Natural fermion mass hierarchy and new signals for the Higgs boson

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We suggest a novel approach towards resolving the fermion mass hierarchy problem within the framework of the standard model. It is shown that the observed masses and mixings can be explained with order one couplings using successive higher dimensional operators involving the SM Higgs doublet field. This scenario predicts a flavor-dependent enhancement in the Higgs boson coupling to the fermions. It also predicts a flavor changing $\overline{t}ch^0$ interaction with a strength comparable to that of $\overline{b}bh^0$, opening up new discovery channels for the Higgs boson at the upgraded Fermilab Tevatron and the CERN LHC. Additional tests of the framework include observable D^0-$\overline{D}^0$ mixing and a host of new phenomena associated with flavor physics at the TeV scale.

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One of the major unresolved puzzles in the standard model (SM) is the observed hierarchy in the masses and mixings of quarks and leptons. In this paper we propose a new approach towards resolving this puzzle. We show that an extremely good fit to all the masses and mixings can be obtained by using higher dimensional operators involving the relevant fermion fields and successive powers of the SM Higgs doublet field. These non-renormalizable operators will have inverse mass dimensions, but the dimensionless couplings which multiply them can all be of order 1. Thus the small Yukawa couplings of the SM will emerge naturally in our approach.

Our proposal differs from most other approaches [1] which address the fermion mass hierarchy puzzle in one crucial way: Since vacuum expectation values (VEVs) of SM singlet scalar fields are not used, the scale of flavor physics cannot be much above the weak scale. In fact, from our fits to the fermion masses and mixings, we find that this scale is in the range 1–2 TeV. While such a scale is high enough to be consistent with all the experimental constraints, it is within reach of planned accelerators in the near future for direct detection.

In our approach, although the spectrum of the theory below 1 TeV is that of the minimal SM, there are a variety of new signals associated with the Higgs boson that are distinct from the SM. The couplings of the Higgs boson to light fermions are enhanced by a flavor-dependent numerical factor relative to that of the SM. The enhanced $\overline{b}bh^0$ and $\overline{\mu}\mu h^0$ couplings would have a significant impact on Higgs boson detection at hadron colliders and at a muon collider [2].

A further consequence of using successive higher dimensional operators is that there are flavor changing neutral current processes mediated by the SM Higgs boson. While all the low energy constraints are well satisfied, the model predicts D^0-$\overline{D}^0$ mixing to be close to the present experimental limit.

The proposed scenario predicts an interesting new discovery channel for the Higgs boson at the upgraded Fermilab Tevatron and the CERN Large Hadron Collider (LHC). It uses the flavor changing Higgs vertex $\overline{t}ch^0$ which has a strength comparable to that of the flavor conserving $\overline{b}bh^0$ vertex. If the mass of h^0 is below 150 GeV, then in the $\overline{t}t$ production at colliders, the decay of the top quark can lead to Higgs boson signals. For $m_{h^0}\approx 200-350$ GeV, the production process $gg\to h^0$ followed by $h^0\to\overline{t}c+\overline{c}t$ can lead to a new signal at LHC.

We assume that some flavor symmetries prevent the direct Yukawa coupling of the SM Higgs doublet (H) to the light fermions. These flavor symmetries are spontaneously broken at a scale $M\approx 1-2$ TeV. The effective theory below M is the SM with one Higgs doublet but with non-renormalizable terms in the Higgs Yukawa couplings. For the quark sector the effective Yukawa Lagrangin is taken to be (with $\tilde{H}\equiv -i\tau_2 H^*$)

$$
\begin{aligned}
\mathcal{L}^{\rm Yuk} = {} & h_{33}^u Q_3 u_3^c \tilde{H} + \left(\frac{H^\dagger H}{M^2}\right)(h_{33}^d Q_3 d_3^c H + h_{22}^u Q_2 u_2^c \tilde{H} \\
& + h_{23}^u Q_2 u_3^c \tilde{H} + h_{32}^u Q_3 u_2^c \tilde{H}) + \left(\frac{H^\dagger H}{M^2}\right)^2 (h_{22}^d Q_2 d_2^c H \\
& + h_{23}^d Q_2 d_3^c H + h_{32}^d Q_3 d_2^c H + h_{12}^u Q_1 u_2^c \tilde{H} + h_{21}^u Q_2 u_1^c \tilde{H} \\
& + h_{13}^u Q_1 u_3^c \tilde{H} + h_{31}^u Q_3 u_1^c \tilde{H}) + \left(\frac{H^\dagger H}{M^2}\right)^3 (h_{11}^u Q_1 u_1^c \tilde{H} \\
& + h_{11}^d Q_1 d_1^c H + h_{12}^d Q_1 d_2^c H + h_{21}^d Q_2 d_1^c H + h_{13}^d Q_1 d_3^c H \\
& + h_{31}^d Q_3 d_1^c H) + \text{H.c.}
\end{aligned}
\tag{1}
$$

In each term of Eq. (1) there is a unique $SU(2)$ group contraction. Only the top quark has a renormalizable Yukawa coupling; other couplings are suppressed by successive powers of $(H^\dagger H/M^2)$. This provides the small expansion parameter needed to explain the light fermion masses. We will show that a good fit to all the fermion masses and mixing angles can be obtained with the dimensionless couplings ($h_{ij}^{u,d}$) in Eq. (1) taking values of order one.

In most other attempts to explain the mass hierarchy, singlet scalar fields (S_i) are employed and the expansion parameter is $\langle S_i\rangle/M$. In such cases, both $\langle S_i\rangle$ and M can be large, for, e.g., near the Planck scale. The low energy theory will be identical to the SM with no modification to the Higgs boson interactions. In contrast, in our approach, the expansion parameter is $\langle H^\dagger H\rangle/M^2$, which implies that M cannot be much above the weak scale. It also results in new interactions of the SM Higgs boson, which can be directly tested.

We will comment on possible ways of deriving the effective Lagrangian of Eq. (1) from renormalizable theories towards the end of the paper. We remark here that the coefficients $h_{ij}^{u,d}/M^2$ could be thought of as background fields which carry flavor quantum numbers. The dimensions of the various operators in Eq. (1) (assumed to have a symmetric form) are allowed to be as low as possible, consistent with the observed masses and mixings and our demand that the coefficients be order 1. The effective Lagrangian inducing the charged lepton masses is taken to have a form identical to that for the down-type quarks [replace the couplings h_{ij}^d by h_{ij}^l in Eq. (1) for the leptons].

Writing $H=(h^0/\sqrt{2}+v, \quad 0)^T$ in unitary gauge with $v \simeq 174$ GeV, and defining a small parameter $\epsilon\equiv v/M$, Eq. (1) leads to the following mass matrices for the up–type and the down–type quarks:

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

$$M_d=\begin{pmatrix} h_{11}^d\epsilon^6 & h_{12}^d\epsilon^6 & h_{13}^d\epsilon^6 \\ h_{21}^d\epsilon^6 & h_{22}^d\epsilon^4 & h_{23}^d\epsilon^4 \\ h_{31}^d\epsilon^6 & h_{32}^d\epsilon^4 & h_{33}^d\epsilon^2 \end{pmatrix} v. \tag{2}$$

The charged lepton mass matrix is obtained from M_d by replacing $h_{ij}^d \to h_{ij}^l$.

The masses of the quarks and leptons can be read off from Eq. (2) in the approximation $\epsilon \ll 1$:

$$\{m_t, m_c, m_u\} \simeq \{|h_{33}^u|,\ \ |h_{22}^u|\epsilon^2,\ \ |h_{11}^u - h_{12}^u h_{21}^u/h_{22}^u|\epsilon^6\}v,$$

$$\{m_b, m_s, m_d\} \simeq \{|h_{33}^d|\epsilon^2,\ \ |h_{22}^d|\epsilon^4,\ \ |h_{11}^d|\epsilon^6\}v,$$

$$\{m_\tau, m_\mu, m_e\} \simeq \{|h_{33}^l|\epsilon^2,\ \ |h_{22}^l|\epsilon^4,\ \ |h_{11}^l|\epsilon^6\}v. \tag{3}$$

The quark mixing angles are found to be $|V_{us}| \simeq |h_{12}^d/h_{22}^d - h_{12}^u/h_{22}^u|\epsilon^2$, $|V_{cb}| \simeq |h_{23}^d/h_{33}^d - h_{23}^u/h_{33}^u|\epsilon^2$, $|V_{ub}| \simeq |h_{13}^d/h_{33}^d - h_{12}^u h_{23}^d/h_{22}^u h_{33}^d - h_{13}^u/h_{33}^u|\epsilon^4$.

To see how the fit works, we choose an illustrative set of input values for the quark masses [3]: $m_u(1\ \text{GeV}) = 5.1$ MeV, $m_c(m_c) = 1.27$ GeV, $m_t^{\text{phys}} = 175$ GeV, $m_d(1\ \text{GeV}) = 8.9$ MeV, $m_s(1\ \text{GeV}) = 175$ MeV, $m_b(m_b) = 4.25$ GeV. We extrapolate all masses to a common scale, conveniently chosen as $m_t(m_t) \simeq 166$ GeV using three loop QCD and one loop QED beta functions. With $\alpha_s(M_Z) = 0.118$, the values of the running masses at $\mu = m_t(m_t)$ are found to be $\{m_t, m_c, m_u\} \simeq \{166,\ 0.60,\ 2.2 \times 10^{-3}\}$ GeV, $\{m_b, m_s, m_d\} \simeq \{2.78,\ 7.5\times10^{-2},\ 3.8\times10^{-3}\}$ GeV, $\{m_\tau, m_\mu, m_e\} \simeq \{1.75,\ 0.104,\ 5.01\times10^{-4}\}$ GeV. An optimal choice of the expansion parameter is $\epsilon = 1/6.5$. ($\epsilon \simeq 1/7 - 1/6$ gives reasonable fits.) This corresponds to a relatively low scale of flavor symmetry breaking: $M \simeq 1.1$ TeV. With $\epsilon = 1/6.5$, we determine the dimensionless coefficients $h_{ij}^{u,d,l}$ from the fermion masses as

$$\{|h_{33}^u|, |h_{22}^u|, |h_{11}^u - h_{12}^u h_{21}^u/h_{22}^u|\} \simeq \{0.96, 0.14, 0.95\},$$

$$\{|h_{33}^d|, |h_{22}^d|, |h_{11}^d|\} \simeq \{0.68, 0.77, 1.65\},$$

$$\{|h_{33}^l|, |h_{22}^l|, |h_{11}^l|\} \simeq \{0.42, 1.06, 0.21\}. \tag{4}$$

We see that, remarkably, all couplings are of order unity. The largest deviation from 1 is the charm Yukawa coupling $|h_{22}^u| \simeq 0.14$. This small fluctuation actually goes in the right direction to explain the magnitude of the Cabibbo angle. From the expression for $|V_{us}|$, with $|h_{22}^u| \simeq 0.14$, and $h_{12}^{u,d}$ of order 1, the correct value of $|V_{us}| \simeq 0.2$ follows naturally. Similarly, $|V_{ub}| \sim 7\epsilon^4 \simeq 0.004$, where the factor of 7 is due to $1/h_{22}^u$ enhancement in the second term of $|V_{ub}|$. Here $|V_{cb}| \simeq 0.04$ can be fit for, e.g., by choosing $|h_{23}^u| \simeq 1.4$ (if that term dominates) or with $|h_{23}^d| \simeq 0.84$ (if it dominates). We thus see that the overall fit of the scheme is quite good and turn to analyze its experimental consequences.

The Yukawa coupling matrices of the SM Higgs boson to the quarks that follow from Eq. (1) are

$$Y_u=\begin{pmatrix} 7h_{11}^u\epsilon^6 & 5h_{12}^u\epsilon^4 & 5h_{13}^u\epsilon^4 \\ 5h_{21}^u\epsilon^4 & 3h_{22}^u\epsilon^2 & 3h_{23}^u\epsilon^2 \\ 5h_{31}^u\epsilon^4 & 3h_{32}^u\epsilon^2 & h_{33}^u \end{pmatrix},$$

$$Y_d=\begin{pmatrix} 7h_{11}^d\epsilon^6 & 7h_{12}^d\epsilon^6 & 7h_{13}^d\epsilon^6 \\ 7h_{21}^d\epsilon^6 & 5h_{22}^d\epsilon^4 & 5h_{23}^d\epsilon^4 \\ 7h_{31}^d\epsilon^6 & 5h_{32}^d\epsilon^4 & 3h_{33}^d\epsilon^2 \end{pmatrix}, \tag{5}$$

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

There are two striking features in Eq. (5). One is that the diagonal couplings are enhanced relative to the respective SM Higgs boson Yukawa couplings by a numerical factor. These enhancement factors are (1,3,7) for (t,c,u) and (3,5,7) for (b,s,d) as well as (τ,μ,e). In two Higgs doublet models or in supersymmetric models, while the couplings of h^0 might be enhanced, they are not flavor dependent; nor do they take these specific values. Second, the Yukawa coupling matrix and the corresponding mass matrix do not diagonalize simultaneously. This will lead to Higgs mediated flavor changing neutral current processes, even though there is only a single electroweak Higgs boson in the theory [4].

The flavor changing neutral current (FCNC) couplings of the SM Higgs boson h^0 to fermions can be obtained from Eqs. (5) and (2). In the quark sector in terms of the mass eigenstates they are given by

$$\mathcal{L}^{\rm FCNC} \simeq \frac{h^0}{\sqrt{2}}(2h_{12}^d\epsilon^6 ds^c + 2h_{21}^d\epsilon^6 sd^c + 4h_{13}^d\epsilon^6 db^c + 4h_{31}^d\epsilon^6 bd^c + 2h_{23}^d\epsilon^4 sb^c + 2h_{32}^d\epsilon^4 bs^c) + \frac{h^0}{\sqrt{2}}(2h_{12}^u\epsilon^4 uc^c + 2h_{21}^u\epsilon^4 cu^c + 4h_{13}^d\epsilon^4 ut^c + 4h_{31}^u\epsilon^4 tu^c + 2h_{23}^u\epsilon^2 ct^c + 2h_{32}^u\epsilon^2 tc^c) + \text{H.c.} \tag{6}$$

For FCNC Higgs couplings in the charged lepton sectors, replace $h_{ij}^d \to h_{ij}^l$ in Eq. (6).

There is a tree–level contribution mediated by the Higgs boson for the $K^0 - \bar{K}^0$ mass difference. We estimate it in the vacuum saturation approximation for the hadronic matrix element [5]. The new contribution, $\Delta m_K^{\rm Higgs}$, is given by

$$\Delta m_K^{\rm Higgs} \simeq \frac{1}{2}\frac{f_K^2 m_K B_K \eta_{QCD}}{m_{h^0}^2}\epsilon^{12}\left\{\left(\frac{1}{6}\frac{m_K^2}{(m_d+m_s)^2}+\frac{1}{6}\right)\times \mathrm{Re}\left[\left(\frac{h_{12}^d + h_{21}^{d*}}{\sqrt{2}}\right)^2\right] - \left(\frac{11}{6}\frac{m_K^2}{(m_d+m_s)^2}+\frac{1}{6}\right)\times \mathrm{Re}\left[\left(\frac{h_{21}^d - h_{12}^{d*}}{\sqrt{2}}\right)^2\right]\right\}. \tag{7}$$

Using $B_K = 0.75$, $f_K \simeq 160$ MeV, $\epsilon \simeq 1/6.5$, $m_s = 175$ MeV, $\eta_{QCD} \simeq 5$, $m_d = 8.9$ MeV, and with $h_{12}^d = 1, h_{21}^d = 0.5$, we obtain $\Delta m_K^{\rm Higgs} \simeq 6 \times 10^{-17}$ GeV, for $m_{h^0} = 100$ GeV. This is a factor of 50 below the experimental value. With $h_{12}^d \simeq h_{21}^{d*} = 1$, $\Delta m_K^{\rm Higgs} \simeq 1 \times 10^{-15}$ GeV, also consistent with experiment.

As for the CP violating parameter ϵ_K, it receives a new contribution from the Higgs boson exchange which can be significant and can even dominate over the the SM Cabibbo-Kobayashi-Maskawa (CKM) contribution. For example, in the choice of parameters with $|h_{12}^d| = 1, |h_{21}^d| = 0.5$, but with their phases being of order 1, ϵ_K arising from the Higgs boson exchange can explain the observed value entirely. This possibility will be tested at the B factory. New contributions to ϵ' are negligible. However, since the prediction for ϵ_K is modified, the standard model fit to the CKM parameter $\mathrm{Im}(V_{td}^* V_{ts})$ will be modified by a factor of order 1 (depending on the relative strength of the CKM and the Higgs boson exchange contribution to ϵ_K). Since ϵ' is directly proportional to $\mathrm{Im}(V_{td}^* V_{ts})$, its prediction will be altered by a factor of order 1 [6].

Electric dipole moments (EDMs) of the neutron and the electron in our scheme are much larger than the SM prediction. The dominant source of the neutron EDM (d_n) is from the two–loop Barr-Zee [7] diagram involving the SM Higgs boson and the Z boson. In our scheme, h^0 has a scalar as well as a pseudoscalar coupling to the u quark. In the physical basis of the u quark the pseudoscalar coupling has a strength of order $\mathrm{Im}(h_{11}^u)(9\epsilon^6)$. From this, we estimate the neutron EDM to be in the range $(10^{-26} - 10^{-27})e$ cm for the phase of order unity and Higgs mass of order 100 GeV. The EDM of the electron (d_e) will arise from an analogous diagram, but its magnitude is suppressed by an additional power of ϵ^2. We estimate $d_e \approx (10^{-27} - 10^{-28})e$ cm. Both d_e and d_n are within reach of future experiments.

All other constraints from low energy flavor changing processes such as $K_L \to \mu^+\mu^-$, $K_L \to \mu e$, $K \to \pi \bar{\nu}\nu$, $\mu \to e\gamma$, $\mu \to 3e$, $B_d - \bar{B}_d$ mixing, etc., are orders of magnitude below the corresponding experimental values and limits. D^0-$\bar{D}^0$ mixing, on the other hand, is predicted to be near the present experimental limit in our scenario. Note that the FCNC $uc^c h^0$ vertex is enhanced by a factor of ϵ^{-2} compared to $ds^c h^0$ vertex [see Eq. (6)]. The new contribution $\Delta m_D^{\rm Higgs}$ is given by an expression analogous to Eq. (7). Using $f_D \simeq 200$ MeV, $B_D \simeq 0.75$, $m_D^2/(m_c + m_u)^2 \simeq 2$, $\eta_{QCD} = 4$, and $h_{12}^u = 1$, $h_{12}^u = 0.5$, we estimate $\Delta m_D^{\rm Higgs} \simeq 7 \times 10^{-14}$ GeV for a Higgs boson mass of 200 GeV, to be compared with the present experimental limit of $\Delta m_D \leqslant 1.6 \times 10^{-13}$ GeV [8]. Allowing for reasonable order-1 uncertainties in the hadronic matrix element and the FCNC couplings, and varying ϵ in the range 1/6−1/7, we conclude that D^0-$\bar{D}^0$ mixing should be not more than a factor of 10−20 below the present limit, provided that the Higgs boson mass is below about 300 GeV. This prediction should be testable in the near future. We should remark that the SM long distance contribution to D^0-$\bar{D}^0$ mixing [9] is expected to be about three orders of magnitude below the current limit, so its discovery should be clear-cut signal for new short-distance physics, such as the one we propose.

There are a variety of new signals associated with the production and decay of the Higgs boson in our scenario. While the tree-level Higgs couplings to the gauge bosons are identical to that of the SM, the Yukawa interactions are modified from that of the SM. The consequences are significant for the strategy to discover the Higgs boson at colliders [10]. We list below the processes where the differences from the SM are most striking.

(1) The Higgs boson couplings to light fermions are enhanced by a flavor-dependent numerical factor. At the Next Linear Collider (NLC) and perhaps the LHC, the Yukawa couplings to (b,τ,t) will be measured [11]. The scenario can thereby be directly tested. The enhanced coupling to μ implies that the Higgs boson production cross section will increase by a factor of 25 at a muon collider [2], relative to that of SM.

(2) The partial width for Higgs boson decay into $\bar{b}b$ will increase by a factor of 9 relative to the SM value. In the SM, the $\bar{b}b$ and the W^*W^* branching ratios become comparable for a Higgs mass of about 135 GeV [12]. Because of the difference in its b-quark coupling, this crossover occurs in our scheme for $m_{h^0} \simeq 155$ GeV. The Higgs boson mass reach via $p\bar{p} \to Wh^0X$ at Tevatron will increase by about 10−20 GeV relative to the SM. This difference should be incorporated into the strategy for discovering the Higgs boson at the upgraded Tevatron [13,12].

(3) The $h^0\gamma\gamma$ coupling, which arises primarily from the $W^{\pm}$ loop, is nearly unaffected in our scheme. Since the $\overline{b}bh^0$ vertex is enhanced by a factor of 9, the branching ratio for $h^0 \to \gamma\gamma$ will decrease by a factor of approximately 9, for $m_{h^0} \leqslant 155$ GeV. This will make the search for an intermediate mass Higgs boson via this process more difficult at the LHC.

(4) The flavor-changing $\overline{t}ch^0$ vertex gives an exciting new discovery channel for the Higgs boson at the upgraded Tevatron. In our scenario this vertex has a strength of order $2\epsilon^2$, which is similar in magnitude to that of the flavor conserving $\overline{b}bh^0$ coupling. If m_{h^0} is less than m_t, the decays $t \to ch^0$ and $\overline{t} \to \overline{c}h^0$ can provide a new channel for Higgs boson discovery. Once produced, h^0 will decay into $\overline{b}b$ with a significant branching ratio. For $m_{h^0}=100$ GeV, and with $|h^u_{23}/h^u_{33}| \simeq |h^u_{32}/h^u_{33}| \simeq 1$, the branching ratio is $Br(t \to ch^0) \simeq 1.1\times 10^{-3}$. With an integrated luminosity of 20 fb^{-1}, about 1.4×10^5 $t\overline{t}$ pairs are expected at Tevatron running at $\sqrt{s}=2$ TeV. This would lead to about 300 Higgs events from t and $\overline{t}$ decays via $t \to ch^0$ and $\overline{t} \to \overline{c}h^0$. The invariant mass of the $b\overline{b}$ jets will provide the Higgs signal. The QCD background can be brought under control by tagging the t (or $\overline{t}$) by its decay into $W+b$. This process could be useful to discover a Higgs boson of mass as large as about 150 GeV at the Tevatron (at which point the kinematic suppression becomes significant). This reach can be as large as 170 GeV at LHC.

For m_{h^0} between m_t and $2m_t$, there is another way to look for the Higgs boson at hadron colliders. The cross section for h^0 production via gg fusion is about 15 pb for $m_{h^0}=200$ GeV at the LHC running at $\sqrt{s}=14$ TeV. The Higgs boson will decay dominantly into W pairs, but the branching fraction into $\overline{t}c+\overline{c}t$ is not negligible: For $h^u_{23} \simeq h^u_{32} \simeq 1$, $Br(h^0 \to \overline{t}c+\overline{c}t) \simeq 1\times 10^{-3}$. With 100 fb^{-1} of data, this will result in 1500 $\overline{t}c+\overline{c}t$ events. The signal will be thus a b jet, a charm and a W. It might be possible to reconstruct the invariant mass of the Higgs boson by studying the leptonic decay of the W from the top quark. (Although there is a neutrino involved, its four-momentum can be reconstructed by measuring the charged lepton momentum, up to a twofold ambiguity.) The background from SM single top quark production can be substantially reduced by the invariant mass requirement. Two jet $+W$ production (where a jet is misidentified as a b) and $\overline{b}bW$ where one b is missed are other dominant backgrounds. The presence of a top quark in the signal but not in these background events can be utilized to provide further cuts.

(5) Since the scale of flavor physics is identified to be $M \approx 1-2$ TeV, new phenomena associated with flavor physics will show up at experiments performed with energies greater than $1-2$ TeV. This will happen at the LHC. One concrete example is the unraveling of the effective vertices in Eq. (1). For example, at $s \sim M^2$, the process $pp \to \overline{b}b(3h^0)$ will proceed without much suppression in the coupling. The relevant dimensionful coupling is $h^d_{33}/M^2 \simeq m_b/v^3$. For $m_{h^0} \sim 100$ GeV, we estimate the cross section for this process at LHC to be in the fb range. The signature will be quiet dramatic: 8 b jets with 3 pairs adding to to the same invariant Higgs boson mass.

We conclude by sketching a concrete model which will induce the desired effective Yukawa Lagrangian on Eq. (1). Although it is conceivable that Eq. (1) has a dynamical origin [14], our explicit construction employs only perturbative physics at the TeV scale. Suppose there are vector-like fermions at the TeV scale in doublet and singlet representations of $SU(2)_L$. Owing to suitable flavor symmetries the usual Yukawa couplings of the SM fermions are forbidden (except for the top quark). These vector fermions also transform under the flavor symmetry in such a way that certain mixed Yukawa couplings of the SM fermions with the vector fermions and the Higgs doublet are allowed. The bare masses of the vector fermions will also take a special form due to flavor symmetries.

Consider the generation of the effective Lagrangian which induces the b quark mass in Eq. (1). Analogous discussions will hold for the c quark and the τ lepton as well. For consistency there must be a minimum of 3 vector fermions that are integrated — otherwise an infinite contribution to the renormalizable Yukawa coupling $Q_3 d_3^c H$ will result. This can be seen by closing an H and $H^\dagger$ line in Eq. (1) to form a loop. With three internal vector fermions, such a diagram will be finite. An example consistent with flavor symmetries is as follows:

$$\begin{aligned} L_b = & a[\overline{Q}_{3L}F_{1R}H + b\overline{G}_{1L}d_{3R}H + c\overline{F}_{3L}G_{3R}H^\dagger + \text{H.c.}] \\ & + m_F(\overline{F}_1F_2+\overline{F}_2F_1) + m(\overline{F}_2F_3+\overline{F}_3F_2) \\ & + m_G(\overline{G}_1G_2+\overline{G}_2G_1) + + m'[\overline{G}_2G_3+\overline{G}_3G_2] + \text{H.c.} \\ & + M_F\overline{F}_3F_3 + M_G\overline{G}_3G_3 . \end{aligned} \tag{8}$$

Here F_i are singlets under $SU(2)_L$ while G_i are doublets. One possible choice of flavor symmetries which yields Eq. (8) is a $U(1)$ with the charges given as $Q[F_L]=-Q[F_R]=(-2,1,0)$, $Q[Q_{3L}]=-Q[d_{3R}]=2$, $Q[H]=0$, $Q[S]=1$, where S is a singlet field that induces the mass terms (m_F, m_G, m, m'). After integrating the G_i and F_i fields, we will arrive at the h^d_{33} term in Eq. (1), given as

$$\frac{h^d_{33}}{M^2} \simeq \frac{abc}{2M_FM_G}\left(\frac{mm'}{m_Fm_G}\right). \tag{9}$$

Here we assumed a hierarchy $m_F \sim m_G \sim m \sim m' \ll M_F \sim M_G$ for simplicity. If the Yukawa couplings a,b,c are of order 1, the mass parameters $M_F \sim M_G \sim 1-2$ TeV from a fit to fermion masses.

The renormalizable Lagrangian of Eq. (8) will also induce at the one-loop level a term $Q_3 d_3^c H$ in the effective Lagrangian. If we denote this induced Yukawa coupling to be Δh^d_{33}, its value can be computed to be

$$\frac{\Delta h_{33}^{d}}{h_{33}^{d}} \simeq \frac{2}{16\pi^2}\left\{ \frac{m_F^2}{M^2}\left[\ln\left(\frac{M_F^2}{m_F^2}\right) -1\right] + \frac{m_G^2}{M^2}\left[\ln\left(\frac{M_G^2}{m_G^2}\right) -1\right]\right\}. \tag{10}$$

This contribution should be compared with $\epsilon^2 \simeq 1/40$ that results from Eq. (1) for the b-quark mass. Note that, in addition to the loop factor, there is a suppression factor of $(m_F/M_F)^2$ in Eq. (10). For $m_F/M_F \simeq 1/4-1/2$, these loop contributions to the fermion masses are much smaller compared with the tree level contributions from Eq. (1). It is straightforward to extend this model to include the lighter fermions where we found a similar behavior.

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[1] See e.g., C.D. Froggatt and H.B. Nielsen, Nucl. Phys. **B147**, 277 (1979); S.M. Barr, Phys. Rev. D **21**, 1424 (1980); H. Georgi, A. Manohar, and A. Nelson, Phys. Lett. **126B**, 169 (1983); B.S. Balakrishna, A.L. Kagan, and R.N. Mohapatra, Phys. Lett. B **205**, 345 (1988).

[2] For a review see V. Barger, M. Berger, J. Gunion, and T. Han, Phys. Rep. **286**, 1 (1997).

[3] G. Gasser and H. Leutwyler, Phys. Rep. **87**, 77 (1982).

[4] For discussions of FCNC processes in two Higgs doublet models, see, e.g., T.P. Cheng and M. Sher, Phys. Rev. D **35**, 3484 (1987); Y.L. Wu and L. Wolfenstein, Phys. Rev. Lett. **73**, 1762 (1994); A. Antaramian, L. Hall, and A. Rasin, *ibid.* **69**, 1871 (1992); L. Hall and S. Weinberg, Phys. Rev. D **48**, 979 (1993).

[5] See, e.g., D. Atwood, L. Reina, and A. Soni, Phys. Rev. D **55**, 3156 (1997).

[6] For a review see S. Bosch *et al.*, Nucl. Phys. **B565**, 3 (2000).

[7] S. Barr and A. Zee, Phys. Rev. Lett. **65**, 21 (1990).

[8] Particle Data Group, C. Caso *et al.*, Eur. Phys. J. C **3**, 1 (1998).

[9] For a review of FCNC in charm physics see S. Pakvasa, hep-ph/9705397.

[10] For reviews see *The Higgs Hunters Guide*, edited by J. Gunion, H. Haber, G. Kane and S. Dawson (Addison-Wesley, Reading, MA, 1994); *Perspectives in Higgs Physics*, edited by G. Kane (World Scientific, Singapore, 1997).

[11] See, e.g., E. Accomando *et al.*, Phys. Rep. **299**, 1 (1998); S. Dawson and L. Reina, hep-ph/9903360.

[12] For recent reviews see M. Spira, hep-ph/9810289; S. Dawson, hep-ph/9901280; Summary talks by M. Carena and H. Haber, Run II Workshop at Fermilab, 1998.

[13] For a recent discussion see T. Han and R-J. Zhang, Phys. Rev. Lett. **82**, 25 (1999).

[14] See, e.g., C. Hill, Phys. Lett. B **345**, 483 (1995) and references therein.