

Higgs-Yukawa coupling constraints on a benchmark one-loop radiative mass model for the bottom, charm and tau

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Measurements of Higgs boson Yukawa couplings to Standard Model (SM) fermions constrain radiative mass models for those species. Such models, motivated by the observed fermion mass hierarchy, act as foils for the SM tree-level mechanism: we cannot claim to have verified the standard mechanism if other possibilities also fit the data well. We construct a benchmark model which generates the top mass at tree level and the bottom, charm, tau, and tau Dirac neutrino masses at one-loop level. Current theoretical and experimental constraints on the model, including from Higgs decays, demonstrate it possesses viable regions of parameter space. We show that future improvements to measurements will not be able to rule out the model, only increase the scale of new physics required, illustrating how difficult it will be to verify the SM fermion mass generation mechanism with great precision. As a bonus, a dark matter candidate is shown to be capable of reproducing the correct relic density within the permitted parameter space.

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I. INTRODUCTION

The observed pattern of the Standard Model (SM) fermion masses, and their apparent hierarchical distribution, is a prominent mystery in modern particle physics. While the SM is compatible with the observed masses, and the associated small Yukawa couplings are technically natural, their hierarchical distribution has no apparent dynamical origin. Radiative mass models are often motivated by this issue and attempt to solve it by generating lighter fermion masses at successively higher loop levels.

These models have been studied extensively (see Refs. [1–18], for example). However, few have considered how radiative mass generation modifies the associated Yukawa couplings, an important phenomenological consideration now that Higgs decays to fermions are being measured. In the SM, the predicted relationship between a fermion's mass and its Yukawa coupling is

$$y_f = \frac{\sqrt{2}m_f}{v}, \quad (1)$$

where v is the Higgs vacuum expectation value (VEV). This relationship is yet to be established at high precision

for any of the SM fermions [19]. Existing data constrain the Yukawa couplings for the top, bottom, and charm quarks, the τ , and the μ . Precision on such measurements is expected to increase in the coming decades for the heavier fermions [20], with upgrades to the LHC, and prospects for experiments such as the future circular collider (FCC), the international linear collider (ILC), and the circular electron positron collider (CEPC). As this occurs, it becomes important to examine how alternative models of mass generation modify (or replicate) this relationship. They can be constrained by such measurements if no disagreement with the SM is found or used to explain such a disagreement if it is measured.

Another aspect that has not been extensively studied is the generation of *all* fermion masses, except for that of the top quark, at loop level; see Refs. [4–6,10,11,17] for existing analyses of such scenarios. The motivation is simply that the top quark mass is the only fermion mass close to the electroweak (EW) scale, so it is naturally expected to arise at tree level. All of the others are suppressed, plausibly because they arise at loop level. However, many radiative models take the bottom quark and tau lepton masses to be generated at tree level along with the top mass.

The purpose of this paper is to progress the study of one-loop radiative models for the bottom, tau, and charm masses in light of Higgs Yukawa coupling constant and other constraints. These species are selected because they are all at the few-GeV level as would be expected from a $1/16\pi^2$ one-loop suppression factor. The benchmark model we present also produces a neutrino Dirac mass at the same loop level. The physical neutrino mass then must arise from

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further suppression, for example, through a seesaw effect, though we do not consider any specific implementation. While we discuss a specific model in this paper, we also intend to present a systematic catalog of related models in a companion paper. We anticipate that our work will form the basis of extensions to higher-loop radiative models for the generation of the smaller fermion masses.

We build on the work of Refs. [15–17] which explored the modified Yukawa couplings when the b quark and τ masses are generated at one loop, seeded by soft-breaking parameters for symmetries that forbid their tree-level masses. As stated above, we extend these frameworks to also produce one-loop charm and neutrino Dirac masses. We do so in light of recent updates to measurements and from the perspective of signals at colliders. We examine the current constraints and future sensitivities on this model for selected slices of parameter space and briefly outline opportunities for future research. We also extend Refs. [15–17] by demonstrating that our benchmark model contains a viable dark matter candidate, a fairly common feature in this class of theories, which adds to their importance.

In Sec. II, we define the model, and we compute current phenomenological constraints and future sensitivities in Sec. III, with the viability of the dark matter candidate established in Sec. IV. We end with discussions of future work and concluding remarks in Secs. V and VI.

II. MODEL

Our model is built using the framework presented in Ref. [17], but extended to treat upper components of electroweak doublets. We introduce four exotic fields to the theory: two vectorlike fermions ψ and χ , a complex scalar η , and a real scalar ϕ . The relevant Lagrangian is

$$\begin{aligned} -\mathcal{L} \supset & y_L^T \overline{L}_L \phi \psi_R + y_L^{Q_i} \overline{Q}_L^i \phi \chi_R + y_R^T \overline{\psi}_L \eta \tau_R \\ & + y_R^\nu \overline{\psi}_L \tilde{\eta} \nu_R + y_R^b \overline{\chi}_L \eta b_R + y_R^c \overline{\chi}_L \tilde{\eta} c_R \\ & + a H \eta^\dagger \phi + m_\psi \overline{\psi}_L \psi_R + m_\chi \overline{\chi}_L \chi_R + \text{H.c.}, \end{aligned} \quad (2)$$

where $\tilde{\eta} = i\sigma_2 \eta^*$, L_L is the third-generation lepton doublet, and $Q_L^{2(3)}$ is the second- (third)-generation quark doublet. We assume η does not acquire a VEV.

The simplest quantum number assignments for the relevant fields given this Lagrangian are shown in Table I. We remain agnostic as to the structure of a full theory treating all three generations of quarks and leptons and thus do not consider any terms coupling exotics to the remaining SM fermions in our analysis.

This Lagrangian admits three softly broken symmetries, denoted by $U(1)_a$, $U(1)_\psi$, and $U(1)_\chi$. The first of these is softly broken by the trilinear term, and $U(1)_\psi$ and $U(1)_\chi$ are softly broken by the ψ and χ mass terms, respectively. The tree-level Yukawa terms for b and c will be hard-breaking terms for $U(1)_{a,\chi}$, while the τ and ν_τ tree-level Yukawa terms will be hard-breaking terms for $U(1)_{a,\psi}$, and thus all four terms are forbidden. None of these three $U(1)$ s can be imposed on the theory as exact symmetries, since this would prevent at least one SM fermion from gaining mass at any order. We do, however, impose exotic particle parity $\mathcal{Z}_2$ (defined in Table I) as an exact symmetry of the theory, to prevent mixing between exotics and SM fields. As a result, the lightest exotic will be stable and thus a dark matter candidate if it is colorless and electrically neutral. We restrict our analysis to regions of parameter space where the lightest exotic satisfies these conditions, and all other exotics can decay into it along with some combination of SM fields.

After electroweak symmetry breaking (EWSB), ϕ will mix with the real part of the lower component η^{low} of η , with mixing angle θ_s , working in the basis in which $\text{Im}\{\eta^{\text{low}}\}$ does not mix with ϕ and $\text{Re}\{\eta^{\text{low}}\}$ at tree level. We denote the mass eigenstate superpositions of ϕ and $\text{Re}\{\eta^{\text{low}}\}$ as $\varphi_{1,2}$, with φ_1 being lighter. The resulting scalar mass eigenvalues are

$$m_{1,2}^2 = \frac{1}{2} \left(m_\phi^2 + m_\eta^2 \pm \sqrt{(m_\eta^2 - m_\phi^2)^2 + 4a^2 v^2} \right). \quad (3)$$

TABLE I. Quantum numbers of relevant fields in our model, including extended lepton number L and extended baryon number B . The $SU(3)_C$ labels are the only allowable ones using fundamental or singlet representations. The $SU(2)_L$ labels are chosen to give the color singlet exotics integer electric charge. See the main text for an explanation of the role played by each of the non-SM symmetries.

	τ_R	ν_R	L_L	b_R	c_R	Q_L	H	ψ_L	ψ_R	χ_L	χ_R	ϕ	η
$SU(3)_C$	1	1	1	3	3	3	1	1	1	3	3	1	1
$SU(2)_L$	1	1	2	1	1	2	2	2	2	2	2	1	2
$U(1)_Y$	-1	0	$-\frac{1}{2}$	$-\frac{1}{3}$	$\frac{2}{3}$	$\frac{1}{6}$	$\frac{1}{2}$	$-\frac{1}{2}$	$-\frac{1}{2}$	$\frac{1}{6}$	$\frac{1}{6}$	0	$\frac{1}{2}$
L	1	1	1	0	0	0	0	1	1	0	0	0	0
$3B$	0	0	0	1	1	1	0	0	0	1	1	0	0
$\mathcal{Z}_2$	1	1	1	1	1	1	1	-1	-1	-1	-1	-1	-1
$U(1)_a$	1	-1	0	1	-1	0	0	0	0	0	0	0	-1
$U(1)_\psi$	1	1	0	0	0	0	0	1	0	0	0	0	0
$U(1)_\chi$	0	0	0	1	1	0	0	0	0	1	0	0	0

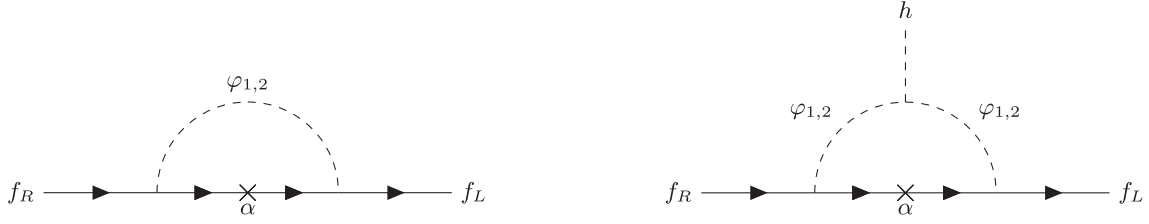


FIG. 1. Mass generation (left) and effective Yukawa coupling (right) diagrams for $f = \tau, \nu_\tau, b, c$ at one loop after EWSB. For τ and ν_τ , the virtual fermion α is ψ , while $\alpha = \chi$ for b and c .

We also obtain the relationship

$$\sin 2\theta_s = \frac{2av}{m_2^2 - m_1^2}. \quad (4)$$

These results differ from those in Ref. [17] only by certain factors of 2, which arise from real rather than complex scalars mixing.

For ease of notation, we also define

$$\alpha = \begin{cases} \psi, & \text{for } f = \tau, \nu_\tau \\ \chi, & \text{for } f = b, c. \end{cases} \quad (5)$$

The Feynman diagrams for SM fermion masses and Yukawa couplings after EWSB are given in Fig. 1. Our expressions for the masses and Yukawa couplings are identical to those in Ref. [17],

$$y_f^{\text{eff}}|_{p_h^2=0} = \frac{y_L^f y_R^f}{16\pi^2} \frac{am_\alpha}{m_1 m_2} \left[\cos^2 2\theta_s F\left(\frac{m_1^2}{m_\alpha^2}, \frac{m_2^2}{m_\alpha^2}\right) + \frac{\sin^2 2\theta_s}{2} \left(\frac{m_2}{m_1} F\left(\frac{m_1^2}{m_\alpha^2}\right) + \frac{m_1}{m_2} F\left(\frac{m_2^2}{m_\alpha^2}\right) \right) \right],$$

$$m_f = \frac{v}{\sqrt{2}} \frac{y_L^f y_R^f}{16\pi^2} \frac{am_\alpha}{m_2 m_1} F\left(\frac{m_1^2}{m_\alpha^2}, \frac{m_2^2}{m_\alpha^2}\right), \quad (6)$$

where we have defined

$$F(x, y) = \frac{\sqrt{xy}}{x-y} \left(\frac{x}{x-1} \ln(x) - \frac{y}{y-1} \ln(y) \right),$$

$$F(x) = \lim_{y \rightarrow x} F(x, y) = \frac{x}{x-1} - \frac{x}{(x-1)^2} \ln(x). \quad (7)$$

The full expression for y_f^{eff} is unwieldy, so we have provided it here in the $m_h \rightarrow 0$ limit but use the full version for our plots. Since $F(x, y) > 0 \forall x, y > 0$, we have $y_f^{\text{eff}} > y_f^{\text{SM}}$ in this model.

The t and c quarks will mix at one loop, since t_L couples to the exotics via the $y_L^{Q_3}$ term. The effective Yukawa and mass terms at one loop are

$$-\mathcal{L} \supset (\overline{t_L} \quad \overline{c_L}) \begin{pmatrix} \frac{y_t v}{\sqrt{2}} & \frac{y_L^{Q_3}}{y_L^{Q_2}} m_c \\ \frac{(y_L^{Q_3} y_R^c)^*}{y_L^{Q_2} y_R^c} m_c & m_c \end{pmatrix} \begin{pmatrix} t_R \\ c_R \end{pmatrix} + \text{H.c.}$$

$$+ (\overline{t_L} \quad \overline{c_L}) \begin{pmatrix} y_t & \frac{y_L^{Q_3}}{y_L^{Q_2}} y_c^{\text{eff}} \\ \frac{(y_L^{Q_3} y_R^c)^*}{y_L^{Q_2} y_R^c} y_c^{\text{eff}} & y_c^{\text{eff}} \end{pmatrix} \begin{pmatrix} t_R \\ c_R \end{pmatrix} \frac{h}{\sqrt{2}} + \text{H.c.}, \quad (8)$$

where m_c, y_c^{eff} are as defined in Eq. (6). In general, these will not be simultaneously diagonalizable, and thus there will exist an interaction between t, c , and h , in the mass basis. Diagonalizing the mass matrix is trivial, so we refrain from displaying the unwieldy formulas for the eigenvalues and mixing angles. In the mass basis, we rewrite the effective Yukawa coupling terms for t and c as

$$-\mathcal{L} \supset (\overline{t'_L} \quad \overline{c'_L}) Y_{\text{eff}} \begin{pmatrix} t'_R \\ c'_R \end{pmatrix} \frac{h}{\sqrt{2}} + \text{H.c.} \quad (9)$$

Taking $y_L^3 = y_R^b$ and $y_L^2 = y_R^c$ to obtain indicative values, we find

$$Y_{\text{eff}} = \begin{pmatrix} 0.991 & 0.014 \\ 0.014 & 0.00018 \end{pmatrix} + \begin{pmatrix} 0.015 & -0.563 \\ -0.563 & 0.297 \end{pmatrix} y_b^{\text{eff}}, \quad (10)$$

where y_b^{eff} is defined in Eq. (6). The partial decay width and branching ratio for the novel $t \rightarrow ch$ decay depends on the off-diagonal elements of Y_{eff} .

III. PHENOMENOLOGY

We considered a variety of constraints to determine whether this model is phenomenologically viable. The theoretical constraints we examine are the perturbative unitarity of the various $y_{L,R}$, and fine-tuning for the Higgs mass parameter μ_H^2 . Specifically, for the former, we can calculate the product $y_L^f y_R^f$ by requiring the observed fermion mass for f be reproduced and then exclude $y_L^f y_R^f > 4\pi$ on our plots. The constraint on the b couplings

is always strongest for the slices of parameter space shown here, so we do not explicitly show constraints on the other couplings. We adopt the fine-tuning measure Δ defined in Ref. [21] and take $\Delta > 100$ to be our benchmark for a high degree of fine-tuning. For consistency with Ref. [17], we take the matching scale to be $\sqrt{m_\phi m_\eta}$. While regions with $\Delta > 100$ may be technically permitted, introducing an unnatural hierarchy between the electroweak and new physics scales undermines the motivation of explaining the hierarchy of fermion masses. We do not show the full perturbative unitarity constraint on the trilinear coupling constant a , instead noting that it results in an upper bound on a roughly of the form $\mathcal{O}(1) \cdot m_1$. Indicative perturbativity bounds on a can be found in Ref. [17].

The experimental constraints we consider are the most recent bounds on the signal strengths for Higgs decays to the relevant fermions, modifications to EW physics, and the novel $t \rightarrow ch$ decay. The signal strength for a process $ii \rightarrow h \rightarrow ff$ is defined in the κ formalism [22,23] as

$$\mu_f^i = \frac{\kappa_i^2 \kappa_f^2}{\kappa_h^2}. \quad (11)$$

All signal strengths are normalized to unity at the SM prediction. In this model, there is no modification (at tree or one-loop level) to the Higgs boson coupling to gluons or photons, and a negligible modification to its coupling to the W and Z bosons, so we consider only the modified Yukawa couplings and Higgs width. The current constraints on the relevant signal strengths are shown in Table II. Note that for $m_\psi = m_\chi$ our model predicts $\kappa_\tau = \kappa_b = \kappa_c$ and thus $\mu_\tau = \mu_b = \mu_c$, so the strongest constraint will always be provided by the signal strength with the most precise measurement.

Corrections to EW physics will primarily occur through contributions to the vacuum polarisations of EW gauge bosons, rather than box or vertex corrections, justifying use of the Peskin-Takeuchi parameters [28,29]. We only consider corrections to the S and T parameters, ignoring any modifications to U since these are typically smaller. Only new physics contributions are considered, so $S = T = U = 0$ is the SM prediction. Under the assumption that $U = 0$, the current 1σ constraints are [19]

TABLE II. Signal strengths for Higgs boson decays into relevant SM fermions [19], normalized to unity at the SM prediction. The dominant contributions to these values are from ATLAS and CMS data [24–27].

Process	$h \rightarrow \tau^+ \tau^-$	$h \rightarrow b\bar{b}$	$h \rightarrow c\bar{c}$
μ_f	0.91 ± 0.09	0.94 ± 0.11	-0.5 ± 3.4

$$\begin{aligned} S &= 0.008 \pm 0.071 \\ T &= 0.021 \pm 0.055. \end{aligned} \quad (12)$$

Neither ψ nor χ will provide any net contribution to these. We find that T provides a stronger bound than S in all regions of parameter space we explore, so we explicitly display the formula and bounds for T only. The one-loop scalar contributions to T are

$$\begin{aligned} T &= \left[\frac{2s_s^2 c_s^2 m_1^2 m_2^2}{m_1^2 - m_2^2} \ln\left(\frac{m_1^2}{m_2^2}\right) + m_1^2 c_s^4 + m_2^2 s_s^4 + m_\eta^2 \right. \\ &\quad \left. + \frac{2c_s^2 m_1^2 m_\eta^2}{m_1^2 - m_\eta^2} \ln\left(\frac{m_\eta^2}{m_1^2}\right) + \frac{2s_s^2 m_2^2 m_\eta^2}{m_2^2 - m_\eta^2} \ln\left(\frac{m_\eta^2}{m_2^2}\right) \right] \\ &\quad \times \frac{1}{16\pi c_W^2 s_W^2 m_Z^2}, \end{aligned} \quad (13)$$

where we have denoted $\sin \theta_i = s_i$ and $\cos \theta_i = \theta_i$.

The rare $t \rightarrow ch$ decay will be induced by mixed interactions of t and c . The partial decay width for $t \rightarrow ch$ is

$$\Gamma_{t \rightarrow ch} = \frac{m_t}{32\pi} \left(1 - \frac{m_h^2}{m_t^2}\right)^2 (0.014 - 0.563 y_b^{\text{eff}})^2 \quad (14)$$

using the values from Eq. (10). We can compare this to the partial decay width

$$\Gamma_{t \rightarrow bW} = \frac{|V_{tb}|^2 m_t^3}{16\pi v^2} \left(1 + 2 \frac{m_W^2}{m_t^2}\right) \left(1 - \frac{m_W^2}{m_t^2}\right)^2. \quad (15)$$

Assuming that $\text{Br}(t \rightarrow bW)$ is approximately unity, we write

$$\text{Br}(t \rightarrow ch) \simeq \frac{\Gamma_{t \rightarrow ch}}{\Gamma_{t \rightarrow bW}} = 0.261 \cdot (0.014 - 0.563 y_b^{\text{eff}})^2. \quad (16)$$

Similar to μ_b , this expression has a quadratic dependence on y_b^{eff} , so these bounds will have a similar form. The current 2σ upper bound on $\text{Br}(t \rightarrow ch)$ is 3.4×10^{-4} [30], though it could be brought as low as 1.6×10^{-5} with the next generation of experiments [31]. We find that the projected bounds on the branching fraction are weaker than the current bound on μ_τ for the parameter space considered, and thus do not show them here.

Current 2σ constraints on this model are shown for two benchmark slices of parameter space in Fig. 2. As can be seen from these plots, viable parameter space with $<1\%$ fine-tuning currently exists for this model. This indicates that it remains sensible to use this model as a foundation for constructing a complete model. The most stringent constraints are generally perturbativity, the μ_τ signal strength, and corrections to T .

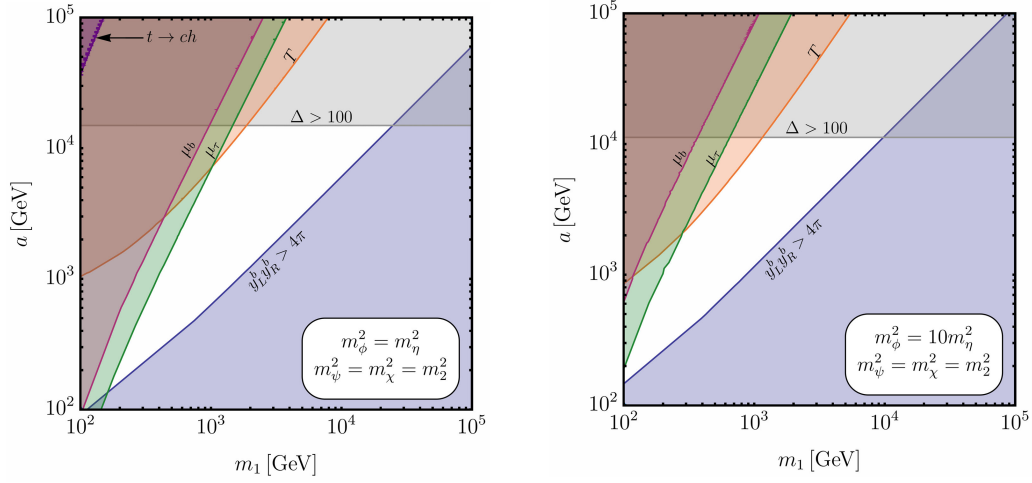


FIG. 2. Theoretical and current 2σ experimental constraints on the model for two benchmark slices of parameter space. Shaded regions are excluded, except for the case of $\Delta > 100$ where they are disfavored.

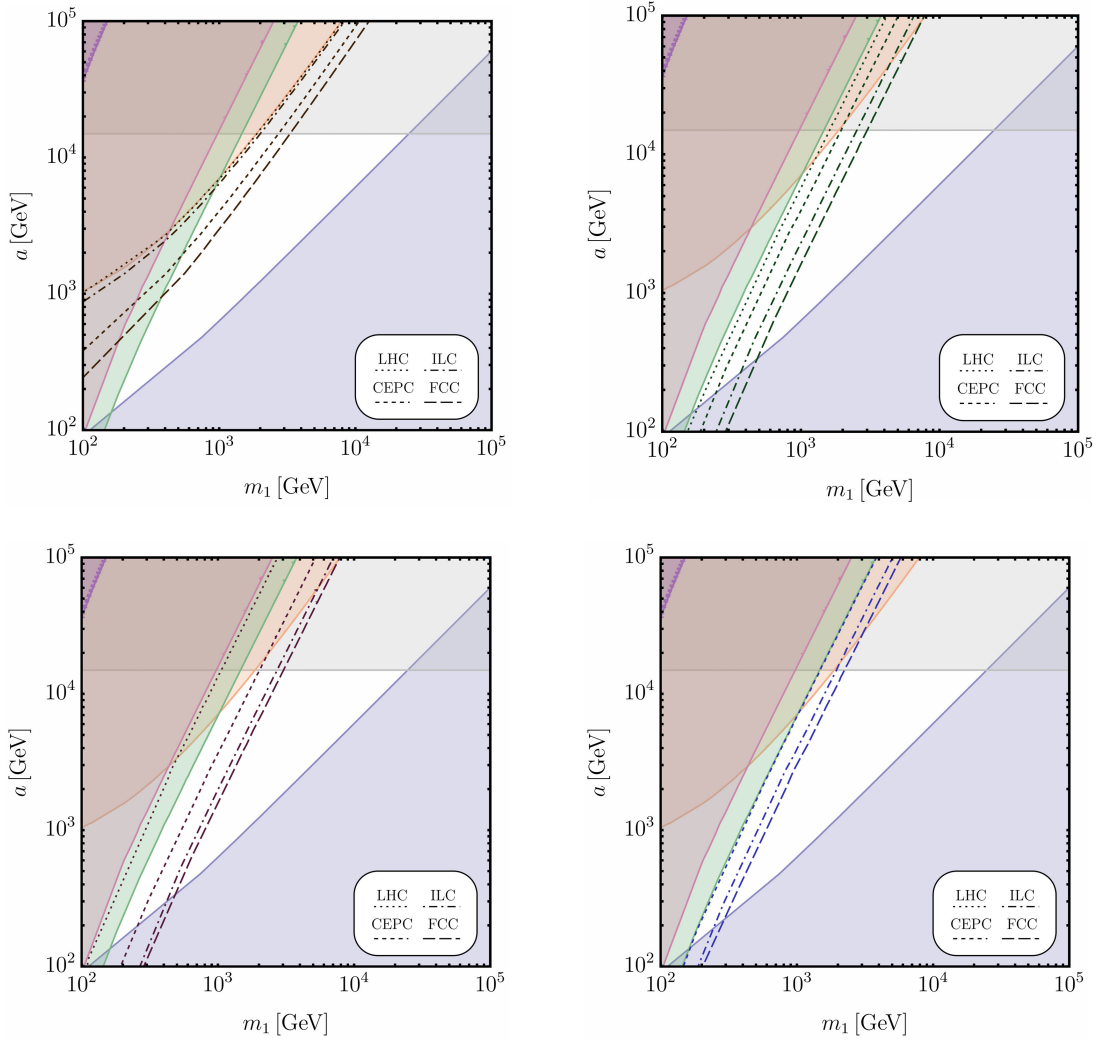


FIG. 3. Projected 2σ experimental sensitivities when $m_\phi^2 = m_\eta^2$ and $m_\psi^2 = m_\chi^2 = m_2^2$. Plots shown are the sensitivities for T (top left), μ_τ (top right), μ_b (bottom left), and μ_c (bottom right). Values for ΔT taken from Ref. [32] and for $\Delta\kappa_\tau$, $\Delta\kappa_b$, and $\Delta\kappa_c$ from Ref. [20]. Shaded regions correspond to the same constraints as in Fig. 2.

Bounds obtained from projected future sensitivities [20,32] for μ_f and T are shown in Fig. 3, assuming agreement with the SM is maintained. Assumptions on the integrated luminosities and center-of-mass energies for these projections are given in Table 1 of Ref. [20], where we have considered the HL-LHC, ILC₅₀₀, the full CEPC program, and the full FCC-ee/eh/hh program in our plots.

Another avenue for constraining such a model is the bounds obtained from direct searches at colliders. Since χ is the only colored exotic, we would expect the bound on its mass to be strongest. The relevant searches are for vector-like quarks (VLQs). However none of the existing VLQ searches can be repurposed, as they assume that the given VLQ primarily decays into SM fields [33], something χ cannot do by the imposition of exotic $\mathcal{Z}_2$. Future searches for VLQs that relax this assumption, or even consider VLQs that cannot decay into exclusively SM fields, would thus be helpful for constraining models that contain fields similar to χ . Existing bounds on squarks and VLQs suggest a lower bound of $m_\chi \gtrsim 1$ TeV would be typical [34–36].

To repurpose bounds from direct searches at colliders for other fields, we note that ϕ resembles a right-handed antineutrino, η resembles a left-handed antilepton doublet, and ψ resembles an anti-Higgsino. Assuming $m_\psi \ll m_\eta$, the lower bound on the η mass is $m_\eta \gtrsim 425$ GeV [37]. This bound gets weaker as m_ψ increases, and there is no constraint for $m_\psi \gtrsim 140$ GeV. This bound thus does not exclude any significant region of parameter space.

IV. DARK MATTER

Our model contains two dark matter candidates, the scalar ϕ_1 and the upper component of ψ . If ψ^{up} is the lightest field, and thus a dark matter candidate, its relic abundance is set by (co)annihilation into τ and ν_L . ψ^{up} is pseudo-Dirac,¹ so neglecting the masses of the SM fields allows us to write the thermally averaged cross section as [38]

$$\langle\sigma v\rangle = \frac{m_\psi^2}{128\pi} \left(\frac{4(y_L^L)^4}{(m_\psi^2 + m_\phi^2)^2} + \frac{(y_R^\tau)^4}{(m_\psi^2 + m_\eta^2)^2} \right). \quad (17)$$

We consider a freeze-out scenario and perform our calculation before EWSB, assuming that this regime constitutes the majority of the temperature region between m_ψ and the freeze-out temperature $T_f \sim m_\psi/20$. Such an assumption is well founded for m_ψ above the TeV level, which turns out to be a highly relevant regime. Requiring that the correct dark matter relic density be reproduced gives the plot in Fig. 4, with blue dots corresponding to parameter values for which this is accomplished. This shows that our model

¹It will receive a Majorana mass at one loop courtesy of ν_R ; however, we expect this mass to be small given large m_{ν_R} .

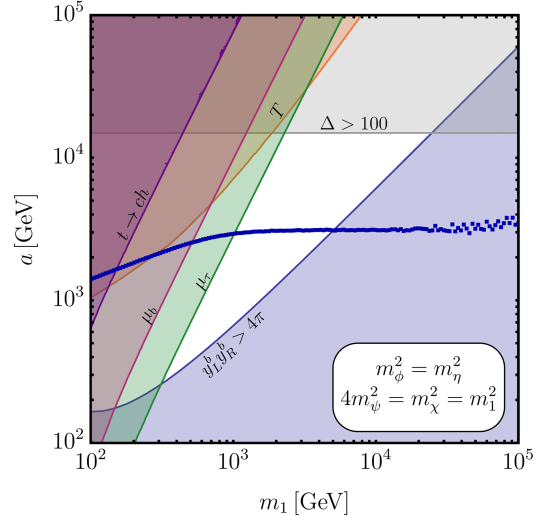


FIG. 4. Exclusion plot with ψ^{up} as a dark matter candidate. Blue dots are points for which the observed relic density is reproduced, given a set value for m_ψ .

contains at least one viable dark matter candidate. A more thorough investigation would be required to fully examine all the possibilities.

V. FUTURE WORK

The most obvious and most important extension of this work is building a model treating all three generations of SM fermions, with lighter masses generated at progressively higher loop levels. This would require some additional symmetry structure, and possibly the introduction of more exotic fields to enable mass generation at higher loops. We note that for a straightforward inclusion of the fermions from other generations, assuming they are charged similarly under the new symmetries, the mass matrices for a fermion type will have the form

$$M^{ij} = y_L^i y_R^j \times \text{const.} \quad (18)$$

This matrix is an outer product of two vectors and is thus rank 1, implying only one fermion of each type can gain mass without further modifications to the theory.

A similar relic abundance analysis could be conducted for the other dark matter candidate, and the dark matter phenomenology for either candidate could also be explored further.

Finally, we note that it is fairly easy to include an axion in this model. If we introduce a scalar that transforms as $S \sim (\mathbf{1}, \mathbf{1}, 0)(1)$ under $\mathcal{G}_{\text{SM}} \otimes U(1)_\chi$, then the UV Lagrangian would include the term $S \overline{\chi} \chi$. When S gets a VEV, the $U(1)_\chi$ symmetry would be spontaneously broken, resulting in an axion field and the aforementioned term becoming the χ mass term.

VI. CONCLUSION

We have extended an earlier framework to present a benchmark model that has one-loop-generated Dirac masses for the τ , b , c , and ν_τ , thus explaining why $m_{\tau,b,c}$ are at the few-GeV level. We demonstrated that this model is currently viable in certain regions of parameter space and showed how the parameter space will be further probed by future experiments, opening the prospect for discovery.

Even if a discovery is not made, ensuing future constraints which maintain agreement with the SM will not be able to rule out such a model, only push the scale of new physics higher and necessitate a greater degree of fine-tuning, making the model less theoretically compelling. It is only in that sense that the SM tree-level fermion mass generation mechanism may become favored.

We briefly outlined avenues for future work in this area, including creating extended models that treat all three generations of SM fermions, exploring the phenomenology of dark matter candidates in the model, and/or introducing an axion to the theory. Note that a full three-family model will likely feature non-SM flavor effects for the lighter quarks and leptons, just as the presented model has the

$t \rightarrow ch$ flavor-changing top-quark decay. These processes are likely to be the best way to constrain complete radiative models. We expect, for example, that the masses of some exotics in some of these models will have to be so large to meet the flavor-violation constraints that consistency with electroweak naturalness will be impossible. A truly successful model would have to avoid this fate, and the search for such a model is a worthy goal of future research.

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DATA AVAILABILITY

There are no publicly available research data or software supporting this manuscript. Requests for further information or data should be sent to the authors.

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