W boson mass shift, dark matter and \((g-2)_\mu\) in ScotoZee model

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We present a singly charged scalar extension of the Scotogenic model, ScotoZee, which resolves the recently reported deviations in \(W\) 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 \(\sim 100\) GeV, required by the shift in \(W\) boson mass, among the inert doublets fields can be evaded by its mixing with the singlet scalar, which is also key to resolving \((g-2)\) anomaly within 1\(\sigma\). We establish the consistency of this framework with dark matter relic abundance while satisfying constraints from charged lepton flavor violation, direct detection as well as collider bounds. The model gives predictions for the lepton flavor violating \(\tau \to \ell\gamma\) processes testable in upcoming experiments.

**Introduction:** The CDF collaboration at Fermilab [1] reported a precision measurement of \(W\) boson mass, \(M_W^{\text{CDF}} = (80.335 \pm 0.0094)\) GeV, which is in tension with the Standard Model (SM) prediction, \(M_W^{\text{SM}} = (80.357 \pm 0.004)\) GeV [2], with an excess at 7\(\sigma\) level, which may be an indication of new physics (NP) beyond the Standard Model (SM). Some possible explanations to the \(W\) boson mass shift can arise at tree level [3–18], or at loop level [19–35], along with the prospect of reconciling the anomalous magnetic moment (AMM) of muon measurement at BNL in 2006 [91] at a combined 4\(\sigma\). The proposed ScotoZee model is a simple charged singlet \(S^+\) \((1,1; -)\) extension of the Scotogenic [97] model, which contains Majorana singlet fermions \(N_{R_1}, (1,0; -)\) and the scalar doublet \((\eta^+, \eta^0) \equiv \eta (2,1/2; -)\), under the gauge group \( SU(2)_L \times U(1)_Y \times Z_2 \). All the new particles are odd under \(Z_2\) while the SM particles are even, guaranteeing the stability of the DM candidate; the lightest among the new neutral \(Z_2\)-odd particles. The charged scalar singlet \(S^+\) not only gives corrections to anomalous magnetic moment of muon and electron through the mixing with charged doublet, but also serves as a portal to generate correct relic abundance for fermionic DM.

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

\[- \mathcal{L}_Y \supset Y_{ij} \bar{t}_L \eta N_{R_i} + f_{ij} \bar{t}_{R_i} S^- N_{R_j} + h.c. \] (1)

The \(Z_2\) symmetry, being exact, prevents \(\eta^0\) from obtaining a non-zero vacuum expectation value (VEV) and neutrinos remain massless at tree level. Moreover, the SM Higgs \(h\) is decoupled from the new \(CP\)-even \((\text{Re}(\eta^0) \approx H)\) and \(-\text{odd} (\text{Im}(\eta^0) \approx A)\) scalars. The charged scalars \(\{\eta^+, S^+\}\) mix giving rise to mass eigenstates \(\{H^+_1, H^+_2\}\). The masses of the scalar fields in the physical basis are given by

\[ m^2_{\eta} = \lambda_1 v^2, \quad m^2_{H(A)} = \mu^2 + \frac{v^2}{2} (\lambda_3 + \lambda_4 \pm \lambda_5), \]

\[ m^2_{H^+} = \frac{1}{2} \left( \mu_2 + \mu_3 \pm \sqrt{(\mu_2 - \mu_3)^2 + 2 \mu_2 v^2} \right), \] (2)

where \(\mu_2 = \mu_1^2 + \frac{\lambda_3}{2} v^2, \quad \mu_3 = \mu_5^2 + \frac{\lambda_5}{2} v^2\). Here \(\mu_1,s, \lambda_i\), and \(\mu\) are the bare-mass terms, quartic couplings, and

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1 The scalar content is the same as the Inert Zee model [98, 99] with only right-handed neutrinos in contrast to vector-like singlets and doublets. The ScotoSinglet model [100] is a neutral scalar extension of the Scotogenic model. Neither models can resolve the discrepancy in \((g-2)_\mu\) [101].
cubic coupling, respectively. The mixing angle between the charged scalar fields is defined as

$$\sin 2\theta = \frac{-\sqrt{2} \mu v}{m_H^2 - m_{H^+}^2}, \quad (3)$$

with the VEV, $v \simeq 246$ GeV. In this work, we comply with the perturbative and vacuum stability conditions [105, 106] constraining the scalar couplings. The Majorana mass term $\frac{1}{2} M_N N_i N_j$ along with the scalar quartic term $\frac{\lambda}{2} (\phi^\dagger \phi)^2 + \text{h.c.}$ breaks the lepton number by two units generating one-loop neutrino mass (c.f. Fig. 1 (top)) $M_\nu$ expressed as

$$(M_\nu)_{ij} = \sum_k Y_{ik} \Lambda_k Y_{kj}^*, \quad \Lambda_k = \frac{M_N}{16\pi^2} \left[ m_H^2 - M_{N_k}^2 \log \frac{m_H^2}{M_{N_k}^2} - (m_H \leftrightarrow m_A) \right]. \quad (4)$$

Here the lightest mass eigenstates $\{H, A\}$ and $N_1$ can serve as a viable bosonic and fermionic DM candidates. It is important to point out that unlike the Scotogenic model, where $M_N$ can be at canonical seesaw scale of $10^9$ GeV or the Yukawa coupling $Y$ arbitrarily small, the $(g - 2)_\mu$ in the model requires the scale to be in the (sub) TeV range along with $O(0.1 - 1.0)$ Yukawa coupling. Thus, a successful explanation of $m_\nu \sim 0.1$ eV would naturally require $m_H$ to be nearly degenerate with $m_A$.

**Correction to $W$ boson mass:** The shift in $W$ boson mass [107] can be evaluated as a function of the oblique parameters, $S, T$ and $U$ [108, 109] 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 [110] except for the $Z_2$ charge preventing the mixing with the SM Higgs doublet. Therefore, we use the expressions for $S, T$ and $U$ given in Ref. [111] under the alignment limit [112]. 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 [4] with the new CDF data to quote the $2\sigma$ allowed ranges of oblique parameters. We confirm the necessity of mass splitting in 2HDM [24, 25] 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 as can be seen from Fig. 2. The splitting $\delta_{H^+} = m_{H^+} - m_H$ depends on the mixing angle, for instance, it can be at most $\sim 140$ GeV for $\sin \theta = 0.2$.

**Anomalous Magnetic Moment:** The charged scalar contributions to anomalous magnetic moment at one-loop [113] as shown in Fig. 1 (bottom) is

$$\Delta a_{H^+}^{m_T} = \frac{m_T^2}{16\pi^2} \left( |Y_{\ell i}|^2 \sin^2 \theta + |f_{\ell i}|^2 \cos^2 \theta \right) \cdot G[m_{H^+}, 2] + \frac{M_N}{m_\ell} \text{Re}(Y_{\ell i} f_{\ell i}^*) \sin 2\theta \cdot G[m_{H^+}, 1], \quad (5)$$
where,

\[
G[M, \varepsilon] = \int_0^1 \frac{x^5(x-1)}{m_\ell^2x^2 + (M^2 - m_\ell^2)x + M^2(1-x)} \, dx
\]  

(6)

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

Note that the Yukawa couplings and the masses of charged scalars are severely restricted by the charged lepton flavor violating (cLFV) processes such as radiative decay \( \ell_i \to \ell_j \gamma \) [114]; 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 the tree level, they arise at the one-loop with large branching ratios [115]. The same is also true for \( \mu \to e \) conversion in the nuclei. We impose these constraints in our parameter scan.

**DM Phenomenology:** In addition to explaining \( W \) boson mass shift and \( \Delta a_\ell \), the proposed model can easily accommodate both the scalar (lightest of \( H \) and \( A \)) and fermionic (lightest among \( N_i \)) dark matter candidates (\( \chi \)). We consider both scenarios and analyze the parameter space by implementing the model in SARAH [119] and numerically evaluating the relic abundance using the software MicrOMEGAs [120]. The relic density of DM is achieved through standard thermal freeze-out mechanism.

For the case of Majorana fermion as a DM (\( \chi \equiv N \)) candidate, the annihilation channel which determines the observed relic abundance is \( \text{DM self-(co-)annihilation into charged leptons} \ell_i \ell_j^\mp \) (light neutrinos \( \nu_\alpha \nu_\beta \)) through \( t \)-channel processes mediated by the \( Z_2 \)-odd scalars, \( H_i^\pm (H, A) \) via Yukawa couplings \( Y \) and/or \( f \). The neutrino oscillation data determines the flavor structure of \( Y \) making it natural to select relatively small \( Y \) and heavy doublet scalar \( \eta \sim O \text{ (TeV)} \) such that the LFV constraints are relaxed. Thus, we choose \( f_{\alpha \beta} \sim O(1) \) \((i = 1, 2)\) and degenerate \( N_i \) to maximize the contribution to annihilation mode \( \chi \chi \to \ell \ell \) via \( s^\pm \); the allowed parameter space in the mass plane can be seen in Fig. 3 (top) along with the region resolving muon AMM for a specific choice of \( \kappa = Y^* f \sin \theta = 0.015 \).

In the case of scalar dark matter, which we choose to be the \( CP \)-even \( H \equiv \chi \) (nearly degenerate\(^2 \) with \( A \) and \( \lambda_5 < 0 \)), pair of DM can annihilate to \( W^+W^- \), \( ZZ \), \( \nu_\alpha \nu_\beta \), \( hh \), \( \ell \ell \), and \( qq \). The low mass regime suffer a strong constraint form LEP [117, 118, 122] which can be satisfied if one assumes \( m_\chi > M_Z/2, m^\pm_H > M_W/2 \) and \( m^+_H + m_\chi > M_W \). For larger DM mass, it predominantly annihilates to a pair of \( W^+W^- \) and \( ZZ \), for which the allowed region is \( m_\chi \gtrsim 500 \text{ GeV} \) and mass splitting \( \delta_{\chi^+} = m^+_H - m_\chi \lesssim 30 \text{ GeV} \) as shown in Fig. 3 (bottom). This can be relaxed by making the Higgs quartic coupling \( \gtrsim 1 \),

\(^2\) The mass splitting is of order \( O(100) \text{ keV} \) [121] to evade direct detection.
oscillation data and efficiently probe the model with LFV neutrino oscillation data. 

The masses of the Majorana fermions and the charged scalars with are constrained by LFV processes. With the choice of nonzero \( f_{ii} \) for \( i = 1, 2 \) to explain both the AMMs as well as to obtain the observed relic density, the mass enhancement to \( \ell_i \to \ell_j \gamma \) severely restricts the parameter space. Such 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 \( Y_{21} f_{11} \approx - Y_{12} f_{22} \) for \( M_{N_i} = M_{N_6} (= m_{\chi}) \). 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 does not play any role in relic abundance, there is more freedom in the choice of parameters and yield no sizeable predictions.

**Conclusions:** In the 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 \( W \) boson it is imperative to investigate new physics contributions for clarification. We propose the ScotoZee 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 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 \( \mathcal{M}_W \) anomaly while evading dangerous LFV processes like \( \mu \to e\gamma \) and \( \mu \to 3e \). In contrast to the IDM/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 \( \Delta \ell \). This model predicts large rates for LFV processes \( \tau \to \ell \gamma \) which can be tested in the future experiments.

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**FIG. 4:** Scattered plot assuming the fermionic DM with the same parameter space given in Fig. 3 (top). Colored shaded regions are the current exclusion limits [127], whereas dashed-dotted lines represent the future projected sensitivities [128]. The orange (green) dots correspond to solutions that resolve \( \Delta \) and \( \Delta \) observed values within their 2σ measured values [129].

a choice strongly constrained by direct detection bound [123–126].

In this work we take the quartic couplings \( \lambda_3 + \lambda_4 + \lambda_5 \ll 1 \) to automatically satisfy direct detection bound obtained from the scalar DM interacting with nucleus at the tree level through the SM Higgs boson. Moreover, it is favored to take the couplings \( Y_{ij} \) small and \( M_N \sim \mathcal{O}(\text{TeV}) \) to be consistent with 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, disfavouring 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 (c.f. Fig 2), thereby admitting the \( \text{CP} \)-even \( H \) to be a viable DM candidate, as shown in Fig. 3 (bottom).

**Neutrino Fit/ Lepton Flavor violation:** The neutrino mass formula of Eq. (4), lepton \( g - 2 \) and the dark matter analysis have close-knot correlation through Yukawa couplings, Majorana fermions, and new scalars. As previously stated, \((g-2)_\mu\) 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 \( m_A \simeq m_H \), 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 [130] to rewrite the Yukawa matrix \( Y \) of Eq. (4) in terms of neutrino mass parameters

\[
Y = \sqrt{\Lambda^{-1}} R \sqrt{M_{\nu}^{\text{diag}}} U_{\text{PMNS}},
\]

where \( R \) is an arbitrary complex orthogonal matrix. The neutrino oscillation parameters are scanned within the 2σ allowed ranges [129] to obtain the Yukawa matrix.

As mentioned earlier, the product of Yukawa couplings \( Y_{ij} f_{ii} \) can explain \((g-2)_\mu\); however these couplings are constrained by LFV processes. With the choice of nonzero \( f_{ii} \) for \( i = 1, 2 \) to explain both the AMMs as well as to obtain the observed relic density, the mass enhancement to \( \ell_i \to \ell_j \gamma \) severely restricts the parameter space. Such 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 \( Y_{21} f_{11} \approx - Y_{12} f_{22} \) for \( M_{N_i} = M_{N_6} (= m_{\chi}) \). 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 does not play any role in relic abundance, there is more freedom in the choice of parameters and yield no sizeable predictions.

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Appendix A: Scalar potential

The most general renormalizable scalar potential of the ScotoZee model is given by:

\[
V = \mu_2^2 \phi^\dagger \phi + \mu_3^2 S^- S^+ + \mu_4^2 \eta^\dagger \eta + \frac{\lambda_1}{2} (\phi^\dagger \phi)^2 + \frac{\lambda_2}{2} (\eta^\dagger \eta)^2 + \lambda_3 (\phi^\dagger \phi)(\eta^\dagger \eta) + \lambda_4 (\phi^\dagger \phi)(\eta^\dagger \eta) + \frac{\lambda_5}{2} (\phi^\dagger \phi)^2 + \text{h.c.}
\]

\[+ \frac{\lambda_6}{2} (S^- S^+)^2 + \lambda_7 (\phi^\dagger \phi)(S^- S^+) + \lambda_8 (\eta^\dagger \eta)(S^- S^+) + \frac{\mu}{2} \{\varepsilon_{\alpha\beta\phi^\dagger \phi^\dagger \eta^\dagger \eta^\dagger} S^- + \text{h.c.}\}.
\]

(8)

Appendix B: Oblique parameters

FIG. 5: Mixing angle $\theta$ as a function of charged scalar mass $m_{H^+}$ for different mass splittings, $\delta_{H^+} = m_{H^+} - m_H$ (left) and the mass splitting $\delta_{H^+}$ as a function of neutral scalars for different choices of charged singlet scalar mass (right) explaining the upward shift in $M_W$ reported by CDF measurement, consistent with the 2$\sigma$ ranges of $S$ and $T$.  

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