

Dark-portal leptogenesis in a nonholomorphic modular scoto-seesaw model

Salah Nasri^{1,*}, Labh Singh^{2,†}, Tapender^{2,‡} and Surender Verma^{2,§}

¹*Department of Physics, United Arab Emirates University, Al-Ain, United Arab Emirates*

²*Department of Physics and Astronomical Science, Central University of Himachal Pradesh, Dharamshala, Himachal Pradesh 176215, India*



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This work explores the neutrino phenomenology of the scoto-seesaw model under nonholomorphic A_4 modular flavor symmetry providing a nonsupersymmetry framework for realization of the modular symmetry. To prevent mixing between the beyond-standard-model fields associated with the tree- and loop-level neutrino mass contributions, we assign even and odd modular weights to these sectors, respectively. The physical allowed ranges of oscillation parameters are used to identify the viable region of modulus parameter τ in its fundamental domain. With the complex modulus τ serving as the unique source of CP violation (all other parameters are real), the framework realizes successful low-scale leptogenesis through CP -violating decays of the lightest righthanded neutrino into Standard model leptons and the Higgs boson. The requisite CP asymmetry arises from one-loop diagrams involving dark-sector states, obviating the need for degenerate mass spectra and thereby circumventing the usual resonant leptogenesis mechanism. The observation of a long-lived charged particle ($\eta^\pm$) in collider experiments would offer compelling evidence for the inert scalar sector of the model and provide a crucial experimental hint as to the dark-sector-assisted generation of neutrino masses and leptogenesis.

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I. INTRODUCTION

The confirmed observation of neutrino oscillations provides definitive evidence that neutrinos are massive and undergo flavor-mixing phenomena that necessitate physics beyond the Standard Model (SM) [1–6]. A fundamental theoretical challenge is the absence of righthanded neutrino counterparts in the SM, which precludes the generation of neutrino mass *via* the conventional Higgs mechanism. While the dimension-five Weinberg operator [7] offers a viable framework for neutrino mass generation, its ultraviolet completion and the specific flavor structure of the neutrino sector remain open questions. Consequently, exploring beyond the standard model (BSM) scenarios is crucial for understanding the origin of neutrino mass. A wide spectrum of models has been proposed to align with oscillation data, including various implementations of the seesaw mechanism, radiative mass generation, and models

involving extra dimensions [8–39]. A common feature in many of these frameworks is the introduction of sterile neutrinos: SM gauge singlets that act as righthanded partners. These sterile neutrinos couple to the active left-handed neutrinos *via* Yukawa interactions. The masses and coupling strengths of these sterile states can span many orders of magnitude, leading to diverse phenomenological implications. For instance, within the canonical type-I seesaw, the observed sub-eV neutrino masses naturally point to a very high mass scale for the righthanded neutrinos, around 10^{15} GeV, which is inaccessible to current experiments. However, alternative realizations of the seesaw mechanism such as the inverse, linear, or extended seesaw, often constructed with discrete or continuous symmetries [40–46] can feature sterile neutrino masses at the TeV scale, rendering them potentially testable at current or future colliders. To date, neutrino oscillation experiments have precisely measured two mass-squared differences associated with atmospheric and solar neutrino oscillations. The existence of these two distinct mass scales may itself be a clue, potentially indicating that the masses for the second and third generations originate from different underlying mechanisms.

To validate our assertion, we employ a recently proposed idea of a nonholomorphic modular symmetry, a non-supersymmetry SI realization of the holomorphic modular symmetry [47]. This approach is rooted in the construction of nonholomorphic modular forms, specifically

*Contact author: snasri@uaeu.ac.ae

†Contact author: sainilabh5@gmail.com

‡Contact author: tapenderphy@gmail.com

§Contact author: s_7verma@hpcu.ac.in

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polyharmonic Maass forms. Unlike their holomorphic counterparts, these forms are invariant under the modular group but are not required to be holomorphic; instead, they satisfy eigenfunction conditions under a Laplace-type differential operator. This expanded functional space incorporates both holomorphic and nonholomorphic components, and accommodates modular forms of both positive and negative weights. Consequently, it enables the construction of more flexible and phenomenologically viable Yukawa coupling structures. Crucially, the rigid constraint of modular invariance is preserved, ensuring the model's predictive power while circumventing the limitations intrinsic to a purely holomorphic framework. Several works pertaining to the neutrino mass model building have already been done in the nonholomorphic modular symmetry [21,48–67].

Further, the observed cosmic matter–antimatter asymmetry remains unexplained in the SM. Leptogenesis is a leading theoretical solution, proposing that this imbalance originated in the lepton sector [68–72]. It typically involves the out-of-equilibrium, CP -violating decay of heavy right-handed neutrinos in the early universe. This lepton asymmetry is subsequently converted into the observed baryon asymmetry *via* sphaleron processes. A notable feature of leptogenesis is its connection to the low-energy neutrino phenomenology through heavy righthanded degrees of freedom in the seesaw framework. The modular symmetry framework offers a highly predictive setting for leptogenesis, as the Yukawa couplings become functions of the complex modulus τ , which in turn acts as a natural source of CP asymmetry. Leptogenesis within modular-invariant theories has been extensively studied in holomorphic formulations [20,22,30–39] and some nonholomorphic extensions [21,65,66].

Motivated by this, we explored the *scoto-seesaw* framework [73–75], where one mass square difference originates at the seesaw level and another at the one-loop (scotogenic) level. To forbid mixing between the tree-level and loop-level BSM sectors, distinct modular weights (κ) are assigned: tree-level fields carry even κ , whereas loop-level fields are odd under κ . In this setup, the heavy righthanded neutrino N_1 , responsible for the tree-level seesaw contribution, decays into the SM leptons and Higgs doublets. The dark sector, consisting of a chiral singlet fermion f , an inert scalar doublet, and a real scalar singlet s , all odd under κ , plays a multifaceted role: (i) it induces the one-loop radiative neutrino mass contribution,

TABLE I. Particle content and charge assignments of the scoto-seesaw model under $SU(2)_L$ and A_4 nonholomorphic modular symmetry.

Fields	$SU(2)_L$	A_4	κ
L	2	3	2
e_R, μ_R, τ_R	1	1, 1', 1''	−2
H	2	1	0
N_1	1	1	2
N_2	1	1'	2
f	1	1	1
η	2	1	1
s	1	1	−1

(ii) it mediates loop-level corrections to the N_1 decay amplitude whose interference with the tree-level process generates a nonzero CP asymmetry, and (iii) it offers a viable dark matter candidate within the model. The interaction $\mu_1 s \eta^\dagger H$, together with the Yukawa structure involving the fermion multiplet f and the singlet s , directly enters the one-loop diagrams responsible for leptogenesis. Consequently, the presence of s and the other BSM fields is not incidental but essential, as they generate the loop-induced CP asymmetry, modify the washout structure, and allow successful leptogenesis at a comparatively lower scale without requiring near-degenerate heavy neutrinos.

The paper is organized as follows. In Sec. II, we present the details of the model framework. Section III is devoted to explaining the procedure involved in numerical estimation of neutrino masses and mixing parameters. In Sec. IV, we discuss the charged lepton flavor violation in the model. A discussion of dark-sector-assisted leptogenesis, together with brief comments on the dark matter implications, is presented in Sec. V. Finally, conclusions are summarized in Sec. VI.

II. MODEL AND FORMALISM

The modular symmetry-based scoto-seesaw model is given in Table I. The basic algebra of the modular symmetry has been given in the Appendix. To ensure there is no mixing between the tree- and loop-level BSM fields, we have assigned even and odd weights to the tree- and loop level fields.

The relevant scalar potential of the model is given by

$$\begin{aligned}
 V_{\text{tree}} = & \mu_H^2 |H|^2 + \mu_\eta^2 |\eta|^2 + \lambda_1 |H|^4 + \lambda_2 |\eta|^4 + \lambda_3 |H|^2 |\eta|^2 + \lambda_4 |\eta^\dagger H|^2 + \tilde{\lambda}_5 [(\eta^\dagger H)^2 + \text{H.c.}] \\
 & + \frac{\mu_s^2}{2} |s|^2 + \lambda_7 |s|^4 + \lambda_8 |s|^2 |H|^2 + \lambda_9 |s|^2 |\eta|^2 + (\mu_1 s H^\dagger \eta + \mu_1^* s \eta^\dagger H).
 \end{aligned} \tag{2.1}$$

Here, H is the SM Higgs doublet¹ and $\tilde{\lambda}_5 = \lambda$. Also, the scalar dark sector is comprised of the Higgs-type scalar doublet

$$\eta = \begin{pmatrix} \eta^+ \\ \frac{(\eta_R + i\eta_I)}{\sqrt{2}} \end{pmatrix}, \quad (2.2)$$

and a singlet scalar s . After spontaneous symmetry breaking, the charged and odd components of the η acquire mass as

$$m_{\eta^+}^2 = \mu_\eta^2 + \frac{1}{2}\lambda_3 v^2, \quad (2.3)$$

$$m_{\eta_I}^2 = \mu_\eta^2 + \frac{1}{2}(\lambda_3 + \lambda_4 - \tilde{\lambda}_5)v^2, \quad (2.4)$$

where v is the vacuum expectation value ($v \approx 246$ GeV) of the SM Higgs. In the (η_R, s) basis, the squared mass matrix is nondiagonal and is given by

$$M^2 = \begin{pmatrix} \frac{\partial^2 V}{\partial \eta_R^2} & \frac{\partial^2 V}{\partial \eta_R \partial s} \\ \frac{\partial^2 V}{\partial s \partial \eta_R} & \frac{\partial^2 V}{\partial s^2} \end{pmatrix} = \begin{pmatrix} \mu_\eta^2 + \frac{1}{2}(\lambda_3 + \lambda_4 - \tilde{\lambda}_5)v^2 & \mu_1 v \\ \mu_1 v & \mu_s^2 + \frac{1}{2}\lambda_8 v^2 \end{pmatrix}, \quad (2.5)$$

which can be diagonalized as

$$\begin{pmatrix} \eta_1 \\ \eta_2 \end{pmatrix} = \begin{pmatrix} \cos \theta & \sin \theta \\ -\sin \theta & \cos \theta \end{pmatrix} \begin{pmatrix} \eta_R \\ s \end{pmatrix}, \quad (2.6)$$

such that

$$\tan 2\theta = \frac{2\mu_1 v}{\mu_\eta^2 - \mu_s^2 + \frac{1}{2}(\lambda_3 + \lambda_4 - \tilde{\lambda}_5 - \lambda_8)v^2}. \quad (2.7)$$

The mass of CP even scalars is given by.² Using a binomial expansion,³ the mass-squared eigenvalues of the inert scalars can be approximated as

$$m_{\eta_1}^2 = \frac{1}{2} \left[\mu_\eta^2 + \mu_s^2 + \frac{1}{2}(\lambda_3 + \lambda_4 - \tilde{\lambda}_5 + \lambda_8)v^2 + \sqrt{\left(\mu_\eta^2 - \mu_s^2 + \frac{1}{2}(\lambda_3 + \lambda_4 - \tilde{\lambda}_5 - \lambda_8)v^2 \right)^2 + 4\mu_1^2 v^2} \right],$$

$$m_{\eta_2}^2 = \frac{1}{2} \left[\mu_\eta^2 + \mu_s^2 + \frac{1}{2}(\lambda_3 + \lambda_4 - \tilde{\lambda}_5 + \lambda_8)v^2 - \sqrt{\left(\mu_\eta^2 - \mu_s^2 + \frac{1}{2}(\lambda_3 + \lambda_4 - \tilde{\lambda}_5 - \lambda_8)v^2 \right)^2 + 4\mu_1^2 v^2} \right].$$

³We considered λ_8 small to eradicate unwanted mixing between SM Higgs and singlet scalar s .

$$m_{\eta_{1,2}}^2 \simeq \frac{1}{4} \left[2(\mu_\eta^2 + \mu_s^2) + v^2(\lambda_3 + \lambda_4 - \tilde{\lambda}_5) \pm \sqrt{(v^2(\lambda_3 + \lambda_4 - \tilde{\lambda}_5) - 2\mu_s^2 + 2\mu_\eta^2)^2 + 16\mu_1^2 v^2} \right]. \quad (2.8)$$

The dark matter candidate in the model can be either a CP even (η_1, η_2) or CP odd (η_I) scalar, depending on the mass spectrum considered. The relevant Yukawa Lagrangian for the charged leptons, based on the field content given in Table I, is given by

$$\mathcal{L} \supset \alpha_1 Y_3^{(0)} (\bar{L} H e_R) + \alpha_2 Y_3^{(0)} (\bar{L} H \mu_R) + \alpha_3 Y_3^{(0)} (\bar{L} H \tau_R) \quad (2.9)$$

where $Y_3^{(0)} = (Y_1^0, Y_2^0, Y_3^0)$ is an A_4 triplet modular form with weight 0. From this Lagrangian, charged lepton mass matrix can be written as

$$M_\ell = v \begin{pmatrix} \alpha_1 Y_1^0 & \alpha_2 Y_3^0 & \alpha_3 Y_2^0 \\ \alpha_1 Y_3^0 & \alpha_2 Y_2^0 & \alpha_3 Y_1^0 \\ \alpha_1 Y_2^0 & \alpha_2 Y_1^0 & \alpha_3 Y_3^0 \end{pmatrix} = v \alpha_1 \begin{pmatrix} Y_1^0 & r_{21} Y_3^0 & r_{31} Y_2^0 \\ Y_3^0 & r_{21} Y_2^0 & r_{31} Y_1^0 \\ Y_2^0 & r_{21} Y_1^0 & r_{31} Y_3^0 \end{pmatrix},$$

where $r_{21} = \alpha_2/\alpha_1$ and $r_{31} = \alpha_3/\alpha_1$. The invariant Yukawa Lagrangian for neutrinos, including the Dirac Yukawa interactions and Majorana mass terms for the righthanded neutrinos, is given by

$$\mathcal{L} \supset \gamma \left(Y_3^{(4)} \bar{L} \tilde{H} N_1 + Y_3^{(4)} \bar{L} \tilde{H} N_2 \right) + \kappa_1 Y_1^{(4)} N_1 N_1 + \kappa_2 Y_{1'}^{(4)} N_2 N_2. \quad (2.10)$$

The righthanded neutrino mass matrix is then given by

$$M_R = \begin{pmatrix} \kappa_1 Y_1^{(4)} & 0 \\ 0 & \kappa_2 Y_{1'}^{(4)} \end{pmatrix}, \quad (2.11)$$

¹Although μ_1 in general can be complex, we take it as real in our study. It is to be noted that most of the terms in the scalar potential remain invariant under modular symmetry as given in [49].

²

and Dirac neutrino mass matrix is

$$M_D = \gamma v \begin{pmatrix} Y_1^4 & Y_3^4 \\ Y_3^4 & Y_2^4 \\ Y_2^4 & Y_1^4 \end{pmatrix}, \quad (2.12)$$

where $Y_3^{(4)} = (Y_1^4, Y_2^4, Y_3^4)$ is A_4 triplet modular form with weight 4. The neutrino masses generated by Type-I seesaw is given by

$$(m^\nu)_{\text{tree}} = -M_D M_R^{-1} M_D^T. \quad (2.13)$$

Additionally, the model exhibits a scotogenic nature. The relevant Yukawa Lagrangian, which is important for the one-loop generation of neutrino mass, can be expressed as

$$L = \beta_L Y_3^{(4)} \bar{L} \tilde{\eta} f + \kappa_S Y_1^{(2)} f f + a Y_1^{(2)} \bar{N}^c f s. \quad (2.14)$$

The loop-induced neutrino mass matrix takes the form

$$(m_{ij}^\nu)_{\text{loop}} = F(m_{\eta_1}, m_{\eta_2}, m_{\eta_l}) M_f h_i h_j, \quad (2.15)$$

where $M_f = \kappa_S Y_1^{(2)}$, and $F(m_{\eta_1}, m_{\eta_2}, m_{\eta_l})$ denotes the loop function

$$F(m_{\eta_1}, m_{\eta_2}, m_{\eta_l}) = \frac{1}{32\pi^2} [\cos \theta \mathcal{F}(m_{\eta_1}) + \sin \theta \mathcal{F}(m_{\eta_2}) - \mathcal{F}(m_{\eta_l})], \quad (2.16)$$

with

$$\mathcal{F}(m_x) = \frac{m_x^2}{M_f^2 - m_{\eta_x}^2} \ln \frac{m_x^2}{M_f^2}. \quad (2.17)$$

The neutrino mass matrix generated at the loop level is

$$(m^\nu)_{\text{loop}} = \beta_L^2 M_f \begin{pmatrix} (Y_1^4)^2 & Y_1^4 Y_2^4 & Y_1^4 Y_3^4 \\ Y_1^4 Y_2^4 & (Y_2^4)^2 & Y_2^4 Y_3^4 \\ Y_1^4 Y_3^4 & Y_2^4 Y_3^4 & (Y_3^4)^2 \end{pmatrix} \times F(m_{\eta_1}, m_{\eta_2}, m_{\eta_l}). \quad (2.18)$$

Considering Eqs. (2.13) and (2.18), the total neutrino mass matrix is given by

$$M^\nu = (m^\nu)_{\text{tree}} + (m^\nu)_{\text{loop}}. \quad (2.19)$$

The next step is to ascertain the viability of the model to reproduce low-energy neutrino phenomenology which will help to identify the allowed parameter space of, initially, free parameters. In order to get the predictions for the neutrino oscillation parameters, we will numerically diagonalize M^ν in the next section.

III. NEUTRINO MASSES AND MIXING

The neutrino mass matrix given by Eq. (2.19) is diagonalized using the relation $U_\nu^T M^\nu U_\nu = \text{diag}(m_{\nu_1}, m_{\nu_2}, m_{\nu_3})$. The mixing matrix $U_{\text{PMNS}} = U_L^\dagger U_\nu$, since the charged lepton mass matrix is not diagonal in the flavor basis. Now, the mixing angle can be extracted from U_{PMNS} as

$$\sin^2 \theta_{13} = |U_{13}|^2, \quad \sin^2 \theta_{12} = \frac{|U_{12}|^2}{1 - |U_{13}|^2},$$

$$\sin^2 \theta_{23} = \frac{|U_{23}|^2}{1 - |U_{13}|^2}, \quad (3.1)$$

where U_{mn} are elements of U_{PMNS} . Further, the CP invariants sensitive to Majorana phases can be written in terms of the elements of the PMNS matrix

$$I_1 = \text{Im}[U_{11}^* U_{12}], \quad (3.2)$$

$$I_2 = \text{Im}[U_{11}^* U_{13}]. \quad (3.3)$$

Using effective Majorana mass

$$m_{\beta\beta} = |U_{11}^2 m_1 + U_{12}^2 m_2 + U_{13}^2 m_3|, \quad (3.4)$$

the model prediction for $m_{\beta\beta}$ is also investigated. We numerically diagonalize the neutrino mass matrix in Eq. (2.19) and elucidate predictions for the neutrino observables. In order to ascertain the allowed parameter space, we define the χ^2 function as

$$\chi^2 = \sum_{i=1}^7 \left(\frac{P_i - P_i^0}{\sigma_i} \right)^2, \quad (3.5)$$

where P_i represents the observables predicted by the model, P_i^0 the central value, and σ_i the error corresponding to 1 σ level as shown in Table III. We have sampled the seven observables in this study, which include three mixing angles, two mass-squared differences (Δm_{21}^2 and Δm_{31}^2), and two lepton mass ratios (m_e/m_μ and m_μ/m_τ) (see Table III). We explore the parameter space (given in Table II) using 10^9 samples in Monte Carlo simulations with mass hierarchy constraint $M_{N_1} < M_f \ll M_{N_2}$. The experimental input values used in our analysis are summarized in Table III. These parameters demonstrate how well the model reproduces the observed neutrino mixing angles in the case of normal hierarchy (NH) yielding a minimum chi-square value of $\chi_{\min}^2 = 0.28$. The model predictions are illustrated through correlation plots obtained for $\chi^2 \leq 25$, where the solid black star denotes the best-fit point corresponding to $\chi_{\min}^2 = 0.28$. Figure 1(a) presents the allowed region of the complex modulus $\tau = \text{Re}[\tau] + i \text{Im}[\tau]$ that remains consistent with the neutrino oscillation data. We find that the data are well fitted with $\text{Re}[\tau] \sim \pm 0.26$ and $\text{Im}[\tau] \sim 0.97$. The corresponding Yukawa couplings are displayed in Fig. 2. The predicted neutrino mixing angles lie within their 3 σ experimental

TABLE II. Ranges of the parameters used for the numerical scan.

Input parameters	Range
$\text{Re}[\tau]$	$\pm[0.0, 0.5]$
$\text{Im}[\tau]$	$[0.8, 1.5]$
γ	$[10^{-6}, 1]$
β_L	$[10^{-6}, 1]$
$\kappa_{1,2}$ (GeV)	$[10^5, 10^{10}]$
κ_s (GeV)	$[10^5, 10^{10}]$
$\nu\alpha_1$ (GeV)	$[10^{-3}, 10^{-2}]$
r_{21}	$[10^2, 1.5 \times 10^3]$
r_{31}	$[10^2, 1.5 \times 10^3]$
m_{η_1} (GeV)	$[500, 10^3]$
$\Delta m_{\eta}^2 = m_{\eta_1}^2 - m_{\eta_2}^2$ (GeV)	$[10^{-3}, 10]$
μ_1 (GeV)	$[10, 10^4]$

ranges as shown in Figs. 1(b) and 1(c). Fig. 1(d) illustrates the correlation between the atmospheric mixing angle θ_{23} and the Dirac-type CP -violating phase δ . We observe that the model exhibits nonzero CP violation with Dirac phase δ being highly constrained to lie in the first and fourth quadrants. While $\text{NO}\nu\text{A}$ data allow normal ordering with $\delta \in [0^\circ, 180^\circ]$ and inverted ordering with $\delta \in [180^\circ, 360^\circ]$ [76], the T2K experiment favors values around $\delta \simeq 270^\circ$ for both hierarchies [77]. This region partially overlaps with the $\text{NO}\nu\text{A}$ and T2K-preferred normal ordering solution. The future precision measurements will therefore provide a decisive test for this framework. The predictions for the effective Majorana mass $m_{\beta\beta}$ as a function of the lightest neutrino mass m_1 are shown in Fig. 1(e). The allowed region corresponds to the parameter space consistent with neutrino oscillation data in the case of NH. The model predicts $m_{\beta\beta}$ values that lie below the current experimental limits from KamLAND-Zen [78] and within the future sensitivities of LEGEND-1000 [79] and nEXO [80]. The observation of neutrinoless double beta decay ($0\nu\beta\beta$) in the upcoming experiments with refined sensitivities will therefore provide a critical test of the framework. The correlation between the invariants I_1 and I_2 (Eqs. (3.2)–(3.3)) is shown in Fig. 1(f). The results indicate small but nonzero values of $I_{1,2}$, implying nonzero Majorana-type CP violation.

TABLE III. The neutrino oscillation observables from NuFIT 6.0 [81]. The values of the mass ratios of charged leptons have been taken from Ref. [82].

Observable	Best-fit $\pm 1\sigma$ (NH)	Best-fit $\pm 1\sigma$ (IH)	3σ range (NH)	3σ range (IH)
$\theta_{12} [^\circ]$	$33.8^{+0.7}_{-0.6}$	$33.8^{+0.7}_{-0.6}$	$31.3 \rightarrow 36.3$	$31.3 \rightarrow 36.3$
$\theta_{23} [^\circ]$	$49.1^{+0.9}_{-1.0}$	$49.2^{+0.9}_{-1.0}$	$39.8 \rightarrow 51.9$	$40.2 \rightarrow 52.1$
$\theta_{13} [^\circ]$	$8.55^{+0.13}_{-0.14}$	$8.57^{+0.13}_{-0.14}$	$8.2 \rightarrow 8.9$	$8.2 \rightarrow 8.9$
$\Delta m_{21}^2 [10^{-5} \text{ eV}^2]$	$7.42^{+0.21}_{-0.20}$	$7.42^{+0.21}_{-0.20}$	$6.82 \rightarrow 8.04$	$6.82 \rightarrow 8.04$
$\Delta m_{31}^2 [10^{-3} \text{ eV}^2]$	$+2.510^{+0.027}_{-0.027}$	$-2.490^{+0.026}_{-0.028}$	$+2.43 \rightarrow +2.58$	$-2.56 \rightarrow -2.42$
m_e/m_μ	0.004737	0.004737	...	...
m_μ/m_τ	0.058823	0.058823	...	...

Furthermore, we have investigated the parameter space for the inverted neutrino mass hierarchy. The model excludes the IH scenario as illustrated in Fig. 3. The figure illustrates the correlation between the mixing angles θ_{12} and θ_{23} , showing that θ_{23} falls outside the 3σ range of the global data, thereby confirming the incompatibility with the IH scenario. While the preference for NH is consistent with global analyses, the prediction of a first-octant solution for θ_{23} is in mild tension with the slight upper-octant preference reported by T2K and $\text{NO}\nu\text{A}$. This discrepancy highlights a key testable prediction of our model. The best-fit values of neutrino observables and corresponding model parameters are given in Tables IV and V, respectively.

IV. LEPTON FLAVOR VIOLATION

One of the key low-energy signatures of physics beyond the SM is the presence of charged lepton flavor-violating (LFV) decays. In Scotogenic frameworks, the presence of new particles involved in the loop can induce sizable contributions to these processes. Among the various LFV channels, the decay $\mu \rightarrow e\gamma$ provides the strongest experimental constraint, with current upper limit $\text{Br}(\mu \rightarrow e\gamma) < 4.2 \times 10^{-13}$ [83]. In our analysis, we impose this bound to restrict the parameter space of the model. The branching ratio for the radiative decay $\mu \rightarrow e\gamma$ is [84]

$$\text{BR}(\mu \rightarrow e\gamma) = \frac{3(4\pi)^3 \alpha_{em}}{4G_F^2} |\mathcal{A}|^2 \text{BR}(\mu \rightarrow e\bar{\nu}_e\nu_\mu), \quad (4.1)$$

where α_{em} is the fine-structure constant, G_F is the Fermi constant, and $\text{BR}(\mu \rightarrow e\bar{\nu}_e\nu_\mu)$ is approximately equal to unity. The dipole amplitude entering Eq. (4.1) is

$$\mathcal{A} = \frac{Y_{11}^* Y_{21}}{32\pi^2 m_{\eta^+}^2} \mathcal{G}(y), \quad (4.2)$$

where m_{η^+} is the mass of the charged scalar field η^+ , and Y_{11} and Y_{21} are the corresponding elements of the rotated Yukawa coupling matrix in the charged-lepton mass basis. This rotated Yukawa matrix is defined as $Y = \beta_L U_L^\dagger (Y_1^4, Y_3^4, Y_2^4)^T$, where U_L is the unitary matrix

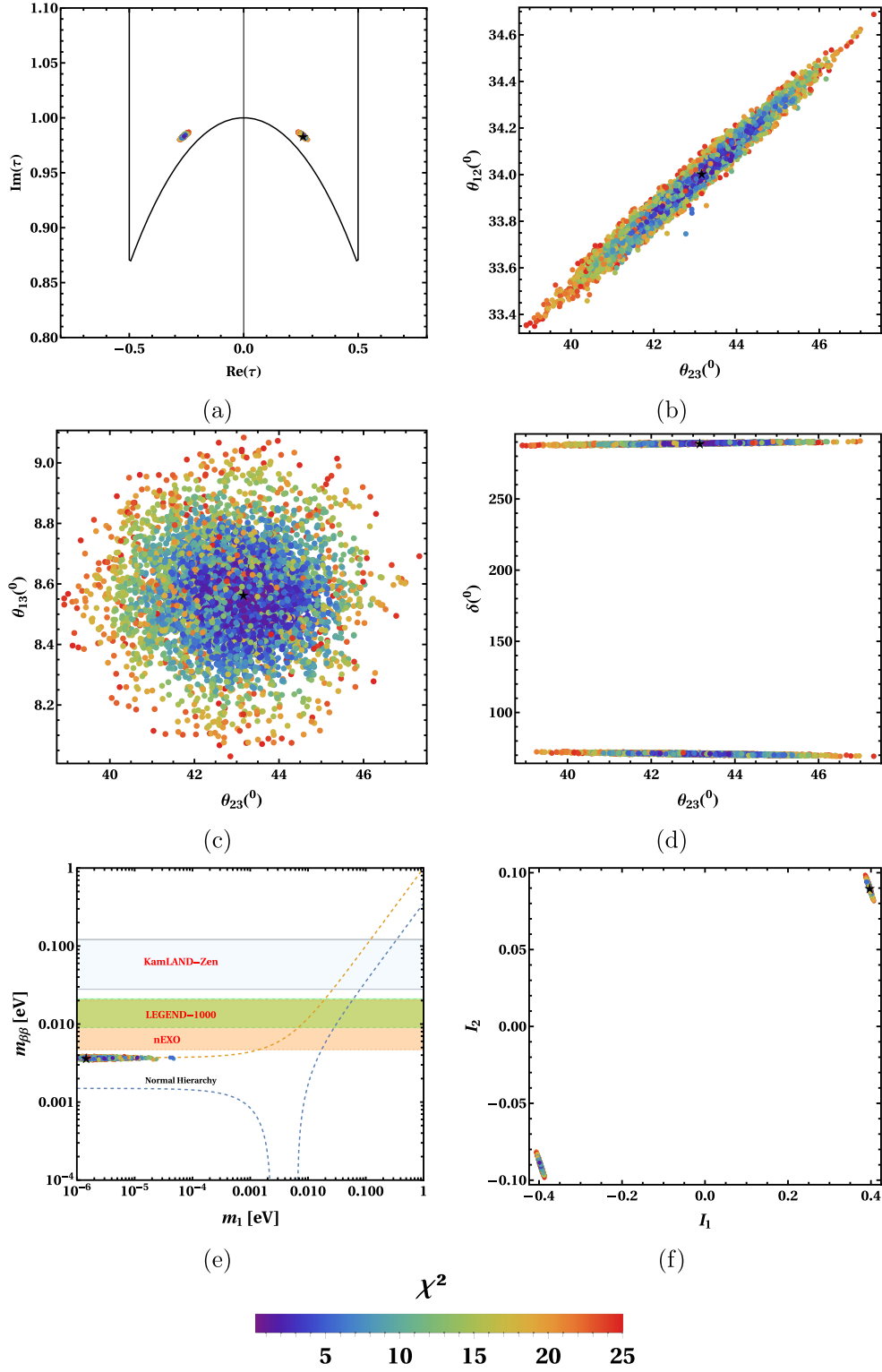


FIG. 1. Allowed parameter space of the complex modulus τ and predictions for neutrino observables: (a) $\text{Im}[\tau]$ vs $\text{Re}[\tau]$, (b) θ_{12} vs θ_{23} , (c) θ_{13} vs θ_{23} , (d) δ vs θ_{23} , (e) $m_{\beta\beta}$ vs m_1 , and (f) I_2 vs I_1 . The black star ($\star$) corresponds to the best-fit value obtained for $\chi^2_{\min} = 0.28$.

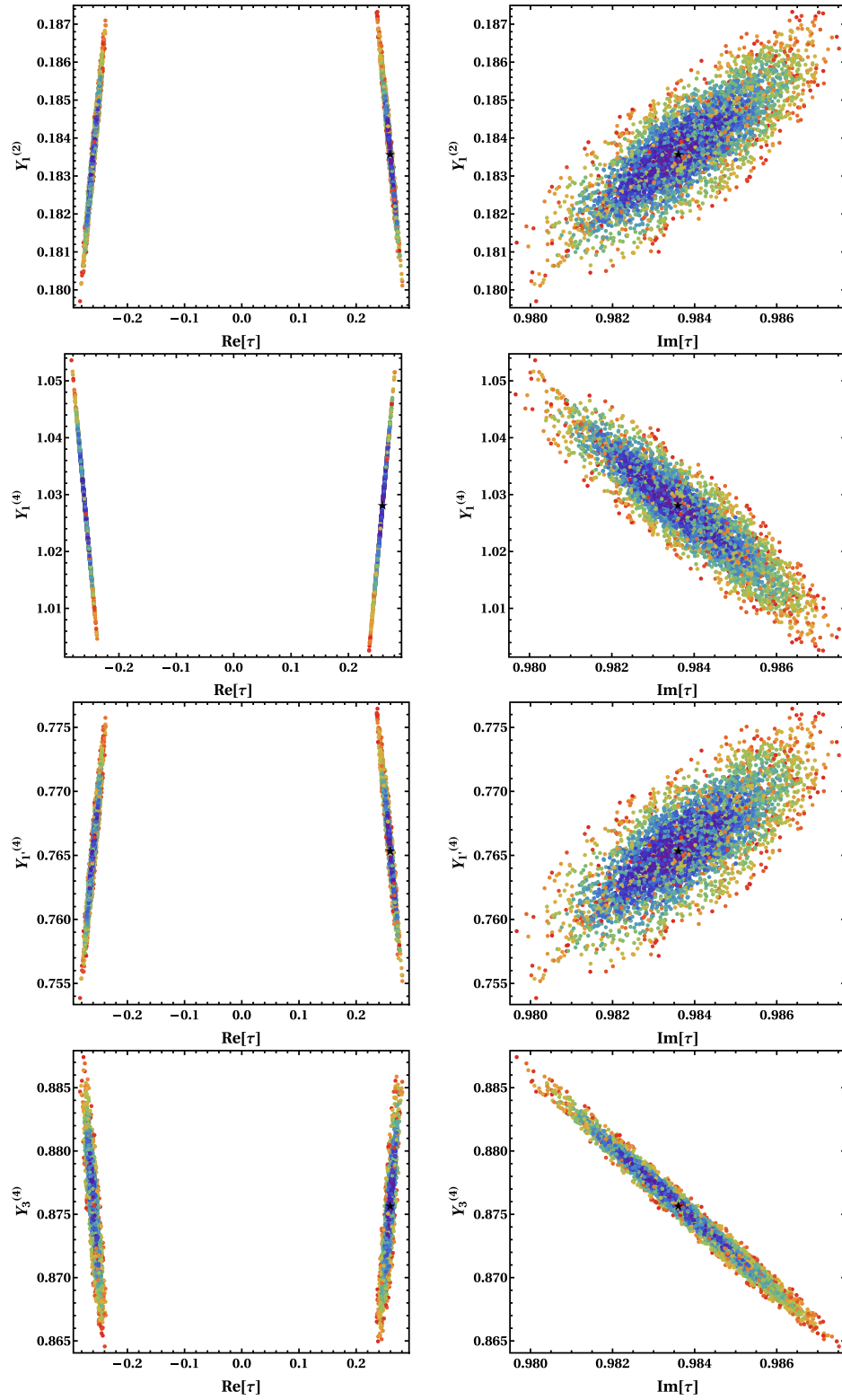


FIG. 2. Variation of Yukawa couplings with real and imaginary part of complex modulus τ . The color code is same as for Fig. 1.

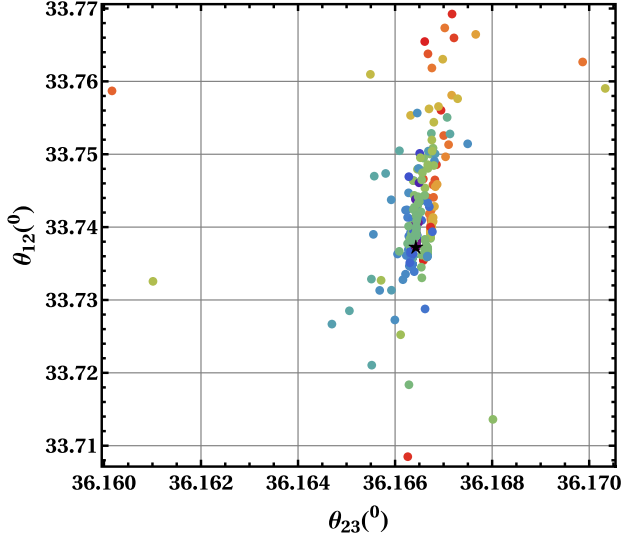


FIG. 3. Correlation between atmospheric (θ_{23}) and solar (θ_{12}) mixing angle for IH.

that diagonalizes the charged-lepton mass matrix. Throughout our numerical analysis, we fix $m_{\eta^+} = m_{\eta_1}$. The loop argument is defined as $y = \frac{M_f^2}{m_{\eta^+}^2}$, where M_f is the mass of the fermion propagating in the loop.

The loop function $\mathcal{G}(y)$ appearing in Eq. (4.2) is

$$\mathcal{G}(y) = \frac{1 - 6y + 3y^2 + 2y^3 - 6y^2 \ln y}{6(1 - y)^4}. \quad (4.3)$$

The value of branching ratio predicted at the best-fit point, corresponding to the minimum of the χ^2 , is presented in Table V. To illustrate how $\text{Br}(\mu \rightarrow e\gamma)$ varies with the mass of the fermion f and the real scalar mass m_{η_R} circulating in the loop, we fix all other parameters at their best-fit values

TABLE IV. Best-fit values for the neutrino oscillation parameters obtained for a minimum χ^2 value of 0.28 for NH.

$\theta_{12}(\circ)$	$\theta_{13}(\circ)$	$\theta_{23}(\circ)$	$\delta(\circ)$	$\Delta m_{21}^2(\text{eV}^2)$	$\Delta m_{31}^2(\text{eV}^2)$	$\text{Re}[\tau]$	$\text{Im}[\tau]$
34.00	8.56	43.15	289.04	7.5×10^{-5}	2.5×10^{-3}	0.26	0.98

TABLE V. Values of the model parameters obtained from the χ^2 analysis corresponding to the $\chi_{\min}^2 = 0.28$ for the NH.

Parameter	Value	Parameter	Value
κ_1 (GeV)	7.90×10^5	κ_2 (GeV)	6.69×10^8
γ	9.52×10^{-6}	κ_S (GeV)	5.22×10^8
β_L	1.00	m_{η_1} (GeV)	6.98×10^2
Δm_{η}^2 (GeV)	9.87×10^{-3}	$v\alpha_1$ (GeV)	4.10×10^{-3}
r_{21}	1.32×10^3	r_{31}	1.67×10^2
M_{N_1} (GeV)	1.21×10^6	M_{N_2} (GeV)	7.24×10^8
M_f (GeV)	3.09×10^7	$\text{Br}(\mu \rightarrow e\gamma)$	5.43×10^{-28}
m_e/m_μ	0.004836	m_μ/m_τ	0.059462

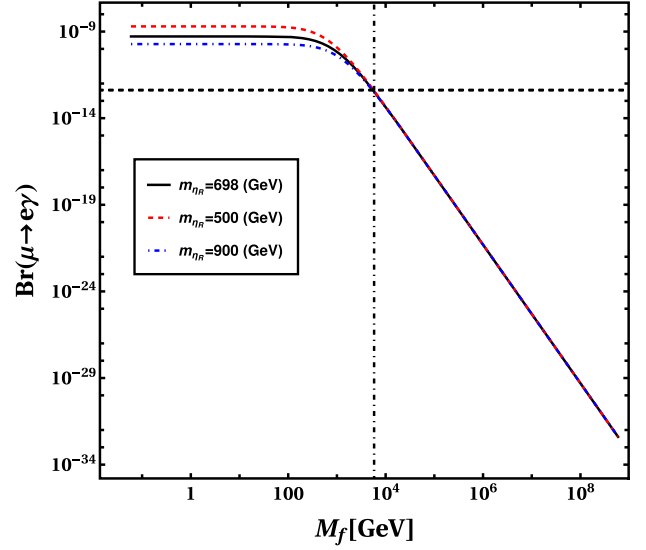


FIG. 4. Correlation between mass of dark-sector fermion (f) and LFV branching ratio $\text{Br}(\mu \rightarrow e\gamma)$ for three different values of scalar mass $m_{\eta_R} \sim m_{\eta_1}$, while keeping all other parameters fixed at their best-fit values given in Table V. The horizontal dashed line shows the current experimental upper bound on $\text{Br}(\mu \rightarrow e\gamma)$. The vertical dot-dashed line indicates the lower bound on the fermion mass M_f required to satisfy the LFV constraint.

and compute the resulting correlation for three representative values of m_{η_R} . The black curve corresponds to the best-fit point, while the other two curves represent the maximum and minimum values of m_{η_R} obtained from our data, as shown in Fig. 4. As seen in the figure, for relatively low fermion masses up to about 10^3 GeV, the suppression is weak and $\text{Br}(\mu \rightarrow e\gamma)$ remains nearly constant and above the current experimental limit. The loop contribution becomes increasingly suppressed, and the branching ratio falls below the experimental upper bound for $M_f \gtrsim 6 \times 10^3$ GeV, irrespective of the mass m_{η_R} . In addition to the radiative decay $\mu \rightarrow e\gamma$, other charged lepton flavor-violating processes can provide complementary constraints on the parameter space of the model. In particular, $\mu - e$ conversion in nuclei, currently constrained by the SINDRUM II experiment [85], can yield bounds comparable to those from $\mu \rightarrow e\gamma$, as both processes are induced by similar dipole operators. Future experiments are expected to significantly improve the sensitivity to charged-lepton flavor violation. The MEG II experiment aims to reach a sensitivity of $\mathcal{O}(6 \times 10^{-14})$ for $\mu \rightarrow e\gamma$ [86], while Mu3e will probe the $\mu \rightarrow eee$ decay down to $\mathcal{O}(2 \times 10^{-15})$ [87]. Searches for $\mu - e$ conversion will also be substantially improved by COMET, with expected sensitivities around $\mathcal{O}(10^{-15})$ and Mu2e [88,89], and by PRISM/PRIME, which aims to probe conversion rates below 10^{-18} [90]. In addition, Belle II will explore charged lepton flavor violation in the τ sector, with projected sensitivities in the range $\mathcal{O}(10^{-9}-10^{-10})$. These experiments will

provide powerful probes of the parameter space of models predicting charged lepton flavor violation.

V. LEPTOGENESIS

In Type-I seesaw framework, successful thermal leptogenesis typically requires a high scale $\mathcal{O}(10^9)$ GeV, known as the Davidson–Ibarra bound, to simultaneously account for neutrino masses and the observed baryon asymmetry [91]. At low leptogenesis scales, flavor effects become crucial and can enhance the generated asymmetry beyond the unflavored approximation [92–95]. To evade this bound and realize leptogenesis at a lower scale, additional CP -violating contributions beyond the minimal seesaw setting are necessary. These contributions often originate *via* new interactions or loop-induced effects involving extra states in nonholomorphic realization of the scoto-seesaw modular framework. To this end, we consider the participation of dark sector particles in the leptogenesis dynamics, assuming a BSM fermion mass hierarchy $M_{N_1} < M_f \ll M_{N_2}$, for which the inverse decays and scatterings involving N_2 are Boltzmann suppressed. Consequently, at temperatures $T \sim M_{N_2}$, the interactions of the heavier fermion N_2 remain in thermal equilibrium, leading to an efficient washout of any lepton asymmetry it may generate. Consequently, successful leptogenesis proceeds solely through the out-of-equilibrium decay of N_1 . For simplicity, we work in the single-flavor approximation, which captures the qualitative behavior of the asymmetry. It is worth noting that the model, also contains a viable candidate for dark matter. In particular, the CP -even neutral scalar originating from the inert doublet, η_1 , can play the role of dark matter provided it is the lightest among κ -odd particles. To ensure that the dark matter is predominantly the doublet component (rather than the singlet s), the mixing between the two CP -even states must remain small. This requirement, in Eq. (2.8), translates into the conditions

$$A < B, \quad |C| \ll |A - B|, \quad (5.1)$$

where

$$A = \mu_\eta^2 + \frac{1}{2}(\lambda_3 + \lambda_4 - \lambda_5)v^2, \quad B = \mu_s^2 \\ (\text{for } \lambda_8 \rightarrow 0), \quad C = \mu_1 v. \quad (5.2)$$

Under these conditions, the lighter mass eigenstate is dominantly composed of η_1 and thus serves as the dark matter candidate. In this context, the dark matter candidate η_1 falls within the class of weakly interacting massive particles (WIMPs). Due to its electroweak gauge interactions, η_1 remains in thermal equilibrium with the Standard Model plasma in the early Universe through annihilation processes such as $\eta_1 \eta_1 \rightarrow W^+ W^-, ZZ$, and $F\bar{F}$ (where F is SM fermion). As the temperature drops

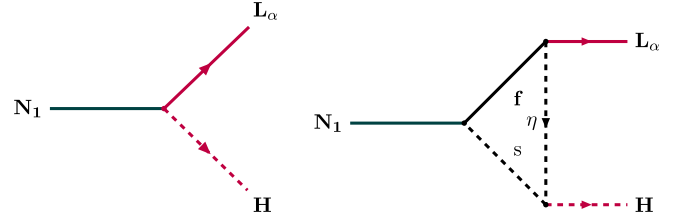


FIG. 5. Tree- and one-loop levels Feynman diagrams for dark-sector portal leptogenesis.

below the dark matter mass, these interactions become inefficient, and the relic abundance is set through the standard thermal freeze-out mechanism.

The CP asymmetry,

$$\epsilon_1 = \frac{\Gamma(N_1 \rightarrow LH) - \Gamma(N_1 \rightarrow \bar{L}H^\dagger)}{\Gamma(N_1 \rightarrow LH) + \Gamma(N_1 \rightarrow \bar{L}H^\dagger)} \\ = \frac{\Gamma(N_1 \rightarrow LH) - \Gamma(N_1 \rightarrow \bar{L}H^\dagger)}{\Gamma_N}, \quad (5.3)$$

generated in the decay of N_1 , shown in Fig. 5, arises from the interference between the tree-level and one-loop diagrams which, further, can be written as [96,97]

$$\epsilon_1 = \frac{1}{8\pi} \frac{\text{Im}[\mathcal{M}_{\text{tree}}^* \mathcal{M}_{\text{loop}}]}{|\mathcal{M}_{\text{tree}}|^2}, \quad (5.4) \\ \epsilon_1 = \frac{1}{8\pi} \frac{\text{Im}(y_D^\dagger y_f y_1 \mu_1)}{(y_D^\dagger y_D) M_{N_1} (1 - \vartheta + \omega) \sqrt{(1 - \vartheta + \omega)^2 - 4\omega}} I_{\text{loop}}, \quad (5.5)$$

where

$$I_{\text{loop}} = 1 + \vartheta - 2\sqrt{\varsigma} + (\kappa - \sqrt{\kappa}(1 - \vartheta + \omega) - \varrho + \omega) \\ \times \ln \left[\frac{\kappa - (1 - \sqrt{\varsigma})^2}{\kappa - \varsigma + \vartheta} \right]$$

is the loop function arises from vertex corrections involving the dark-sector fermion f , inert scalars, and $y_D = U_L^\dagger M_D/v$. The mass ratios are defined as

$$\kappa = \frac{M_f^2}{M_{N_1}^2}, \quad \varrho = \frac{m_\eta^2}{M_{N_1}^2}, \quad \vartheta = \frac{m_h^2}{M_{N_1}^2}, \\ \omega = \frac{m_l^2}{M_{N_1}^2}, \quad \varsigma = \frac{m_s^2}{M_{N_1}^2}. \quad (5.6)$$

All CP -violating phases entering ϵ_1 originate from the complex modular forms. With the modular symmetry constraints, the Yukawa couplings in our model are determined by modular forms

$$y_f = \beta_L Y_3^{(4)}, \quad y_1 = a Y_3^{(2)}. \quad (5.7)$$

In the limit of vanishing SM Higgs and lepton masses, the CP asymmetry simplifies to

$$\epsilon_1 = \frac{1}{8\pi} \frac{\text{Im}(y_D^\dagger y_f y_1 \mu_1)}{(y_D^\dagger y_D) M_{N_1}} \left(1 - 2\sqrt{\zeta} + (\kappa - \sqrt{\kappa} - \varrho) \ln \left[\frac{\kappa - (1 - \sqrt{\zeta})^2}{\kappa - \zeta} \right] \right). \quad (5.8)$$

The corresponding Boltzmann equations for comoving densities of N_1 and $B - L$ can be written as

$$\begin{aligned} \frac{dY_{N_1}}{dz} = & -D_N(Y_{N_1} - Y_{N_1}^{\text{eq}}) - \frac{s}{H(z)Z} (Y_{N_1}^2 - (Y_{N_1}^{\text{eq}})^2) \left[\langle \sigma v \rangle_{N_1 N_1 \rightarrow ss} + \langle \sigma v \rangle_{N_1 N_1 \rightarrow HH^\dagger} \right. \\ & + \langle \sigma v \rangle_{N_1 N_1 \rightarrow L_a \bar{L}_\beta} \left. \right] - \frac{s}{H(z)Z} (Y_{N_1} - Y_{N_1}^{\text{eq}}) \left[2Y_l^{\text{eq}} \langle \sigma v \rangle_{\bar{L} N_1 \rightarrow \bar{q} t} + 4Y_t^{\text{eq}} \langle \sigma v \rangle_{N_1 t \rightarrow \bar{L} q} \right. \\ & + 2Y_s^{\text{eq}} \langle \sigma v \rangle_{N_1 s \rightarrow \bar{L} \eta^\dagger} + 2Y_H^{\text{eq}} \langle \sigma v \rangle_{N_1 H \rightarrow L_a \nu_\mu} + 2Y_f^{\text{eq}} \langle \sigma v \rangle_{N_1 f \rightarrow H \eta^\dagger} \left. \right], \end{aligned} \quad (5.9)$$

$$\begin{aligned} \frac{dY_{B-L}}{dz} = & -\epsilon_1 D_N(Y_{N_1} - Y_{N_1}^{\text{eq}}) - W_{\text{ID}} Y_{B-L} - \frac{s}{H(z)Z} Y_{B-L} \left[2Y_H^{\text{eq}} \langle \sigma v \rangle_{LH^\dagger \rightarrow \bar{L} H} + Y_{N_1} \langle \sigma v \rangle_{\bar{L} N_1 \rightarrow \bar{q} t} \right. \\ & + 2Y_q^{\text{eq}} \langle \sigma v \rangle_{\bar{L} q \rightarrow N_1 t} + 2Y_l^{\text{eq}} \langle \sigma v \rangle_{LL \rightarrow H^\dagger H^\dagger} + Y_\eta^{\text{eq}} \langle \sigma v \rangle_{\bar{L} \eta^\dagger \rightarrow N_1 s} + Y_V^{\text{eq}} \langle \sigma v \rangle_{\bar{L} \nu_\mu \rightarrow H N_1} \left. \right]. \end{aligned} \quad (5.10)$$

Here, $Y_i = n_i/s$ denotes comoving density with n_i being number density of species ' i ' and $\rho_s = \frac{2\pi^2}{45} g_* T^3$ being entropy density of the Universe. $z = M_{N_1}/T$ and $H(z) = \sqrt{\frac{4\pi^3 g_*}{45} \frac{T^2}{M_{\text{Pl}}}}$ are the Hubble parameters at high temperatures. The decay term D_N is defined as

$$D_N = \frac{\langle \Gamma_N \rangle}{Hz} = K_N z \frac{\kappa_1(z)}{\kappa_2(z)}, \quad (5.11)$$

where $K_N = \Gamma_N/H(z=1)$, with $\kappa_i(z)$ being the modified Bessel function of the i^{th} kind. The washout due to inverse decay (ID) is

$$W_{\text{ID}} = \frac{1}{4} K_N z^3 \kappa_1(z). \quad (5.12)$$

The thermal averaged cross-section is defined as

$$\begin{aligned} \langle \sigma v \rangle_{ij \rightarrow kl} = & \frac{1}{8T m_i^2 m_j^2 \kappa_2(z_i) \kappa_2(z_j)} \\ & \times \int_{(m_i+m_j)^2}^{\infty} ds \frac{\lambda(\rho_s, m_i^2, m_j^2)}{\sqrt{\rho_s}} \kappa_1(\sqrt{\rho_s}/T) \sigma, \end{aligned} \quad (5.13)$$

with $z_i = m_i/T$ and $\lambda(\rho_s, m_i^2, m_j^2) = [\rho_s - (m_i + m_j)^2][\rho_s - (m_i - m_j)^2]$. The lepton asymmetry at the sphaleron decoupling epoch $T_{\text{sph}} \sim 130$ GeV gets converted into baryon asymmetry as

$$Y_B \simeq a_{\text{sph}} Y_{B-L} = \frac{8N_F + 4N_H}{22N_F + 13N_H} Y_{B-L} = 0.34 Y_{B-L}, \quad (5.14)$$

with $N_F = 3$, $N_H = 2$ being the fermion generations and the number of scalar doublets in our model, respectively.

The dependence of the CP asymmetry parameter ϵ_1 on the trilinear coupling μ_1 is shown in Fig. 6, where all other parameters are fixed according to Table IV. We observe that increasing μ_1 leads to a substantial enhancement of the asymmetry parameter ϵ_1 as $\epsilon_1 \propto \mu_1$ Eq. (5.8). Though increasing μ_1 enhances the CP asymmetry, for

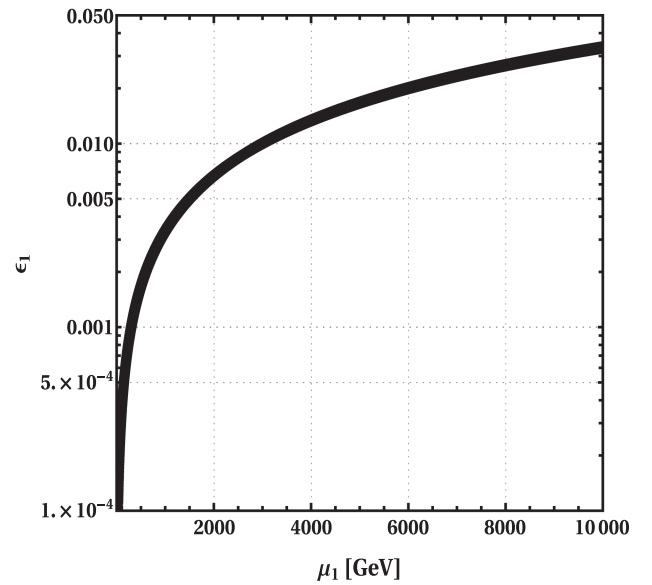


FIG. 6. Correlation between CP asymmetry (ϵ_1) as a function of dimensionful coupling μ_1 , keeping other parameters fixed using Table IV.

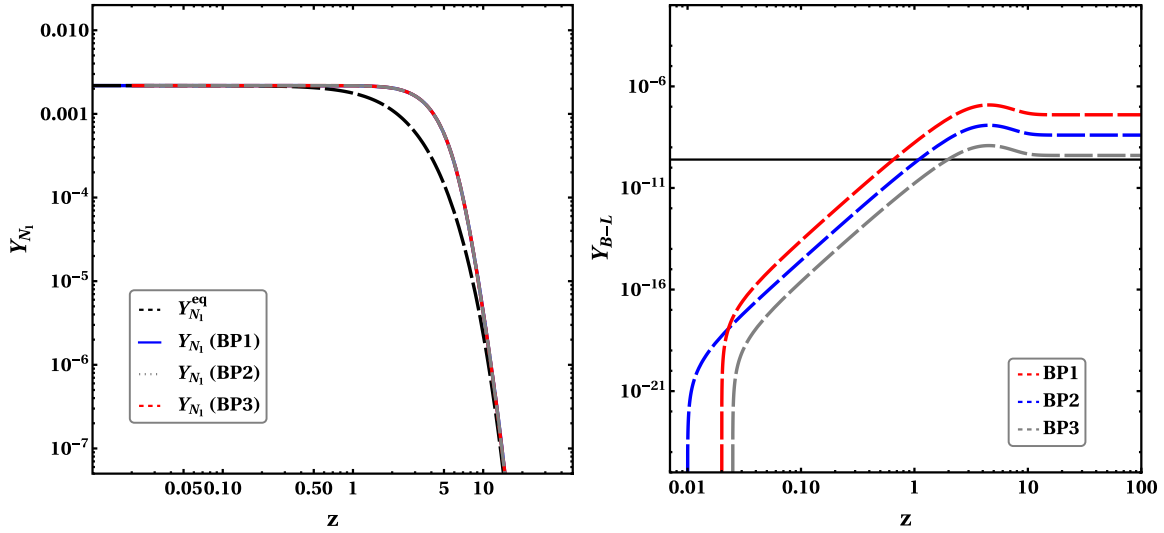


FIG. 7. *Left panel:* Evolution of comoving number density of N_1 for different benchmark points given in Table VI. *Right panel:* Evolution of comoving number density of $B - L$ asymmetry for different benchmark points. The solid horizontal line corresponds to the experimentally required value of the lepton asymmetry (using Eq. (5.14)).

sufficiently large values of μ_1 (e.g., $\mu_1 \sim 10^4$ GeV), the resulting ϵ_1 can become as large as $\mathcal{O}(1)$, which in turn leads to a $B - L$ asymmetry exceeding the value required to reproduce the observed baryon asymmetry. This indicates that μ_1 cannot be arbitrarily increased, and successful leptogenesis imposes an upper bound on the trilinear coupling μ_1 to prevent the overproduction of the baryon asymmetry. The numerical results for the comoving number density of N_1 (left panel) and the $B - L$ asymmetry (right panel) are presented in Fig. 7. With the evolution of the Universe or increase in z , the N_1 starts decaying to the SM particles (Fig. 5), thereby reducing their respective comoving number densities. Figure 7 (left panel) shows the evolution of the comoving number density of N_1 as a function of $z = M_{N_1}/T$ for the three benchmark points considered. At early times ($z \ll 1$), all solutions closely track the equilibrium distribution $Y_{N_1}^{eq}$, as expected for a relativistic species efficiently coupled to the thermal bath. As the temperature drops below M_{N_1} , the equilibrium abundance decreases, while the actual abundance departs from equilibrium and begins to freeze out.

For all benchmark points, the number density Y_{N_1} remains slightly above the equilibrium value during the freeze-out phase. This mild departure from equilibrium

reflects the fact that the decay and scattering interactions of N_1 are still moderately efficient around $z \sim \mathcal{O}(1-10)$, delaying the freeze-out relative to the purely equilibrium expectation. The three benchmark curves overlap throughout the evolution, indicating that the corresponding variations in the Yukawa couplings or dark-sector parameters have only a subleading impact on the dynamics of N_1 depletion. As a consequence, the CP -asymmetry-producing decays take place in a regime where N_1 is slightly overabundant compared to equilibrium, ensuring a nonzero departure from equilibrium.

The evolution of $B - L$ asymmetry for three benchmark points (Table VI) is shown in the right panel of Fig. 7. In all cases, the asymmetry begins to build up once the decays of N_1 start to dominate over inverse decay and scattering washout processes. Since the Dirac Yukawa coupling is fixed to a small value $y_D \sim 10^{-6}$ by neutrino oscillation data, the inverse-decay washout term, which scales as $\Gamma_{ID} \propto y_D^2$, remains weak throughout the relevant temperature range. Consequently, the washout rate Γ_{ID}/H becomes inefficient at relatively early times, allowing the generated CP asymmetry to accumulate smoothly.

A distinctive feature of this framework is that the dark sector responsible for CP violation is controlled by the trilinear parameter μ_1 while the effective washout involving

TABLE VI. Best-fit values for the neutrino oscillation parameters obtained from the χ^2 analysis correspond to a minimum χ^2 value of 0.28 for NH.

Benchmark point	M_{N_1} (GeV)	M_η (GeV)	Δm^2 (GeV)	μ_s (GeV)	μ_1 (GeV)	M_f (GeV)	$\text{Re}[\tau]$	$\text{Im}[\tau]$
BP1	1.21×10^6	6.98×10^2	9.87×10^{-3}	10^3	10^4	3.09×10^7	0.26	0.98
BP2	1.28×10^6	6.02×10^2	9.97×10^{-3}	10^3	5×10^3	3.40×10^7	0.25	0.98
BP3	1.23×10^6	6.23×10^2	9.96×10^{-3}	10^3	10^3	3.24×10^7	0.25	0.98

dark-sector fields depends on the fixed parameter $\mu_s = 1$ TeV. For simplicity, μ_s is kept the same for all benchmark points, such that the s -mediated washout channels remain moderately suppressed. Moreover, the kinematic hierarchy $M_{N_1} \sim 10^6$ GeV $< M_f \sim 10^7$ GeV forbids on-shell $N_1 N_1 \rightarrow ff$ scatterings, thereby eliminating an otherwise potentially strong source of thermal washout. The remaining scattering channels, such as $N_1 N_1 \rightarrow HH^\dagger$ and $N_1 N_1 \rightarrow L\bar{L}$, are further suppressed due to their dependence on the small Yukawa coupling y_D and the relatively small inert-scalar mass splitting Δm^2 .

The benchmark points in Table VI correspond to progressively smaller values of μ_1 , leading to different magnitudes of the loop-induced CP asymmetry. As a result, BP1 produces the largest final asymmetry, followed by BP2 and BP3, as reflected in Fig. 7 (right panel). In all benchmarks, the $B-L$ asymmetry saturates around $z \sim 10$, corresponding to the epoch at which N_1 fully departs from equilibrium and washout effects become negligible. The resulting asymmetries are compatible with the value required to reproduce the observed baryon

asymmetry of the Universe after sphaleron conversion, demonstrating that the model can successfully realize leptogenesis without invoking resonant enhancement. Further, we revisit leptogenesis in our framework by incorporating flavor effects, which are essential at the relevant scale. For the benchmark points presented in Table VI, the mass of the lightest righthanded neutrino is $M_{N_1} \sim 10^6$ GeV, placing the scenario well within the fully flavored regime ($T \lesssim 10^9$ GeV). In this temperature range, charged lepton Yukawa interactions are in thermal equilibrium, and the individual lepton flavors $\alpha = e, \mu, \tau$ must be treated separately. The CP asymmetry is given by

$$\epsilon_1^\alpha = \frac{1}{8\pi (y_D^\dagger y_D)_{11}} \frac{\text{Im}[(y_D^\dagger)_{1\alpha} (\sum_\beta (y_f)_\beta (y_D)_{\beta 1}) y_1 \mu_1]}{M_{N_1} (1 - \vartheta + \omega) \sqrt{(1 - \vartheta + \omega)^2 - 4\omega}} I_{\text{loop}} \quad (5.15)$$

with $\alpha, \beta = e, \mu, \tau$. The Boltzmann equation describing the generation of $B-L$ asymmetry for each lepton flavor is given by

$$\begin{aligned} \frac{dY_{B-L_\alpha}}{dz} = & -\epsilon_1^\alpha D_N (Y_{N_1} - Y_{N_1}^{\text{eq}}) - W_{\text{ID}}^\alpha Y_{B-L_\alpha} - \frac{s}{H(z)z} Y_{B-L_\alpha} \left[2Y_H^{\text{eq}} \langle \sigma v \rangle_{L_\alpha H^\dagger \rightarrow \bar{L}_\beta H} + Y_{N_1} \langle \sigma v \rangle_{\bar{L}_\alpha N_1 \rightarrow \bar{q} t} \right. \\ & \left. + 2Y_q^{\text{eq}} \langle \sigma v \rangle_{\bar{L}_\alpha q \rightarrow N_1 t} + 2Y_l^{\text{eq}} \langle \sigma v \rangle_{L_\alpha L_\beta \rightarrow H^\dagger H^\dagger} + Y_\eta^{\text{eq}} \langle \sigma v \rangle_{\bar{L}_\alpha \eta^\dagger \rightarrow N_1 s} + Y_V^{\text{eq}} \langle \sigma v \rangle_{\bar{L}_\alpha V_\mu \rightarrow H N_1} \right]. \end{aligned} \quad (5.16)$$

For BP3 in Table VI, the Dirac Yukawa matrix $y_D = U_L^\dagger M_D / v$ and the dark-sector Yukawa vector y_f (corresponding to N_1 coupling only in Eq. (2.12) are numerically given by

$$y_D = \begin{pmatrix} -1.16 \times 10^{-6} - 1.13 \times 10^{-6}i & 2.43 \times 10^{-6} - 2.81 \times 10^{-6}i \\ -2.60 \times 10^{-6} + 3.27 \times 10^{-6}i & 4.68 \times 10^{-6} - 7.26 \times 10^{-6}i \\ -1.22 \times 10^{-5} - 7.01 \times 10^{-6}i & -1.02 \times 10^{-5} - 5.03 \times 10^{-6}i \end{pmatrix}, \quad y_f = \begin{pmatrix} 1.015 - 0.173i \\ 0.385 + 0.788i \\ 0.589 + 0.491i \end{pmatrix}.$$

The CP asymmetry is computed separately for each flavor where all CP -violating phases originate from the complex modulus τ through the modular Yukawa couplings. Using the above numerical inputs, we obtain the flavor-dependent CP asymmetries:

$$\epsilon_1^e \simeq -2.3 \times 10^{-9}, \quad \epsilon_1^\mu \simeq -1.0 \times 10^{-7}, \quad \epsilon_1^\tau \simeq -8.5 \times 10^{-7}. \quad (5.17)$$

A pronounced hierarchy is observed $|\epsilon_1^\tau| > |\epsilon_1^\mu| \gg |\epsilon_1^e|$, indicating that the CP asymmetry is dominantly generated in the τ -flavor. This behavior can be directly understood from the structure of y_D , where the coupling of N_1 to the τ -flavor is significantly larger than to the electron and muon flavors. In addition, the complex structure of y_f (from dark-sector), arising from modular forms, introduces nontrivial CP phases that enhance the asymmetry without being aligned with y_D , thereby preventing complete washout.

Since all flavor asymmetries carry the same sign, their contributions add constructively to the total lepton asymmetry. The total CP asymmetry is therefore given by

$$\epsilon_1 = \sum_\alpha \epsilon_1^\alpha \simeq -9.6 \times 10^{-7},$$

which is clearly dominated by the τ -flavor contribution. These results demonstrate that leptogenesis in this framework is effectively τ -dominated as shown in the Fig. 8, with subleading contributions from the muon flavor and negligible contributions from the electron flavor. This feature arises from the interplay between the hierarchical structure of the Dirac Yukawa matrix and the complex modular Yukawa couplings in the dark sector, and it plays a crucial role in determining the final baryon asymmetry after solving the flavored Boltzmann equations. Further, the lightest dark sector particle, *i.e.*, real part of inert scalar, is our viable dark matter particle (WIMP type) in the model which has been widely explored in the literature [97,98].

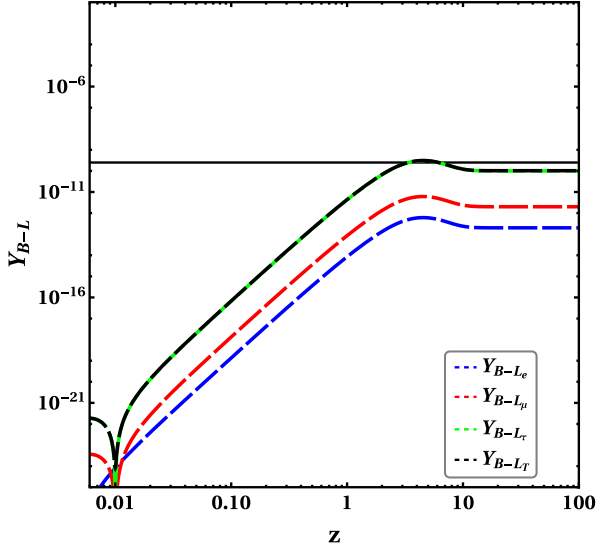


FIG. 8. Boltzmann evolution of the flavor asymmetries Y_{B-L_e} , Y_{B-L_μ} , and Y_{B-L_τ} as functions of $z = M_{N_1}/T$ for the BP3 in Table VI. The total asymmetry $Y_{B-L_T} = \sum_\alpha Y_{B-L_\alpha}$ is shown in black. The horizontal line represents the observed lepton asymmetry of the Universe.

A distinctive prediction of our construction arises from the inert scalar sector, in which the charged scalar $\eta^\pm$ is nearly degenerate with the dark matter candidate η_1 . For the benchmark point, with $\delta m = m_{\eta^\pm} - m_{\eta_1} \sim 0.1$ GeV, the decay $\eta^\pm \rightarrow \eta_1 + W^{\pm*}$ proceeds only through an off-shell W boson and is therefore phase-space suppressed. Since $m_{\eta^\pm} - m_{\eta_1} \ll m_W$, the resulting decay width scales as $\Gamma \propto (\delta m)^5$, leading to a macroscopic decay length of $c\tau_{\eta^\pm} \sim \mathcal{O}(1-10)$ cm. Consequently, $\eta^\pm$ behaves as a long-lived charged particle inside ATLAS/CMS experiments. The ATLAS Collaboration has recently performed a dedicated search for long-lived charged states using the full 140 fb^{-1} dataset of pp collisions at $\sqrt{s} = 13$ TeV. The analysis exploits signatures characterized by high specific ionization energy loss and time-of-flight measurements [99]. Interpreting their results⁴ in the context of stau-to-gravitino decays, ATLAS excludes masses up to 560 GeV for such long-lived charged particles. Observation of such a long-lived charged particle signal would provide strong support for the inert scalar sector of the model and offer an important experimental handle on the dark-sector-assisted origin of neutrino masses and leptogenesis.

VI. CONCLUSION

This work investigates the neutrino phenomenology and dark-sector-assisted leptogenesis in a scoto-seesaw framework constructed under a nonholomorphic A_4 modular flavor symmetry. To ensure a clean separation between

the fields responsible for the tree-level (seesaw) and loop-induced (scotogenic) contributions to the neutrino mass matrix, we assign even modular weights to the fields in the seesaw sector and odd modular weights to those in the radiative sector. This modular-weight arrangement forbids undesired mixing terms and enforces a controlled hierarchical structure between these two neutrino-mass-generating sectors. By restricting the complex modulus τ to lie within its fundamental domain and performing a systematic numerical scan over the remaining model parameters, we identify regions of parameter space that are fully consistent with the latest neutrino global-fit data. We find that the data is satisfied for $\text{Re}[\tau] \sim \pm 0.26$ and $\text{Im}[\tau] \sim 0.97$. Also, the CP -violating phase δ is lying at $\sim 70^\circ$ and $\sim 290^\circ$. The model is consistent only with a normal hierarchical neutrino mass pattern. In this setup, the complex modulus τ acts as the unique source of CP violation, since all Yukawa couplings and scalar-sector parameters are taken to be real. Consequently, the predicted low-energy neutrino mixing phases and the high-energy CP -violating interactions relevant for leptogenesis are entirely determined by the modular origin of the flavor sector. Within this construction, successful leptogenesis emerges from the CP -violating out-of-equilibrium decays of the lightest righthanded neutrino N_1 into SM leptons and the Higgs doublet at $M_{N_1} \sim 10^6$ GeV. The loop-induced CP asymmetry is generated through one-loop diagrams mediated by the dark-sector fields intrinsic to the scotogenic mechanism, notably without relying on any mass degeneracy or resonance enhancement between the heavy neutrino states. This demonstrates that the interplay of modular flavor structures and radiative dynamics can naturally accommodate the observed baryon asymmetry of the Universe while simultaneously reproducing the correct neutrino mass and mixing pattern. Finally, we have also incorporated a comment on the unique collider signature of $\eta^\pm$, which may behave as a long-lived charged particle at ATLAS/CMS experiments.

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DATA AVAILABILITY

The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

⁴The same interpretation has been inferred in Ref. [100].

APPENDIX: NONHOLOMORPHIC MODULAR FLAVOR SYMMETRY

The conventional modular flavor symmetry employs holomorphic modular forms as Yukawa couplings, a structure naturally emerging in supersymmetric settings. However, modular invariance alone does not necessitate holomorphicity. By relaxing this condition, we enter the broader framework of *non-holomorphic modular symmetry*, where Yukawa couplings are polyharmonic Maass forms—automorphic functions satisfying the harmonic condition

$$\Delta_k Y(\tau) = 0, \quad (\text{A1})$$

with $\Delta_k = -4y^2 \partial_\tau \partial_{\bar{\tau}} + 2iky \partial_{\bar{\tau}}$ being the weight- k hyperbolic Laplacian, rather than requiring holomorphicity. These functions transform covariantly under the modular group

$$Y(\gamma\tau) = (c\tau + d)^k \rho(\gamma) Y(\tau), \quad \gamma \in \text{SL}(2, \mathbb{Z}), \quad (\text{A2})$$

where ρ denotes a representation of the finite modular group $\Gamma'_N = \text{SL}(2, \mathbb{Z})/\Gamma(N)$ (or Γ_N for even weight). Their Fourier expansion reveals the characteristic nonholomorphic structure

$$Y(\tau) = \sum_{n \geq 0} c^+(n) q^n + c^-(0) y^{1-k} + \sum_{n < 0} c^-(n) \Gamma(1-k, -4\pi n y) q^n, \quad (\text{A3})$$

where the terms proportional to $c^-(0)$ and $c^-(n)$ explicitly break holomorphicity.

A crucial connection to ordinary modular forms exists via differential operators: the Maass raising operator D^{1-k} and the ξ -operator $\xi_k = 2iy^k \partial_{\bar{\tau}}$ map weight- k polyharmonic

Maass forms to weight- $(2-k)$ modular forms in conjugate and original representations, respectively. This enables explicit construction of polyharmonic forms from known modular forms through the relations

$$\xi_k Y_{\mathbf{r}}^{(k)}(\tau) = \alpha \Omega Y_{\mathbf{r}'}^{(2-k)}(\tau), \quad (\text{A4})$$

$$D^{1-k} Y_{\mathbf{r}}^{(k)}(\tau) = \beta Y_{\mathbf{r}}^{(2-k)}(\tau), \quad (\text{A5})$$

where Ω relates conjugate representations and α, β are normalization constants.

In flavor model building, matter fields ψ_i are assigned modular weights $-k_i$ and transform in irreducible representations ρ_i of Γ'_N . Yukawa couplings are now polyharmonic Maass forms $Y_{\mathbf{r}}^{(k_Y)}(\tau)$ of weight k_Y and representation $\mathbf{r}$. Modular invariance imposes the conditions

$$k_Y = k_{\psi^c} + k_{\psi} + k_H, \quad (\text{A6})$$

$$\rho_Y \otimes \rho_{\psi^c} \otimes \rho_{\psi} \otimes \rho_H \ni \mathbf{1}. \quad (\text{A7})$$

This nonholomorphic extension significantly expands the Yukawa coupling space while preserving modular invariance. Negative modular weights become admissible, and nonholomorphic terms introduce new parameters that enhance phenomenological flexibility. The vacuum expectation value of τ simultaneously breaks modular and CP symmetries, offering a unified origin for flavor structures and CP violation. As an illustrative example, at level $N = 3$ (finite group A_4), polyharmonic Maass forms exist for weights $k = -4, -2, 0, 2, 4, 6, \dots$, organizing into A_4 multiplets, thereby enabling novel texture constructions beyond the holomorphic paradigm.

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