I. INTRODUCTION

Supersymmetry is probably the most appealing idea for solving the hierarchy problem of particle physics. Large corrections to the Higgs mass are eliminated due to a symmetry between fermions and bosons, which becomes manifest at energies of order the weak scale $v_{\text{EW}} = 174$ GeV. The problem is that many supersymmetric models, in particular, the minimal supersymmetric standard model (MSSM), imply that some superpartners and the Higgs boson should have already been discovered at LEP. The absence of such discoveries pushes these models into corners with a relatively large tuning among its basic parameters, of order 1% or worse. Naturalness suggests that supersymmetry, if present at the weak scale, should be combined with additional ingredients.

One possibility is to make the Higgs a pseudo-Goldstone boson (pGB) of a global symmetry broken at the scale $f \simeq v_{\text{EW}}$ [1,2]. This idea is usually referred to as double protection or the superlittle Higgs. In this class of models, the Higgs and the fermionic sectors are realized in a similar vein as in the little Higgs models [3], in particular, collective breaking of the global symmetry is implemented.

Combining supersymmetry with the little Higgs mechanism further softens quantum corrections to the Higgs mass. In particular, the one-loop corrections to the Higgs mass are completely finite, making electroweak symmetry breaking technically natural and alleviating the fine-tuning problem of minimal supersymmetry. The downside of softening quantum corrections are the reduced one-loop contributions to the physical Higgs mass [4]. As a consequence, it is very difficult to make the Higgs mass larger than the LEP bound of 114 GeV without reintroducing fine-tuning or complicating the models by additional structures. The main point of this paper is to show that in this class of models the Higgs mass can actually be below 114 GeV without conflicting the experimental constraints from LEP (for a related model see [5]). This is possible because the model predicts nonstandard decays for the Higgs. As advocated, for example, in [6], the LEP bounds can be relaxed if the Higgs boson decays to a final state with four light standard model (SM) states.

The global symmetry breaking pattern of our model is $SU(3)_C \times SU(3)_W \times U(1)_X$, which is embedded into an $SU(3)_C \times SU(3)_W \times U(1)_X$ gauge symmetry. This symmetry breaking pattern will ensure the existence of a fifth pseudo-Goldstone boson $\eta$ (in addition to the SM Higgs doublet). This $\eta$ is a singlet under the SM gauge symmetries, and is naturally light: lighter than the Higgs boson, and has tree-level derivative couplings to the Higgs suppressed by $v_{\text{EW}}/f$. As long as $f$ is not much larger than the electroweak scale, the decay $h \rightarrow \eta \eta$ dominates, while the branching ratio for the standard $h \rightarrow b \bar{b}$ channel can be suppressed below 20%, which is consistent with LEP bounds for a Higgs mass of order $m_2$. If this is the case, then the dominant Higgs decay channel involves at least four SM states. The composition of the final state crucially depends on the mass and couplings of the intermediate state $\eta$, which in turn depends on the embedding of the SM fermions into representations of the enlarged gauge symmetry. We will find that with our choice of matter representations the couplings of $\eta$ to the down type quarks and charged leptons can be very strongly suppressed. Thus, unlike in all previous models of the light hidden Higgs, the pseudoscalar mass does not have to be squeezed into a small window (few GeV $< m_\eta < 2 m_h$) in order to avoid the stringent LEP bounds on the $4b$ final state; instead, all the parameter space up to half the Higgs mass is available.

As the decay modes $\eta \rightarrow b \bar{b}$ and $\eta \rightarrow \tau \bar{\tau}$ are suppressed, the branching ratio for $\eta$ decaying into two charm quarks is by far the largest. The dominant Higgs decay channel is then $h \rightarrow 2 \eta \rightarrow 4 c$ for which the LEP bounds are very similar as for the $4g$ final state [7]. The branching ratio for $h \rightarrow 2 \eta \rightarrow 2 c 2 g$ (where the decay of $\eta$ into gluons now proceeds mainly via a loop of charm quarks and its symmetry partners) is at the level $10^{-2} - 10^{-1}$, while the branching ratio for decays with two photons in the final state is even more suppressed, at the level of $10^{-5} - 10^{-4}$. Since charm tagging is difficult in hadron colliders such as...
the LHC and the Tevatron, the *charming Higgs* may well be buried under the QCD background unless dedicated search strategies are devised.

### II. GAUGE SECTOR, SYMMETRY BREAKING, AND GOLDSTONE BOSONS

We consider a supersymmetric model with the Higgs arising as a pseudo-Goldstone boson of an approximate $SU(3)$ global symmetry spontaneously broken to $SU(2)$. The global $SU(3)$ is a residue of an extended gauge symmetry broken at higher energies of order $10 \text{ TeV}$. In our model, the SM gauge symmetry is extended to $SU(3)_C \times SU(3)_W \times U(1)_Y$ which is then broken by two pairs of triplet Higgses $H_u, \Phi_u = (1, \bar{3})_{1/3}$ and $H_d, \Phi_d = (1, 3)_{-1/3}$. We assume that the $\Phi$'s and $\mathcal{H}$'s do not mix in the superpotential. This leads to an enlarged $SU(3) \Phi \times SU(3) \mathcal{H}$ approximate global symmetry where the two group factors independently rotate the respective triplet pair. The $\Phi$'s are assumed to have a supersymmetric vacuum expectation value (VEV):

$$
\langle \Phi_a \rangle = \langle \Phi_d \rangle^T = (0, 0, F/\sqrt{2}),
$$

with $F \sim 10 \text{ TeV}$. This breaks the gauge group down to $SU(3)_C \times SU(2)_W \times U(1)_Y$ with the hypercharge realized as $Y = -T_3/\sqrt{3} + X$. On the other hand, $SU(3)_G$ survives down to lower energies. Ultimately, loops involving the top quark and its symmetry partners generate a negative mass squared for $H_{u,d}$ (and also the quartic term), which induces a VEV of $H_{u,d}$ of order $M_{soft}$. Then the approximate global $SU(3)_G$ symmetry is spontaneously broken to $SU(2)$ and produces five pGBs. Four of them corresponds to the Higgs doublet whose three components are nonphysical and eaten by the $W$ and $Z$ bosons. This leaves two physical pGBs. It is convenient to use the following embedding of these two pGBs into the Higgs triplets:

$$
\mathcal{H}_{u,d} = f_{u,d} (\frac{\sin(\tilde{h} / \sqrt{2} f)}{e^{\frac{\pi i}{2}} \cos(\tilde{h} / \sqrt{2} f)}),
$$

where $f_u = f s_b, f_d = f c_b, c_b = \sqrt{1 - s_b^2}$ and $t_b = s_b / c_b$ is the analogue of the MSSM $\tan \beta$. The field $\tilde{h}$ is the pGB Higgs whose VEV will break the electroweak symmetry. The other physical pGB $\tilde{\eta}$ is a singlet under the SM gauge interactions. The Higgs boson field $h$ is obtained by the shift $h = \tilde{h} + \sqrt{2} \tilde{v}$, while the canonically normalized singlet is $\eta = \eta \cos(\tilde{v} / f)$. Once the fermions are introduced (see the next section) the nonlinear sigma model scale $f$ is generated dynamically by loops of the top quark and its symmetry partners, in close analogy to generation of the Higgs VEV in the MSSM: the radiative corrections lead to negative contributions to the soft mass of $\mathcal{H}_u$ triggering the global $SU(3)_G$ symmetry breaking. We are interested in the case where $f$ is not too large, of order 350–400 GeV, which requires some mild tuning among the model parameters. The radial mode corresponding to the oscillations of $f$ (which is not a pGB) has a mass of order 200–300 GeV. The top/stop loops also generate the VEV $(\tilde{h}) = \sqrt{2} \tilde{v}$ of the Higgs field. The electroweak scale is related to the Higgs VEV by

$$
v_{EW} = f \sin(\tilde{v} / f),
$$

and the Higgs mass ends up in the range 80–90 GeV for a generic point in the parameter space.

### III. MATTER FIELDS

The third generation quarks and leptons are embedded into the following anomaly free representations:

$$
Q = (t^0, b^0, \bar{t}^0)
$$

$$
SU(3)_C \quad SU(3)_W \quad U(1)_Y
$$

$$
t_i^c = \frac{3}{2} \quad 1 \quad -2/3
$$

$$
b_c = \frac{3}{2} \quad 1 \quad 1/3
$$

$$
L_{1,2} = (\tau_{1,2}, \nu_{1,2}, \tilde{\tau}_{1,2})
$$

$$
L_c = (\nu^c_i, \tau^c_i, \tilde{\nu}^c_i)
$$

The third generation quark sector is fairly simple; in fact, it coincides with the extended quark sector of common little Higgs models. Compared to the SM, only one heavy vectorlike top quark is added. This matter content can be obtained from an underlying $SU(6) \times U(1)_Y$ symmetry, see the Appendix. The masses for all quarks can be obtained from the superpotential

$$
y_1 t_i^c \Phi_a Q + y_2 t_i^c H_u Q + \frac{y_b}{\mu_y} b_i Q \Phi_d \mathcal{H}_d.
$$

The SM top mass follows

$$
m_t = \frac{s_b y_1 y_2 F}{\sqrt{(y_1 F)^2 + 2 (s_b y_2 F)^2}} v_{EW},
$$

while the mass of the heavy top partner is

$$
m_{t'} = \sqrt{(y_1 F)^2 / 2 + (s_b y_2 F)^2}.
$$

Note that the couplings in Eq. (5) are nongeneric, as the gauge symmetry also allows for $\tilde{y}_1 \tilde{t}_i^c H_u Q + \tilde{y}_2 \tilde{t}_i^c \Phi_a Q$, which we assume to be absent or small (which in practice means $10^{-3}$ or less). In this case the approximate global $SU(3)_G$ symmetry acting on $\mathcal{H}_{u,d}$ is only collectively broken by the renormalizable couplings. That is to say, $SU(3)_G$ would be unbroken by (5) if $y_2 = 0$, and also if $y_1 = 0$ (in which case the global rotation of $\mathcal{H}_u$ could be compensated by a global rotation of $Q$). Thus, the order parameter for the explicit breaking of the global $SU(3)_G$ symmetry is proportional to $\max(y_1 F, y_2 F)$. When the collective breaking is implemented the top loop contributions to the pGB Higgs mass are finite at one-loop in a supersymmetric theory because they have to be proportional to $M_{soft}^2$, and there is simply no room for another mass parameter. In order to keep the size of loop corrections under control both $y_2 F$ and $y_1 F$ have to
be below 1 TeV (since \( F \sim 10 \text{ TeV} \) the latter requires \( y_1 \leq 0.1 \)). Given these assumptions, electroweak symmetry breaking has no fine-tuning at all.

In Eq. (5), the last term giving rise to the bottom quark mass is nonrenormalizable. This can be easily cured by integrating in a vectorlike pair of \( SU(3)_W \) and color triplets \( V, V_c \) with zero \( U(1)_\chi \) charge, and adding the couplings \( \mu_V V_c V + y b V_c Q \Phi_d + y h b V \mathcal{H}_d \) to the superpotential. For large \( \mu_V \) this leads to the effective nonrenormalizable interaction in (5) with \( y_b = y h_b y_2 \) and the resulting bottom mass is \( m_b = y h_b v_{\text{EW}} F / \sqrt{2} \mu_V \). Because of the fact that the bottom quark is much lighter than the electroweak scale we are free to choose large \( \mu_V \), as large as 10–100 TeV. Without including \( V, V_c \) the one-loop \( \beta \) function for QCD vanishes, while with these fields added it is positive which leads to a Landau pole below the grand unified theory scale. Typically, there is also a Landau pole for the \( y_2 \) Yukawa coupling which should be large in order to minimize fine-tuning. The location of the Landau pole \( \Lambda = \min \{ \Lambda_{\text{QCD}}, \Lambda_{y_2} \} \) varies between 10\(^3\) \((y_2 \approx 2.1)\) and 10\(^9\) TeV \((y_2 \leq 1.8)\).

For the second generation quarks we assume that the same matter structure is repeated, just replacing the Yukawa couplings \( y_i \rightarrow y_{ci} \). In particular, the charm mass is given by

\[
m_c = \frac{s_y c_y c_f F}{\sqrt{(y_c f)^2 + 2(s_y c_y c_f)^2}} v_{\text{EW}}. \tag{7}
\]

One can see that there are two ways to make the charm quark mass hierarchically smaller than \( v_{\text{EW}} \): either \( y_c f \ll y_{c2} f \) (where \( y_{c2} f \approx 400 \text{ GeV} \) to make the charm partner heavy enough) or the other way around. In the following we choose the former possibility; the latter leads to very suppressed couplings of \( \eta \) to the charm quarks and similar phenomenology as in [5].

The charged lepton sector is somewhat more complicated and dictated by anomaly cancellation. The SM tau lepton has three heavy partner states, while the neutrino has one; note that Majorana masses are not allowed. We have multiple options for Yukawa couplings consistent with collective symmetry breaking. Here, we concentrate on the following set:

\[
\alpha_{1j} \tau^i_1 L_1 \Phi_d + \alpha_{2j} \tau^i_2 L_2 \Phi_d + \beta_{3j} \Phi_d L_2 + \alpha_{13} \tau^i_3 L_1 \mathcal{H}_d. \tag{8}
\]

This is a collective Yukawa, since a global \( SU(3) \) emerges if one sets either \( \alpha_{13} \) or \( \alpha_{1j} \) to zero. In the absence of the \( \alpha_{13} \) term all tau partners pick up a mass proportional to \( F \) and are pushed into the TeV range. Including \( \alpha_{13} \) provides the mass for the SM tau.

**IV. HIGGS DECAYS: \( h \rightarrow 2\eta \) VS \( h \rightarrow b \bar{b} \)**

We first discuss the Higgs decay modes, and argue that the \( h \rightarrow \eta \eta \) mode dominates if the global symmetry scale \( f \) is not much larger than the electroweak scale. Even though \( \eta \) is a \( SU(2) \) singlet, it does have a tree-level derivative coupling to the Higgs field \( h \) due to \( h \) partly living in the third component of \( \mathcal{H}_{d,c} \). The symmetry preserving derivative coupling (characteristic to exact Goldstone bosons) originates from the Higgs kinetic terms,

\[
\mathcal{L}_{\text{pGB}} = \frac{1}{2} (\partial \mu \tilde{h})^2 + \frac{2}{3} \cos^2(\tilde{h}/\sqrt{2} f) (\partial \mu \tilde{h})^2. \tag{9}
\]

After canonical normalization of the pGB fields this leads to the following vertex of the Higgs boson with two singlets:

\[
\mathcal{L}_{\eta \eta} = -h (\partial \mu \eta)^2 \tan(\tilde{h}/f) / \sqrt{2} f. \tag{10}
\]

The decay width of the Higgs boson into two singlets is given by

\[
\Gamma_{h \rightarrow \eta \eta} = \frac{1}{64 \pi} \sqrt{1 - \frac{4 m_h^2}{m_h^2}} \left( \frac{1}{f^2} \right) \frac{m_h}{f^2} \Gamma_{\text{SM}}. \tag{11}
\]

The coupling of the Higgs boson to the SM quarks and leptons is the same as in the SM, up to an additional factor \( \cos(\tilde{h}/f) \) that arises due to its pGB nature [cf. the sines and cosines in the parametrization Eq. (2)]. Thus, the decay width into a pair of SM fermions is given by \( \Gamma_{h \rightarrow f \bar{f}} = \cos^2(\tilde{h}/f) \Gamma_{\text{SM}} / f \). The relevant quantity for LEP searches, customarily denoted as \( \xi_{h \rightarrow b \bar{b}}^2 \), is the branching ratio for a decay into \( b \) quarks multiplied by the suppression of the Higgs production cross section. The latter is also relevant here because the coupling of the pGB Higgs to the Z boson is multiplied, much as the Higgs-fermion coupling, by the factor \( \cos(\tilde{h}/f) < 1 \) as compared to the SM coupling. It then follows

\[
\xi_{h \rightarrow b \bar{b}}^2 = \frac{\Gamma_{\text{SM}}}{\Gamma_{h \rightarrow \eta \eta}} \frac{1 - \frac{4 m_h^2}{m_h^2}}{f^2} \left( \frac{m_h}{f^2} \right) \Gamma_{h \rightarrow f \bar{f}}. \tag{12}
\]

If \( f \) is as small as 350–400 GeV, the \( b \bar{b} \) branching ratio is at the level of 10–20% for the Higgs mass of order \( m_Z \). That is small enough to avoid exclusion by LEP. Once \( f \) is raised to around 450 GeV or higher, the generic 114.4 GeV limit from LEP cannot be significantly relaxed—the \( b \bar{b} \) branching ratio becomes large enough to have been observable at LEP.

**V. \( \eta \) COUPLINGS AND DECAYS**

The pGB pseudoscalar \( \eta \) has linear couplings to the SM fermions of the form \( i \bar{c} \gamma_j \eta \gamma_5 c \). The partial width for \( \eta \rightarrow f \bar{f} \) is given by

\[
\Gamma(\eta \rightarrow f \bar{f}) = \frac{N_c}{12 \pi} \left( \frac{4 m_f^2}{m_\eta^2} \right) \frac{m_f^2}{8 \pi^2} f. \tag{13}
\]
At tree level, the bottom quark does not couple to $\eta$ at all in the effective theory below $\mu_V$, cf. Eq. (5) [a tiny coupling suppressed by $\mu_V$ is generated when we integrate in the vectorlike pair $V, V^\dagger$.] The leading coupling generated by a penguin diagram involving two top quarks and the $W$ boson. One can estimate $\tilde{y}_b \sim 1/16\pi v \left( \frac{m_b m_{\tiny_{\text{t}}}}{v^2} \right) \log(m_{\tiny_{\text{t}}}^2/m_b^2) \sim 10^{-4}$, where $\tilde{y}_b \sim 0.2$ is the coupling of $\eta$ to the top quark. The loop factor provides enough suppression of $\tilde{y}_b$, so that $\eta$ does not decay into two bottom quarks even when it is kinematically allowed! This interesting feature distinguishes our setup from all previous hidden Higgs models in the literature where decays to $b$ quarks could be avoided only for $m_\eta < 2m_b$.

For the charm quarks, the coupling to $\eta$ originates from the term $s_b \gamma_2 f \cos(\tilde{v}/f)e^{i\tilde{e}/\sqrt{2f}}(c_\pi^2 \bar{c} \bar{c}^0)$, while for the tau lepton the relevant term is $c_\tau \bar{c}_\tau f \cos(\tilde{v}/f)e^{-i\tilde{e}/\sqrt{2f}}(\tau_2^3 \bar{c} \bar{c}^0)$. Expressing the original fields in terms of mass eigenstates one finds

$$\tilde{y}_c = \frac{m_c}{\sqrt{2f}}, \quad \tilde{y}_\tau = \frac{m_\tau}{\sqrt{2f}} \frac{m_\tau^2}{M^2}, \quad (14)$$

For $f \sim 350$ GeV we find $\tilde{y}_c \sim 10^{-3}$, while $\tilde{y}_\tau$ is additionally suppressed by the ratio of the tau mass to its heavy partner mass, $(m_\tau/M_{\tiny_{\text{t}}})^2 \sim 10^{-5}$. Given the loop suppression of $\tilde{y}_b$, the charm coupling $\tilde{y}_c$ remains by far the largest coupling. Therefore, the dominant decay channel of the pseudoscalar is $\eta \to c\bar{c}$, with the total width of order keV.

In Fig. 1 we present the branching ratios of $\eta$ for a typical point in the Yukawa space. In the entire range of $\eta$ masses, the decay into charm quarks dominates, with the next-to-leading decay into two gluons suppressed by a factor of 100.

![FIG. 1 (color online). The branching ratios of $\eta$ into bottom (blue), tau (green), charm (orange), gluons (dashed black), and photons (dashed red) as a function of its mass. We picked the following generic point in Yukawa coupling space: $y_1 = 0.109$, $y_2 = 1.8$, $y_{b1} = 0.3$, $y_{b2} = 0.365$, $y_{c1} = 0.0003$, $y_{c2} = 0.1$, $\alpha_{11} = \alpha_{22} = \beta_2 = 0.1$, $\alpha_{13} = 0.102$, while the remaining Yukawas are set to zero.](image)

VI. CONCLUSIONS

We presented a supersymmetric model where the lightest Higgs boson decays dominantly into four charm quarks. This decay is mediated by two on-shell pseudoscalars $\eta$, each subsequently decaying into two charm quarks. Besides the interesting phenomenology, our model is motivated by solving the fine-tuning problem of minimal supersymmetric theories. The Higgs is a pGB of a spontaneously broken global symmetry, and it is protected against divergent quantum corrections at one-loop, including logarithmic divergences. This opens up the possibility of completely natural electroweak symmetry breaking. The softening of the quantum correction implies that the Higgs boson cannot be much heavier than 80–90 GeV without reintroducing fine-tuning. This mass range is however perfectly allowed by all existing constraints thanks to the fact that the $h \to 4c$ decay channel is poorly constrained by the existing LEP analyses.

To our knowledge, this peculiar pattern of the Higgs branching ratios is not available in any other model in the literature. Signatures of Higgs cascade decays are currently searched for at the Tevatron [8] in the $4\tau$ and $2\tau 2\mu$ channels. The $h \to 4\tau$ topologically has also been studied in the very recent re-analysis of the LEP data [9]. These channels are motivated by the hidden Higgs models based on the next to MSSM [6]. The existence of well-motivated hidden Higgs models with suppressed decays to bottom quarks and tau leptons prompts extending the scope of collider searches. The feasibility of detecting fully hadronic final states like $4c$ and $4g$ should be assessed. Moreover, the searches should cover a larger range of the intermediate pseudoscalar masses, including $m_\eta > 2m_b$. $B$ factories unfortunately will not be able to search in the latter range [10].

Detection of this charming Higgs is challenging at hadron colliders, since four light jets will be difficult to discriminate from the QCD background, while the traditional $h \to 2\gamma$ channel will be suppressed by about a factor of 10. However it should be possible to observe a “fake” Higgs at the LHC: the radial mode associated with the oscillations around the VEV $f$ [5]. The mass of this fake Higgs is in the 200–400 GeV range, with couplings to SM fermions and gauge bosons suppressed by $\sin(\tilde{v}/f) \sim 1/2$ with respect to the usual SM Higgs coupling, so that its decay to four leptons should be observable at the LHC. Besides this heavy Higgs-like scalar, many of the super- and little-partners of the SM particles should also be observable at the LHC. In particular, naturalness dictates that the heavy (fermionic) top partner $T$ is expected below 1 TeV, in which case it will be QCD produced with a sizable cross section, and will be observable via its decay into $Zt, h t$ and $Wb$. Similarly, the stop $\tilde{t}$ should be below 1 TeV and easily available at the LHC.
ACKNOWLEDGMENTS

We would like to thank T. Han, I. Yavin, A. Schwimmer, and S. Lee for useful discussions. The work of A. F. was supported in part by the Department of Energy Grant No. DE-FG02-96ER40949. The research of B. B. and C. C. has been supported in part by NSF Grant No. PHY-0757868. C. C. was also supported in part by a U.S.-Israeli BSF grant.

APPENDIX: THE FLIPPED MATTER CONTENT

Inspired by the well-known embedding of the SM matter content into the flipped $SU(5)/C_2 \times U(1)$, we found a generalization to any $SU(N)$ group. An anomaly free representation containing the chiral matter content of one family is given by

$$\Box_{N-4} \times \Box_{(N-2)} + \frac{(N-4)(N-3)}{2} \times 1_N$$

The particular case of $SU(6) \times U(1)_X$ (with matter

$$\Box_1 + 2 \times \Box_{-2} + 3 \times 1_3$$

) contains the $SU(3)_{QCD} \times SU(3)_W \times U(1)_Z \times U(1)_X$ subgroup, where hypercharge is identified as $Y = X/3 - T_8/6$ with $T_8 = \text{diag}(1, 1, -2)$ of $SU(3)_W$.