Asymptotically Free Natural Supersymmetric Twin Higgs Model

Marcin Badziak

Institute of Theoretical Physics, Faculty of Physics, University of Warsaw, ul. Pasteura 5, PL–02–093 Warsaw, Poland;
Department of Physics, University of California, Berkeley, California 94720, USA;
and Theoretical Physics Group, Lawrence Berkeley National Laboratory, Berkeley, California 94720, USA

Keisuke Harigaya

Department of Physics, University of California, Berkeley, California 94720, USA
and Theoretical Physics Group, Lawrence Berkeley National Laboratory, Berkeley, California 94720, USA

Twin Higgs (TH) models explain the absence of new colored particles responsible for natural electroweak symmetry breaking (EWSB). All known ultraviolet completions of TH models require some nonperturbative dynamics below the Planck scale. We propose a supersymmetric model in which the TH mechanism is introduced by a new asymptotically free gauge interaction. The model features natural EWSB for squarks and gluino heavier than 2 TeV even if supersymmetry breaking is mediated around the Planck scale, and has interesting flavor phenomenology including the top quark decay into the Higgs boson and the up quark which may be discovered at the LHC.

Introduction.—Models of natural electroweak symmetry breaking (EWSB), e.g., supersymmetric (SUSY) models [1–4] and composite Higgs models [5,6], generically predict new light colored particles, called top partners, so that the quantum correction to the Higgs mass is suppressed. Null results of the LHC searches, however, show that new colored particles are heavy, which calls for fine-tuning of the parameters of the theories; this is known as the little hierarchy problem. In light of this fact the idea that the light top partners are not charged under the standard model (SM) $SU(3)_c$ gauge group has become increasingly attractive. Twin Higgs (TH) models [7] are one of the most studied realizations of the idea.

A crucial ingredient of TH models is an approximate global $SU(4)$ symmetry under which the SM Higgs boson and its mirror (or twin) partner transform as a fundamental representation. The Higgs boson is a pseudo-Nambu-Goldstone boson associated with the spontaneous breakdown of the $SU(4)$ symmetry. The $SU(4)$ symmetry of the Higgs mass term emerges from a $\mathbb{Z}_2$ symmetry exchanging the SM fields with their mirror counterparts. The light top partners are then charged under the mirror gauge group rather than the SM one. Standard lore says that ultraviolet (UV) completion of TH models involves some nonperturbative dynamics. This is because the quality of the $SU(4)$

symmetry requires a large $SU(4)$ invariant quartic term which points to UV completions based on composite Higgs models [8–11]. SUSY UV completions of the TH model also exist [12–17]. Acceptable tuning of the electroweak (EW) scale at the level of 5%–10% can be, however, obtained only with a low Landau pole scale, which requires UV completion by some strong dynamics. SUSY models that are able to keep the tuning at the level of 5%–10% without resorting to the TH mechanism also require a low cutoff scale [18].

In this Letter we propose a SUSY twin Higgs model with an asymptotically free $SU(4)$ invariant quartic coupling. The model remains perturbative up to around the Planck scale, and does not require any further UV completion below the energy scale of gravity. As a result the Yukawa couplings of the SM particles are given by renormalizable interactions.

Setup.—It was proposed in Ref. [16] that an $SU(4)$ invariant quartic coupling may be obtained from a D-term potential of a new $U(1)_X$ gauge symmetry. The model suffers from a low Landau pole scale of the $U(1)_X$ gauge interaction. A model with a non-Abelian $SU(2)_X$ gauge symmetry was proposed in Ref. [17], so that the Landau pole scale is far above the TeV scale. Still the gauge interaction is asymptotically nonfree. In order for the gauge interaction to be perturbative up to a high energy scale of $10^{16–18}$ GeV, the $SU(4)$ invariant quartic coupling at the TeV scale must be small, and the TH mechanism does not work perfectly well; fine-tuning of order 1% is required to obtain a correct EWSB scale.

In this Letter, we present an extension of the model such that the new gauge interaction is asymptotically free. In the

model presented in Ref. [17], the new gauge symmetry $SU(2)_X$ is assumed to be $\mathbb{Z}_2$ neutral, and mirror particles are charged under $SU(2)_X$. We instead assume that $SU(2)_X$ has a mirror partner $SU(2)'_X$, under which mirror particles are charged. As a result the number of $SU(2)_X$ charged fields is reduced, so that the $SU(2)_X$ gauge interaction is asymptotically free. A similar group structure in a nontwin SUSY model was introduced in Ref. [19] to achieve asymptotically free gauge theory.

The charged matter content of the model is shown in Table I. The up-type SM and mirror Higgs bosons are embedded into $\mathcal{H}$ and $\mathcal{H}'$, respectively. The resultant D-term potentials of the gauge symmetries are not $SU(4)$ invariant. Once $SU(2)_X \times SU(2)'_X$ symmetry is broken down to a diagonal subgroup $SU(2)_D$ by a nonzero vacuum expectation value (VEV) of a bifundamental Σ, both the SM and mirror Higgs bosons are fundamental representations of $SU(2)_D$, and the D-term potential is approximately $SU(4)$ invariant below the symmetry breaking scale. The $SU(2)_D$ symmetry is completely broken down by the VEVs of S, $\bar{S}$, S', $\bar{S}'$.

The right-handed top quark is embedded into $\bar{Q}_R$, so that a large enough top Yukawa coupling is obtained via the superpotential $W \sim \mathcal{H}\bar{Q}_R Q_3$, where Q_i is the ith generation of left-handed quarks. The right-handed up quark is also embedded into $\bar{Q}_R$. The VEV of ϕ_u gives a mass to the charm quark via $W \sim \phi_u \bar{u}_2 Q_2$. We assume that Yukawa couplings $\mathcal{H}\bar{Q}_R Q_{1,2}$ and $\phi_u \bar{u}_2 Q_{1,3}$ are small so that tree

level flavor changing neutral currents (FCNCs) are suppressed. H_d gives masses to down-type quarks and charged leptons via $W \sim H_d \bar{d} Q + H_d L \bar{e}$. We assume that the Yukawa couplings involving $\phi_{d,1,2}$ are suppressed; otherwise large FCNCs are induced. For details of the extended Higgs sector see the Supplemental Material [20]. Because of the $SU(2)_X$ invariance, after $\mathcal{H}$ and ϕ_u obtain their VEVs, one linear combination of the two components in $\bar{Q}_R$ remains massless at the tree level. The one loop quantum correction with a charged wino, charged Higgsinos in $\mathcal{H}$, and down-type left-handed squarks inside the loop generates the up-quark mass. The mass of the Higgsinos in $\mathcal{H}$ is given by the $SU(2)_X$ symmetry breaking and hence the loop mediates the breaking. One cannot embed a charm quark instead of an up quark into $\bar{Q}_R$ because the above loop correction is too small to generate the charm quark mass. The field $\bar{E}$ cancels the anomaly of $U(1)_Y$-$SU(2)_X^2$, while $E_{1,2}$ cancels that of $U(1)_Y^3$. The charged lepton is, in general, the mixture of $\bar{E}$, E, $\bar{e}$, and the charged component of L due to possible mixing $W \sim \bar{e}E$. The neutrino mass can be obtained by the seesaw mechanism [21–23] with the superpotential $W \sim \phi_u L N$.

In Fig. 1, we show the renormalization group (RG) running of the gauge coupling constants and the top Yukawa coupling, where we use the NSVZ beta function [24] with the anomalous dimension evaluated at the one-loop level. We see that the new gauge interaction is asymptotically free. Here and hereafter, we approximate the RG running above the $SU(2)_D$ symmetry breaking scale by that of the $SU(2)_X \times SU(2)'_X$ symmetric theory. This is a good approximation as long as the $SU(2)_X \times SU(2)'_X$ breaking scale is within the same order of magnitude as the $SU(2)_D$ breaking scale.
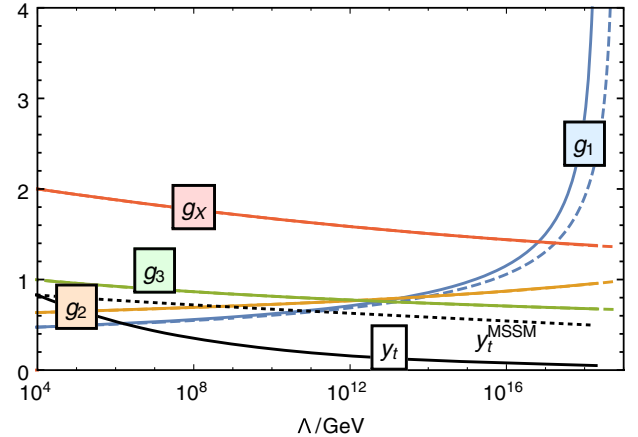
TABLE I. The charged matter content of the model. In addition to the fields shown in the table, the model contains $SU(2)_X$-neutral left-handed quarks $Q_{1,2,3}$, a right-handed charm $\bar{u}_2$, right-handed leptons $\bar{e}_{1,2,3}$, left-handed leptons $L_{1,2,3}$, and right-handed neutrino $N_{1,2,3}$ as well as their mirror partners.

	$SU(2)_X$	$SU(2)'_X$	3-2-1	3'-2'-1'
$\mathcal{H}$	2		$(\mathbf{1},\mathbf{2},1/2)$	
$\mathcal{H}'$		2		$(\mathbf{1},\mathbf{2},1/2)$
Σ	2	2		
S	2			
$\bar{S}$	2			
S'		2		
$\bar{S}'$		2		
$\bar{Q}_R$	2		$(\bar{\mathbf{3}},\mathbf{1},-2/3)$	
$\bar{Q}_R{}'$		2		$(\mathbf{3},\mathbf{1},-2/3)$
$\bar{E}$	2		$(\mathbf{1},\mathbf{1},1)$	
$\bar{E}'$		2		$(\mathbf{1},\mathbf{1},1)$
$E_{1,2}$			$((\mathbf{1},\mathbf{1},-1))$	
$E'_{1,2}$				$(\mathbf{1},\mathbf{1},-1)$
ϕ_u			$(\mathbf{1},\mathbf{2},1/2)$	
ϕ'_u				$(\mathbf{1},\mathbf{2},1/2)$
$H_d, \phi_{d,1,2}$			$(\mathbf{1},\mathbf{2},-1/2)$	
$H'_d, \phi'_{d,1,2}$				$(\mathbf{1},\mathbf{2},-1/2)$

FIG. 1. RG running of g_X (red), g_1 (blue), g_2 (yellow), g_3 (green) and the top Yukawa coupling y_t (black) for $m_X = 10$ TeV, $m_{\text{stop}} = 2$ TeV, $g_X(m_X) = 2$, and $\tan\beta = 3$. Solid lines correspond to the case where all states beyond the Minimal Supersymmetric Standard Model (MSSM) have masses around m_X. Dashed lines assume $M_{E_1} = 10^7$, $M_{E_2} = 10^9$ GeV, see text for details. Dotted black line corresponds to the running of y_t in the MSSM.

The model possesses many new states with nonzero hypercharge, which make the Landau pole scale of $U(1)_Y$ much lower than in the SM. Nevertheless, this Landau pole appears around 10^{18} GeV, as seen from Fig. 1, which is rather close to the Planck scale. The Landau pole scale is pushed up if some of the new states are much heavier than the TeV scale. Actually, we can give a large Dirac mass $M_{E_1}\bar{e}_1 E_1 + M_{E_2}\bar{e}_2 E_2$. After integrating them out, the electron and muon masses are given by a dimension-5 term $W \sim (S\phi_d\bar{E}L_i + \bar{S}\phi_d\bar{E}L_i)/M_{E_1}$. For $\mathcal{O}(1)$ coupling of $W \sim S\bar{E}E + \phi_d\bar{E}L$, the Dirac masses may be as large as $M_{E_1} \approx 10^7$, $M_{E_2} \approx 10^9$ GeV. The RG running in such a case is also shown in Fig. 1.

Let us evaluate the magnitude of the $SU(4)$ invariant coupling. We assume that Σ obtains its VEV v_Σ in a SUSY way, e.g., by a superpotential $W \sim Y(\Sigma^2 - v_\Sigma^2)$, where Y is a chiral multiplet, and that v_Σ is much larger than the TeV scale, say few tens of TeV. Then below the scale v_Σ the theory is well described by a SUSY theory with an $SU(2)_D$ gauge symmetry. The symmetry breaking of $SU(2)_D$ should involve the SUSY breaking effect, so that the D-term potential of $SU(2)_D$ does not decouple after the symmetry breaking. We introduce chiral multiplets Ξ, Ξ', and the superpotential

$$W = \kappa\Xi(S\bar{S} - M^2) + \kappa\Xi'(S'\bar{S}' - M^2), \quad (1)$$

where κ, M are constants and soft masses

$$V_{\text{soft}} = m_S^2(|S|^2 + |\bar{S}|^2 + |S'|^2 + |\bar{S}'|^2). \quad (2)$$

Here we assume that the soft masses of S and $\bar{S}$ are the same. Otherwise, the asymmetric VEVs of S and $\bar{S}$ give a large soft mass to the Higgs doublet through the D-term potential of $SU(2)_D$. Splitting between the soft masses by quantum corrections does not introduce a large soft Higgs mass [17]. In the Supplemental Material [20] more details on the masses of the $SU(2)_X \times SU(2)'_X$ sector are provided. Assuming that all Higgs bosons apart from the SM-like and twin Higgs are heavy, negligible VEV of ϕ_u and integrating out S fields, the $SU(4)$ invariant quartic coupling of the SM Higgs H and the mirror Higgs H' is given by

$$V = \frac{g_X^2}{8}\sin^4\beta(1 - \epsilon^2)(|H|^2 + |H'|^2)^2,$$

$$\epsilon^2 \equiv \frac{m_X^2}{2m_S^2 + m_X^2}, \quad (3)$$

where $\tan\beta$ is the ratio of the up-type Higgs component to the down-type Higgs component in H.

Natural electroweak symmetry breaking.—Asymptotic freedom of the new gauge interactions allows an $SU(4)$ invariant coupling of $\mathcal{O}(1)$, which enforces TH mechanism. Moreover, large g_X strongly suppresses the top Yukawa

coupling at high energy scales, as seen from Fig. 1, which results in additional suppression of the correction to the Higgs mass parameter from stops and gluino. However, for very large values of g_X, the tuning of the EW scale is dominated by a finite threshold correction from the gauge bosons of the new interaction:

$$(\delta m_{H_u}^2)_X = 3\frac{g_X^2}{64\pi^2}m_X^2\ln(\epsilon^{-2}). \quad (4)$$

For large values of g_X, which we are most interested in, the strongest lower mass limit on the new gauge boson mass of $m_X \gtrsim g_X \times 4$ TeV originates from the mixing between the Z boson and the $SU(2)_D$ gauge bosons which breaks custodial symmetry; see Ref. [17] for a detailed derivation of this bound using the EW precision observables. The threshold correction in Eq. (4) is smaller for larger ϵ which leads also to smaller $SU(4)$ invariant coupling; some intermediate value of ϵ minimizes the tuning of the EW scale. Not too small ϵ, i.e., not too heavy S fields, is also preferred to avoid a large two-loop correction to $m_{H_u}^2$ proportional to $g_X^4 m_S^2$.

In order to quantify the tuning we use the measure [14]

$$\Delta_v \equiv \Delta_f \times \Delta_{v/f}, \quad (5)$$

where the tuning in percent is $100\%/\Delta_v$ and

$$\Delta_{v/f} = \frac{1}{2}\left(\frac{f^2}{v^2} - 2\right),$$

$$\Delta_f = \max_i\left(\left|\frac{\partial\ln f^2}{\partial\ln x_i(\Lambda)}\right|, 1\right). \quad (6)$$

Here $\langle H\rangle \equiv v$, $\langle H'\rangle \equiv v'$, and $f \equiv \sqrt{v^2 + v'^2}$ is the decay constant of the spontaneous $SU(4)$ breaking. $\Delta_{v/f}$ measures the tuning to obtain $v < f$ via explicit soft $\mathbb{Z}_2$ symmetry breaking, which is required by the Higgs coupling measurements [25], implying $f \gtrsim 2.8v$ [26,27]. In our numerical analysis we fix $f = 3v$. Δ_f measures the tuning to obtain the scale f from the soft SUSY breaking. $x_i(\Lambda)$ are the parameters of the theory evaluated at the mediation scale of the SUSY breaking Λ including $m_{H_u}^2$, $m_{Q_3}^2$, $m_{\bar{u}_3}^2$, M_1^2, M_2^2, M_3^2, μ^2, m_S^2, and m_Ξ^2, where m_Ξ^2 is the soft mass of Ξ. In the following numerical analysis we assume $m_S^2 = 0$ at the mediation scale and a value of m_Ξ^2 such that $m_S = m_X$ at the $SU(2)_D$ breaking scale, corresponding to $\epsilon^2 = 1/3$, is generated via the RG running with $\kappa = 0.2$ at the mediation scale; see Ref. [17] for more details of the calculation of Δ_v.

An intriguing feature of the SUSY twin Higgs models is that the tree-level Higgs mass squared can be about twice that in the MSSM [14] so large quantum corrections are not required to obtain 125 GeV [16]. The tuning does not depend strongly on $\tan\beta$ so in the following numerical
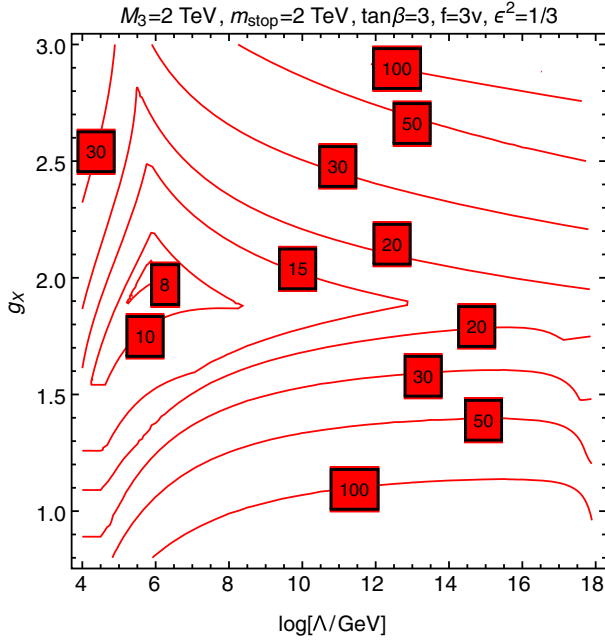
FIG. 2. Fine-tuning Δ_v of the model in the plane Λ-g_X for $m_{\text{stop}} = 2$ TeV, $\tan\beta = 3$, $f = 3v$, $\mu = 500$ GeV, $M_1 = M_2 = 200$ GeV and the soft gluino mass $M_3 = 2$ TeV. We fix $\epsilon^2 = 1/3$ which corresponds to $m_S = m_X$.
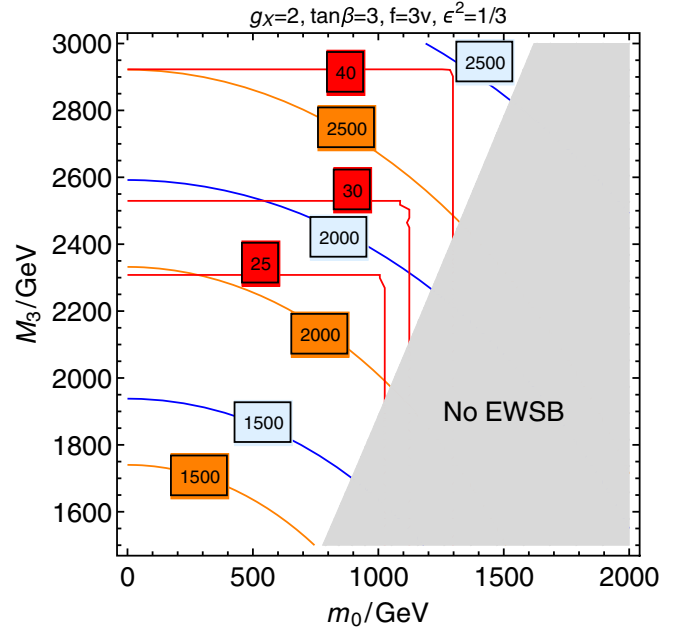
FIG. 3. Masses of the lightest (right-handed) stop and the right-handed up squark (blue) and the lightest squark of the first generation other than the right-handed up squark (orange) as a function of gluino soft mass parameter M_3 defined at the stop mass scale and universal soft scalar mass m_0 at the UV scale of $\Lambda = 10^{16}$ GeV for $g_X = 2$ and $\epsilon^2 = 1/3$. Red contours depict Δ_v.

analysis we fix $\tan\beta = 3$ which reproduces the observed Higgs mass within theoretical uncertainties for the stop masses of about 2 TeV; see Ref. [16] for more details. The tuning in the plane $\Lambda - g_X$ for $m_{\text{stop}} = M_3 = 2$ TeV at the TeV scale is shown in Fig. 2. We see that the tuning decreases with increasing g_X due to the TH mechanism as long as $g_X \lesssim 2$. For larger g_X the tuning becomes dominated by the threshold correction in Eq. (4) and the two-loop correction from the soft masses of S fields, so further increasing g_X worsens the tuning. For the optimal value of $g_X \approx 2$ the tuning is only at the level of 5%–10% even for very large mediation scales. This allows us to employ gravitational interactions as a source of SUSY breaking mediation without excessive fine-tuning, in contrast to the MSSM and previously proposed SUSY TH models.

The above discussion of tuning, similarly to all previous papers on SUSY TH models, assumed the soft stop masses at the low scale as an input without paying attention to the question of what kind of SUSY breaking mechanism can realize the spectrum. Since in this model the TH mechanism is at work also for high mediation scales, we calculate the spectrum using simple UV boundary conditions. We assume a universal soft scalar masses m_0 for the SM charged fields at the mediation scale, which explains the smallness of the flavor violation from SUSY particles. m_S^2 and m_Ξ^2 are determined in the same way as before. We fix all soft trilinear terms $A_0 = 0$ at the mediation scale. On the other hand, there is no well-motivated choice for gaugino masses since in this model the gauge couplings do not

unify. Thus, similarly as before we take gaugino masses at the low scale as input, with $M_1 = M_2 = 200$ GeV and M_3 as a variable, and include their RG running.

Using the above assumptions we show in Fig. 3 the contours of masses for the lightest stop and the lightest first-generation squark other than the right-handed up squark in the plane m_0-M_3. The lightest stop is mostly right handed and roughly degenerate with the right-handed up squark. An important constraint on the parameter space is provided by the condition of correct EWSB since the top Yukawa coupling is much smaller during the RG evolution than in MSSM, so the negative corrections from stops and gluino to $m_{H_u}^2$ [28,29] are smaller. This suppression is only partly compensated by the negative correction from the S fields. In consequence, for too large m_0, $m_{H_u}^2$ is positive at the low scale. This can be easily circumvented by assuming m_{H_u} smaller than m_0 at the mediation scale. Even without this assumption there are parts of parameter space that give EWSB as well as the viable sparticle spectrum. In this example, gluino is slightly heavier than squarks and the tuning at the level of 5% can be achieved with this simple UV boundary condition leading to squarks and gluino masses that comfortably satisfy the LHC constraints. In the Supplemental Material [20] we present a typical mass spectrum.

Flavor and collider phenomenology.—Asymptotic freedom for g_X is obtained thanks to a small number of SM

fermions charged under $SU(2)_X$. This implies a nontrivial flavor structure of the model which may impact flavor observables. As explained before, we have assumed a flavor structure in Yukawa couplings to suppress most of the tree-level FCNCs. Tree-level FCNCs are, however, unavoidable in the top sector as we embed the right-handed up quark in $\bar{Q}_R$.

The heavy Higgs in $\mathcal{H}$ which we call H_2 couples to quarks via $\mathcal{L} = y_t H_2 \bar{u}_R Q_3$. We expect non-negligible $t \to hu$ decays through mixing between the SM-like Higgs h and the neutral component in H_2 which we call H_2^0. The resultant h-t-u coupling λ_{htu} is as large as $m_Z^2/m_{H_2}^2$. The current upper limit on BR$(t \to hu)$ is 2×10^{-3} corresponding to λ_{htu} of about 0.1 [30,31], which implies a lower bound on m_{H_2} of few hundred GeV. The future sensitivity of the high-luminosity LHC to BR$(t \to hu)$ is around 10^{-4} [32], so this process will serve as an important probe of the model.

Flavor violation in the top sector has also impact on the rare decays of mesons. We find that the strongest constraint comes from a possible deviation in BR$(b \to s\gamma)$ due to one-loop corrections involving the charged component of H_2, that we refer to as $H_2^\pm$, and the up quark, which is not suppressed by the GIM mechanism [33]. Translating the bound obtained in Ref. [34] for a type-II two-Higgs-doublet model using the loop function in Ref. [35], we obtain the lower bound $m_{H_2} \gtrsim 200$ GeV.

The heavy Higgs H_2 is produced in proton colliders via the process $u + g \to H_2^+ b, H_2^0 t$ involving the strong interaction and the top Yukawa coupling, with a dominant decay mode $H_2^+ \to u\bar{b}, H_2^0 \to u\bar{t}, \bar{u}t$. None of the existing searches give relevant constraints on the masses of these new Higgs bosons.

The right-handed up squark is almost degenerate in mass with the right-handed stop and decays mainly to the top or bottom quark and a Higgsino. The signal resembles that of the right-handed stop but with a much larger cross section.

Discussion.—We have presented the first SUSY model that accommodates tuning of the EW scale 5%–10% for stops and gluino heavier than 2 TeV, even if SUSY breaking mediation occurs close to the Planck scale. This is achieved in the novel UV completion of the TH mechanism which is at work thanks to a new asymptotically free $SU(2)_X$ gauge interaction. The EW scale is much more naturally obtained than in the MSSM. The unnecessity of large stop masses nor large A terms to obtain the Higgs mass of 125 GeV improves the naturalness by a factor of $\mathcal{O}(10$–$100)$ alone. The TH mechanism and the suppression of the top Yukawa coupling further improve the naturalness by a factor of $\mathcal{O}(10)$. This scenario may be tested via flavor-violating top decays which are generically correlated with deviations from the SM prediction for the $b \to s\gamma$ decay.

[1] L. Maiani, Conf. Proc. C **7909031**, 1 (1979).
[2] M. J. G. Veltman, Acta Phys. Pol. B **12**, 437 (1981).
[3] E. Witten, Nucl. Phys. **B188**, 513 (1981).
[4] R. K. Kaul, Phys. Lett. B **109B**, 19 (1982).
[5] D. B. Kaplan and H. Georgi, Phys. Lett. **136B**, 183 (1984).
[6] D. B. Kaplan, H. Georgi, and S. Dimopoulos, Phys. Lett. **136B**, 187 (1984).
[7] Z. Chacko, H. S. Goh, and R. Harnik, Phys. Rev. Lett. **96** (2006) 231802.
[8] P. Batra and Z. Chacko, Phys. Rev. D **79** (2009) 095012.
[9] M. Geller and O. Telem, Phys. Rev. Lett. **114**, 191801 (2015).
[10] R. Barbieri, D. Greco, R. Rattazzi, and A. Wulzer, J. High Energy Phys. 08 (2015) 161.
[11] M. Low, A. Tesi, and L. T. Wang, Phys. Rev. D **91**, 095012 (2015).
[12] A. Falkowski, S. Pokorski, and M. Schmaltz, Phys. Rev. D **74**, 035003 (2006).
[13] S. Chang, L. J. Hall, and N. Weiner, Phys. Rev. D **75**, 035009 (2007).
[14] N. Craig and K. Howe, J. High Energy Phys. 03 (2014) 140.
[15] A. Katz, A. Mariotti, S. Pokorski, D. Redigolo, and R. Ziegler, J. High Energy Phys. 01 (2017) 142.
[16] M. Badziak and K. Harigaya, J. High Energy Phys. 06 (2017) 065.
[17] M. Badziak and K. Harigaya, J. High Energy Phys. 10 (2017) 109.
[18] S. Dimopoulos, K. Howe, and J. March-Russell, Phys. Rev. Lett. **113**, 111802 (2014).
[19] P. Batra, A. Delgado, D. E. Kaplan, and T. M. P. Tait, J. High Energy Phys. 02 (2004) 043.
[20] See Supplemental Material at http://link.aps.org/supplemental/10.1103/PhysRevLett.120.211803 for more details of the mass spectrum of the model.
[21] T. Yanagida, Conf. Proc. C **7902131**, 95 (1979).
[22] M. Gell-Mann, P. Ramond, and R. Slansky, Conf. Proc. C **790927**, 315 (1979).
[23] P. Minkowski, Phys. Lett. **67B**, 421 (1977).
[24] V. A. Novikov, M. A. Shifman, A. I. Vainshtein, and V. I. Zakharov, Nucl. Phys. **B229**, 381 (1983).
[25] G. Aad *et al.* (ATLAS and CMS Collaborations), J. High Energy Phys. 08 (2016) 045.
[26] D. Buttazzo, F. Sala, and A. Tesi, J. High Energy Phys. 11 (2015) 158.
[27] A. Ahmed, J. High Energy Phys. 02 (2018) 048.
[28] L. E. Ibanez and G. G. Ross, Phys. Lett. **110B**, 215 (1982).
[29] K. Inoue, A. Kakuto, H. Komatsu, and S. Takeshita, Prog. Theor. Phys. **68**, 927 (1982); **70**, 330(E) (1983).

[30] M. Aaboud *et al.* (ATLAS Collaboration), J. High Energy Phys. 10 (2017) 129.

[31] C. W. Chiang, H. Fukuda, M. Takeuchi, and T. T. Yanagida, J. High Energy Phys. 11 (2015) 057.

[32] ATLAS Collaboration, Report No. ATL-PHYS-PUB-2016-019.

[33] S. L. Glashow, J. Iliopoulos, and L. Maiani, Phys. Rev. D **2** (1970) 1285.

[34] M. Misiak and M. Steinhauser, Eur. Phys. J. C **77**, 201 (2017).

[35] A. J. Buras, M. Misiak, M. Munz, and S. Pokorski, Nucl. Phys. **B424**, 374 (1994).