3 \bar{D}_{1L} d_{3R} S^\dagger \tag{18}$$

produces an effective coupling below the TeV scale proportional to

$$f_2 f_3 f_9 f_{10} \frac{F}{M} \frac{S^\dagger S}{M^2} \bar{q}_{3L} d_{3R} H + \text{H.c.} \tag{19}$$

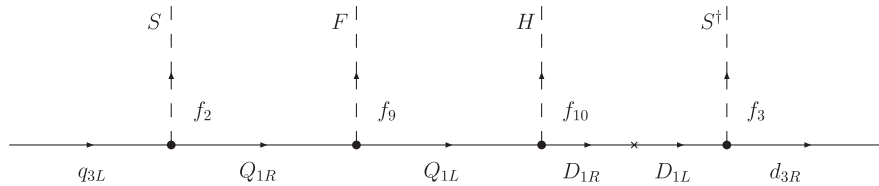Integrating out the heavy fermions in the tree level diagram composed from the couplings

$$f_4 \bar{q}_{2L} Q_{2R} S^\dagger + f_{11} \bar{Q}_{2R} Q_{2L} F + f_{12} \bar{Q}_{2L} U_{1R} \bar{H}$$

$$+ f_{13} \bar{U}_{1R} U_{1L} F + f_5 \bar{U}_{1L} u_{3R} S \tag{20}$$

produces an effective coupling below the TeV scale proportional to

$$f_4 f_5 f_{11} f_{12} f_{13} \frac{F^2}{M^2} \frac{S^\dagger S}{M^2} \bar{q}_{2L} u_{3R} \bar{H} + \text{H.c.} \tag{21}$$

Integrating out the heavy fermions in the tree level diagram composed from the couplings

$$f_6 \bar{q}_{2L} Q_{3R} S + f_{14} \bar{Q}_{3R} Q_{3L} F^\dagger + f_{15} \bar{Q}_{3L} U_{2R} \bar{H}$$

$$+ M \bar{U}_{2R} U_{2L} + f_7 \bar{U}_{2L} u_{2R} S^\dagger \tag{22}$$

produces an effective coupling below the TeV scale proportional to

$$f_6 f_7 f_{14} f_{15} \frac{F^\dagger}{M} \frac{S^\dagger S}{M^2} \bar{q}_{2L} u_{2R} \bar{H} + \text{H.c.} \tag{23}$$

Finally, integrating out the heavy fermions in the tree level diagram composed from the couplings

$$f_4 \bar{q}_{2L} Q_{2R} S^\dagger + f_{17}^* \bar{Q}_{2R} Q_{4L} S^\dagger + f_{19} \bar{Q}_{4L} D_{2R} H$$

$$+ f_{20} \bar{D}_{2R} D_{2L} F^\dagger + f_{21} \bar{D}_{2L} D_{3R} S + M \bar{D}_{3R} D_{3L}$$

$$+ f_8 \bar{D}_{3L} d_{2R} S \tag{24}$$

produces an effective coupling below the TeV scale proportional to

$$f_4 f_8 f_{17}^* f_{19} f_{20} f_{21} \frac{F^\dagger}{M} \frac{(S^\dagger S)^2}{M^4} \bar{q}_{2L} d_{2R} H + \text{H.c.} \tag{25}$$

There is an additional very similar tree level diagram contributing to h_{22}^d composed from the couplings

$$f_4 \bar{q}_{2L} Q_{2R} S^\dagger + f_{11} \bar{Q}_{2R} Q_{2L} F + f_{16} \bar{Q}_{2L} Q_{4R} S^\dagger$$

$$+ f_{18} \bar{Q}_{4R} Q_{4L} F^\dagger + f_{19} \bar{Q}_{4L} D_{2R} H + f_{20} \bar{D}_{2R} D_{2L} F^\dagger$$

$$+ f_{21} \bar{D}_{2L} D_{3R} S + M \bar{D}_{3R} D_{3L} + f_8 \bar{D}_{3L} d_{2R} S, \tag{26}$$

FIG. 7. The Feynman diagram associated with Eq. (18).

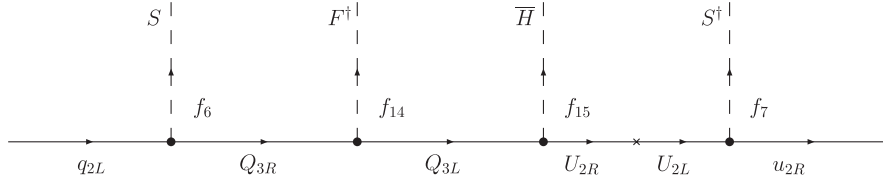FIG. 8. The Feynman diagram associated with Eq. (20).

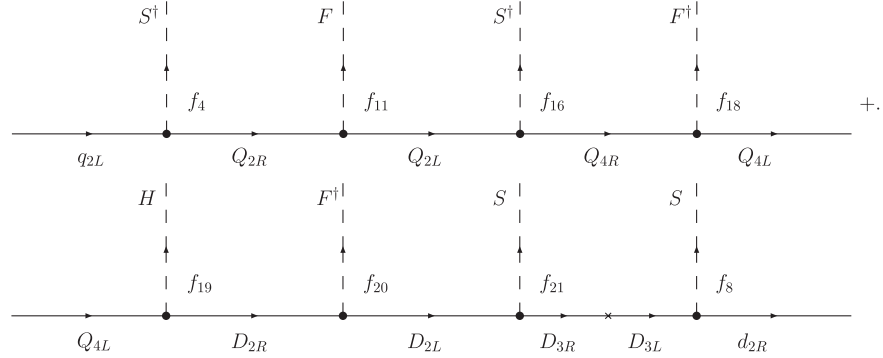FIG. 9.   The Feynman diagram associated with Eq. (22).

FIG. 10.   The Feynman diagram associated with Eq. (26).

which produces an effective coupling below the TeV scale proportional to

$$f_4 f_8 f_{11} f_{16} f_{18} f_{19} f_{20} f_{21} \frac{F(F^\dagger)^2}{M^3} \frac{(S^\dagger S)^2}{M^4} \bar{q}_{2L} d_{2R} H + \text{H.c.}$$

(27)

The above equations can be understood through the Feynman diagrams shown in Figs. 7–10.

B. Three generation model

Here we present a concrete example of a full three generation TeV scale model that reproduces an effective action like Eq. (1) at the electroweak scale. This model uses a single electroweak messenger scalar S, but employs three TeV scale flavon scalars F_1, F_2, and F_3, each corresponding to a different broken $U(1)_{F_i}$ flavon symmetry. As before the SM quarks are neutral under $U(1)_S$. The third generation up-type quark fields also carry no $U(1)_{F_i}$ charges, while the other quark fields have flavor-dependent nonzero $U(1)_{F_i}$ charges.

The model has a rather large number of new heavy fermions: seven pairs of new color triplet weak doublet fermion fields Q_{iL}, Q_{iR}, six pairs of color triplet up-type weak singlets U_{iL}, U_{iR}, and eight pairs of color triplet down-type weak singlets D_{iL}, D_{iR}. Each pair is vectorlike with respect to the SM gauge group and $U(1)_S$, thus no anomalies are introduced with respect to these gauge groups, and each vectorlike pair naturally acquires a Dirac mass of order M [when they have the same $U(1)_{F_i}$ charges] or of order the vev of some F_i [when one of their $U(1)_{F_i}$ charges differs by one]. We assume that both the F_i vevs and M are of order a TeV. Any residual anomaly in the $U(1)_{F_i}$ symmetries can be handled either by introducing more heavy fermions or using the Green-Schwarz mechanism at the TeV scale.

We do not suggest that this model is the most efficient one implementing the basic concepts of our proposal. We have made an explicit tradeoff, in some sense, of maximizing the number of the new heavy fermions required in order to minimize the complexity of the messenger sector and the charge assignments.

With the charge assignments listed in Table III the only dimension 4 couplings of fermions to the Higgs scalar are

$$f_1 \bar{q}_{3L} u_{3R} \bar{H} + f_2 \bar{Q}_{2L} U_{2R} \bar{H} + f_3 \bar{Q}_{4R} U_{4L} \bar{H} + f_4 \bar{Q}_{6L} U_{3R} \bar{H} + f_5 \bar{Q}_{7L} U_{6R} \bar{H} + f_6 \bar{Q}_{3L} D_{3R} H + f_7 \bar{Q}_{4L} D_{4R} H$$
$$+ f_8 \bar{Q}_{7L} D_{7R} H + \text{H.c.}$$

(28)

The only dimension 4 couplings of fermions to the S messenger scalar are

TABLE III. Charge assignments in the three generation model for the scalar fields H, S, F_i, the SM quark fields q_{iL}, u_{iR}, d_{iR}, and the heavy quark pairs Q_{iL}, Q_{iR}, U_{iL}, U_{iR}, D_{iL}, D_{iR}.

Field	$U(1)_Y$	$U(1)_S$	$U(1)_{F1}$	$U(1)_{F2}$	$U(1)_{F3}$	Field	$U(1)_Y$	$U(1)_S$	$U(1)_{F1}$	$U(1)_{F2}$	$U(1)_{F3}$
q_{1L}	1/6	0	1	2	1	U_{1L}	2/3	1	0	1	1
q_{2L}	1/6	0	0	1	0	U_{1R}	2/3	1	0	1	0
q_{3L}	1/6	0	0	0	0	U_{2L}	2/3	1	1	0	0
u_{1R}	2/3	0	0	1	1	U_{2R}	2/3	1	1	1	0
U_{2R}	2/3	0	1	0	0	U_{3L}	2/3	2	2	1	−1
U_{3R}	2/3	0	0	0	0	U_{3R}	2/3	2	2	1	−1
d_{1R}	−1/3	0	1	2	0	U_{4L}	2/3	2	1	1	0
d_{2R}	−1/3	0	0	1	1	U_{4R}	2/3	2	2	1	0
d_{3R}	−1/3	0	1	0	0	U_{5L}	2/3	2	0	1	0
Q_{1L}	1/6	1	1	2	0	U_{5R}	2/3	2	0	2	0
Q_{1R}	1/6	1	1	2	1	U_{6L}	2/3	3	0	2	0
Q_{2L}	1/6	1	1	1	0	U_{6R}	2/3	3	1	2	0
Q_{2R}	1/6	1	0	1	0	D_{1L}	−1/3	1	1	2	0
Q_{3L}	1/6	1	1	0	0	D_{1R}	−1/3	1	1	2	1
Q_{3R}	1/6	1	0	0	0	D_{2L}	−1/3	1	0	1	1
Q_{4L}	1/6	2	1	0	0	D_{2R}	−1/3	1	1	1	1
Q_{4R}	1/6	2	1	1	0	D_{3L}	−1/3	1	1	0	0
Q_{5L}	1/6	2	2	2	0	D_{3R}	−1/3	1	1	0	0
Q_{5R}	1/6	2	1	2	0	D_{4L}	−1/3	2	1	0	0
Q_{6L}	1/6	2	2	1	−1	D_{4R}	−1/3	2	1	0	0
Q_{6R}	1/6	2	2	1	0	D_{5L}	−1/3	2	1	1	1
Q_{7L}	1/6	3	1	2	0	D_{5R}	−1/3	2	1	1	0
Q_{7R}	1/6	3	2	2	0	D_{6L}	−1/3	2	1	2	1
H	1/2	0	0	0	0	D_{6R}	−1/3	2	1	2	0
S	0	1	0	0	0	D_{7L}	−1/3	3	1	1	0
F_1	0	0	1	0	0	D_{7R}	−1/3	3	1	2	0
F_2	0	0	0	1	0	D_{8L}	−1/3	3	1	2	0
F_3	0	0	0	0	1	D_{8R}	−1/3	3	1	2	1

$$f_9\bar{q}_{1L}Q_{1R}S^\dagger + f_{10}\bar{q}_{2L}Q_{2R}S^\dagger + f_{11}\bar{q}_{3L}Q_{3R}S^\dagger + f_{12}\bar{U}_{1L}u_{1R}S + f_{13}\bar{U}_{2L}u_{2R}S + f_{14}\bar{D}_{1L}d_{1R}S + f_{15}\bar{D}_{2L}d_{2R}S$$
$$+ f_{16}\bar{D}_{3L}d_{3R}S + f_{17}\bar{Q}_{2L}Q_{4R}S^\dagger + f_{18}\bar{Q}_{1L}Q_{5R}S^\dagger + f_{19}\bar{Q}_{7L}Q_{5R}S^\dagger + f_{20}\bar{Q}_{5L}Q_{7R}S^\dagger + f_{21}\bar{U}_{4L}U_{2R}S + f_{22}\bar{U}_{5L}U_{1R}S$$
$$+ f_{23}\bar{U}_{6L}U_{5R}S + f_{24}\bar{D}_{4L}D_{3R}S + f_{25}\bar{D}_{3L}D_{4R}S^\dagger + f_{26}\bar{D}_{5L}D_{2R}S + f_{27}\bar{D}_{6L}D_{1R}S + f_{28}\bar{D}_{1L}D_{6R}S^\dagger + f_{29}\bar{D}_{7L}D_{5R}S$$
$$+ f_{30}\bar{D}_{8L}D_{6R}S + f_{31}\bar{D}_{6L}D_{8R}S^\dagger + \text{H.c.} \tag{29}$$

The direct fermion mass terms and mixings consistent with the flavon symmetries and SM gauge symmetries generated by operators of dimension 4 or less are

$$f_{32}\bar{Q}_{1L}Q_{1R}F_3^\dagger + f_{33}\bar{Q}_{2L}Q_{2R}F_1 + f_{34}\bar{Q}_{3L}Q_{3R}F_1 + f_{35}\bar{Q}_{3L}Q_{3R}F_2^\dagger + f_{36}\bar{Q}_{4L}Q_{4R}F_2^\dagger + f_{37}\bar{Q}_{5L}Q_{5R}F_1 + f_{38}\bar{Q}_{5L}Q_{6R}F_2$$
$$+ f_{39}\bar{Q}_{6L}Q_{6R}F_3^\dagger + f_{40}\bar{Q}_{7L}Q_{7R}F_1^\dagger + f_{41}\bar{U}_{1L}U_{1R}F_3 + f_{42}\bar{U}_{2L}U_{2R}F_2^\dagger + M\bar{U}_{3L}U_{3R} + f_{43}\bar{U}_{3L}U_{4R}F_3^\dagger + f_{44}\bar{U}_{4L}U_{4R}F_1^\dagger$$
$$+ f_{45}\bar{U}_{5L}U_{5R}F_2^\dagger + f_{46}\bar{U}_{6L}U_{6R}F_1^\dagger + f_{47}\bar{D}_{1L}D_{1R}F_3^\dagger + f_{48}\bar{D}_{2L}D_{2R}F_1^\dagger + M\bar{D}_{3L}D_{3R} + M\bar{D}_{4L}D_{4R} + f_{49}\bar{D}_{5L}D_{5R}F_3$$
$$+ f_{50}\bar{D}_{4L}D_{5R}F_2^\dagger + f_{51}\bar{D}_{6L}D_{6R}F_3 + f_{52}\bar{D}_{7L}D_{7R}F_2^\dagger + M\bar{D}_{8L}D_{7R} + f_{53}\bar{D}_{8L}D_{8R}F_3^\dagger + \text{H.c.}, \tag{30}$$

where for simplicity of notation we have used M to denote all the TeV scale mass parameters.

Thus, integrating out the heavy fermions in the tree level diagram composed from the couplings

$$f_{11}\bar{q}_{3L}Q_{3R}S^\dagger + f_{34}^*\bar{Q}_{3R}Q_{3L}F_1^\dagger + f_{10}\bar{Q}_{3L}D_{3R}H$$
$$+ M\bar{D}_{3R}D_{3L} + f_3\bar{D}_{3L}d_{3R}S \tag{31}$$

produces an effective coupling below the TeV scale pro-

portional to

$$f_{11}f_{34}^*f_{10}f_3 \frac{F_1^\dagger}{M}\frac{S^\dagger S}{M^2}\bar{q}_{3L}d_{3R}H + \text{H.c.} \tag{32}$$

Integrating out the heavy fermions in the tree level diagram composed from the couplings

$$f_{10}\bar{q}_{2L}Q_{2R}S^\dagger + f_{33}^*\bar{Q}_{2R}Q_{2L}F_1^\dagger + f_2\bar{Q}_{2L}U_{2R}\bar{H}$$
$$+ f_{42}^*\bar{U}_{2R}U_{2L}F_2 + f_{13}\bar{U}_{2L}u_{2R}S \tag{33}$$

produces an effective coupling below the TeV scale proportional to

$$f_{10}f_{33}^*f_2f_{42}^*f_{13}\frac{F_1^\dagger F_2}{M^2}\frac{S^\dagger S}{M^2}\bar{q}_{2L}u_{2R}\bar{H} + \text{H.c.} \tag{34}$$

Integrating out the heavy fermions in the tree level diagram composed from the couplings

$$f_{10}\bar{q}_{2L}Q_{2R}S^\dagger + f_{33}^*\bar{Q}_{2R}Q_{2L}F_1^\dagger + f_{17}\bar{Q}_{2L}Q_{4R}S^\dagger$$
$$+ f_{36}^*\bar{Q}_{4R}Q_{4L}F2 + f_7\bar{Q}_{4L}D_{4R}H + M\bar{D}_{4R}D_{4L}$$
$$+ f_{24}^*\bar{D}_{4L}D_{3R}S + M\bar{D}_{3R}D_{3L} + f_3\bar{D}_{3L}d_{3R}S \tag{35}$$

produces an effective coupling below the TeV scale proportional to

$$f_{10}f_{33}^*f_{17}f_{36}^*f_7f_{24}^*f_3\frac{F_1^\dagger F_2}{M^2}\frac{(S^\dagger S)^2}{M^4}\bar{q}_{2L}d_{3R}H + \text{H.c.} \tag{36}$$

Integrating out the heavy fermions in the tree level diagram composed from the couplings

$$f_{11}\bar{q}_{3L}Q_{3R}S^\dagger + f_{34}^*\bar{Q}_{3R}Q_{3L}F_1^\dagger + f_{10}\bar{Q}_{3L}D_{3R}H$$
$$+ f_{24}^*\bar{D}_{3R}D_{4L}S^\dagger + f_{50}\bar{D}_{4L}D_{5R}F_2^\dagger + f_{49}^*\bar{D}_{5R}D_{5L}F_3^\dagger$$
$$+ f_{26}\bar{D}_{5L}D_{2R}S + f_{48}^*\bar{D}_{2R}D_{2L}F_1 + f_{15}\bar{D}_{2L}d_{2R}S \tag{37}$$

produces an effective coupling below the TeV scale proportional to

$$f_{11}f_{34}^*f_{10}f_{24}^*f_{50}f_{49}^*f_{26}f_{48}^*f_{15}\frac{F_1^\dagger F_1 F_2^\dagger F_3^\dagger}{M^4}\frac{(S^\dagger S)^2}{M^4}\bar{q}_{3L}d_{2R}H$$
$$+ \text{H.c.} \tag{38}$$

There is also another tree level contribution to h_{32}^d, proportional to

$$f_{11}f_{34}^*f_{10}f_{24}f_{50}f_{49}^*f_{26}f_{48}^*f_{15}\frac{F_1^\dagger F_1 F_2^\dagger F_3^\dagger}{M^4}\frac{(S^\dagger S)^2}{M^4}\bar{q}_{3L}d_{2R}H$$
$$+ \text{H.c.} \tag{39}$$

Integrating out the heavy fermions in the tree level diagram composed from the couplings

$$f_{10}\bar{q}_{2L}Q_{2R}S^\dagger + f_{33}^*\bar{Q}_{2R}Q_{2L}F_1^\dagger + f_{17}\bar{Q}_{2L}Q_{4R}S^\dagger$$
$$+ f_{36}^*\bar{Q}_{4R}Q_{4L}F2 + f_7\bar{Q}_{4L}D_{4R}H + M\bar{D}_{4R}D_{4L}$$
$$+ f_{50}\bar{D}_{4L}D_{5R}F_2^\dagger + f_{49}^*\bar{D}_{5R}D_{5L}F_3^\dagger + f_{26}\bar{D}_{5L}D_{2R}S$$
$$+ f_{48}^*\bar{D}_{2R}D_{2L}F_1 + f_{15}\bar{D}_{2L}d_{2R}S \tag{40}$$

produces an effective coupling below the TeV scale proportional to

$$f_{10}f_{33}^*f_{17}f_{36}^*f_7f_{50}f_{49}^*f_{26}f_{48}^*f_{15}\frac{F_1^\dagger F_1 F_2^\dagger F_2 F_3^\dagger}{M^5}$$
$$\times \frac{(S^\dagger S)^2}{M^4}\bar{q}_{3L}d_{2R}H + \text{H.c.} \tag{41}$$

Integrating out the heavy fermions in the tree level diagram composed from the couplings

$$f_9\bar{q}_{1L}Q_{1R}S^\dagger + f_{32}^*\bar{Q}_{1R}Q_{1L}F_3 + f_{18}^*\bar{Q}_{1L}Q_{5R}S^\dagger$$
$$+ f_{19}^*\bar{Q}_{5R}Q_{7L}S^\dagger + f_8\bar{Q}_{7L}D_{7R}H + M\bar{D}_{7R}D_{8L}$$
$$+ f_{30}\bar{D}_{8L}D_{6R}S + f_{28}\bar{D}_{6R}D_{1L}S + f_{14}\bar{D}_{1L}d_{1R}S \tag{42}$$

produces an effective coupling below the TeV scale proportional to

$$f_9f_{32}^*f_{18}^*f_{19}^*f_8f_{30}f_{28}f_{14}\frac{F_3}{M}\frac{(S^\dagger S)^3}{M^6}\bar{q}_{1L}d_{1R}H + \text{H.c.} \tag{43}$$

There are four other very similar tree level contributions to h_{11}^d.

Integrating out the heavy fermions in the tree level diagram composed from the couplings

$$f_9\bar{q}_{1L}Q_{1R}S^\dagger + f_{32}^*\bar{Q}_{1R}Q_{1L}F_3 + f_{18}^*\bar{Q}_{1L}Q_{5R}S^\dagger$$
$$+ f_{19}^*\bar{Q}_{5R}Q_{7L}S^\dagger + f_5\bar{Q}_{7L}U_{6R}\bar{H} + f_{46}^*\bar{U}_{6R}U_{6L}F_1$$
$$+ f_{23}\bar{U}_{6L}U_{5R}S + f_{45}^*\bar{U}_{5R}U_{5L}F_2 + f_{22}\bar{U}_{5L}U_{1R}S$$
$$+ f_{41}^*\bar{U}_{1R}U_{1L}F_3^\dagger + f_{12}\bar{U}_{1L}u_{1R}S \tag{44}$$

produces an effective coupling below the TeV scale proportional to

$$f_9f_{32}^*f_{18}^*f_{19}^*f_5f_{46}^*f_{23}f_{45}^*f_{22}f_{41}^*f_{12}\frac{F_1F_2F_3^\dagger F_3}{M^4}$$
$$\times \frac{(S^\dagger S)^3}{M^6}\bar{q}_{1L}u_{1R}\bar{H} + \text{H.c.} \tag{45}$$

Integrating out the heavy fermions in the tree level diagram composed from the couplings

$$f_9\bar{q}_{1L}Q_{1R}S^\dagger + f_{32}^*\bar{Q}_{1R}Q_{1L}F_3 + f_{18}^*\bar{Q}_{1L}Q_{5R}S^\dagger$$
$$+ f_{19}^*\bar{Q}_{5R}Q_{7L}S^\dagger + f_8\bar{Q}_{7L}D_{7R}H + f_{52}^*\bar{D}_{7R}D_{7L}F_2$$
$$+ f_{29}\bar{D}_{7L}D_{5R}S + f_{49}^*\bar{D}_{5R}D_{5L}F_3^\dagger + f_{26}\bar{D}_{5L}D_{2R}S$$
$$+ f_{48}^*\bar{D}_{2R}D_{2L}F_1 + f_{15}\bar{D}_{2L}d_{2R}S \tag{46}$$

produces an effective coupling below the TeV scale pro-

portional to

$$f_9 f_{32}^* f_{18}^* f_{19}^* f_8 f_{52}^* f_{29} f_{49}^* f_{26} f_{48}^* f_{15} \frac{F_1 F_2 F_3^\dagger F_3}{M^4}$$

$$\times \frac{(S^\dagger S)^3}{M^6} \bar{q}_{1L} d_{2R} H + \text{H.c.} \tag{47}$$

Integrating out the heavy fermions in the tree level diagram composed from the couplings

$$f_9 \bar{q}_{1L} Q_{1R} S^\dagger + f_{32}^* \bar{Q}_{1R} Q_{1L} F_3 + f_{18}^* \bar{Q}_{1L} Q_{5R} S^\dagger$$

$$+ f_{19}^* \bar{Q}_{5R} Q_{7L} S^\dagger + f_8 \bar{Q}_{7L} D_{7R} H + f_{52}^* \bar{D}_{7R} D_{7L} F_2$$

$$+ f_{29} \bar{D}_{7L} D_{5R} S + f_{50}^* \bar{D}_{5R} D_{4L} F_2 + f_{24}^* \bar{D}_{4L} D_{3R} S$$

$$+ M \bar{D}_{3R} D_{3L} + f_3 \bar{D}_{3L} d_{3R} S \tag{48}$$

produces an effective coupling below the TeV scale proportional to

$$f_9 f_{32}^* f_{18}^* f_{19}^* f_8 f_{52}^* f_{29} f_{50}^* f_{24}^* f_3 \frac{(F_2)^2 F_3}{M^3} \frac{(S^\dagger S)^3}{M^6} \bar{q}_{1L} d_{3R} H$$

$$+ \text{H.c.} \tag{49}$$

There is one other very similar tree level contribution to h_{13}^d.

Integrating out the heavy fermions in the tree level diagram composed from the couplings

$$f_{11} \bar{q}_{3L} Q_{3R} S^\dagger + f_{34}^* \bar{Q}_{3R} Q_{3L} F_1^\dagger + f_{10} \bar{Q}_{3L} D_{3R} H$$

$$+ f_{24}^* \bar{D}_{3R} D_{4L} S^\dagger + f_{50} \bar{D}_{4L} D_{5R} F_2^\dagger + f_{29}^* \bar{D}_{5R} D_{7L} S^\dagger$$

$$+ f_{52}^* \bar{D}_{7L} D_{7R} F_2^\dagger + M \bar{D}_{7R} D_{8L} + f_{30} \bar{D}_{8L} D_{6R} S$$

$$+ f_{28} \bar{D}_{6R} D_{1L} S + f_{14} \bar{D}_{1L} d_{1R} S \tag{50}$$

produces an effective coupling below the TeV scale proportional to

$$f_{11} f_{34}^* f_{10} f_{24}^* f_{50} f_{29}^* f_{52}^* f_{30} f_{28} f_{14} \frac{F_1^\dagger (F_2^\dagger)^2}{M^3} \frac{(S^\dagger S)^3}{M^6} \bar{q}_{3L} d_{1R} H$$

$$+ \text{H.c.} \tag{51}$$

The following effective couplings are not generated or are generated at higher order in ϵ / β: h_{23}^u, h_{32}^u, h_{12}^u, h_{21}^u, h_{13}^u, h_{31}^u, and h_{21}^d. As already indicated these couplings are not needed to reproduce the observed SM quark masses and mixings. For illustration, h_{32}^u arises from the effective coupling

$$f_{11} f_{34}^* f_6 f_{25} f_7^* f_{36} f_{17}^* f_2 f_{42}^* f_{13} \frac{F_1^\dagger F_2 F_2^\dagger}{M^3} \frac{(S^\dagger S)^2}{M^4}$$

$$\times \frac{H^\dagger H}{M^2} \bar{q}_{3L} u_{2R} \bar{H} + \text{H.c.}, \tag{52}$$

so the extra suppression relative to Eq. (1) is by an additional factor of β as well as an additional factor of ϵ.

Since h_{12}^u and h_{21}^u have extra suppression in this model, $D^0 - \bar{D}^0$ mixing also has extra suppression. This weakens

the lower bound on m_s derived in Sec. III B. Similarly since h_{23}^u and h_{32}^u have extra suppression the relatively large BR for $t \to cs$ discussed in Sec. IV B will not occur for this particular realization.

VI. CONCLUSION

We have presented a framework in which only the top quark obtains its mass from the Yukawa interaction with the SM Higgs boson via dimension 4 operators. All the other quarks receive their masses from operators of dimension 6 or higher involving a complex scalar S that is part of an extended Higgs sector and whose vev is at the electroweak scale. The successive hierarchy of light quark masses is generated via the expansion parameter $(\frac{S^\dagger S}{M^2}) \sim \epsilon^2$, where $\epsilon \equiv \frac{v_s}{M} \sim 0.15$. All the couplings of the higher dimensional operators are of order 1. We are able to generate the appropriate hierarchy of fermion masses with this small parameter ϵ. Since v_s is at the EW scale, the physics of the new scale M is not far above a TeV. We predict a neutral scalar s, which gives rise to signals that could be detected at the LHC or at the Tevatron. We make new predictions for Higgs decays and for top decays. The model has a light Z' that has very weak couplings to SM fermions, but could be light enough to be produced via mixing in Higgs decays at the LHC; this could give rise to invisible Higgs decays, displaced vertices from the Z' decays, or multilepton final states, depending on the mass and lifetime of the Z'.

We have presented an explicit model in which the effective interaction given in (1) is realized. This involves extending the SM gauge symmetry by an abelian gauge symmetry $U(1)_S$ and a local flavon symmetry group $U(1)_{F_1} \times U(1)_{F_2} \times U(1)_{F_3}$. The flavon symmetry is spontaneously broken at the TeV scale by complex flavon scalars F_1, F_2, F_3, whereas the $U(1)_S$ symmetry is broken at the electroweak scale by the complex scalar S, which is a SM singlet extension of the SM Higgs sector. S acts as the messenger of both flavor and electroweak symmetry breaking. The model requires the existence of vectorlike quarks and leptons, both EW doublets and singlets, at the TeV scale. These can be probed at the LHC. Their decays will be a new source for Higgs production and give rise to final states with four Z's or four Z''s and other interesting new physics signals at the LHC.

We have restricted ourselves to models where all of the hierarchies of the SM quark and charged lepton masses and mixings arise from powers of the vev of a single messenger field. In [6], a framework was suggested in which all of these hierarchies arise from powers of $\beta = (\frac{H^\dagger H}{M^2})$. As we saw in the previous section, in explicit models it is natural to generate powers of both ϵ and β. Thus the model presented here and the framework of [6] are two extremes of a more general class of models. Obviously one could also generalize by introducing a more complicated messenger sector, i.e. further extending the Higgs sector.

A truly viable model should have fewer species of heavy fermions than were required in our example, ameliorating what is otherwise a dramatic worsening of the little hierarchy problem of the standard model. This could be achieved by a more efficient construction of the messenger sector and its interplay with the flavon sector. Another interesting direction is to attempt to generate some of the higher order effective couplings from the top quark Yukawa, as was done successfully with leptoquark-generated loop diagrams in [4].

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