

Dirac neutrinos from a second Higgs doublet

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We propose a minimal extension of the standard model in which neutrinos are Dirac particles and their tiny masses are explained without requiring tiny Yukawa couplings. A second Higgs doublet with a tiny vacuum expectation value provides neutrino masses while simultaneously improving the naturalness of the model by allowing a heavier standard-model-like Higgs boson consistent with electroweak precision data. The model predicts a $\mu \rightarrow e\gamma$ rate potentially detectable in the current round of experiments, as well as distinctive signatures in the production and decay of the charged Higgs H^+ of the second doublet which can be tested at future colliders. Neutrinoless double beta decay is absent.

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I. INTRODUCTION

Since the discovery of neutrino oscillations, many models of neutrino mass have been constructed. The most straightforward is to incorporate Dirac neutrino masses into the standard model (SM) by introducing three (or two [1]) right-handed neutrinos ν_R coupled to the SM Higgs doublet Φ_1 analogously to the SM quarks and charged leptons. The Yukawa Lagrangian becomes

$$\begin{aligned} \mathcal{L}_{\text{Yuk}} = & -y_{ij}^d \bar{d}_{R_i} \Phi_1^\dagger Q_{L_j} - y_{ij}^u \bar{u}_{R_i} \tilde{\Phi}_1^\dagger Q_{L_j} - y_{ij}^e \bar{e}_{R_i} \Phi_1^\dagger L_{L_j} \\ & - y_{ij}^{\nu} \bar{\nu}_{R_i} \tilde{\Phi}_1^\dagger L_{L_j} + \text{H.c.}, \end{aligned} \quad (1)$$

where $\tilde{\Phi}_1 \equiv i\sigma_2 \Phi_1^*$ is the conjugate of the Higgs doublet, $Q_L \equiv (u_L, d_L)^T$ and $L_L \equiv (\nu_L, e_L)^T$ are the left-handed quark and lepton doublets, respectively, and a sum over the generation indices i, j is implied.

This implementation has two problems. First, realistic Dirac neutrino masses below ~ 1 eV require nine independent dimensionless Yukawa couplings $|y_{ij}^{\nu}| \lesssim 10^{-11}$. Second, the right-handed neutrinos ν_R are uncharged under the SM gauge group, so that Majorana mass terms $M_{ij} \nu_{R_i} \nu_{R_j}$ are allowed by the gauge symmetry. The Majorana mass terms can be eliminated by imposing a global symmetry such as lepton number. Alternatively, a large Majorana mass for the right-handed neutrinos leads to naturally light left-handed Majorana neutrinos via the type-1 seesaw [2], $m_\nu \sim (y^\nu v_1)^2/M$, where $v_1 \simeq 246$ GeV is the SM Higgs vacuum expectation value (vev). This possibility motivates the experimental search for neutrinoless double beta decay, which can happen only if the neutrino is a Majorana particle. Most other neutrino mass models also yield Majorana neutrinos.

In this paper we introduce a minimal model for Dirac neutrino masses that does not require tiny neutrino Yukawa couplings. Our motivation is to provide a viable, renormalizable model with minimal new field content which appears entirely below the TeV scale. The smallness of the

neutrino masses relative to those of the quarks and charged leptons is explained by sourcing them from a second Higgs doublet Φ_2 with a tiny vev $v_2 \sim$ eV. The second Higgs doublet yields two neutral scalars and a charged scalar pair at the electroweak scale, providing signatures at the CERN Large Hadron Collider (LHC) that can be used to discriminate the model from other neutrino mass models and to extract new information about the neutrino masses. The charged scalar also contributes to the lepton flavor violating decay $\mu \rightarrow e\gamma$ at a rate potentially within reach of the currently running MEG experiment, which would provide additional sensitivity to the model parameters. The model can be made consistent with all existing experimental constraints, including standard big bang nucleosynthesis (BBN) which generally constrains models with new light degrees of freedom. The second doublet has the additional benefit of allowing a heavier SM-like Higgs boson to be consistent with the precision electroweak data, thereby easing the fine-tuning problem of a light Higgs boson [3].

In the next two sections we describe the model and show that consistency with BBN can be achieved by keeping the neutrino Yukawa couplings y_i^ν below about 1/30. In Sec. IV we discuss the phenomenology. After considering the decay modes of the new Higgs bosons, we derive a constraint on the charged Higgs mass from existing data from the CERN Large Electron-Positron (LEP) collider. We then address predictions for $\mu \rightarrow e\gamma$. We also show that the model is consistent with constraints from the muon anomalous magnetic moment and tree-level muon and tau decay. We finish with a discussion of LHC search prospects for the charged Higgs and possible effects on the phenomenology of the SM-like Higgs. Finally we summarize our conclusions in Sec. V.

II. THE MODEL

The field content is that of the SM with the addition of a new scalar doublet Φ_2 —with the same gauge quantum numbers as the SM Higgs doublet Φ_1 —and three gauge-singlet right-handed neutrino fields ν_{R_i} which will pair up with the three left-handed neutrinos of the SM to form

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Dirac particles. We impose a global U(1) symmetry under which the new fields Φ_2 and ν_{R_i} carry charge +1 while all SM fields are uncharged. This U(1) symmetry is needed to forbid Majorana mass terms for the ν_{R_i} while simultaneously enforcing a Yukawa coupling structure in which only Φ_2 couples to right-handed neutrinos. The fourth term in Eq. (1) is then replaced according to

$$-y_{ij}^{\nu} \bar{\nu}_{R_i} \tilde{\Phi}_1^{\dagger} L_{L_j} \rightarrow -y_{ij}^{\nu} \bar{\nu}_{R_i} \tilde{\Phi}_2^{\dagger} L_{L_j}. \quad (2)$$

If the U(1) symmetry is unbroken, Φ_2 has zero vev and the neutrinos are strictly massless [4]. An identical Yukawa structure can be obtained using a Z_2 symmetry, as in the models of Refs. [5,6]; however, this does not forbid Majorana mass terms for ν_{R_i} .

In order to generate a vev for Φ_2 , we break the global U(1) explicitly using a dimension-2 term in the Higgs potential of the form $m_{12}^2 \Phi_1^{\dagger} \Phi_2$. This results in a seesaw-like relation [5],

$$v_2 = m_{12}^2 v_1 / M_A^2, \quad (3)$$

where M_A is the mass of the neutral pseudoscalar (defined below). For $M_A \sim 100$ GeV, $v_2 \sim \text{eV}$ is achieved for m_{12}^2 of order (a few hundred keV)². By breaking the global U(1) explicitly we avoid an extremely light scalar as is present in the model of Ref. [6]; this allows us to satisfy BBN constraints without resorting to nonstandard cosmology.

With these considerations the scalar potential is

$$\begin{aligned} V = & m_{11}^2 \Phi_1^{\dagger} \Phi_1 + m_{22}^2 \Phi_2^{\dagger} \Phi_2 - [m_{12}^2 \Phi_1^{\dagger} \Phi_2 + \text{H.c.}] \\ & + \frac{\lambda_1}{2} (\Phi_1^{\dagger} \Phi_1)^2 + \frac{\lambda_2}{2} (\Phi_2^{\dagger} \Phi_2)^2 \\ & + \lambda_3 (\Phi_1^{\dagger} \Phi_1) (\Phi_2^{\dagger} \Phi_2) + \lambda_4 (\Phi_1^{\dagger} \Phi_2) (\Phi_2^{\dagger} \Phi_1). \end{aligned} \quad (4)$$

We can choose m_{12}^2 real and positive without loss of generality by rephasing Φ_2 and putting the excess phase into y_{ij}^{ν} , which is already a general complex matrix. Stability of the potential requires $\lambda_1, \lambda_2 > 0$, $\lambda_3 > -\sqrt{\lambda_1 \lambda_2}$, and $\lambda_4 > -\sqrt{\lambda_1 \lambda_2} - \lambda_3$. Note that even after the global U(1) is broken, conventional lepton number survives as an accidental symmetry of the model. Neutrinoless double beta decay is thus absent.¹

¹As in the SM, higher-dimensional operators can violate lepton number; we assume such operators are Planck-suppressed. The leading contribution is from $(L_L \Phi_1^{\dagger})^2 / M_{\text{Pl}}$, which yields a Majorana mass term for the left-handed neutrinos of order 3×10^{-3} meV—compare effective electron neutrino Majorana masses $|m_{ee}| > 14$ meV for the inverted neutrino mass hierarchy and $|m_{ee}| > 2$ meV for the normal hierarchy (except for a possible cancellation region) if the neutrinos are Majorana particles [7]. Because the rate for neutrinoless double beta decay is proportional to the square of the effective electron neutrino Majorana mass, a Planck-suppressed Majorana mass term would yield a neutrinoless double beta-decay rate of order 10^6 times smaller than the ultimate reach of the 100-ton-scale experiments [7] proposed to probe the normal hierarchy in Majorana neutrino scenarios.

The mass-squared parameters m_{11}^2 and m_{22}^2 suffer from large radiative corrections with quadratic sensitivity to the high-scale cutoff of the theory, just as the SM Higgs mass-squared parameter does. The resulting hierarchy problem could be solved as in the SM by embedding our model into a supersymmetric or strong-dynamics theory at the TeV scale. Note however that because m_{12}^2 is the only source of breaking of the global U(1), its size is technically natural—radiative corrections to m_{12}^2 are proportional to m_{12}^2 itself and are only logarithmically sensitive to the cutoff. The smallness of m_{12}^2 required in our model could thus be explained by higher-scale physics, e.g., through spontaneous breaking of the U(1) in a hidden sector which is then communicated to the Higgs sector by heavy messenger particles or at high loop order.

We minimize the potential with the following considerations. Φ_1 must obtain the usual SM Higgs vev through spontaneous symmetry breaking with $m_{11}^2 < 0$; neglecting m_{12}^2 and v_2 we obtain $v_1^2 = -2m_{11}^2/\lambda_1$. To avoid a pseudo-Nambu-Goldstone boson with mass $\sim v_2$, the global U(1) is not also broken spontaneously; this is achieved for $m_{22}^2 + (\lambda_3 + \lambda_4)v_1^2/2 > 0$. Defining $\Phi_i = (\phi_i^+, (v_i + \phi_i^{0,r} + i\phi_i^{0,i})/\sqrt{2})^T$ and neglecting mixing terms suppressed by v_2/v_1 , the mass eigenstates are two neutral scalars $h^0 \simeq \phi_1^{0,r}$ (SM-like) and $H^0 \simeq -\phi_2^{0,r}$, a charged scalar $H^+ \simeq -\phi_2^+$, and a neutral pseudoscalar $A^0 \simeq -\phi_2^{0,i}$. Mixing between Φ_1 and Φ_2 can be ignored when $y_i^{\nu} \gg v_2/v_1$, as we assume here. The physical masses are

$$\begin{aligned} M_h^2 = & \lambda_1 v_1^2, & M_{H^+}^2 = & m_{22}^2 + \lambda_3 v_1^2/2, \\ M_{A,H}^2 = & M_{H^+}^2 + \lambda_4 v_1^2/2, \end{aligned} \quad (5)$$

where we again neglect terms suppressed by v_2/v_1 . In particular, H^0 and A^0 are degenerate and can be lighter or heavier than H^+ depending on the sign of λ_4 . Note that a mass splitting $(M_{H^+} - M_A)$ of either sign in this model yields a positive contribution to the ρ parameter [8], which serves to increase the best-fit value for the SM-like Higgs mass M_h in the electroweak fit. This eases the tension [9] between the standard electroweak fit, which prefers a light Higgs in the SM, and the lower bound on the SM Higgs mass from LEP. It also eases the “little hierarchy problem,” i.e., the fine-tuning of the Higgs mass-squared parameter against radiative corrections required for a cutoff above 1 TeV [3].

In the limit $v_2 \ll v_1$ the Yukawa couplings of the new scalars are given by

$$\begin{aligned} \mathcal{L}_{\text{Yuk}} = & (m_{\nu_i}/v_2) H^0 \bar{\nu}_i \nu_i - (im_{\nu_i}/v_2) A^0 \bar{\nu}_i \gamma_5 \nu_i \\ & - (\sqrt{2} m_{\nu_i}/v_2) [U_{\ell i}^* H^+ \bar{\nu}_i P_L e_{\ell} + \text{H.c.}], \end{aligned} \quad (6)$$

where $U_{\ell i}$ is the Pontecorvo-Maki-Nakagawa-Sakata matrix, for which we use the convention of Ref. [10]. In particular, the scalar Yukawa couplings are entirely fixed

by neutrino-sector parameters and the vev v_2 of the second doublet.

The scalars also couple among themselves; neglecting couplings suppressed by v_2 , the Feynman rules for all triple-scalar couplings in the model are $-i\lambda_3 v_1$ for the $h^0 H^+ H^-$ vertex, $-i(\lambda_3 + \lambda_4)v_1$ for the $h^0 H^0 H^0$ and $h^0 A^0 A^0$ vertices, and $-3i\lambda_1 v_1$ for the $h^0 h^0 h^0$ vertex.

III. BIG BANG NUCLEOSYNTHESIS

In our model the right-handed neutrino degrees of freedom will be populated in the early universe due, e.g., to $\ell^+ \ell^- \leftrightarrow \nu_R \bar{\nu}_R$ via t -channel H^+ exchange. The model is thus constrained by the limit on new relativistic degrees of freedom during BBN, $\delta N_{\nu, \max} = 1.44$ [11] at 95% confidence level (CL). To evade this bound, the right-handed neutrinos must be colder than the left-handed neutrinos, $T_{\nu_R}/T_{\nu_L} \leq (\delta N_{\nu, \max}/3)^{1/4}$ [12]. This can be achieved if the right-handed neutrinos drop out of thermal equilibrium early enough; in terms of the effective number of relativistic degrees of freedom g_* at the times of decoupling of the left- and right-handed neutrinos we have $T_{\nu_R}/T_{\nu_L} = (g_{*L}/g_{*R})^{1/3}$. Inserting $g_{*L} = N_B + 7N_F/8 = 43/4$ (where $N_B = 2$ photon and $N_F = 10$ e^\pm and ν_L spin degrees of freedom), we find $g_{*R} \geq 43/4 + 7.9$. We therefore require that the ν_R decouple from the thermal bath above the quark-hadron transition at 200–400 MeV, yielding 51 extra effective degrees of freedom from muons, u , d , and s quarks, and gluons.

This puts an upper bound on the H^+ -mediated $\ell^+ \ell^- \leftrightarrow \nu_R \bar{\nu}_R$ cross section via [12] $T_{d, \nu_R}/T_{d, \nu_L} \approx (\sigma_R/\sigma_L)^{-1/3} = [4v_2^4 M_{H^+}^4 / v_1^4 m_{\nu_i}^4 |U_{\ell i}|^4]^{1/3}$. Using $T_{d, \nu_L} \approx 3$ MeV [13] and imposing $T_{d, \nu_R} \gtrsim 300$ MeV we obtain an upper bound on the neutrino Yukawa couplings,

$$y_i^\nu \equiv \sqrt{2} m_{\nu_i} / v_2 \lesssim \frac{1}{30} \left[\frac{M_{H^+}}{100 \text{ GeV}} \right] \left[\frac{1/\sqrt{2}}{|U_{\ell i}|} \right], \quad (7)$$

which for $M_{H^+} \sim 100$ GeV is comparable to the SM bottom quark Yukawa coupling. If the lightest neutrino is nearly massless, we may take, for the normal neutrino mass hierarchy, $m_{\nu_3} \sim \sqrt{\Delta m_{32}^2} \sim 0.05$ eV and $U_{\mu 3} \approx 1/\sqrt{2}$, which yields $v_2 \gtrsim 2$ eV. In the inverted hierarchy the limit due to $y_{1,2}^\nu$ is comparable. The limit on v_2 scales linearly with the heaviest neutrino mass.

IV. PHENOMENOLOGY

We now consider the phenomenology of the scalar sector. The decays of the new scalars are controlled by the underlying U(1) symmetry—for $M_{A,H} > M_{H^+}$ ($M_{A,H} < M_{H^+}$), the decay modes are $H^+ \rightarrow \ell^+ \nu$ and $A^0, H^0 \rightarrow \nu \bar{\nu}$, $W^\pm H^\mp$ ($H^+ \rightarrow \ell^+ \nu$, $W^+ A^0$, $W^+ H^0$ and $A^0, H^0 \rightarrow \nu \bar{\nu}$). All other decays are suppressed by the tiny U(1) breaking v_2 . In particular, the tree-level couplings of H^0 ,

A^0 to W and Z bosons, quarks, and charged leptons are suppressed by v_2/v_1 . Decays to gluons (via a quark loop) are thus also suppressed. The decays $A^0, H^0 \rightarrow \gamma\gamma$ through an H^+ loop are suppressed by the tiny $A^0 H^+ H^-$, $H^0 H^+ H^-$ couplings $\sim \lambda_2 v_2$. The decay $H^0 \rightarrow h^0 h^0$ is also suppressed by v_2 . Loop-induced decays of $H^0 \rightarrow ZZ, WW$ or $H^+ \rightarrow W^+ Z, W^+ \gamma$ through a lepton triangle are suppressed by a neutrino mass insertion.

In what follows we assume $M_{A,H} \geq M_{H^+}$ (i.e., $\lambda_4 \geq 0$) and focus on the decays of the charged Higgs. First, note that H^- decays via the neutrino Yukawa couplings into a left-handed charged lepton, in contrast to the usual Type-I or II two Higgs doublet model [14] in which H^- decays into a right-handed charged lepton. In particular, $H^- \rightarrow \tau^- \bar{\nu}$ produces a left-handed τ^- , so that usual charged Higgs searches that take advantage of the τ helicity to suppress W backgrounds are not applicable in this model.

The charged Higgs decays into all nine combinations of $\ell_i \nu_j$. Summing over final-state neutrinos we obtain the decay width into a particular charged lepton species ℓ ,

$$\Gamma(H^+ \rightarrow \ell^+ \nu) = \frac{M_{H^+} \langle m_\nu^2 \rangle_\ell}{8\pi v_2^2}, \quad (8)$$

where we define the expectation value of the neutrino mass-squared in a flavor eigenstate, $\langle m_\nu^2 \rangle_\ell = \sum_i m_{\nu_i}^2 |U_{\ell i}|^2$ [15] (here $\langle m_\nu^2 \rangle_e$ is the same observable that is measured in tritium beta-decay endpoint experiments like KATRIN [16]). Imposing the BBN constraint $y_i^\nu \lesssim 1/30$ yields an upper bound on the H^+ total width, $\Gamma_{H^+}^{\text{tot}} \lesssim 1.3 \times 10^{-4} M_{H^+}$, or for $M_{H^+} = 100$ GeV, $\Gamma_{H^+}^{\text{tot}} \lesssim 13$ MeV. The charged Higgs is thus narrow but not long-lived. Similarly, the new neutral Higgs widths into neutrinos are given by $\Gamma(H^0 \rightarrow \nu_i \bar{\nu}_i) = \Gamma(A^0 \rightarrow \nu_i \bar{\nu}_i) = M_A m_{\nu_i}^2 / 8\pi v_2^2$, yielding an upper bound on the width to neutrinos of $\Gamma(H^0, A^0 \rightarrow \nu \bar{\nu}) \lesssim 6.6 \times 10^{-5} M_A$.

The charged Higgs branching ratios are given by $\text{BR}(H^+ \rightarrow \ell^+ \nu) = \langle m_\nu^2 \rangle_\ell / \sum_e \langle m_\nu^2 \rangle_e$ and are shown in Fig. 1 where we scan over the 2σ neutrino parameter ranges from Ref. [10]. These branching ratios are the

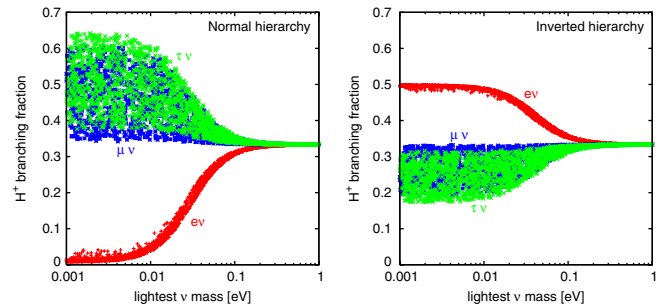


FIG. 1 (color online). Charged Higgs decay branching fractions to $e\nu$, $\mu\nu$, and $\tau\nu$ as a function of the lightest neutrino mass.

same as those of the singly charged triplet Higgs state Φ^+ in the type-2 seesaw model [17] as given in Ref. [18].

A. Constraints from LEP

Searches for leptons plus missing energy at LEP can be used to set limits on the charged Higgs in this model. Charged Higgs pair production is dominantly through s -channel Z and γ exchange; the t -channel neutrino exchange amplitude is suppressed by two powers of $y_i^{\nu} \lesssim 1/30$ and we thus neglect it. The usual LEP charged Higgs search [19] relies on H^+ decays to $\tau\nu$ or $q\bar{q}'$, as are expected in the Type-I or II two Higgs doublet model [14]. In our model, charged Higgs decays to quarks are absent and $\text{BR}(H^+ \rightarrow \tau^+\nu)$ reaches at most 0.65 for the normal hierarchy and 0.33 for the inverted hierarchy (Fig. 1); because of the sizable branching fractions of H^\pm into $e\nu$ or $\mu\nu$, we find that these channels will provide a stronger exclusion than the usual $\tau\nu$ channel.

LEP studied $e^+e^-p_T^{\text{miss}}$ and $\mu^+\mu^-p_T^{\text{miss}}$ in the context of searches for the supersymmetric scalar partners of e and μ (selectrons and smuons) with decays to the corresponding lepton plus a neutralino which escapes the detector [20,21]. For a massless neutralino the kinematics reproduce those of H^+H^- pair production. Using the smallest allowed branching fraction into the relevant decay mode from Fig. 1, we translate the LEP-combined 95% CL cross section limits [21] into a lower bound on M_{H^+} as shown in Fig. 2. We find $M_{H^+} \gtrsim 65\text{--}83$ GeV, depending on the hierarchy and the lightest neutrino mass. These bounds could be improved by a combined analysis of the $e^+e^-p_T^{\text{miss}}$, $\mu^+\mu^-p_T^{\text{miss}}$, and $\tau^+\tau^-p_T^{\text{miss}}$ channels and the inclusion of mixed-flavor channels.

For $M_A < M_{H^+}$, decays of $H^+ \rightarrow W^+A^0$, W^+H^0 could compete with the leptonic decay and potentially invalidate our LEP limits quoted above. This bosonic decay proceeds via the SU(2) gauge coupling; partial widths for off-shell

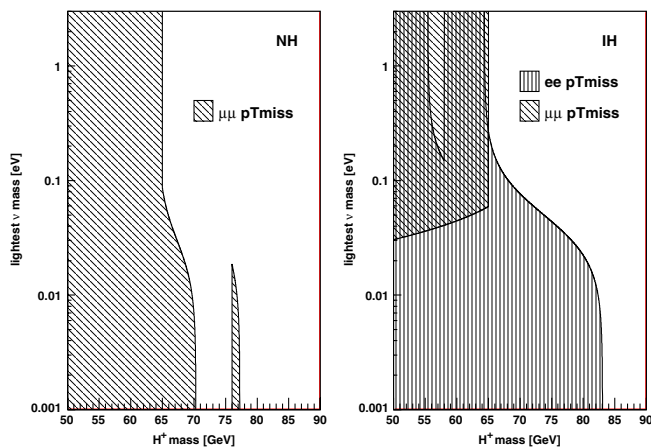


FIG. 2 (color online). Charged Higgs mass lower bound as a function of the lightest neutrino mass for the normal (NH) and inverted (IH) neutrino mass hierarchies, based on LEP selectron and smuon cross section limits from Ref. [21].

decays $H^+ \rightarrow W^+A^0/H^0$ were computed in Ref. [22] and are implemented in the public FORTRAN code HDECAY [23]. We note however that the Z boson invisible width [24] constrains $M_A \geq 43.7$ GeV at 95% CL. Taking $M_{H^+} = 83$ GeV, the largest possible partial width for off-shell $H^+ \rightarrow W^+A^0/H^0$ relevant to the LEP limits presented here is then 0.12 MeV. Taking y_i^{ν} at the BBN bound $\sim 1/30$, off-shell decays constitute a branching ratio of less than 1%, so that our LEP limits remain valid. The branching ratio for off-shell decays can reach 10% for $y_i^{\nu} \sim 1/100$.

B. $\mu \rightarrow e\gamma$ and $\tau \rightarrow \ell\gamma$

At one loop, H^+ mediates lepton flavor violating decays $L \rightarrow \ell\gamma$ with a branching fraction [15,25]

$$\text{BR}(L \rightarrow \ell\gamma) = \text{BR}(L \rightarrow e\nu\bar{\nu}) \frac{\alpha_{\text{em}}}{96\pi} \frac{|\sum_i m_{\nu_i}^2 U_{Li}^* U_{\ell i}|^2}{2G_F^2 M_{H^+}^4 v_2^4}, \quad (9)$$

where $(L, \ell) = (\mu, e)$ or $(\tau, e \text{ or } \mu)$.² Unitarity of the Pontecorvo-Maki-Nakagawa-Sakata matrix yields $\sum_i m_{\nu_i}^2 U_{Li}^* U_{\ell i} = -\Delta m_{21}^2 U_{L1}^* U_{\ell 1} + \Delta m_{32}^2 U_{L3}^* U_{\ell 3}$. Scanning over the allowed neutrino parameter ranges [10] with $M_{H^+} = 100$ GeV and $v_2 = 2$ eV yields $\text{BR}(\mu \rightarrow e\gamma) \leq 8 \times 10^{-12}$, with significant dependence on $\sin\theta_{13} \equiv |U_{e3}|$ as shown in Fig. 3.³ This branching fraction is below the current experimental 90% CL upper limit of $\text{BR}(\mu \rightarrow e\gamma) \leq 1.2 \times 10^{-11}$ [28], but for $\sin\theta_{13} \gtrsim 0.01$ is within reach of the MEG experiment which expects a sensitivity of about 10^{-13} after running to the end of 2011 [29]. In particular, in the MEG sensitivity range the $\Delta m_{32}^2 U_{\mu 3}^* U_{e3}$ term dominates, so that $\text{BR}(\mu \rightarrow e\gamma) \propto \sin^2\theta_{13}/v_2^4 M_{H^+}^4$. This provides a measurement of v_2 if $\sin\theta_{13}(M_{H^+})$ can be obtained from neutrino oscillation (collider) experiments.

Similarly, for $M_{H^+} = 100$ GeV and $v_2 = 2$ eV we find $\text{BR}(\tau \rightarrow \mu\gamma) = (0.7\text{--}2.1) \times 10^{-11}$ and $\text{BR}(\tau \rightarrow e\gamma) \leq 1.5 \times 10^{-12}$. These are well below the current 90% CL bounds of 6.8×10^{-8} [30] and 1.1×10^{-7} [31], respectively, as well as the expected reach at the SuperB next-generation flavor factory of 2×10^{-9} in either channel [32].

C. Muon anomalous magnetic moment

The charged Higgs also contributes to the muon anomalous magnetic moment $a_\mu \equiv (g - 2)/2$. Adapting the cal-

²The type-2 seesaw yields an analogous formula [26] with the same dependence on the neutrino parameters but dominated by loops involving the doubly charged triplet state Φ^{++} .

³Note that the rate for $\mu \rightarrow e\gamma$ ($\tau \rightarrow e\gamma$) vanishes when $U_{e3} = [\Delta m_{21}^2/\Delta m_{32}^2][U_{L1}^* U_{e1}/U_{L3}]$, with $L = \mu$ (τ) [27]. The rate for $\tau \rightarrow \mu\gamma$ cannot vanish given the known neutrino parameter values.

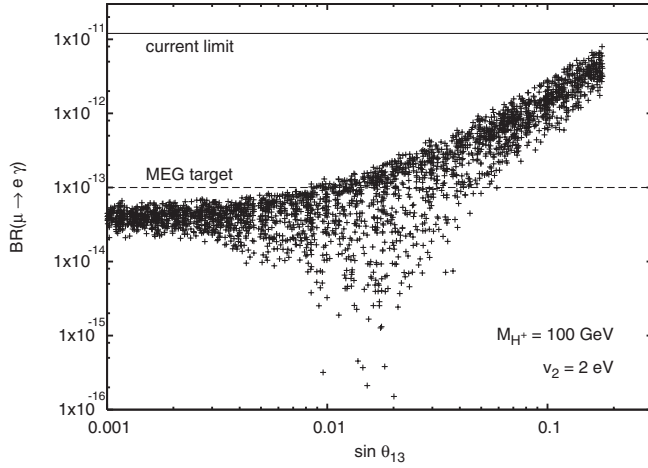


FIG. 3. Branching ratio of $\mu \rightarrow e\gamma$ versus $\sin\theta_{13}$. The normal and inverted hierarchies populate the same region.

culaton of Ref. [33] we find the one-loop H^+ contribution,

$$\delta a_\mu^{H^+} = -\frac{m_\mu^2 \langle m_\nu^2 \rangle_\mu}{48\pi^2 M_{H^+}^2 v_2^2}. \quad (10)$$

Taking $M_{H^+} \sim 100$ GeV and $y_i^\nu \lesssim 1/30$ results in $\delta a_\mu^{H^+}$ two to three orders of magnitude smaller than the current experimental uncertainty on a_μ . The two-loop H^+ contribution [34] is also well below the current sensitivity.

D. Tree-level muon and tau decay

In our model the decays $L \rightarrow \ell \nu \bar{\nu}$ receive contributions from tree-level charged Higgs exchange. Neglecting m_ν in the kinematics, the differential cross section for the H^+ -mediated process is identical to that for the W -mediated process but with the ν and $\bar{\nu}$ momenta interchanged, and the interference term is zero. The decay widths become $\Gamma(L \rightarrow \ell \nu \bar{\nu}) = \Gamma^{\text{SM}}(L \rightarrow \ell \nu \bar{\nu}) \times [1 + \langle m_\nu^2 \rangle_L \langle m_\nu^2 \rangle_\ell / 8G_F^2 M_{H^+}^4 v_2^4]$ [15], where the neutrino mass dependence in the H^+ contribution violates lepton flavor universality. In the standard parametrization (see, e.g., Ref. [35]) we obtain $g_\mu/g_e \simeq 1 + \langle m_\nu^2 \rangle_\tau \times (\langle m_\nu^2 \rangle_\mu - \langle m_\nu^2 \rangle_e) / 16G_F^2 M_{H^+}^4 v_2^4$ and $g_\mu/g_\tau \simeq 1 + \langle m_\nu^2 \rangle_e \times (\langle m_\nu^2 \rangle_\mu - \langle m_\nu^2 \rangle_\tau) / 16G_F^2 M_{H^+}^4 v_2^4$. Taking $M_{H^+} \simeq 100$ GeV and imposing the BBN constraint $y_i^\nu \lesssim 1/30$ results in deviations at the 10^{-6} level, well within the current experimental constraints $g_\mu/g_\tau = 0.9982 \pm 0.0021$ and $g_\mu/g_e = 0.9999 \pm 0.0020$ [35] as well as the expected SuperB reach of ± 0.0005 in either quantity [35].

E. H^+H^- production at LHC

We now consider charged Higgs search prospects at the LHC. The charged Higgs can be pair produced via $pp \rightarrow \gamma^*, Z^* \rightarrow H^+H^-$; we compute the cross section including next-to-leading-order QCD corrections using PROSPINO [36,37]. For $M_{H^+} = 100(500)$ GeV we find a cross section

of 300 (0.60) fb, with a theoretical uncertainty of $\sim 25\%$ [37]. This cross section provides direct access to the isospin of H^+ , allowing our doublet model (in which $T_{H^+}^3 = 1/2$) to be distinguished from the type-2 seesaw (in which the triplet state Φ^+ has $T_{\Phi^+}^3 = 0$). We find that because of this isospin difference, the cross section for H^+H^- is 2.7 (2.6) times larger than that for $\Phi^+\Phi^-$ for $M_{H^+}, \Phi^+ = 100(500)$ GeV.⁴

We note also that the branching fraction of H^+ to $\mu\nu$ or $e\nu$ is always at least 1/3 (Fig. 1), which provides distinctive search channels for H^+ at the LHC. Details will be given in a forthcoming paper [39]. Measurement of the characteristic pattern of H^+ branching fractions would provide strong evidence for the connection of H^+ to the neutrino sector, as well as allowing a determination of the neutrino mass hierarchy and providing some sensitivity to the lightest neutrino mass.

F. Effects on the SM-like Higgs

Finally we comment on the phenomenology of the SM-like Higgs h^0 . The charged Higgs will contribute to the one-loop amplitude for $h^0 \rightarrow \gamma\gamma$, yielding

$$\Gamma(h^0 \rightarrow \gamma\gamma) = \Gamma^{\text{SM}}(h^0 \rightarrow \gamma\gamma) \left[1 - \lambda_3 \delta \left[\frac{100 \text{ GeV}}{M_{H^+}} \right]^2 \right]^2, \quad (11)$$

where for $M_h = 120$ GeV and $M_{H^+} = 100(200, 1000)$ GeV, $\delta = 0.20(0.17, 0.16)$. This provides experimental access to the coupling λ_3 , which appears in the $h^0 H^+ H^-$, $h^0 A^0 A^0$, and $h^0 H^0 H^0$ vertices. Again for $M_h = 120$ GeV, the experimental precision with which $\Gamma(h^0 \rightarrow \gamma\gamma)$ can be extracted at the LHC has been estimated at $\sim 15\% - 30\%$ [40]; at the International Linear e^+e^- Collider with 500 GeV center-of-mass energy the precision remains at the $\sim 25\%$ level due to limited statistics [41]. This would provide sensitivity to $|\lambda_3| \sim 1$ only at the 1σ level. This precision could be improved to $\sim 5\%$ at a linear collider with 1 TeV center-of-mass energy [42] or $\sim 2\%$ at a photon collider running on the Higgs resonance [43], providing 4–10 σ sensitivity to $|\lambda_3| \sim 1$, respectively.

More importantly, our model impacts the range of SM-like Higgs masses allowed by the standard electroweak fit. In particular, the new scalars can increase the allowed range for M_h by giving a positive contribution to the ρ parameter. Using the result for a generic two Higgs doublet model from Ref. [8], the new scalars yield

$$\Delta\rho = \alpha_{\text{em}} \Delta T = \frac{\alpha_{\text{em}}}{8\pi M_W^2 s_W^2} F(M_{H^+}^2, M_A^2), \quad (12)$$

where M_W is the W boson mass, s_W denotes the sine of the

⁴For recent LHC phenomenology studies of the triplet states in the type-2 seesaw, see Refs. [18,38]. The signatures are dominated by $\Phi^{++}\Phi^{--}$ and $\Phi^{\pm\pm}\Phi^\mp$ production.

weak mixing angle, and $F(m_1^2, m_2^2) \equiv (m_1^2 + m_2^2)/2 - [m_1^2 m_2^2 / (m_1^2 - m_2^2)] \ln(m_1^2 / m_2^2)$. Here the shift in the oblique parameter ΔT is defined relative to a SM reference point with SM Higgs mass set equal to M_h . It was shown in Ref. [3] that a SM-like Higgs mass in the range of 400–600 GeV can be made consistent with the electroweak fit if $\Delta T \approx 0.25 \pm 0.1$. This can be achieved in our model for $M_{H^+} = 100$ GeV with $M_A \approx 200$ –250 GeV, corresponding to $\lambda_4 = 2\sqrt{2}G_F(M_A^2 - M_{H^+}^2) \approx 1.0$ –1.7.

For $M_h > 2M_{H^+}$ or $2M_A$, the additional decays $h^0 \rightarrow H^+ H^- \rightarrow \ell^+ \ell'^- \nu \bar{\nu}$ or $h^0 \rightarrow H^0 H^0$, $A^0 A^0 \rightarrow 4\nu$ appear. The partial widths for these decays are given above threshold by

$$\begin{aligned} \Gamma(h^0 \rightarrow H^+ H^-) &= \frac{\lambda_3^2}{16\pi\sqrt{2}G_F M_h} \sqrt{1 - \frac{4M_{H^+}^2}{M_h^2}}, \\ \Gamma(h^0 \rightarrow A^0 A^0) &= \Gamma(h^0 \rightarrow H^0 H^0) \\ &= \frac{(\lambda_3 + \lambda_4)^2}{32\pi\sqrt{2}G_F M_h} \sqrt{1 - \frac{4M_A^2}{M_h^2}}. \end{aligned} \quad (13)$$

For example, for $M_{H^+} = 100$ GeV and $M_h = 300$ GeV we obtain $\Gamma(h^0 \rightarrow H^+ H^-) = 3.0\lambda_3^2$ GeV; for comparison the total width of a 300 GeV SM Higgs boson is 8.5 GeV [23].

Below threshold, these Higgs-to-Higgs decays are suppressed by the small y_i' unless $A^0, H^0 \rightarrow H^\pm W^\mp$ or $H^+ \rightarrow W^+ A^0 / H^0$ is open.

V. CONCLUSIONS

We introduced a simple new TeV-scale model for Dirac neutrinos which explains the smallness of neutrino masses by sourcing them from a second Higgs doublet with tiny vev \sim eV. The model predicts distinctive decay patterns of H^+ controlled by the neutrino mass spectrum and mixing matrix, which can be tested at the LHC. The isospin of H^+ can be measured at the LHC, allowing the model to be distinguished from the type-2 seesaw for Majorana neutrino mass generation. The model also predicts a signal in $\mu \rightarrow e\gamma$ at the currently running MEG experiment if $\sin\theta_{13} \gtrsim 0.01$ and $\nu_2 \lesssim 6$ eV. Because the model conserves lepton number, neutrinoless double beta decay is absent.

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