W boson mass shift, dark matter, and $(g-2)_\ell$ in a scotogenic-Zee model

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We present a singly charged scalar extension of the scotogenic model, the scotogenic-Zee model, which resolves the recently reported deviations in the W boson mass as well as lepton $g-2$. The model admits a scalar or a fermionic dark matter while realizing naturally small radiative neutrino masses. The mass splitting of ∼100 GeV, which is required by the shift in W boson mass, among the inert doublet fields can be evaded by its mixing with the singlet scalar, which is also key to resolving the $(g-2)_\ell$ anomaly within 1σ. We establish the consistency of this framework with dark matter relic abundance while satisfying constraints from charged lepton flavor violation, direct detection, and collider bounds. The model gives predictions for the lepton flavor violating $\tau \to \ell\gamma$ processes that will be testable in upcoming experiments.

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I. INTRODUCTION

The CDF Collaboration at Fermilab [1] reported a precision measurement of W boson mass, $M_W^{\mathrm{CDF}} = (80.4335 \pm 0.0094)$ GeV, which is in tension with the Standard Model (SM) prediction, $M_W^{\mathrm{SM}} = (80.357 \pm 0.004)$ GeV [2], with an excess at the 7σ level that may be an indication of new physics beyond the SM. The new result from the CDF Collaboration, with a much reduced uncertainty, has a higher precision than the Particle Data Group (PDG) world average of $M_W^{\mathrm{PDG}} = (80.377 \pm 0.012)$ GeV [3], which takes into account the W mass measurements from LEP [4], Tevatron [5] (CDF [6] and D0 [7]), and the LHCb Collaboration [8]. The PDG average, which is in agreement with the SM prediction, disagrees with the CDF Run-II result. It has been shown that the improvement in parton density functions [9] and perturbative matrix elements [10,11] cannot account for this discrepancy. However, the discrepancy may be due to high-twist power corrections within the SM that are not normally considered in perturbative calculations [12,13]. This work proceeds under the assumption that the new CDF measurement will be validated as the correct result for the W boson mass.

Some possible explanations to the W boson mass shift can arise at tree level [14–29] or at loop level [30–46], along with the prospect of reconciling one or more discrepancies [47–69], such as flavor anomalies and dark matter. Various other papers [11,70–99] also examined the consequences of the CDF M_W anomaly on new physics scenarios.

Independently, the Muon $(g-2)$ Collaboration at Fermilab [100] confirmed the long-standing discrepancy in the anomalous magnetic moment (AMM) of muon measurement at BNL in 2006 [101] at a combined 4.2σ deviation,[1] $\Delta a_\mu^{\mathrm{exp}} = (2.51 \pm 0.59) \times 10^{-9}$, from the SM prediction (see Ref. [103] and the references therein). In addition to these recent anomalies, astrophysical and cosmological observations [104–106] present compelling evidence for the existence of dark matter (DM), for which the SM fails to provide an explanation. Moreover, one of the major shortcomings of the SM is its inability to explain the origin of nonzero neutrino mass substantiated by several experiments [107].

In this work we show that a simple extension of the scotogenic model [108] with a charged singlet (the scotogenic-Zee model) can simultaneously address all the previously mentioned puzzles. Our novel scotogenic-Zee model[2] is the simplest model that furnishes a direct link

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[1]Recent results from the BMW Collaboration [102] agree with the experimental measurement within 1.6σ.

[2]The scalar content is the same as the inert Zee model [109,110], with only right-handed neutrinos, in contrast to vectorlike singlets and doublets. The scotogenic-singlet model [111] is a neutral scalar extension of the scotogenic model. Neither model can resolve the discrepancy in $(g-2)_\mu$ [112].

between neutrino mass generation, dark matter, and the AMM of the muon and also provides an upward mass shift in the W boson that is in agreement with the CDF measurement. Additionally, the presumed anomaly in the AMM of the electron [113–115] can also be addressed within the same framework. We explore the parameter space of the scotogenic-Zee model spanned by both the bosonic and fermionic DM candidates while being consistent with the current experimental constraints.

The rest of the paper is organized as follows. In Sec. II we give a brief description of the scotogenic-Zee model while discussing the neutrino mass generation and the scalar sector. In Secs. III and IV we respectively introduce and examine a resolution to the W mass shift and AMM phenomenology in the model. Section V discusses DM phenomenology for both the scalar and fermionic DM candidates. Lastly, we integrate the three puzzles (W-mass shift, lepton $g-2$, and DM) with neutrino mass generation and lepton flavor violating (LFV) constraints and invoke a highly predictive flavor texture in Sec. VI before concluding in Sec. VII.

II. MODEL

The proposed scotogenic-Zee model is a simple charged singlet $S^+(1,1;-)$ extension of the scotogenic [108] model that contains Majorana singlet fermions $N_{R_i}(1,0;-)$ and the scalar doublet $(\eta^+,\eta^0)\equiv\eta(2,1/2;-)$ under the gauge group $SU(2)_L\times U(1)_Y\times\mathbb{Z}_2$. All the new particles are odd under $\mathbb{Z}_2$, while the SM particles are even, guaranteeing the stability of the DM candidate; the lightest among the new neutral $\mathbb{Z}_2$-odd particles. The charged scalar singlet S^+ not only gives corrections to the anomalous magnetic moment of the muon and electron through mixing with the charged doublet but also serves as a portal to generate the correct relic abundance for fermionic DM.

The effective Yukawa Lagrangian in the extended model can be written as

$$-\mathcal{L}_Y\supset Y_{ij}\bar{L}_{L_i}\tilde{\eta}N_{R_j}+f_{ij}\bar{\ell}_{R_i}S^-\bar{N}_{R_j}+\text{H.c.} \tag{1}$$

The $\mathbb{Z}_2$ symmetry, being exact, prevents η^0 from obtaining a nonzero vacuum expectation value (VEV), and neutrinos remain massless at tree level. Moreover, the SM Higgs h is decoupled from the new CP-even $[\mathrm{Re}(\eta^0)\approx H]$ and -odd $[\mathrm{Im}(\eta^0)\approx A]$ scalars. The tree-level scalar potential of the model is given by

$$\begin{aligned}V&=\mu_h^2\phi^\dagger\phi+\mu_S^2S^-S^++\mu_\eta^2\eta^\dagger\eta+\frac{\lambda_1}{2}(\phi^\dagger\phi)^2+\frac{\lambda_2}{2}(\eta^\dagger\eta)^2\\&+\lambda_3(\phi^\dagger\phi)(\eta^\dagger\eta)+\lambda_4(\phi^\dagger\eta)(\eta^\dagger\phi)+\frac{\lambda_5}{2}\{(\phi^\dagger\eta)^2+\text{H.c.}\}\\&+\frac{\lambda_6}{2}(S^-S^+)^2+\lambda_7(\phi^\dagger\phi)(S^-S^+)+\lambda_8(\eta^\dagger\eta)(S^-S^+)\\&+\frac{\mu}{2}\{\varepsilon_{\alpha\beta}\phi^\alpha\eta^\beta S^-+\text{H.c.}\}.\end{aligned} \tag{2}$$

The charged scalars $\{\eta^+,S^+\}$ mix, giving rise to the mass eigenstates $\{H_1^+,H_2^+\}$. The masses of the scalar fields in the physical basis are given by

$$\begin{aligned}m_h^2&=\lambda_1v^2,\qquad m_{H(A)}^2=\mu_\eta^2+\frac{v^2}{2}(\lambda_3+\lambda_4\pm\lambda_5),\\m_{H_i^+}^2&=\frac{1}{2}\left(\mu_2+\mu_3\pm\sqrt{(\mu_2-\mu_3)^2+2\mu^2v^2}\right),\end{aligned} \tag{3}$$

where $\mu_2=\mu_\eta^2+\frac{\lambda_3}{2}v^2$ and $\mu_3=\mu_S^2+\frac{\lambda_7}{2}v^2$. Here $\mu_{\eta,S}$, λ_i, and μ are the bare-mass terms, quartic couplings, and cubic coupling, respectively. The mixing angle between the charged scalar fields is defined as

$$\sin2\theta=\frac{-\sqrt{2}\mu v}{m_{H_1^+}^2-m_{H_2^+}^2}, \tag{4}$$

with the VEV $v\simeq246$ GeV. In this work, we comply with the perturbative and vacuum stability conditions [116,117] constraining the scalar couplings. We also have ensured that our parameter space does not drive the mass parameters μ_η^2 and μ_S^2 to negative values; negative bare-mass terms can lead the inert doublet or the charged singlet to attain a nonzero VEV, thereby breaking the $\mathbb{Z}_2$ symmetry [118]. Such $\mathbb{Z}_2$ parity violation not only destabilizes the DM candidates but also allows unacceptably large tree-level neutrino masses. These arise from the renormalization group (RG) evolution of the bare-mass terms and might pose a serious problem when the coupling of the inert scalars to heavy neutrinos are of $\mathcal{O}(1)$. As will be discussed in the following sections, the coupling of η to N_{R_k} is small enough to avoid breaking the $\mathbb{Z}_2$ symmetry. On the other hand, the coupling between S^+ and N_{R_k} is taken to be $\mathcal{O}(1)$. Thus, in order for the model to be valid above the electroweak (EW) scale (such that the $\mathbb{Z}_2$ symmetry is preserved), it is important to check that the RG evolution preserves not only $\mu_\eta^2>0$ but also $\mu_S^2>0$. The trilinear coupling μ gives positive contributions to the RG evolution of the new scalar mass-squared parameters, thereby helping prevent the breaking of the model symmetries. Note that it is also essential to ensure that μ is not much larger than the scalar masses, as it can result in a deeper minimum than the SM one [119,120]. The Majorana mass term $\frac{1}{2}M_{N_i}N_iN_i$ along with the scalar quartic term $\frac{\lambda_5}{2}\{(\phi^\dagger\eta)^2+\text{H.c.}\}$ breaks the lepton number by two units while generating the one-loop neutrino mass [see Fig. 1 (top sketch)] $\mathcal{M}_\nu$, which is expressed as

$$\begin{aligned}(\mathcal{M}_\nu)_{ij}&=\sum_kY_{ik}\Lambda_kY_{kj}^*,\\\Lambda_k&=\frac{M_{N_k}}{16\pi^2}\left[\frac{m_H^2}{m_H^2-M_{N_k}^2}\log\frac{m_H^2}{M_{N_k}^2}-(m_H\leftrightarrow m_A)\right].\end{aligned} \tag{5}$$

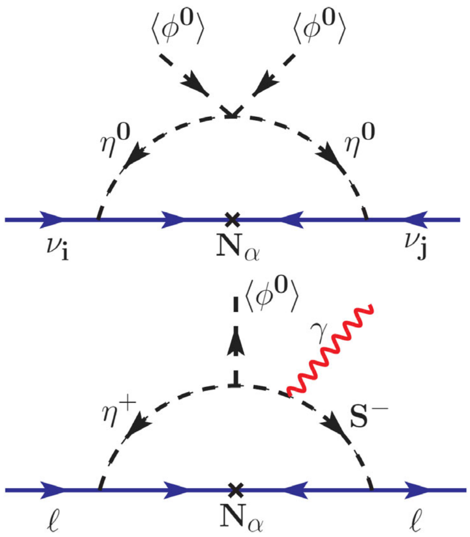

FIG. 1. Top sketch: radiative neutrino mass generation at one loop. Bottom sketch: dominant correction to AMMs arising through the chiral enhancement. The cross represents the mass insertion, whereas $\ell = e(\mu)$ for the electron (muon) AMM.

Here the lightest mass eigenstates $\{H, A\}$ and N_i can serve as viable bosonic and fermionic DM candidates, respectively. It is important to point out that unlike the scotogenic model, where M_N can be arbitrarily small at the canonical seesaw scale of 10^9 GeV or the Yukawa coupling Y, the $(g-2)_\mu$ in the model requires the scale to be in the (sub-)TeV range along with $\mathcal{O}(0.1\text{–}1.0)$ Yukawa coupling. Thus, a successful explanation of $m_\nu \sim 0.1$ eV would naturally require m_H to be nearly degenerate with m_A.

III. CORRECTION TO THE *W* BOSON MASS

The shift in W boson mass [121] can be evaluated as a function of the oblique parameters S, T, and U [122,123] that quantify the deviation of a new physics model from the SM through radiative corrections arising from shifts in gauge boson self-energies. The oblique parameters in our model get corrections from the extended Higgs sector which is same as in the Zee model [124] except for the $\mathbb{Z}_2$ charge preventing the mixing with the SM Higgs doublet. Therefore, we use the expressions for S, T and U given in Ref. [125] under the alignment limit [126]. Note that the corrections to U at one-loop level is suppressed compared to S and T.

With the new precision measurement of M_W by CDF, some electroweak (EW) observables are expected to suffer from this deviation. We incorporate the global EW fit [15] with the new CDF data to quote the 2σ allowed ranges of oblique parameters. We confirm the necessity of mass splitting in 2HDM [35,36] to accommodate the recent CDF results and show that the introduction of the charged singlet scalar allows the components of the doublet field to be degenerate in certain regions for a specific choice of $\sin\theta$

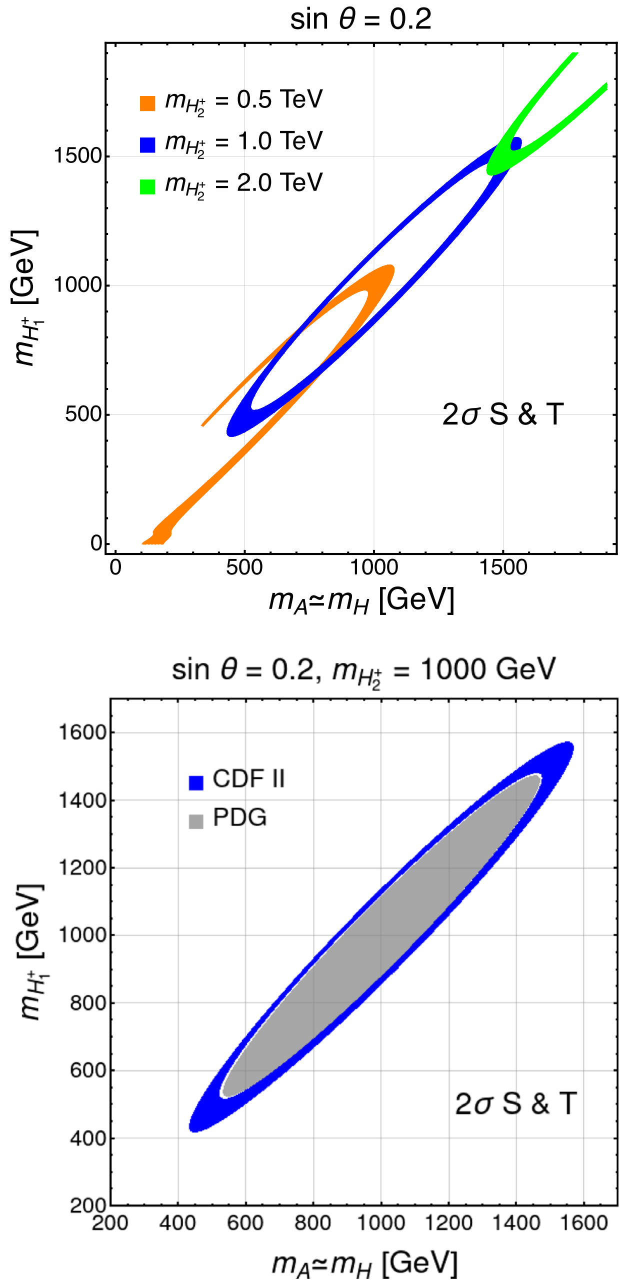

FIG. 2. Top panel: the mass splitting between the components of the doublet scalar required by the new CDF measurement of W boson mass using 2σ ranges for S and T from Ref. [15]. Bottom panel: the 2σ band allowed by the CDF measurement, in contrast to the PDG world average [3,40].

and $m_{H_2^+}$, as can be seen from Fig. 2 (top) (for more detail see Fig. 5). The splitting $\delta_{H+} = m_{H_1^+} - m_H$ depends on the mixing angle, for instance, it can be at most ~140 GeV for $\sin\theta = 0.2$. In spite of explaining the CDF W mass shift, the proposed model remains consistent with the previous experiments resulting in the PDG world average of the W boson mass. It can be seen in Fig. 2 (bottom panel) that a region consistent with the older experiments opens up more parameter space in the model, even allowing for degeneracy between the doublet fields in the entire range of allowed parameter space for particular values of $\sin\theta$ and $m_{H_2^+}$.

IV. ANOMALOUS MAGNETIC MOMENT

The charged scalar contributions to the anomalous magnetic moment at one loop [127] shown in Fig. 1 (bottom sketch) are

$$\Delta a_\ell^{H_1^+} = \frac{m_\ell^2}{16\pi^2}\Bigg((|Y_{\ell i}|^2 \sin^2\theta + |f_{\ell i}|^2 \cos^2\theta) G[m_{H_1^+}, 2] + \frac{M_{N_i}}{m_\ell}\mathrm{Re}(Y_{\ell i} f^*_{\ell i}) \sin 2\theta G[m_{H_1^+}, 1]\Bigg), \tag{6}$$

where

$$G[M, \epsilon] = \int_0^1 \frac{x^\epsilon (x-1)dx}{m_\ell^2 x^2 + (M^2 - m_\ell^2)x + M_{N_i}^2(1-x)} \tag{7}$$

and $\Delta a_\ell^{H_2^+} = \Delta a_\ell^{H_1^+}(\theta \to \frac{\pi}{2} + \theta)$. The dominant contribution to Δa_μ comes from the Majorana neutrino mass enhancement, aided by the mixing of the charged scalar mediators as shown in Fig. 1 (bottom sketch). The sign of the product of the Yukawa couplings and the mixing angle can be chosen independently. This in turn allows for the simultaneous explanations of Δa_ℓ ($\ell = e, \mu$). Moreover, Δa_μ provides an upper limit on the mass of the Majorana neutrino (charged scalar) on the order of 15 (6.5) TeV with $f, Y \sim \mathcal{O}(1)$. The mass limit is relaxed in the case of Δa_e.

Note that the Yukawa couplings and the masses of charged scalars are severely restricted by the charged lepton flavor violating processes such as radiative decay $\ell_i \to \ell_j \gamma$ [128]; such processes are enhanced in our model by the mass insertion of Majorana neutrinos. Moreover, although trilepton decays such as $\mu \to 3e$ do not occur at tree level, they arise at one loop with large branching ratios [129]. The same is also true for $\mu - e$ conversion in the nuclei. We impose these constraints in our parameter scan.

V. DM PHENOMENOLOGY

In addition to explaining the W boson mass shift and Δa_ℓ, the proposed model can easily accommodate both the scalar (lightest of H and A) and fermionic (lightest among N_i) dark matter candidates (χ). We consider both scenarios and analyze the parameter space by implementing the model in SARAH [130] and numerically evaluating the relic abundance using the software `micrOMEGAs` [131]. The relic density of DM is achieved through the standard thermal freeze-out mechanism.

For the case of the Majorana fermion as a DM ($\chi \equiv N$) candidate, the annihilation channel which determines the observed relic abundance is DM self-(co)annihilation into charged leptons $\ell_\alpha^- \ell_\beta^+$ (light neutrinos $\nu_\alpha \bar{\nu}_\beta$) through t-channel processes mediated by the $\mathbb{Z}_2$-odd scalars H_i^+ (H, A) via the Yukawa coupling Y and/or f. The neutrino oscillation data determine the flavor structure of Y, making it natural to select a relatively small Y and a heavy doublet scalar $\eta \sim \mathcal{O}(\mathrm{TeV})$ such that the LFV constraints are relaxed. Thus, we choose $f_{ii} \sim \mathcal{O}(1)$ $(i = 1, 2)$ and degenerate N_i to maximize the contribution to annihilation mode $\chi\chi \to \ell\ell$ via S^+; the allowed parameter space in the mass plane can be seen in Fig. 3 (top panel) along with the region resolving muon AMM for a specific choice of $\kappa = Y^* f \sin\theta = 0.015$.

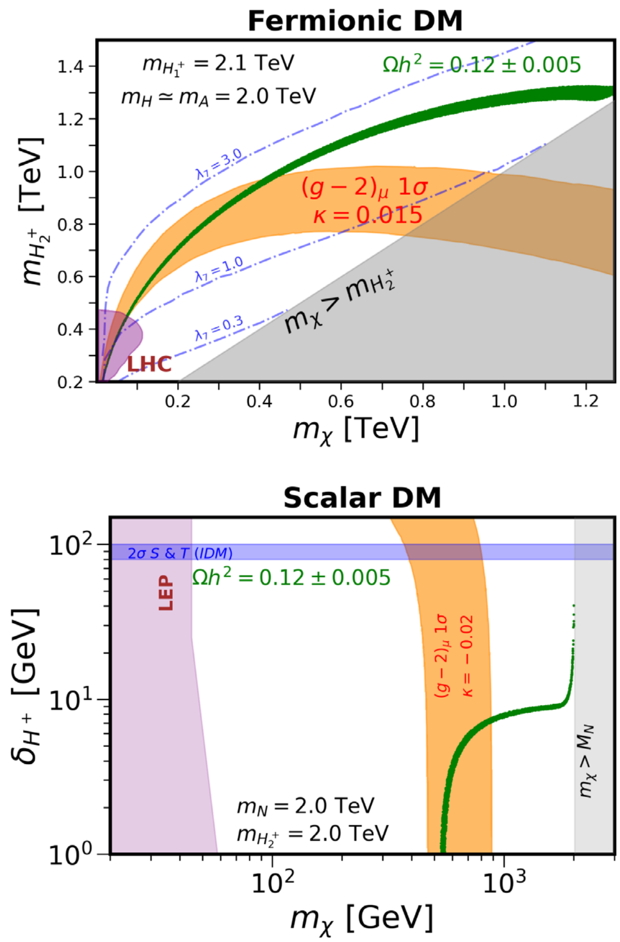

FIG. 3. The allowed parameter space for fermion (top panel) and neutral scalar (bottom panel) dark matter candidate to simultaneously explain the AMM of the muon and the W boson mass shift. The green (orange) band satisfies the correct relic abundance $[(g-2)_\mu]$ for a fixed $\kappa = Y^* f \sin\theta$. The choice of scalar mass provides an upward W boson mass shift, as expected from the CDF measurements. The purple band (top panel) is the exclusion region from the $pp \to \ell^+\ell^- + \not{E}_T$ signature obtained by recasting supersymmetric slepton searches at the LHC [132] and from chargino production searches at LEP [133,134] (bottom panel). The blue dash-dotted lines (top panel) show the direct detection cross section lower bounds [135,136] from the LZ Collaboration experiment [137] for different quartic couplings λ_7 and our choice of Yukawa coupling $f = 1.5$.

Note that in Fig. 3 the effects of coannihilation are not taken into account for fermionic DM. The coannihilation between fermionic DM and the charged scalar $H_2^{\pm}$ becomes important when the mass difference is small: $(m_{H_2^{\pm}} - m_N)/m_N < 0.1$ [138]. This dominates over the annihilation processes mainly when $f < 1$, a parameter space not favorable for our work to incorporate $(g-2)_\ell$ and satisfy LFV. Thus, we simply avoid coannihilation by taking a larger mass splitting.

Although fermionic DM does not contribute to tree-level direct detection cross sections, it can arise at one-loop order via photon Z and Higgs penguin diagrams involving charged scalars and charged leptons in the loop [136]. The relevant dimension-6 operators for Majorana DM direct detection are the ones with the bilinears $\bar{\chi}\chi$, $\bar{\chi}\gamma_5\chi$, and $\bar{\chi}\gamma^\mu\gamma_5\chi$. In the model presented here, both Yukawa couplings f and Y give rise to such diagrams. As previously mentioned, the Yukawa coupling Y needs to be small. Thus, the dominant contribution comes from the coupling $f \sim \mathcal{O}(1)$ of Eq. (1) with the right-handed charged lepton ℓ_R and the charged scalar S^+ in the loop. In this case, the Z penguin contribution is suppressed due to the absence of axial-vector coupling to ℓ_R. Photon mediated processes can lead to an anapole operator [139] of the form $\mathcal{O}^q_{AV} = (\bar{\chi}\gamma^\mu\gamma_5\chi)(\bar{q}\gamma_\mu q)$, where q represents the quark flavor eigenstates. Anapole contributions, however, are momentum and velocity suppressed in the nonrelativistic limit [140]. We find that the cross section from the anapole contribution for the DM masses in our framework is at most $\mathcal{O}(10^{-55})$ cm^2 [136,141], which is well beyond the current sensitivity. The dominant contribution, therefore, appears from the Higgs penguin diagram leading to spin-independent direct detection from the operator $\mathcal{O}^q_{SS} = m_q(\bar{\chi}\chi)(\bar{q}q)$. The corresponding Wilson coefficient is dependent not only on the Yukawa coupling f but also on the quartic coupling λ_7. This contribution can therefore be suppressed by choosing a small λ_7, allowing for f coupling to be $\gg$1. The allowed parameter space with the current bounds [137] for direct detection cross sections [135,136] is shown in Fig. 3 (top panel) for different choices of λ_7 and a fixed Yukawa coupling $f = 1.5$.

In the case of scalar dark matter, which we choose to be the CP-even $H \equiv \chi$ (nearly degenerate[3] with A and $\lambda_5 < 0$), DM can annihilate to W^+W^-, ZZ, $\nu_\alpha\nu_\beta$, hh, $\bar{\ell}\ell$, and $\bar{q}q$. The low mass regime suffers a strong constraint from LEP [133,134,143] which can be satisfied if one assumes that $m_\chi > M_Z/2$, $m_{H_1^+} > M_W/2$, and $m_{H_1^+} + m_\chi > M_W$. For larger DM mass, it predominantly annihilates to a pair of W^+W^- and ZZ, for which the allowed region is $m_\chi \gtrsim 500$ GeV and the mass splitting $\delta_{H^+} = m_{H_1^+} - m_\chi \lesssim 30$ GeV, as shown in Fig. 3 (bottom panel). This can be relaxed by making the Higgs quartic coupling $\gtrsim$1, a choice strongly constrained by the direct detection bound [137,144–147].

In this work we take the quartic couplings $\lambda_3 + \lambda_4 + \lambda_5 \ll 1$ to automatically satisfy the direct detection bound obtained from the scalar DM interacting with the nucleus at tree level through the SM Higgs boson. Moreover, it is favored to take the couplings Y_{ij} small and $M_N \sim \mathcal{O}(\mathrm{TeV})$ to be consistent with the neutrino fit, which implies that the DM analysis is indistinguishable from the known inert doublet model (IDM). It turns out that the CDF measurement requires mass splitting among the inert doublet fields of $\mathcal{O}(100)$ GeV, thus disfavoring the scalar DM candidates in the scotogenic IDM. However, the mixing between the charged scalars in this model allows the components of the doublet field to be degenerate (cf. Fig. 2), thereby admitting the CP-even H to be a viable DM candidate, as shown in Fig. 3 (bottom panel).

The direct detection cross section can arise at one loop from vertices like $HHZZ$, $HHWW$, HZA, and $HWH_1^{\pm}$ driven by the gauge coupling. These contributions can be as important as the tree-level processes. In the model presented here, the quartic couplings that give rise to tree-level direct detection are chosen to be small; thus, its contribution in comparison to the loop effect is negligible. For the masses $m_\chi \geq 500$ GeV required to explain the total observed relic density within the model [see Fig. 3 (bottom panel)], we find that the one-loop cross section is $\sim 1.1 \times 10^{-46}$ cm^2 (see Refs. [148–150] for details). This evades the current experimental bound [137,144–147] but will be sensitive to upcoming experiments [151,152].

VI. NEUTRINO FIT/ LEPTON FLAVOR VIOLATION

The neutrino mass formula of Eq. (5), lepton $g-2$, and the dark matter analysis have close-knit correlation through Yukawa couplings, Majorana fermions, and new scalars. As previously stated, $(g-2)_\ell$ sets an upper bound on the masses of the Majorana fermions and the charged scalars with $f, Y \sim \mathcal{O}(1)$. Moreover, the maximum splitting among the doublet fields is restricted by the shift in W boson mass, thereby forcing the parameter space to the region $m_A \simeq m_H$, which is crucial in explaining the observed neutrino oscillation data.

In order to check the consistency with the neutrino oscillation data and efficiently probe the model with LFV observables, we adopt Casas-Ibarra parametrization [153] to rewrite the Yukawa matrix Y of Eq. (5) in terms of neutrino mass parameters,

$$Y = U\sqrt{\mathcal{M}_\nu^{\mathrm{diag}}}R^\dagger\sqrt{\Lambda}^{-1}, \tag{8}$$

where R is an arbitrary complex orthogonal matrix. The neutrino oscillation parameters are scanned within the 2σ allowed ranges [154] to obtain the Yukawa matrix.

[3]The mass splitting is on the order of $\mathcal{O}(100)$ keV [142] to evade direct detection.

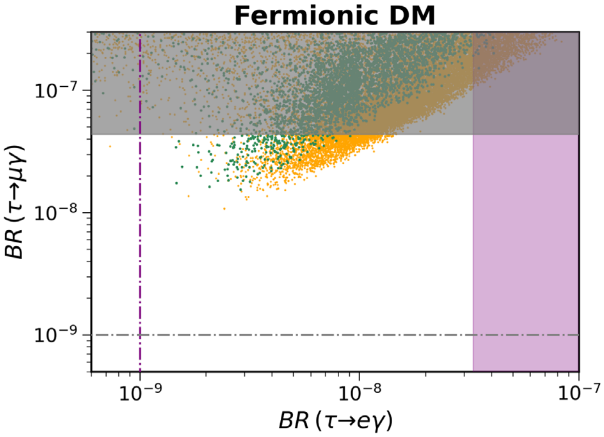

FIG. 4. Scattered plot assuming the fermionic DM with the same parameter space as given in Fig. 3 (top panel). Colored shaded regions are the current exclusion limits [155], whereas dash-dotted lines represent the future projected sensitivities [156]. The orange (green) dots correspond to solutions that resolve Δa_μ ($\Delta a_{e,\mu}$) and satisfy the observed relic density as well as neutrino oscillation observables within their 2σ measured values [154].

As mentioned earlier, the product of Yukawa couplings $Y_{\ell i} f^*_{\ell i}$ can explain $(g-2)_\ell$; however, these couplings are constrained by LFV processes. To facilitate a direct correlation between the observed relic density and both the muon and the electron $(g-2)$, we consider two (degenerate) stable DM candidates with their respective Yukawa couplings of $\mathcal{O}(1)$. This also allows more parameter space for the DM mass m_χ and mediator mass $m_{H_2^+}$, as shown in Fig. 3 (top panel). With such a choice of a large f_{ii} ($i = 1, 2$), the mass enhancement to $\ell_i \to \ell_j \gamma$ severely restricts the parameter space. Such a chirally enhanced contribution can be suppressed with a suitable choice of Yukawa couplings and masses of the Majorana fermions. For instance, chirally enhanced $\mu \to e\gamma$ can be evaded with the choice of $Y_{12} = Y_{21} \sim 0$ or $Y_{21} f_{11} \simeq -Y_{12} f_{22}$ for $M_{N_1} = M_{N_2} (= m_\chi)$. We then check the consistency of our fit by computing the branching fractions of $\ell_i \to \ell_j \gamma$ and $\ell_i \to 3\ell_j$ process at one-loop level and make testable predictions for fermionic DM (see Fig. 4). In the case of scalar DM, since the Yukawa coupling f does not play any role in relic abundance, there is more freedom in the choice of parameters and it yields no sizable predictions. It is important to note that a single Majorana fermionic DM candidate by itself can successfully resolve $(g-2)_\mu$ with $f_{\mu i} \sim \mathcal{O}(1)$. This would open up the parameter space with much weaker constraints arising from LFV processes. However, such a choice would not lead to direct correlation between $(g-2)_e$ and DM, and also does not lead to a LFV prediction.

VII. CONCLUSIONS

In light of recent experimental results confirming a 4.2σ discrepancy in the measurement of $(g-2)_\mu$ and a possible 7σ excess in the mass of the W boson, it is imperative to investigate new physics contributions for clarification. We propose the scotogenic-Zee model, a simple charged singlet extension of the scotogenic model, to show a direct correlation between these anomalies and the observed neutrino oscillation data, as well as the dark matter relic abundance. We explore the parameter space spanned by both the bosonic and fermionic dark matter candidates and provide a coherent resolution to electron and muon AMMs and M_W anomaly while evading dangerous LFV processes like $\mu \to e\gamma$ and $\mu \to 3e$. In contrast to the inert doublet/scotogenic models, where the small mass splitting among the doublet fields required for the observed relic density is disfavored by the CDF measurement, the scalar DM candidate in our model survives due to the presence of the extra charged singlet, which is essential in resolving Δa_ℓ. This model predicts large rates for LFV processes $\tau \to \ell\gamma$ which can be tested in future experiments.

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APPENDIX: OBLIQUE PARAMETERS

Here we provide additional plots in the model to show the allowed parameter space consistent with the upward CDF W mass shift.

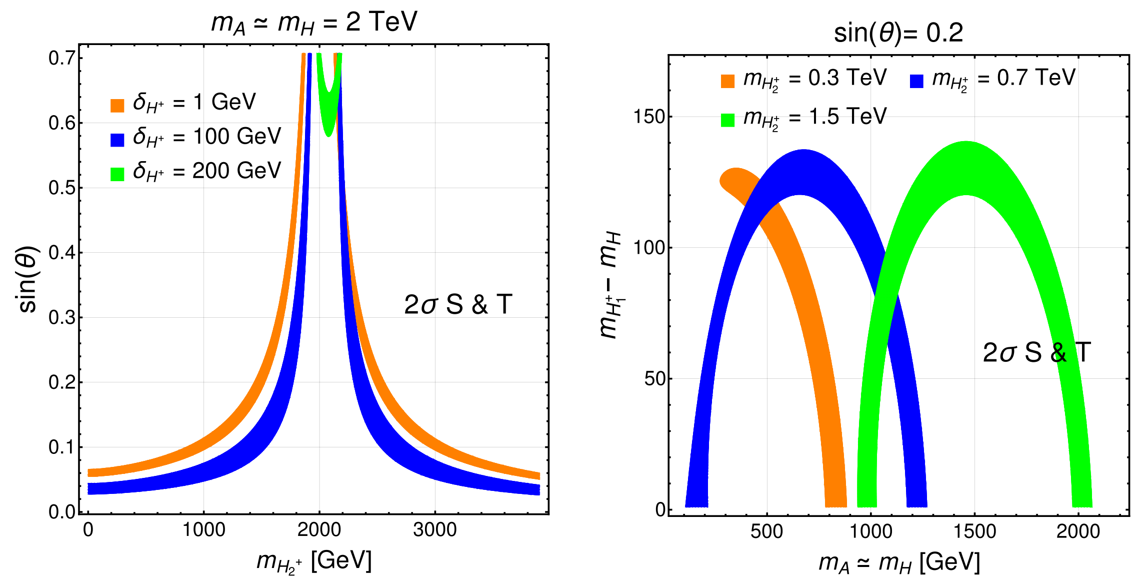

FIG. 5. Mixing angle θ as a function of the charged scalar mass $m_{H_2^+}$ for different mass splittings (left panel) $\delta_{H^+} = m_{H_1^+} - m_H$ and (right panel) the mass splitting δ_{H^+} as a function of neutral scalars for different choices of charged singlet scalar mass explaining the upward shift in M_W reported in the CDF measurements, which is consistent with the 2σ ranges of S and T.

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