

Dark matter-driven low-scale leptogenesis via neutrino portal

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We propose a novel framework for low-scale leptogenesis within an extension of the Standard Model that includes three $SU(2)$ singlet right-handed neutrinos, a singlet vectorlike neutral fermion, and a real scalar field. In this setup, the CP asymmetry arises through a rich interplay of mechanisms, including two-body decays of the lightest right-handed neutrino into leptons and Higgs boson or into dark-sector particles, as well as multiple $2 \rightarrow 2$ scattering processes involving visible and dark sector particles. Crucially, the CP -violating phases originate not only from conventional vertex and self-energy corrections of the decay of the right-handed neutrinos but also from novel interference effects in the scattering processes mediated by the dark sector, which significantly enrich the sources of asymmetry. A distinctive feature of our framework is the direct connection between the dark sector and leptogenesis, providing a unified explanation for both the matter–antimatter asymmetry and Dark Matter abundance. Moreover, the viable parameter range including the mass spectrum is within the range that Large Hadron Collider can explore.

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I. INTRODUCTION

The baryon asymmetry of the Universe (BAU) is one of the most puzzling problems of modern cosmology. While the Standard Model (SM) satisfies the qualitative requirements set through Sakharov conditions [1], the experimentally measured Higgs mass renders the electroweak phase transition a smooth crossover [2–5], and the CP violation present in the SM is far too small to account for the amount of the BAU adequately. There are many beyond the Standard Model (BSM) frameworks that have been proposed that invoke a first-order phase transition associated with electroweak symmetry breaking [6–8] to fulfill the required out-of-equilibrium condition. However, within the SM gauge symmetry, an alternative approach is considered. Rather than directly producing a baryon asymmetry, this approach first generates an asymmetry in the leptonic sector through lepton number (L)-violating decays. Subsequently, this lepton asymmetry can be converted into the observed baryon asymmetry through $(B + L)$ -violating electroweak (EW) sphaleron processes [6]. This widely studied approach, known as *leptogenesis* [9], was proposed

by Fukugita and Yanagida. A comprehensive review of leptogenesis can be availed at Ref. [10]. An appealing aspect of the standard leptogenesis is that the CP -violating, out-of-equilibrium decays responsible for generating lepton asymmetry can originate from the same heavy neutrinos involved in the seesaw mechanism [11–14], which explains the smallness of neutrino masses [15].

While matter–antimatter asymmetry remains one of the long-standing mysteries of the Universe, another intriguing puzzle lies in the composition of matter itself in the Universe. Observations from the spectral analysis of the CMB radiation [16], and many other independent observations [12–14,17] have shown that about 85% of the matter content of the Universe is nonbaryonic and non-luminous dark matter (DM). While dark matter plays a central role in cosmic structure formation through its gravitational interactions, it is usually expected to couple only very weakly to the Standard Model particles, as stronger interactions would have resulted in observable signals. Although neutrinos share several characteristics required for DM particles, their tiny mass and a large free-streaming length make them insufficient to fully account for the dark matter content of the Universe [18]. Since none of the other SM particles possess the required properties of a viable DM candidate, a plethora of BSM frameworks have been proposed to account for its existence. Among the various DM paradigms, the weakly interacting massive particle (WIMP) scenario is one of the most extensively studied frameworks [19]. In this picture, a DM particle with a mass around the electroweak scale and weak-scale interactions is thermally produced in the early Universe.

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As the Universe expands and cools, the interaction rate drops below the Hubble expansion rate, causing the dark matter to decouple from the thermal bath. The resulting frozen-out abundance naturally accounts for the observed dark matter relic density. However, conventional WIMP scenarios are being increasingly challenged by the difficulty of simultaneously satisfying the observed dark matter relic density and the stringent bounds from direct detection experiments [20]. In direct detection searches, WIMPs are expected to elastically scatter off atomic nuclei, producing measurable nuclear recoil signals. In standard WIMP frameworks, the same interactions that determine the thermal relic abundance also control the DM-nucleus scattering cross section. Consequently, the null results from direct detection experiments impose severe constraints on the parameter space consistent with the observed relic density, significantly limiting the viability of such models. Motivated by these tensions, alternative paradigms have been proposed. One attractive scenario among these is the so-called feebly interacting massive particles (FIMPs). In a FIMP scenario, the dark matter interacts so weakly with the Standard Model that it is produced nonthermally in the early Universe and naturally evades direct detection bounds. This production mechanism, known as *freeze-in*, leads to a relic density that asymptotically approaches the observed DM relic density [21–23]. Extensive analyses of the freeze-in mechanism have been explored in Refs. [23–28], and specifically within the context of baryogenesis to explain the BAU in Ref. [29].

In this work, we propose a unified framework that simultaneously addresses the origin of the baryon asymmetry of the Universe, the nature of dark matter, and the generation of neutrino masses. The model constitutes a minimal extension of the SM particle content, introducing a singlet vectorlike fermion ψ , a neutral real scalar singlet ϕ , and three generations of heavy right-handed neutrinos $N_{1,2,3}$. The dark sector is stabilized by a discrete Z_2 symmetry, under which ψ and ϕ are odd, while the right-handed neutrinos and all the SM fields are even. In conventional type-I seesaw scenarios with hierarchical right-handed neutrinos, successful thermal leptogenesis requires the lightest heavy neutrino mass to satisfy the Davidson-Ibarra bound, $m_{N_1} \geq 10^9$ GeV [30]. Various models allowing low-scale leptogenesis with hierarchical neutrinos have been explored in the literature [31–44], while scenarios involving quasidegenerate right-handed neutrinos, including resonant leptogenesis with enhanced CP asymmetry are studied in Refs. [45–50]. In contrast to these approaches, in the present setup, we demonstrate that a sufficient leptogenesis can be achieved with the lightest heavy neutrino mass being at the TeV scale without invoking resonant enhancement. This becomes possible through the neutrino portal interactions linking the right-handed neutrinos to the dark sector fields. Some selected literature in neutrino portal dark matter studies includes

Refs. [51–55]. In our setup, these interactions open additional decay channels, loop contributions, and scattering processes, leading to sizable dark-sector-induced CP violation. Moreover, the new processes involving leptons and antileptons in both initial and final states modify the reaction rates entering the Boltzmann equations, thereby significantly impacting the evolution of the lepton asymmetry. A key feature of this scenario is the intimate interplay between dark matter dynamics and leptogenesis: the same couplings that govern the dark sector phenomenology also play a central role in enhancing the CP asymmetry required for successful BAU generation. We show that the region of parameter space consistent with the observed dark matter relic density (DMRD) and other phenomenological constraints naturally accommodates the required low-scale leptogenesis. Furthermore, the presence of CP -violating interactions in the dark sector can also lead to an asymmetry within the dark sector itself, opening up additional avenues for phenomenological exploration.

The structure of the article is as follows. In Sec. II, we present the theoretical framework of the model, including the particle content, symmetries, and relevant interactions. Section III is devoted to the dark matter phenomenology, where the WIMP realizations are discussed and the parameter space consistent with the observed relic abundance is identified. In Sec. IV, we analyze the mechanism of leptogenesis, highlighting the novel sources of CP violation arising from neutrino-portal interactions with the dark sector, including decay and scattering contributions. The numerical solutions of the coupled Boltzmann equations and a detailed discussion of the generated asymmetries are presented in Sec. V. Finally, we summarize our results and outline future directions in Sec. VI.

II. THE MODEL

The model framework under consideration extends the SM by introducing three right-handed gauge-singlet neutrinos, denoted by N_i , a singlet vectorlike fermion ψ , and a singlet real scalar field ϕ . Both ϕ and ψ are assigned to be odd under a discrete Z_2 symmetry, while all other fields remain even. This discrete symmetry is introduced to ensure the stability of the DM candidate. The gauge symmetry of the model is given by the SM gauge group supplemented with an additional discrete Z_2 symmetry. Depending on the mass hierarchy between ϕ and ψ , the lighter of the two becomes stable and can serve as a viable DM candidate. In addition, we assign a generalized lepton number [global $U(1)_{L'}$ symmetry] to the fermion ψ , the heavy neutrinos N_i , the SM lepton doublet L and the SM lepton singlet e_R , while the scalar ϕ and rest of the Standard Model particles are taken to be lepton-number neutral. This symmetry governs the generation of asymmetries by distinguishing ψ from its antiparticle $\bar{\psi}$. The relevant particle content of this extended model, along with their transformation properties under the SM gauge symmetry,

$U(1)_{L'}$ symmetry, and the Z_2 symmetry, is summarized in Table I.

The Lagrangian for this model framework is given by

$$\begin{aligned}\mathcal{L}_m = & \mathcal{L}_{\text{SM}} + (\partial^\mu \phi)^\dagger (\partial_\mu \phi) + \bar{\psi} i \gamma^\mu \partial_\mu \psi + \bar{N}_i i \gamma^\mu \partial_\mu N_i \\ & - m_\psi \bar{\psi} \psi - m_{N_{ij}} \bar{N}_i N_j \\ & - (y_i \bar{\psi} \phi N_i + Y_{ij}^N \bar{L}_i \tilde{H} N_j + \text{H.c.}) - \lambda' \phi^2 H^\dagger H \\ & - \frac{1}{2} m^2 \phi^2 - \zeta \phi^4,\end{aligned}\quad (1)$$

where $L_i = (\nu_{e_i}, e_i)_L^T$ are the left-handed lepton doublets with $e_i = e, \mu, \tau$, and H denotes the SM Higgs doublet field. After electroweak symmetry breaking (EWSB), the Higgs field acquires a vacuum expectation value (VEV),

$$\langle H \rangle = \left(0, \frac{v}{\sqrt{2}} \right)^T, \quad (2)$$

while the scalar field ϕ does not develop a VEV, ensuring that the Z_2 symmetry remains an exact symmetry. The Higgs VEV contributes to the mass of the scalar field ϕ , yielding its physical mass square

$$m_\phi^2 = m^2 + \lambda' v^2. \quad (3)$$

The exact Z_2 symmetry prevents any mixing between ϕ and the SM Higgs boson. The only interaction connecting the dark scalar to the SM Higgs sector arises from the quartic term $\lambda' \phi^2 H^\dagger H$, which can open an additional Higgs decay channel $H \rightarrow \phi\phi$ for sufficiently light ϕ . However, the coupling λ' can be chosen to be very small, independently of the scalar mass, thereby evading constraints from the invisible Higgs decay searches.

The mass of the vectorlike fermion ψ is not generated through symmetry breaking and is entirely determined by the parameter m_ψ . Although ψ does not couple directly to the Standard Model fields, it communicates with the visible sector through interactions mediated by the heavy right-handed neutrinos N_i and the scalar field ϕ . As we shall demonstrate in the subsequent Sections, these neutrino-portal interactions play a crucial role in linking the dynamics of the dark sector to the generation of the

matter-antimatter asymmetry of the Universe. This interplay between dark matter and leptogenesis constitutes the central focus of the present work.

The model addresses the origin of neutrino masses via the standard type-I seesaw mechanism [56], induced by the Yukawa interaction governed by the couplings Y_{ij}^N . For simplicity, we assume the Majorana mass matrix of the right-handed neutrinos to be diagonal, $m_{N_{ij}} = m_{N_i} \delta_{ij}$, so that the physical masses of the heavy neutrinos are given by m_{N_i} . The seesaw Yukawa couplings Y_{ij}^N are fixed using the Casas-Ibarra parametrization [57].

The right-handed neutrinos N_i , being Majorana fermions, naturally provide the necessary ingredients for leptogenesis through their out-of-equilibrium decays. In the standard type-I seesaw framework, leptogenesis arises solely from the Yukawa interactions among N_i , the SM leptons, and the Higgs doublet. In the present model, this conventional picture is extended by the additional interaction term $\bar{\psi} N_i \phi$, which directly connects the right-handed neutrinos to the dark sector. The associated Yukawa couplings y_i , being unconstrained by low-energy experiments, introduce new sources of CP violation in addition to those present in the standard leptogenesis scenario. As a result, the decay of N_i can generate asymmetries simultaneously in the visible and dark sectors. Moreover, this interaction opens up a variety of scattering processes involving leptons, dark-sector particles, and heavy neutrinos, which can act both as sources of lepton asymmetry and as channels for transferring asymmetry between the two sectors. In addition, when kinematically allowed, the decay of the scalar field ϕ can itself generate a dark-sector asymmetry, which may subsequently influence the evolution of the lepton asymmetry through these scattering processes. The combined effect of these decay and scattering channels establishes a nontrivial interplay between dark matter dynamics and leptogenesis, which constitutes the central focus of this work.

A. Collider constraints

Major collider constraints on the present framework arise from searches for heavy neutral leptons (HNLs), corresponding to the right-handed neutrinos N_i , at the Large Hadron Collider (LHC). At the LHC, HNLs can be produced through their mixing with Standard Model leptons, typically via processes such as $pp \rightarrow W^* \rightarrow \ell N$, followed by N decays into leptons and electroweak gauge bosons. Both of the ATLAS and the CMS collaborations have carried out extensive searches for HNLs in prompt and displaced decay channels [58–62]. These searches are sensitive to the active-sterile mixing along the masses of HNLs. For HNL masses between 100 GeV to 20 TeV, the mixing is limited to $< (0.1-1)$ [60]. We shall restrict our study to heavy neutrinos with masses $m_{N_i} \gtrsim 2$ TeV and active-sterile mixing angles smaller than $|U_{\ell N}|^2 \lesssim 1.2 \times 10^{-5}$, well below the current

TABLE I. Relevant particle content of the model. All the SM particles are Z_2 even. The SM leptons have $U(1)_{L'}$ charge equal to $+1$, while all the other SM particles have zero $U(1)_{L'}$ charge.

Particles	Spin	$SU(3)_c$	$SU(2)_L$	$U(1)_Y$	$U(1)_{L'}$	Z_2
$N_{1,2,3}$	$\frac{1}{2}$	1	1	0	1	1
ψ	$\frac{1}{2}$	1	1	0	1	-1
ϕ	0	1	1	0	0	-1
L	$\frac{1}{2}$	1	2	$-\frac{1}{2}$	1	1
e_R	$\frac{1}{2}$	1	1	-1	1	1

sensitivity of LHC searches. Within this regime, collider constraints are naturally evaded, allowing us to focus on the cosmological implications of the model.

We first analyze the framework to determine the regions of parameter space consistent with the observed dark matter relic density. Subsequently, we investigate the viability of leptogenesis within this restricted parameter space, with particular emphasis on the novel contributions arising from neutrino-portal interactions. In this setup, the interplay between the dynamics of the dark sector and lepton number generation plays a central role, enabling successful low-scale leptogenesis while simultaneously satisfying constraints on dark matter.

III. DARK MATTER

Both ϕ and ψ are odd under an exact discrete Z_2 symmetry, which guarantees their stability and renders them viable DM candidates. When kinematically allowed, the heavier of the two particles can decay into the lighter one, effectively yielding a single-component DM scenario. On the other hand, if such decays are kinematically forbidden or sufficiently suppressed, a two-component DM framework arises naturally, with both ϕ and ψ contributing to the total dark matter relic abundance.

As discussed in Sec. II, the quartic coupling between ϕ and the standard Higgs field, λ' is taken to be very small, so as to avoid collider constraints. Thus, the only interaction ϕ and ψ have with the visible sector is through the Yukawa interaction $\bar{\psi}\phi N_i$, controlled by the respective couplings, y_i . For this coupling being in the range of $\mathcal{O}(y_i) < 10^{-10}$, we have FIMP dark matter, which follows a nonthermal genesis of the dark matter relic abundance. We have explored this case and found that for a large range of dark matter masses in the GeV to TeV range, the compatible parameters lead to $y_i \sim 10^{-12}$. However, such small couplings do not lead to any significant effect in the leptogenesis. Therefore, we shall not discuss this scenario any further in this work.

On the other hand, the WIMP scenario is of particular interest in this framework, as the presence of sizable couplings can lead to a significant impact of dark matter interactions on leptogenesis. The viable dark matter candidates in this case are dictated by the mass hierarchy among the exotic particles. Accordingly, we consider the different possibilities separately. In the following, we first demonstrate why a single-component dark matter scenario is not viable within this framework. We then systematically analyze the two-component scenarios, identify the corresponding viable regions of parameter space, and discuss their phenomenological consequences.

A. Single-component dark matter

Depending on the mass hierarchy, we may have either a scalar DM scenario or a scotino (fermionic) DM scenario.

In the first case of scalar DM scenario, the dark scalar, ϕ is lighter than the new fermion, ψ . The scalar DM candidate, ϕ will remain in thermal equilibrium with the SM particles in the early Universe through its Higgs portal interactions. This makes it a potential WIMP candidate. The annihilation of DM will be controlled by the quartic coupling, $\lambda'\phi^2 H^\dagger H$. However, the same coupling also decides the DM-nucleus scattering in the direct detection experiments. Therefore, it turned out to be quite challenging to deplete the DM sufficiently to avoid an overabundance in DMRD while satisfying the measurements at the direct detection experiments. This is a typical issue with scalar DM and Higgs portal interactions, and the present model framework is no exception.

In the scotino DM scenario, the dark fermion ψ remains in thermal equilibrium with the Standard Model (SM) particles through processes such as $\psi\psi \rightarrow N_i N_j, \phi\phi$. However, for $m_\psi < (m_\phi, m_N)$, no tree-level annihilation channels are available that can efficiently deplete the ψ abundance to reproduce the observed dark matter relic density during thermal freeze-out, as the annihilation channels, $\psi\psi \rightarrow N_i N_j / \phi\phi$ are kinematically forbidden. Annihilation processes of the type $\psi\psi \rightarrow \text{SM}$ may arise at higher orders; however, these contributions are strongly suppressed by loop factors, the smallness of the Higgs-portal coupling λ' required by collider constraints, and the bounds on the seesaw Yukawa couplings.

In view of the above constraints, we do not pursue the single-component WIMP scenario any further in this work.

B. Coexisting fermion-scalar two-component dark matter

In kinematic regions where ϕ and ψ are mass degenerate or nearly degenerate, the coannihilation channel $\phi\psi \rightarrow \text{SM}$ becomes significant. In contrast, the annihilation processes $\phi\phi \rightarrow \text{SM}$ and $\psi\psi \rightarrow \text{SM}$ are highly suppressed, such that the dark matter relic abundance is predominantly determined by the coannihilation channel. Notice that the masses of ψ and ϕ are controlled independently by the bare mass parameters m_ψ and m_ϕ , respectively, which enter their physical masses, apart from the Higgs VEV-induced contribution in the case of m_ϕ and quantum corrections in both sectors. As a result, the mass degeneracy between ψ and ϕ can be tuned to the desired level and remains technically stable against radiative corrections without being destabilized by quantum effects.

The process mediated by the right-handed neutrinos, $\phi\psi \rightarrow \ell H$, is the main coannihilation channel. The Feynman diagram corresponding to these processes is depicted in Fig. 1.

The temperature evolution of the yields Y_ϕ and Y_ψ corresponding to the fields ϕ and ψ , respectively; is governed by the coupled Boltzmann equations,

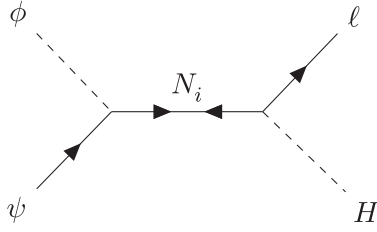


FIG. 1. Feynman diagram showing the tree-level coannihilation process $\phi\psi \rightarrow \ell H$.

$$\begin{aligned} \mathcal{H}Sx \frac{dY_\phi}{dx} &= \gamma^{\text{sc}} \left(1 - \frac{Y_\phi}{Y_\phi^{\text{eq}}} \frac{Y_\psi}{Y_\psi^{\text{eq}}} \right), \quad \text{and} \\ \mathcal{H}Sx \frac{dY_\psi}{dx} &= \gamma^{\text{sc}} \left(1 - \frac{Y_\phi}{Y_\phi^{\text{eq}}} \frac{Y_\psi}{Y_\psi^{\text{eq}}} \right), \end{aligned} \quad (4)$$

where γ^{sc} denotes the thermally averaged cross section for the processes $\phi\psi \rightarrow \ell H$, $\phi\bar{\psi} \rightarrow \bar{\ell} H^\dagger$, $\phi\bar{\psi} \rightarrow \ell H$, and $\phi\bar{\psi} \rightarrow \bar{\ell} H^\dagger$ (see Appendix C). We have assumed that the scotino ψ and its antiparticle $\bar{\psi}$ have equal number densities. The slight asymmetry, while important in the leptogenesis, is not relevant in the dark matter dynamics.

As discussed later in the context of leptogenesis, the complex Yukawa coupling y_i associated with the interaction term $\bar{\psi}\phi N_i$ can generate a small asymmetry between the number densities of ψ and $\bar{\psi}$. However, this effect is subdominant for the dark matter dynamics, which is primarily governed by the WIMP freeze-out process. Nevertheless, such an asymmetry in the fermionic dark sector can have important implications for leptogenesis, and this constitutes one of the key features of our framework. A detailed discussion of this aspect is presented in the leptogenesis Section.

The remaining parameters in Eq. (4) have their standard meanings: $x = m_\phi/T$ denotes the scaled inverse temperature, $\mathcal{H}$ is the Hubble parameter, S is the entropy density of the Universe, and $Y_{\phi,\psi}^{\text{eq}}$ are the equilibrium yields. Since the equilibrium number densities of ϕ and ψ differ primarily by their internal degrees of freedom, the ratios Y/Y^{eq} are approximately equal for the two species. Nevertheless, for completeness, we solve the coupled equations in Eq. (4) explicitly.

While the Boltzmann equations determine the late-time comoving yields $Y_{\phi,\psi}$, cosmological observations constrain the present-day dark matter abundance in terms of the dimensionless relic density parameter Ωh^2 . Accordingly, the relic density contributions are expressed as

$$\Omega_{\psi,\phi} h^2 = \frac{Sh^2 m_{\psi,\phi}}{\rho_c} Y_{\psi,\phi}, \quad (5)$$

where $\rho_c = \frac{3\mathcal{H}^2}{8\pi G}$ is the critical density, with $\mathcal{H} = 100 h \text{ km s}^{-1} \text{ Mpc}^{-1}$ and G denoting Newton's constant.

The total dark matter relic abundance is then given by

$$\Omega_{\text{tot}} h^2 = 2\Omega_\psi h^2 + \Omega_\phi h^2. \quad (6)$$

Here $\Omega_\psi h^2$ denotes the contribution from either ψ or $\bar{\psi}$ separately. Since the Boltzmann equation is solved for a single scotino species with $Y_\psi = Y_{\bar{\psi}}$, the total fermionic dark matter contribution carries an explicit factor of two in Eq. (6).

Notice that, for a fixed m_{N_i} , the relevant seesaw Yukawa couplings Y_{ij}^N are fixed. Consequently, the coannihilation cross section for the process $\phi\psi \rightarrow \ell H$ depends only on the common dark matter mass $m_\phi = m_\psi$ and the Yukawa coupling y_i . For a fixed relic abundance (or, equivalently, for a fixed thermally averaged effective annihilation rate) an increase in the dark matter mass must be compensated by a decrease in the coupling y_i , leading to an anticorrelation between $m_{\phi,\psi}$ and y_i . As a result, the perturbative limit $y_i = \sqrt{4\pi}$ corresponds to the minimum value of $m_{\phi,\psi}$ that reproduces the observed dark matter relic abundance. Further, with m_{N_1} in the TeV range and $m_{N_{2,3}} \gg m_{N_1}$, the effect of N_2 and N_3 mediated processes are negligible in the coannihilation process. Our analysis indicates that the DM candidate must be in the range of $\mathcal{O}(10^2)$ GeV to satisfy the observed relic density in this scenario. To remind the reader, large couplings y_i required here will not influence the detection prospects at the direct detection experiments, as the process is mediated by heavy neutrinos, which do not have any direct interaction with the nucleus.

Considering benchmark scenarios with m_{N_1} in the TeV range, Table II presents the lower bound on the dark matter mass corresponding to the perturbative limit of the coupling y_1 . We keep the value of the quartic coupling, λ' fixed to 1.0×10^{-10} all through this study. As mentioned above, this evades the dark matter direct detection constraints, as well as the Higgs sector collider constraints. As an illustrative example, for $m_{N_1} = 2 \text{ TeV}$ one finds $m_{\phi,\psi} \geq 200 \text{ GeV}$. The sensitivity of the relic abundance to variations in m_{N_1} is also illustrated in the Table II. A decrease in m_{N_1} enhances

TABLE II. Benchmark values of model parameters $m_\phi = m_\psi$, m_{N_1} , and y_1 and the corresponding DM relic abundance $\Omega_{\text{tot}} h^2$. The heavier right-handed neutrino masses are fixed at $m_{N_2} = 10^9 \text{ GeV}$ and $m_{N_3} = 2.6 \times 10^{13} \text{ GeV}$, and the quartic coupling is fixed at $\lambda' = 10^{-10}$.

BPs	$m_\phi = m_\psi$ (GeV)	m_{N_1} (GeV)	y_1	$\Omega_{\text{tot}} h^2$
BP 1	200	2000	$2\sqrt{\pi}$	0.12
BP 2	200	1000	$2\sqrt{\pi}$	0.04
BP 3	200	3000	$2\sqrt{\pi}$	0.21
BP 4	200	1000	$\sqrt{1.56}$	0.12
BP 5	187	1000	$2\sqrt{\pi}$	0.12
BP 6	209	3000	$2\sqrt{\pi}$	0.12

the coannihilation cross section for the process $\phi\psi \rightarrow \ell h$, resulting in a smaller relic density. This effect can be compensated either by reducing the $\bar{\psi}\phi N_1$ Yukawa coupling y_1 or by lowering the common dark matter mass $m_\phi = m_\psi$. Conversely, increasing m_{N_1} suppresses the coannihilation cross section. To maintain the observed relic abundance in this case, the minimum allowed value of $m_\phi = m_\psi$ shifts to larger values, since the coupling cannot exceed the perturbative limit. As mentioned above, for $m_{N_1} \ll m_{N_{2,3}}$, the contributions of N_2 and N_3 to the dark matter dynamics are negligible, and their precise masses are therefore not relevant for the present discussion. However, these masses play an important role in leptogenesis, since the corresponding seesaw Yukawa couplings enter the generation of the lepton asymmetry. Accordingly, we fix $m_{N_2} = 10^9$ GeV and $m_{N_3} = 2.6 \times 10^{13}$ GeV.

In the subsequent analysis, we fix $m_{N_1} = 2$ TeV as an illustrative benchmark point, keeping in mind the sensitivity of both the dark matter dynamics and leptogenesis, as discussed later, to this parameter space. As shown above, this choice requires a minimum dark matter mass of $m_{\phi,\psi} \simeq 200$ GeV when the coupling saturates its perturbative limit, $y_1 = \sqrt{4\pi}$. We find that such large couplings generically lead to a lepton asymmetry that is much below the observed value.

Our analysis further indicates that the generated lepton asymmetry is more sensitive to the coupling y_1 than to the dark matter mass, since y_1 enters both the CP asymmetry and the relevant reaction rates. In practice, viable leptogenesis typically requires y_1 in the range 0.01–0.1. Ideally, one would perform a comprehensive scan over the dark matter mass and coupling to identify the region consistent with the relic density, and independently carry out a similar scan for leptogenesis. The overlap of these regions would then determine the parameter space compatible with both dark matter constraints and successful leptogenesis. However, such a global scan is computationally prohibitive.

We therefore adopt a more pragmatic strategy. We fix the dark matter mass to $m_{\phi,\psi} = 250$ GeV, motivated by its proximity to the minimum mass required for the chosen neutrino-sector parameters. The coupling y_1 is then scanned over a coarse range to identify regions corresponding to underabundant and overabundant dark matter. For these values of y_1 , the resulting lepton asymmetry is subsequently computed to assess the viability of this parameter region. The results are presented in Table III. Let us now briefly examine the sensitivity of the above discussion to variations in m_{N_1} . As shown in Table II, increasing m_{N_1} while keeping it within the few-TeV range modifies the minimum allowed dark matter mass by no more than $\mathcal{O}(10\%)$. For dark matter masses above these minima, one can always identify a corresponding value of the coupling y_1 that reproduces the observed relic abundance. We also note that an underabundant dark matter contribution is phenomenologically acceptable, provided

TABLE III. Dark matter relic abundance $\Omega_{\text{tot}} h^2$ as a function of the coupling y_1 for a fixed DM mass of 250 GeV. The heavier right-handed neutrino masses are fixed at $m_{N_2} = 10^9$ GeV and $m_{N_3} = 2.6 \times 10^{13}$ GeV.

BPs	$m_\phi = m_\psi$ (GeV)	m_{N_1} (GeV)	y_1	$\Omega_{\text{tot}} h^2$
BP 7	250	2000	0.0042	0.12
BP 8	250	2000	0.01	0.032
BP 9	250	2000	0.1	0.0033
BP 10	250	2000	0.5	0.0028

the total dark matter relic density is not required to be saturated by this model alone.

IV. LEPTOGENESIS

We now turn our attention to the details of leptogenesis within the framework proposed in this study, where the required lepton number violation and CP asymmetry arise from the decay of heavy Majorana neutrinos $N_{1,2,3}$. In the conventional thermal leptogenesis scenario, such asymmetry is generated via the out-of-equilibrium decays of right-handed Majorana neutrinos into the SM leptons and the Higgs doublet. However, CP violation arises here through the quantum corrections to the decay process, both self-energy corrections and vertex corrections. In the standard scenario, these corrections are decided by the type-I seesaw interaction [9,63]. In our framework, $y_i \bar{\psi} \phi N_i$ provides the additional self-energy and vertex correction contributions. This additional interaction modifies the dynamics of leptogenesis in several important ways. Notably, the vertex and self-energy loop corrections contributing to CP violation in the decay processes receive new contributions from diagrams involving the fields ϕ and ψ . These additional contributions can potentially enhance the CP asymmetry compared to the standard scenario, and allow a viable leptogenesis at lower energy scales. Notice that, apart from the standard decay channel, $N_1 \rightarrow \ell H$, the heavy neutrino N_1 decays to the dark sector, $N_1 \rightarrow \phi\psi$ when kinematically allowed. In this study, learning from the dark matter analysis, we restrict ourselves to a degenerate $\phi - \psi$ dark matter, which is always considerably lighter than N_1 . The decay of a heavy neutrino into dark sector particles also provides the sources of CP asymmetry at one loop, very similar to the standard scenario. This asymmetry, however, does not induce any lepton number asymmetry directly. At the same time, this can generate significant asymmetry in the dark sector, which in turn can influence the thermal evolution of the lepton number asymmetry. While strict mass degeneracy of ψ with ϕ disallows ϕ decay, making it part of the dark matter spectrum, sufficient nonzero mass splitting could allow ϕ decay through three-body or four-body final states. In principle, these three-body and four-body decays of ϕ could also generate additional CP violation. However, in the rest of this study, we impose strict mass degeneracy of ϕ and ψ to

prohibit such three- or four-body decays of ϕ . Other sources of CP violation arise from scattering processes. Scattering channels mediated by right-handed neutrinos can also violate CP symmetry through the interference of amplitudes involving different heavy neutrinos, N_i and N_j with $i \neq j$. The relevant scattering processes in this case are $\phi\psi \rightarrow \ell H$, $\phi\bar{\psi} \rightarrow \ell H$, $\ell\phi \rightarrow H\psi$, $\bar{\ell}\phi \rightarrow H^\dagger\psi$, $\bar{\ell}\psi \rightarrow H^\dagger\phi$, and $\ell\psi \rightarrow H\phi$.

In the standard leptogenesis, scattering processes typically contribute to washout effects, often significantly reducing the lepton asymmetry generated by CP -violating decays. In the present framework, however, scattering processes play three distinct roles: they act as source terms that generate lepton asymmetry through their intrinsic CP violation, as washout terms that suppress the generated asymmetry, and as asymmetry-transfer processes that redistribute the asymmetry between the dark sector and the leptonic sector, and vice versa.

Before proceeding to the thermal evolution of the yields of the various particle species, we first analyze the CP asymmetries arising in the different processes.

A. CP asymmetry

The CP asymmetry arising from the decay of right-handed neutrinos N_i , as discussed in Refs. [9,64,65], originates from the interference between tree-level and loop-level diagrams. In the standard case, the loop-level contributions to the decay of right-handed neutrinos arise from diagrams with leptons and the Higgs boson running in the loop. In contrast, within the framework considered here, the CP asymmetry in the decay of N_1 receives contributions not only from the standard lepton-Higgs loops but also from additional quantum corrections involving the dark-sector fields ψ and ϕ . While the general idea of linking dark sector dynamics with leptogenesis has been explored previously, notably in Refs. [32,40,44], the scenario studied here is substantially different. Our framework does not rely on asymmetric dark matter, but instead considers a multicomponent WIMP scenario in which the dark matter abundance is set by thermal freeze-out. Furthermore, scattering processes play a qualitatively new role in our analysis, contributing not only to washout but also directly to the generation and transfer of lepton asymmetry, hence an important distinction of our study in comparison to Ref. [32]. We discuss the various contributions to the CP asymmetry in the following.

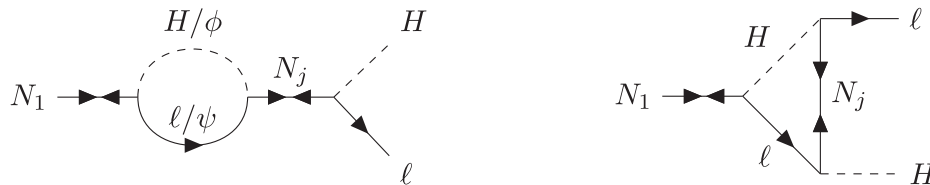


FIG. 2. Feynman diagrams representing the one-loop corrections to the decay $N_1 \rightarrow \ell H$ including the dark sector contributions to the self-energy.

1. CP asymmetry from the decay of $N_1 \rightarrow \ell H$

The CP asymmetry arising from the decay of N_1 is defined as

$$\epsilon_1 = \frac{\Gamma(N_1 \rightarrow \ell H) - \Gamma(N_1 \rightarrow \bar{\ell} H^\dagger)}{\Gamma(N_1 \rightarrow \ell H) + \Gamma(N_1 \rightarrow \bar{\ell} H^\dagger)}, \quad (7)$$

where $\Gamma(N_1 \rightarrow \ell H)$, and $\Gamma(N_1 \rightarrow \bar{\ell} H^\dagger)$ are the decay widths of N_1 into leptons and antileptons, respectively, given explicitly by

$$\Gamma_{N_1 \rightarrow \ell H} = \Gamma_{N_1 \rightarrow \ell H}^{\text{tree}}(1 + \epsilon_1), \quad (8)$$

$$\Gamma_{N_1 \rightarrow \bar{\ell} H^\dagger} = \Gamma_{N_1 \rightarrow \bar{\ell} H^\dagger}^{\text{tree}}(1 - \epsilon_1) \quad (9)$$

where the tree-level decay width is given by

$$\Gamma_{N_1 \rightarrow \ell H}^{\text{tree}} = \Gamma_{N_1 \rightarrow \bar{\ell} H^\dagger}^{\text{tree}} = \frac{(Y^{N^\dagger} Y^N)_{11}}{8\pi m_{N_1}^3} (m_{N_1}^2 - m_H^2)^2, \quad (10)$$

where m_{N_1} and m_H are the masses of the heavy neutrino and the SM Higgs boson, respectively. The quantum corrections to the decay process at one loop level leading to CP violation are shown in Fig. 2. Notice that, in addition to the standard contributions arising through the presence of H and ℓ in the loop, in the present case, the quantum corrections get contributions from the interaction of N_i with the dark sector particles, ϕ and ψ . These additional contributions are proportional to the new Yukawa couplings, y_i . While in the DM dynamics, only the magnitudes of these couplings $|y_i|$ play a role; the complex nature of the couplings y_i is vital for leptogenesis. The CP asymmetry including the contributions arising from the dark sector dynamics is given below with total CP asymmetry ϵ_1 getting contributions from the self-energy (ϵ^s) and vertex correction (ϵ^v),

$$\epsilon_1 = \epsilon^s + \epsilon^v, \quad \text{where} \quad (11)$$

$$\epsilon^s = \frac{1}{8\pi K_{11}} \sum_{j=2}^3 \frac{m_{N_1}}{m_{N_1}^2 - m_{N_j}^2} \times \text{Im}(m_{N_j} K_{1j}^2 + m_{N_1} \kappa_{j1} K_{1j} + m_{N_j} \kappa_{1j} K_{1j}), \quad \text{and} \quad (12)$$

TABLE IV. Numerical values of ϵ^s for some of the parameter values used in our analysis, illustrating that the inclusion of dark matter contributions in the loop significantly enhances the magnitude of ϵ^s .

m_{N_1}	m_{N_2}	m_{N_3}	y_1	y_2	y_3	ϵ^s
2000	10^9	26×10^{12}	0	0	0	2.1×10^{-8}
2000	10^9	26×10^{12}	$0.0059(1+i)$	$0.2(1-i)$	$0.3\sqrt{2\pi}(1+i)$	1.2×10^{-7}
2000	10^9	26×10^{12}	$\frac{0.1(1+i)}{\sqrt{2}}$	$0.2(1-i)$	$0.3\sqrt{2\pi}(1+i)$	1.2×10^{-6}
2000	1.5×10^7	26×10^{12}	0	0	0	1.05×10^{-7}
2000	1.5×10^7	26×10^{12}	$0.12(1+i)$	$0.2(1-i)$	$0.3\sqrt{2\pi}(1+i)$	2.1×10^{-5}

$$\epsilon^v = \frac{1}{8\pi K_{11}} \sum_{j=2}^3 \text{Im}(K_{1j}^2) \mathcal{F}\left(\frac{m_{N_j}^2}{m_{N_1}^2}\right), \quad (13)$$

where $K_{ij} = (Y^{N^\dagger} Y^N)_{ij}$, and $\kappa_{ij} = (y_i^\dagger y_j)$, and the function

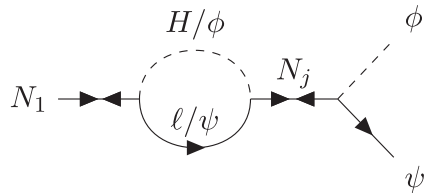
$$\mathcal{F}(x) = \sqrt{x} \left[1 + (1+x) \ln\left(\frac{x}{1+x}\right) \right]$$

is related to vertex loop factor. In the self-energy correction contribution, ϵ^s given in Eq. (12) has three different terms. The first term is related to the standard self-energy part of CP asymmetry, whereas the second and the third terms are the novel contributions due to ϕ and ψ in the loop. These additional terms are crucial for low-scale leptogenesis, without requiring mass degeneracy of right handed neutrinos. The importance of the dark sector contributions to the self-energy corrections is evident from Table IV. In the absence of dark matter effects, corresponding to setting $y_i = 0$, the resulting CP -violation parameter is suppressed by one or more orders of magnitude, depending on the choice of parameters.

2. CP asymmetry from decay of $N_1 \rightarrow \phi\psi$

As discussed before in Sec. IVA 1, the CP asymmetry generated by the decay of N_1 into the visible sector is influenced by its interaction with the dark sector. At the same time, the decay of N_1 into the dark sector also generates the CP asymmetry, leading to asymmetric generation of ψ and its antiparticle, $\bar{\psi}$. The CP asymmetry in this case is defined as

$$\epsilon_{\text{DM}} = \frac{\Gamma(N_1 \rightarrow \phi\psi) - \Gamma(N_1 \rightarrow \phi\bar{\psi})}{\Gamma(N_1 \rightarrow \phi\psi) + \Gamma(N_1 \rightarrow \phi\bar{\psi})},$$



where ϕ is its own antiparticle. The decay widths are explicitly given by

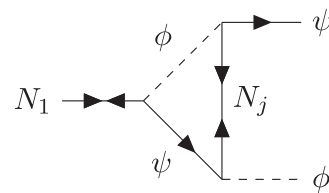
$$\Gamma_{N_1 \rightarrow \phi\psi} = \Gamma_{N_1 \rightarrow \phi\psi}^{\text{tree}} (1 + \epsilon_{\text{DM}}), \quad (14)$$

$$\Gamma_{N_1 \rightarrow \phi\bar{\psi}} = \Gamma_{N_1 \rightarrow \phi\bar{\psi}}^{\text{tree}} (1 - \epsilon_{\text{DM}}) \quad (15)$$

where the tree-level contribution is

$$\begin{aligned} \Gamma_{N_1 \rightarrow \phi\psi}^{\text{tree}} &= \Gamma_{N_1 \rightarrow \phi\bar{\psi}}^{\text{tree}} = \frac{|y_1|^2}{16\pi m_{N_1}^3} [m_{N_1}^2 + m_\psi^2 - m_\phi^2] \\ &\times \sqrt{(m_{N_1}^2 - m_\phi^2 - m_\psi^2)^2 - 4m_\phi^2 m_\psi^2}, \end{aligned}$$

while the interference between the tree-level amplitude and the one-loop contributions gives rise to the CP asymmetry parameter ϵ_{DM} . The Feynman diagrams associated with the one-loop processes are shown in Fig. 3. As is clear, the CP asymmetry in this decay channel potentially induces asymmetry in the number density of ψ and $\bar{\psi}$. Since the lepton number asymmetry is related to the number density of N_1 , it naturally raises a very interesting question: Does dark sector asymmetry affect the visible sector (lepton-antilepton) asymmetry? We will address it in the numerical discussion in Sec. V where we shall demonstrate that the dark sector asymmetry feeding into the lepton-antilepton asymmetry has a significant role in certain regions of the parameter space. At the same time, we assert that the effect of this asymmetry in dark matter relic density is negligible (in the relevant parameter region) as illustrated in the previous Section. Thus, unlike the asymmetric dark matter scenario discussed in the literature [32], we have a thermal WIMP dark matter framework.

FIG. 3. Feynman diagrams representing the one-loop corrections to the decay $N_1 \rightarrow \phi\psi$.

The CP asymmetry arising in the decay $N_1 \rightarrow \phi\psi$ is given by

$$\epsilon_{\text{DM}} = \epsilon_{\text{DM}}^s + \epsilon_{\text{DM}}^v, \quad \text{where} \quad (16)$$

$$\epsilon_{\text{DM}}^s = \frac{1}{8\pi\kappa_{11}} \sum_{j=2}^3 \frac{m_{N_1}}{m_{N_1}^2 - m_{N_j}^2} \times \text{Im}(m_{N_j}\kappa_{1j}^2 + m_{N_1}K_{j1}\kappa_{1j} + m_{N_j}K_{1j}\kappa_{1j}), \quad \text{and} \quad (17)$$

$$\epsilon_{\text{DM}}^v = \frac{1}{8\pi\kappa_{11}} \sum_{j=2}^3 \text{Im}(\kappa_{1j}^2) \mathcal{F}\left(\frac{m_{N_j}^2}{m_{N_1}^2}\right), \quad (18)$$

respectively.

3. CP asymmetry from $2 \rightarrow 2$ scattering process $\phi\psi \rightarrow \ell H$

A new contribution to the total lepton asymmetry comes from the scattering process, $\psi\phi \rightarrow H\ell$ mediated by heavy neutrinos $N_{1,2,3}$. Notice that, while Ref. [44] discuss the asymmetry in scattering processes induced through self-energy corrections of the propagator, in our framework the CP asymmetry is generated at the tree level through interference of two different propagators. The corresponding Feynman diagram is shown in Fig. 4. The importance of these CP asymmetry sources is that they affect the dark sector as well as visible sector asymmetries both. Therefore, these channels turn into a bridge between the dark sector and the visible sector asymmetries. An asymmetry generated in the dark sector is transferred to the visible sector and vice versa. The CP asymmetry generated

by scattering process $\psi\phi \rightarrow H\ell$ is parametrized as

$$\epsilon_2 = \frac{\sigma(\psi\phi \rightarrow H\ell) - \sigma(\bar{\psi}\phi \rightarrow H^\dagger\bar{\ell})}{\sigma(\psi\phi \rightarrow H\ell) + \sigma(\bar{\psi}\phi \rightarrow H^\dagger\bar{\ell})}, \quad (19)$$

where σ 's denote the cross sections of corresponding processes.

At the tree level, interference between two different scattering $\phi\psi \rightarrow \ell H$ processes that are mediated by different heavy right-handed neutrinos provides the additional sources of CP asymmetry. This phenomenon illustrates how multiple pathways in particle scattering can lead to significant effects in terms of CP violation. Here, we take an approximation where $m_{N_1} \ll m_{N_2} \ll m_{N_3}$. In this scenario, the dominant contribution to the denominator arises from the process that is mediated by the lightest heavy right-handed neutrino N_1 , whereas ϵ_2 itself gets contributions from the interference between N_1 and N_2 mediated channels. In the center of mass frame of the process (with the center of mass energy $\sqrt{s}$), the generated asymmetry is given by

$$\epsilon_2 = \frac{4 \text{Im}[y_1^* y_2 Y_{1\alpha}^N (Y_{2\alpha}^N)^*] \text{Im}[D_1^* D_2]}{|y_1|^2 K_{11} |D_2|^2}, \quad (20)$$

where $K_{11} = (Y^{N\dagger} Y^N)_{11}$. The parameters $y_{1,2}$ and $Y_{(1,2)\alpha}^N$, with $\alpha = 1, 2, 3$, denote the Yukawa couplings of N_1 to the dark sector and to the visible sector, respectively. Besides that $D_i = s - m_{N_i}^2 + im_{N_i}\Gamma_i$, with $i = 1, 2$, denotes the denominators associated with the propagators of different right-handed neutrinos with Γ_i being the total decay width of the mediating heavy neutrino N_i . After substituting D_i into Eq. (20), ϵ_2 becomes,

$$\epsilon_2 = \frac{4 \text{Im}[y_1^* y_2 Y_{1\alpha}^N (Y_{2\alpha}^N)^*] [(s - m_{N_1}^2)m_{N_2}\Gamma_2 - (s - m_{N_2}^2)m_{N_1}\Gamma_1]}{|y_1|^2 K_{11} [(s - m_{N_2}^2)^2 + m_{N_2}^2 \Gamma_2^2]}. \quad (21)$$

It is clear from Eq. (21) that ϵ_2 depends on the center of mass energy $\sqrt{s}$, the Yukawa couplings y and Y , the masses of right-handed neutrinos, m_{N_i} , and the total decay widths of right-handed neutrinos. Here, for m_{N_3} being very heavy in comparison to m_{N_1} and m_{N_2} , the contribution from N_3 is highly suppressed, and therefore has been neglected in our analysis.

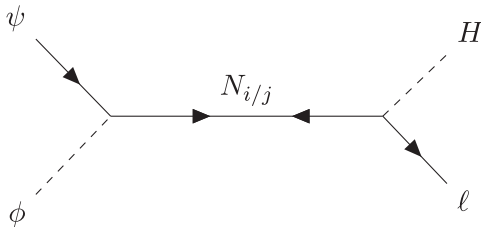


FIG. 4. Feynman diagrams of the scattering process $\phi\psi \rightarrow \ell H$ ($i \neq j$).

4. CP asymmetry from $2 \rightarrow 2$ scattering process $\phi\bar{\psi} \rightarrow \ell H$

Another scattering process $\phi\bar{\psi} \rightarrow \ell H$ also violates the CP asymmetry. The relevant Feynman diagram is shown in Fig. 5. Similar to the scattering process $\phi\psi \rightarrow \ell H$, this channel also works as a bridge between the dark sector

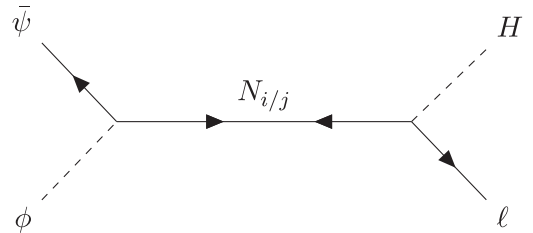


FIG. 5. Feynman diagrams for the scattering process $\phi\bar{\psi} \rightarrow \ell H$ ($i \neq j$).

and the visible sector asymmetries. The CP asymmetry generated by the scattering process $\phi\bar{\psi} \rightarrow \ell H$ is parametrized as

$$\epsilon'_2 = \frac{\sigma(\bar{\psi}\phi \rightarrow H\ell) - \sigma(\psi\phi \rightarrow H^\dagger\bar{\ell})}{\sigma(\bar{\psi}\phi \rightarrow H\ell) + \sigma(\psi\phi \rightarrow H^\dagger\bar{\ell})}, \quad (22)$$

where σ denotes the cross section of the associated process. The CP asymmetry associated with this process ϵ'_2 originates from the interference of two different N_i mediated processes shown in Fig. 5. For being insignificant, the

contributions from N_3 has been ignored in this analysis. After putting the cross-sections of relevant processes in Eq. (22), ϵ'_2 in the center of mass frame of the process becomes,

$$\epsilon'_2 = \frac{4 \operatorname{Im}[y_1 y_2^* Y_{1\alpha}^N (Y_{2\alpha}^N)^*] \operatorname{Im}[D_1^* D_2]}{|y_1|^2 K_{11} |D_2|^2}, \quad (23)$$

where $\alpha = 1, 2, 3$ is summed over, and D_i denotes the denominator of the corresponding propagator. Substituting the explicit expression for this D_i factor, ϵ'_2 is given by

$$\epsilon'_2 = \frac{4 \operatorname{Im}[y_1 y_2^* Y_{1\alpha}^N (Y_{2\alpha}^N)^*] [(s - m_{N_1}^2) m_2 \Gamma_2 - (s - m_{N_2}^2) m_{N_1} \Gamma_1]}{|y_1|^2 K_{11} [(s - m_{N_2}^2)^2 + m_{N_2}^2 \Gamma_2^2]}. \quad (24)$$

Notice that the total decay width of N_1 depends on the mass spectrum of the dark sector, therefore ϵ_2 and ϵ'_2 indirectly depend on the masses of DM candidates.

5. CP asymmetry from $2 \rightarrow 2$ scattering t channels

Apart from the above CP asymmetry generating $2 \rightarrow 2$ s-channel scattering processes, there are several $2 \rightarrow 2$ t-channel scattering processes that violate CP symmetry. The relevant Feynman diagrams are listed in Fig. 6. Notice that both ϕ and ψ are DM candidates. Certain t-channel scattering processes allow the transfer of asymmetry generated in the dark sector to the visible sector and vice versa, whereas in a few other t-channel processes, this asymmetry transfer does not occur.

The CP asymmetry parameter generated by scattering processes $\ell\phi \rightarrow H\psi$, $\bar{\ell}\phi \rightarrow H^\dagger\bar{\psi}$, $\bar{\ell}\psi \rightarrow H^\dagger\phi$ and $\ell\psi \rightarrow H\phi$ are given by

$$\epsilon_1^t = \frac{\Gamma(\ell\phi \xrightarrow{t} H\psi) - \Gamma(\bar{\ell}\phi \xrightarrow{t} H^\dagger\bar{\psi})}{\Gamma(\ell\phi \xrightarrow{t} H\psi) + \Gamma(\bar{\ell}\phi \xrightarrow{t} H^\dagger\bar{\psi})} = -\frac{\Gamma(\bar{\ell}\psi \xrightarrow{t} H^\dagger\phi) - \Gamma(\ell\bar{\psi} \xrightarrow{t} H\phi)}{\Gamma(\bar{\ell}\psi \xrightarrow{t} H^\dagger\phi) + \Gamma(\ell\bar{\psi} \xrightarrow{t} H\phi)}, \quad (25)$$

$$\epsilon_2^t = \frac{\Gamma(\bar{\ell}\phi \xrightarrow{t} H^\dagger\bar{\psi}) - \Gamma(\ell\phi \xrightarrow{t} H\bar{\psi})}{\Gamma(\bar{\ell}\phi \xrightarrow{t} H^\dagger\bar{\psi}) + \Gamma(\ell\phi \xrightarrow{t} H\bar{\psi})} = -\frac{\Gamma(\ell\psi \xrightarrow{t} H\phi) - \Gamma(\bar{\ell}\bar{\psi} \xrightarrow{t} H^\dagger\phi)}{\Gamma(\ell\psi \xrightarrow{t} H\phi) + \Gamma(\bar{\ell}\bar{\psi} \xrightarrow{t} H^\dagger\phi)}, \quad (26)$$

respectively. The CP asymmetry parameters ϵ_1^t and ϵ_2^t in the center of mass frame of the process are given as follows

$$\epsilon_1^t = \frac{4 \operatorname{Im}[y_1^* y_2 Y_{1\alpha}^N (Y_{2\alpha}^N)^*] \operatorname{Im}[D_1^{t*} D_2^t]}{|y_1|^2 K_{11} |D_2^t|^2}, \quad \text{and} \quad (27)$$

$$\epsilon_2^t = \frac{4 \operatorname{Im}[y_1^* y_2 (Y_{1\alpha}^N)^* Y_{2\alpha}^N] \operatorname{Im}[D_1^{t*} D_2^t]}{|y_1|^2 K_{11} |D_2^t|^2}, \quad (28)$$

where $y_{1,2}$ and $Y_{(1,2)\alpha}^N$ with $\alpha = 1, 2, 3$ are Yukawa couplings. With the propagator factor, $D_i^t = t - m_{N_i}^2 + im_{N_i}\Gamma_i$, the expressions for ϵ_1^t and ϵ_2^t are given by

$$\epsilon_1^t = \frac{4 \operatorname{Im}[y_1^* y_2 Y_{1\alpha}^N (Y_{2\alpha}^N)^*] [(t - m_{N_1}^2) m_{N_2} \Gamma_2 - (t - m_{N_2}^2) m_{N_1} \Gamma_1]}{|y_1|^2 K_{11} [(t - m_{N_1}^2)^2 + m_{N_1}^2 \Gamma_1^2]}, \quad \text{and} \quad (29)$$

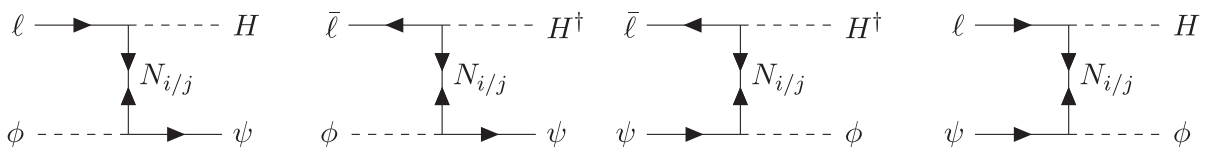


FIG. 6. Feynman diagrams for the scattering process $\phi\ell \rightarrow \psi H$, $\phi\bar{\ell} \rightarrow \psi H^\dagger$, $\psi\bar{\ell} \rightarrow \phi H^\dagger$, and $\psi\ell \rightarrow \phi H$, where $i \neq j$.

$$\epsilon_2' = \frac{4 \operatorname{Im} [y_1^* y_2 (Y_{1\alpha}^N)^* Y_{2\alpha}^N] [(t - m_{N_1}^2) m_{N_2} \Gamma_2 - (t - m_{N_2}^2) m_{N_1} \Gamma_1]}{|y_1|^2 K_{11} [(t - m_{N_2}^2)^2 + m_{N_2}^2 \Gamma_2^2]}, \quad (30)$$

respectively, where m_{N_i} are the right handed neutrino masses of N_i , and Γ_i are the total decay widths of N_i . The imaginary parts of coupling products and denominator of propagator products are responsible for generating the CP violating phase. It is clear that the complex couplings, mass splittings, and decay widths of intermediate particles play a crucial role in generating CP violating phases. The dependency of ϵ_1' on right-handed neutrino masses and couplings is the same as that in the case of ϵ_2 .

In total, there are ten CP asymmetry sources in this framework. Two are associated with the decays of N_1 . Another six CP asymmetry sources are connected to the scattering processes (two from s-channel processes and remaining four from t-channel processes). The last two CP asymmetry sources are associated with three-body decay and four-body decay of ϕ . These last two sources of CP asymmetry are not relevant in our analysis due to mass degeneracy of m_ϕ and m_ψ .

B. Time evolution of the number densities of particles

The time evolution of the number densities of N_1 , ϕ , ψ , and $\bar{\psi}$ is crucial for studying the asymmetries in both the dark sector and the visible sector. The evolution of the particle number densities when the system is out of thermal

equilibrium is governed by the Boltzmann equations. In scenarios where the masses of the right-handed neutrino (RHN) N_i satisfy the hierarchy $m_{N_3} \gg m_{N_2} \gg m_{N_1}$, any lepton asymmetry generated by the decays of the heavier neutrinos N_3 and N_2 at high temperatures is efficiently washed out before the EWSB. As a result, the lepton asymmetry that survives to lower temperatures is predominantly produced by the decay of the lightest right-handed neutrinos, N_1 . The evolution of the number density of N_1 is influenced by several processes, including its decay and scattering processes. Similarly, the number density of the leptons is decided by processes involving their production (like N_1 decay and scattering processes), and annihilation (like the inverse decay and scattering processes). The processes relevant to the visible sector asymmetry and the dark sector asymmetry in our case are presented in Appendix A. These interactions collectively determine the abundance of N_1 , ϕ , ψ , $\bar{\psi}$, and the resulting lepton asymmetry. The generated lepton asymmetry can subsequently be changed into the observed baryon asymmetry of the Universe via sphaleron processes [66]. The relevant Boltzmann equations for yields $Y_{N_1} = \frac{n_{N_1}}{s}$, $Y_L = \frac{n_l - n_{\bar{l}}}{s}$, and $Y_{\Delta\psi} = \frac{n_\psi - n_{\bar{\psi}}}{s}$ are as follows:

$$\mathcal{H}S_z \frac{dY_{N_1}}{dz} = \left(1 - \frac{Y_{N_1}}{Y_{N_1}^{\text{eq}}}\right) (\gamma_{D_1} + \gamma_{D_2} + 2\gamma_s + 2\gamma_{t_1} + 2\gamma_{t_2} + 2\gamma_V), \quad (31)$$

$$\begin{aligned} \mathcal{H}S_z \frac{dY_L}{dz} = & \left(\frac{Y_{N_1}}{Y_{N_1}^{\text{eq}}} - 1\right) \epsilon_1 \gamma_{D_1} + 2\Theta_{\epsilon_2}^{\phi\psi} + 2\Theta_{\epsilon_2'}^{\phi\bar{\psi}} - 2\Theta_{\epsilon_1'}^{\phi\ell} - 2\Theta_{\epsilon_1'}^{\psi\bar{\ell}} + 2\Theta_{\epsilon_2'}^{\phi\bar{\ell}} + 2\Theta_{\epsilon_2'}^{\psi\ell} - \gamma_D^{\text{os}} \epsilon_1 \\ & + \frac{Y_{\Delta\psi}}{Y_{\psi}^{\text{eq}}} \left(\frac{\gamma_{\phi\psi}}{2} + \frac{\gamma'_{\phi\psi}}{2} + \frac{\gamma_{\phi\ell}^t}{2} - \frac{\gamma_{\phi\bar{\ell}}^t}{2}\right) \\ & - \frac{Y_L}{Y_L^{\text{eq}}} \left(\frac{\gamma_{D_1}}{2} + \gamma_{t_1} + \gamma_{t_2} + \gamma_{V_1} + \gamma_{V_3} + \gamma_{V_4} + \gamma_{V_5} + \frac{\gamma_{\phi\psi}}{2} + \frac{\gamma'_{\phi\psi}}{2} + \frac{\gamma_{\phi\ell}^t}{2} + \frac{\gamma_{\phi\bar{\ell}}^t}{2} + \gamma_{\psi\bar{\ell}}^t\right. \\ & \left. + \frac{Y_{N_1}}{Y_{N_1}^{\text{eq}}} (\gamma_s + \gamma_{V_2} + \gamma_{V_6}) - \gamma_D^{\text{os}}\right), \quad \text{and} \end{aligned} \quad (32)$$

$$\begin{aligned} \mathcal{H}S_z \frac{dY_{\Delta\psi}}{dz} = & \left(\frac{Y_N}{Y_N^{\text{eq}}} - 1\right) \epsilon_{\text{DM}} \gamma_{D_2} - 2\Theta_{\epsilon_2}^{\phi\psi} + 2\Theta_{\epsilon_2'}^{\phi\bar{\psi}} + 2\Theta_{\epsilon_1'}^{\phi\ell} + 2\Theta_{\epsilon_1'}^{\psi\bar{\ell}} + 2\Theta_{\epsilon_2'}^{\phi\bar{\ell}} + 2\Theta_{\epsilon_2'}^{\psi\ell} + \gamma_D^{\text{os}} \epsilon_1 \\ & + \frac{Y_L}{Y_L^{\text{eq}}} \left(\frac{\gamma_{\phi\psi}}{2} - \frac{\gamma'_{\phi\psi}}{2} + \frac{\gamma_{\phi\ell}^t}{2} - \gamma_{\phi\bar{\ell}}^t - \gamma_D^{\text{os}}\right) \\ & - \frac{Y_{\Delta\psi}}{Y_{\psi}^{\text{eq}}} \left(\frac{\gamma_{D_2}}{2} + \frac{\gamma_{\phi\psi}}{2} + \frac{\gamma'_{\phi\psi}}{2} + \frac{\gamma_{\phi\ell}^t}{2} + \frac{\gamma_{\psi\bar{\ell}}^t}{2} + \frac{\gamma_{\phi\bar{\ell}}^t}{2} + \frac{\gamma_{\psi\ell}^t}{2}\right), \end{aligned} \quad (33)$$

where γ_{D_1} and γ_{D_2} are the thermal averaged decay widths of $N_1 \rightarrow \ell H$ and $N_1 \rightarrow \phi\psi$, respectively. γ_{t_1} is the thermal average of the processes $N_1 u \leftrightarrow d\bar{\ell}$ and $N_1 \bar{u} \leftrightarrow \bar{d}\ell$, and γ_{t_2} is the same for the processes $N_1 d \leftrightarrow u\bar{\ell}$ and $N_1 \bar{d} \leftrightarrow \bar{u}\ell$. We denote the sum of these contributions as $\gamma_t = \gamma_{t_1} + \gamma_{t_2}$. Similarly, γ_s corresponds to the s -channel processes $\ell N_1 \leftrightarrow d\bar{u}$, and $\bar{\ell} N_1 \leftrightarrow \bar{d}u$, while γ_{V_i} , where $i = 1, \dots, 6$, are the thermal averaged cross sections of scattering processes involving the gauge bosons: $N_1 H \xleftrightarrow{s} V_\mu \ell$, $N_1 \ell \xleftrightarrow{s} V_\mu H$, $N_1 V_\mu \xleftrightarrow{t} H\ell$, $N_1 V_\mu \xleftrightarrow{u} H\bar{\ell}$, $N_1 H^\dagger \xleftrightarrow{t} V_\mu \ell$, $N_1 \ell \xleftrightarrow{t} V_\mu H^\dagger$, and their charged conjugate processes in that order, with $V_\mu = W_\mu, Z_\mu$. The quantities $\gamma_{\phi\psi}$, $\gamma'_{\phi\psi}$, $\gamma'_{\phi\ell}$, $\gamma'_{\phi\bar{\ell}}$, $\gamma'_{\psi\bar{\ell}}$, γ_{LL} , γ_{HL} , and $\gamma'_{\psi\ell}$ represent the thermally averaged scattering cross-sections of the processes $\phi\psi \rightarrow \ell H$, $\phi\bar{\psi} \rightarrow \bar{\ell} H$, $\phi\ell \rightarrow \psi H$, $\phi\bar{\ell} \rightarrow \bar{\psi} H^\dagger$, $\psi\bar{\ell} \rightarrow \phi H^\dagger$, $\ell\ell \rightarrow H^\dagger H^\dagger$, $H\ell \rightarrow H^\dagger \bar{\ell}$, and $\psi\ell \rightarrow \phi H$, respectively (as given in the Appendix (C1)), while $\Theta_{\epsilon_a}^{ab}$, given in Appendix (C3), represent the thermal average of $\epsilon_a \sigma^{ab}$ corresponding to the respective scattering process. In addition, $\gamma_D^{os} = \gamma_D \mathcal{B}_1 \mathcal{B}_2$, where γ_D is the thermal average of the total decay width of N_1 , whereas $\mathcal{B}_1$ and $\mathcal{B}_2$ correspond to the branching fractions to $N_1 \rightarrow \ell H$ and $N_1 \rightarrow \phi\psi$, respectively. These terms represent the on-shell contributions explicitly subtracted from the scattering process, $\phi\psi \rightarrow \ell H$ and $\phi\bar{\psi} \rightarrow \bar{\ell} H$, and the conjugate channels, in order to avoid double counting of the real intermediate N_1 state already included through decay and inverse-decay contributions. For the range of parameters considered in this work, these on-shell subtraction terms are numerically negligible compared to the other reaction rates. In Eq. (32), the terms proportional to $Y_{\Delta\psi}$ indicate the transfer of dark-sector asymmetry into the lepton asymmetry. Similarly, in Eq. (33),

the terms involving Y_L represent the transfer of visible sector asymmetry into the dark sector. The strength of the transfer of asymmetry from the dark sector to the visible sector and vice versa depends on the choice of parameter values. In a simplified scenario, if the scattering processes involving dark matter occur in such a way that the numbers of dark sector particles ψ and $\bar{\psi}$ remain equal during the evolution, then no dark sector asymmetry is generated, nor can any dark-sector asymmetry be transferred to the visible sector. Under this condition, only two equations are required to analyze the resulting lepton asymmetry. As described below, this simplified case can be solved analytically in a schematic manner.

1. Case when asymmetry in dark sector is absent

We shall explore the specific case when dark sector asymmetry is negligible. Consequently, the abundances of ψ and $\bar{\psi}$ remain equal [$Y_\psi(z) = Y_{\bar{\psi}}(z)$]. Under this condition, Eqs. (31) and (32) simplify to

$$\frac{dY_{N_1}}{dz} + P_N(z)Y_{N_1} = Q_N(z), \quad (34)$$

$$\frac{dY_L}{dz} + P_L(z)Y_L = Q_L(z), \quad (35)$$

where $P_N(z)$, $Q_N(z)$, $P_L(z)$, and $Q_L(z)$ are defined in Eqs. (36)–(39), respectively.

$$P_N(z) = \frac{\gamma_{D_1} + \gamma_{D_2} + 2\gamma_s + 2\gamma_{t_1} + 2\gamma_{t_2} + 2\gamma_V}{\mathcal{H}S_Z Y_{N_1}^{\text{eq}}}, \quad (36)$$

$$Q_N(z) = Y_{N_1}^{\text{eq}} P_N(z), \quad (37)$$

$$P_L(z) = \frac{1}{\mathcal{H}S_Z Y_L^{\text{eq}}} \left[\frac{\gamma_{D_1}}{2} + \gamma_{t_1} + \gamma_{t_2} + \gamma_{V_1} + \gamma_{V_3} + \gamma_{V_4} + \gamma_{V_5} + \frac{\gamma_{\phi\psi}}{2} + \frac{\gamma'_{\phi\psi}}{2} + \frac{\gamma_{\phi\ell}}{2} + \frac{\gamma'_{\phi\bar{\ell}}}{2} + \gamma'_{\psi\bar{\ell}} + \frac{Y_{N_1}}{Y_{N_1}^{\text{eq}}} (\gamma_s + \gamma_{V_2} + \gamma_{V_6}) \right], \quad \text{and} \quad (38)$$

$$Q_L(z) = \frac{\epsilon_1 \gamma_{D_1}}{\mathcal{H}S_Z} \left(\frac{Y_{N_1}}{Y_{N_1}^{\text{eq}}} - 1 \right) + \frac{2}{\mathcal{H}S_Z} [\Theta_{\epsilon_2}^{\phi\psi} + \Theta_{\epsilon_2'}^{\phi\bar{\psi}} - \Theta_{\epsilon_1'}^{\phi\ell} - \Theta_{\epsilon_1'}^{\psi\bar{\ell}} + \Theta_{\epsilon_2'}^{\phi\bar{\ell}} + \Theta_{\epsilon_2'}^{\psi\ell}] \\ = \epsilon_1 Q_1(z) + Q_{\epsilon_2}(z) + Q_{\epsilon_2'}(z) + Q_{\epsilon_1'}(z) + Q_{\epsilon_1'}'(z) + Q_{\epsilon_2'}(z) + Q_{\epsilon_2'}'(z). \quad (39)$$

In Eq. (39), the individual source terms $Q_i(z)$ are defined implicitly through decomposition. This simple form, allows us to solve Eqs. (34) and (35) formally with the solutions

$$Y_{N_1}(z) = \frac{1}{\mu(z)} \left[Y_{N_1}^{\text{eq}}(z_0) + \int_{z_0}^z \mu(s) Q_N(s) ds \right], \quad \text{and} \quad (40)$$

$$Y_L(z) = \epsilon_1 \kappa_1 + \kappa_2 + \kappa_3 + \kappa_4 + \kappa_5 + \kappa_6 + \kappa_7, \quad (41)$$

respectively, where $\mu(z) = \exp(\int_{z_0}^z P_N(z')dz')$. The definitions of the various efficiency factors κ_i are given by

$$\begin{aligned}\kappa_1 &= \frac{1}{\nu(z)} \int_{z_0}^z \nu(z') Q_1(z') dz', & \kappa_2 &= \frac{1}{\nu(z)} \int_{z_0}^z \nu(z') Q_{\epsilon_2}(z') dz', \\ \kappa_3 &= \frac{1}{\nu(z)} \int_{z_0}^z \nu(z') Q_{\epsilon'_2}(z') dz', & \kappa_4 &= \frac{1}{\nu(z)} \int_{z_0}^z \nu(z') Q_{\epsilon'_1}(z') dz', \\ \kappa_5 &= \frac{1}{\nu(z)} \int_{z_0}^z \nu(z') Q'_{\epsilon'_1}(z') dz', & \kappa_6 &= \frac{1}{\nu(z)} \int_{z_0}^z \nu(z') Q_{\epsilon'_2}(z') dz', & \kappa_7 &= \frac{1}{\nu(z)} \int_{z_0}^z \nu(z') Q'_{\epsilon'_2}(z') dz',\end{aligned}\quad (42)$$

where each κ_i represents the accumulated contribution of the corresponding source term $Q_i(z)$, while being weighted by the integrating factor $\nu(z)$ given by

$$\nu(z) = \exp\left(\int_{z_0}^z P_L(z') dz'\right). \quad (43)$$

These efficiency factors encode the different sources of matter-antimatter asymmetry. Specifically, κ_1 contributes to the asymmetry only when Y_{N_1} deviates from its equilibrium value $Y_{N_1}^{\text{eq}}$. The factors κ_2 and κ_3 arise from the s -channel $2 \rightarrow 2$ scattering processes $\phi\psi \rightarrow \ell H$ and $\phi\bar{\psi} \rightarrow \ell H$, respectively; their contributions are maximized when the CP -violating combinations $\epsilon_2 \sigma^{\phi\psi}$ and $\epsilon'_2 \sigma^{\phi\bar{\psi}}$ are largest. Similarly, κ_4 and κ_5 originate from the t -channel processes $\ell\phi \rightarrow H\psi$ and $\bar{\ell}\psi \rightarrow \phi H^\dagger$, respectively, while κ_6 and κ_7 stem from the related t -channel processes $\bar{\ell}\phi \rightarrow H^\dagger\psi$ and $\ell\psi \rightarrow H\phi$, respectively. The contributions from κ_4 , κ_5 , κ_6 , and κ_7 are maximized when the respective CP -violating scattering rates $\epsilon_1^t \sigma^{\phi\ell}$, $\epsilon_1^t \sigma^{\bar{\ell}\psi}$, $\epsilon_2^t \sigma^{\phi\bar{\ell}}$, and $\epsilon_2^t \sigma^{\ell\psi}$ are the largest. Finally, the CP asymmetry parameter ϵ_1 and all efficiency factors $\kappa_{1,2,\dots,7}$ depend on the Yukawa couplings $y_{1,2,3}$, the heavy neutrino masses $m_{N_{1,2,3}}$, and the DM masses m_ϕ and m_ψ . Consequently, the lepton asymmetry Y_L depends solely on these free parameters.

In Eqs. (32) and (33), the parameters ϵ_2 , ϵ'_2 , ϵ_1^t , and ϵ_2^t appear linearly. The CP asymmetry parameters ϵ_1 and ϵ_{DM} are associated with the decays of RHNs, which behave as a source term for the visible sector and dark sector asymmetry, respectively. While the scattering processes serves as sources to generate asymmetry in both the dark sector and the visible sector, they also act as a washout process, as well as transferring asymmetry from the dark sector to the visible sector and vice versa. The CP asymmetry sources ϵ_2 and ϵ'_2 contribute inversely to each other in the leptonic sector.

V. NUMERICAL ANALYSIS AND RESULTS

In this Section, we perform a numerical analysis of the lepton asymmetry (visible sector) and DM asymmetry generated by decays and scattering processes involving the heavy right-handed neutrinos. These asymmetries depend on the free parameters y_i , m_{N_i} (for $i = 1, 2, 3$),

m_ϕ , and m_ψ . The CP asymmetry parameters discussed in Sec. IV A play a crucial role in the generation of these asymmetries. In this framework, there are eight such parameters and we have considered only ϵ_1 , ϵ_{DM} , ϵ_2 , ϵ'_2 , ϵ_1^t , and ϵ_2^t in our investigation as the CP asymmetries from three-body and four-body decays of ϕ are not allowed here due to mass degeneracy in the dark sector ($m_\phi = m_\psi$). The CP asymmetry parameter ϵ_1 and ϵ_{DM} depend on y_i and m_{N_i} , while other CP asymmetry parameters ϵ_2 , ϵ'_2 , ϵ_1^t and ϵ_2^t depend on the decay width of N_1 in addition to y_i , and m_{N_i} . We discuss the visible sector asymmetry study in Subsection V A and dark sector asymmetry study in Subsection V B below.

A. Numerical analysis of visible sector asymmetry

As explained in the previous sections, in the present framework, the lepton asymmetry is influenced by the dark sector in multiple ways. Our numerical study elaborated in this section will categorically establish that this influence is significant enough to have sufficiently large leptogenesis with TeV-scale N_1 . We shall illustrate this with the results presented in Fig. 7, where the baryon number asymmetry, η is given as a function of $z = \frac{m_{N_1}}{T}$. Parameter values used for this study are given in Table V. The couplings, y_2 and y_3 as well as the masses of the heavy neutrinos, N_2 and N_3 are fixed in both the cases described in the first two plots. In Fig. 7(a), the observed baryon asymmetry is achievable for $m_\phi = m_\psi = 200$ GeV for a corresponding $y_1 = 0.0059(1-i)$. However, notice that this case corresponds to an overabundant dark matter scenario. The red curve corresponding to $m_\phi = m_\psi = 250$ GeV for the same coupling produces larger than observed baryon asymmetry, while being consistent with the dark matter abundance. Figure 7(a) also demonstrates the significance of dark matter asymmetry (η_{DM}) in leptogenesis. As expected, when $\eta_{\text{DM}} = 0$, the baryon asymmetry is insensitive to the dark matter mass. One can conclude that, for $m_{N_1} = 2$ TeV, and for the other right-handed neutrino masses and seesaw Yukawa's fixed as in this case, the dark matter mass around or below 200 GeV is not consistent with both dark matter observations and required baryon asymmetry. Keeping the dark matter mass a little higher, at 250 GeV, it is possible to adjust the coupling y_1 to get the right value of η , while being consistent with the

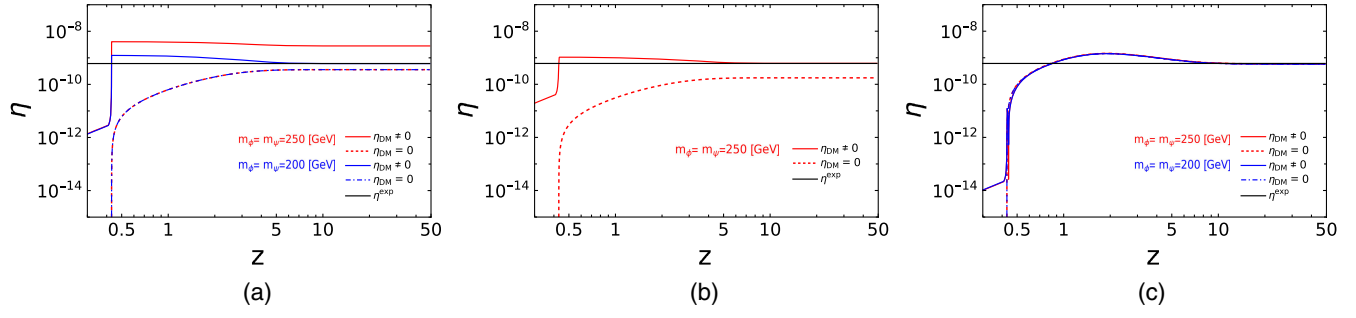


FIG. 7. Baryon asymmetry parameter as a function of the dimensionless variable $z = m_{N_1}/T$. The three different scenarios, (a), (b), and (c), correspond to different sets of parameter values as given in Table V.

TABLE V. Fixed parameter points for η vs z plot Fig. 7.

B.P.	m_{N_1} (GeV)	m_{N_2} (GeV)	m_{N_3} (GeV)	y_1	y_2	y_3
(a)	2000	10^9	26×10^{12}	$0.0059(1-i)$	$0.20(1+i)$	$0.3\sqrt{2\pi}(1-i)$
(b)	2000	10^9	26×10^{12}	$\frac{0.1}{\sqrt{2}}(1-i)$	$0.20(1+i)$	$0.3\sqrt{2\pi}(1-i)$
(c)	2000	1.5×10^7	26×10^{12}	$0.12(1-i)$	$0.20(1+i)$	$0.3\sqrt{2\pi}(1-i)$

dark matter observations, as illustrated in Fig. 7(b). The coupling required in this case is $y_1 = \frac{0.1}{\sqrt{2}}(1-i)$, whereas, Fig. 7(c) demonstrates that the dark matter asymmetry is quite sensitive to the value of m_{N_2} . At $m_{N_2} = \mathcal{O}(10^7)$ GeV, the effect of η_{DM} is negligible. However, suitable couplings provide the required baryon asymmetry, while being consistent with the dark matter observations for dark matter mass of 250 GeV.

Although leptogenesis depends on the parameters $y_1, y_2, y_3, m_\phi, m_\psi, m_{N_1}, m_{N_2}, m_{N_3}$, the resulting asymmetry η is particularly sensitive to just two of them: y_1 and m_{N_2} . Consequently, a broad range of (y_1, m_{N_2}) values can reproduce the observed η . This is explicitly shown in the scans over the parameter space as presented in Figs. 8–10.

The scans were carried out over the ranges $\frac{10^{-3}(1-i)}{\sqrt{2}} < y_1 < \frac{(1-i)}{\sqrt{2}}$ and $8 \times 10^4 \text{ GeV} < m_{N_2} < 10^9 \text{ GeV}$, while other parameters were fixed at values given in Table VI. For each point in this parameter space, we solved the Boltzmann Eqs. (31)–(33) and extracted the corresponding value of η at $z = 45$. The left panel in Fig. 8 shows η as a function of m_{N_2} with the color scale encoding $|y_1|$. The right panel in Fig. 8 shows η as a function of $|y_1|$, with the color scale encoding m_{N_2} . Both panels include horizontal lines marking the observed value of the BAU η_{obs} in 3σ limit. Corresponding to a large value of m_{N_2} , a combination with a small value of y_1 can provide the observed value of η . It is evident that $\eta \propto \frac{1}{m_{N_2}}$ with $\eta > 0$. This behavior

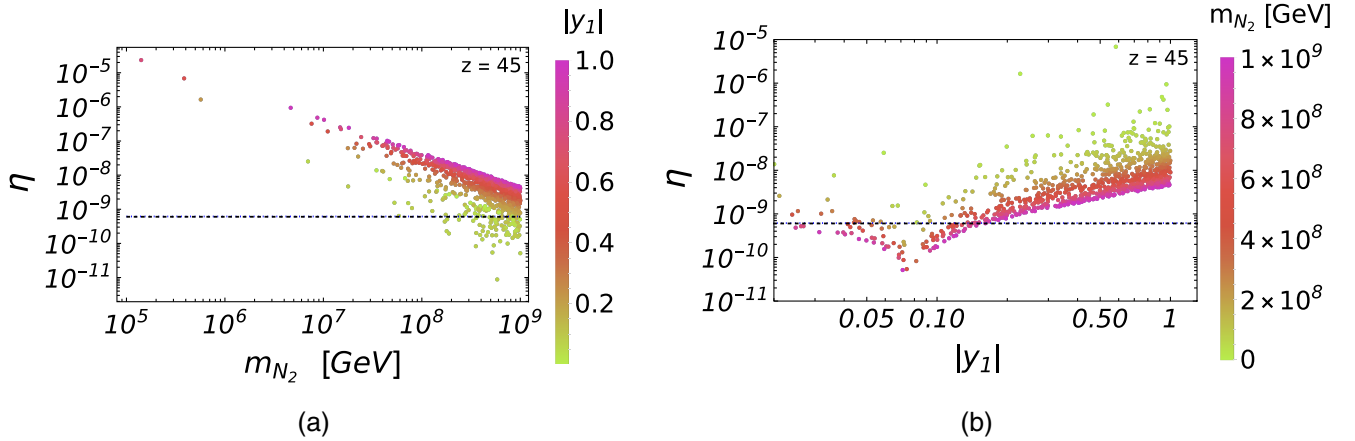


FIG. 8. Scan of the baryon asymmetry parameter $\eta[z = 45]$, (a) as a function of the right-handed neutrino mass m_{N_2} , with color variation representing the Yukawa coupling $|y_1|$, (b) as a function of $|y_1|$, with color variation representing m_{N_2} . The parameter values considered for the scan are given in Table VI.

TABLE VI. Parameter values considered for the scan results presented in Figs. 8–10.

m_{N_1} (GeV)	m_{N_3} (GeV)	m_ϕ (GeV)	m_ψ (GeV)	y_2	y_3
2000	26×10^{12}	250	250	$0.200781(1+i)$	$0.3\sqrt{2}\pi(1-i)$
	$8 \times 10^4 < m_{N_2} < 10^9$			$\frac{10^{-3}}{\sqrt{2}}(1-i) < y_1 < \frac{1}{\sqrt{2}}(1-i)$	

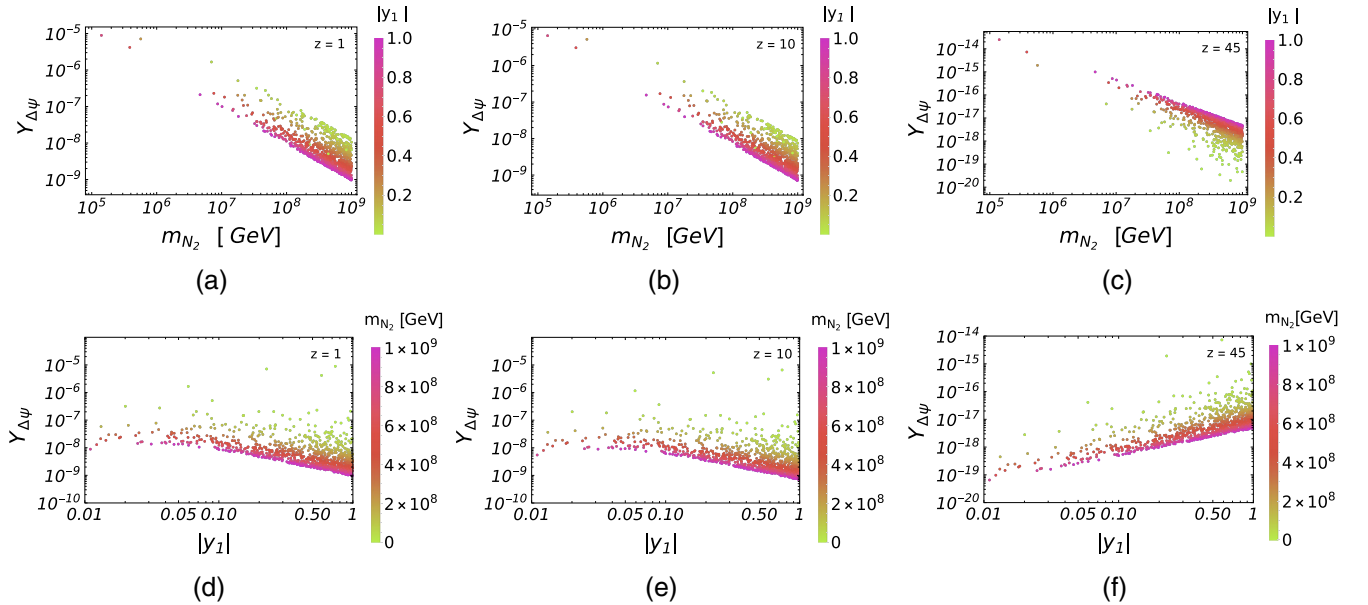
typically arises from the CP asymmetry parameters. For a large value of m_{N_2} , a smaller value of y_1 leads to a smaller value of η , whereas a larger value of y_1 results in a larger value of η . For small values of $|y_1|$, η remains suppressed, whereas it increases significantly at larger values of y_1 .

B. Numerical analysis of dark sector asymmetry

We investigate the numerical behavior of the dark sector asymmetry and its possible impact on the visible-sector asymmetry. Our analysis indicates that, for certain choices of parameters, the dark sector asymmetry can significantly affect the visible sector asymmetry. This is particularly so when the dark sector asymmetry is sizable at early times (small values of z). For other regions of parameter space, however, the visible sector asymmetry remains largely insensitive to the dark-sector contribution. To study this interplay, we compute the dark sector asymmetry $Y_{\Delta\psi}$ at three representative values of the evolution parameter z : $z = 1$, $z = 10$, and $z = 45$, as shown in Fig. 9. The parameter values used in this analysis are listed in Table VI. For each point in parameter space, we numerically solve Eq. (33) to obtain the corresponding values of $Y_{\Delta\psi}$ at the three chosen values of z . We then perform parameter scans of $Y_{\Delta\psi}$ as a function of m_{N_2} and $|y_1|$ at $z = 1$, $z = 10$, and $z = 45$, as illustrated in Fig. 9.

The dark-sector asymmetry at $z = 1$ and $z = 10$ is both sizable and nearly identical. In contrast, at later times ($z = 45$) the asymmetry becomes strongly suppressed. For all values of z , the asymmetry $Y_{\Delta\psi}$ increases as the mass of the right-handed neutrino m_{N_2} decreases. The dependence of $Y_{\Delta\psi}$ on the Yukawa coupling $|y_1|$, however, varies with z and exhibits distinct behaviors at different stages of the evolution. In this analysis, the mass of the lightest right-handed neutrino is fixed at $m_{N_1} = 2000$ GeV. When the mass m_{N_2} approaches m_{N_1} , the CP -asymmetry parameter ϵ_{DM} increases sharply, leading to an enhancement of the dark sector asymmetry. As m_{N_2} moves farther away from m_{N_1} , the dark sector asymmetry decreases steadily, as shown in subfigures (a), (b), and (e) of Fig. 9. An increase in the Yukawa coupling y_1 enhances both the decay rate of the process $N_1 \rightarrow \phi\psi$ and the associated scattering cross sections. As a result, the dark sector asymmetry initially grows because the parameter ϵ_{DM} itself increases with y_1 . Beyond a critical value of y_1 , however, washout processes become dominant and suppress the asymmetry. This leads to a characteristic turnover, after which the dark sector asymmetry decreases, as seen in subfigures (d), and (e) of Fig. 9.

Although the dark sector asymmetry can affect the visible sector asymmetry, as illustrated in Fig. 7 and further supported by the scatter plot in Fig. 10, the transfer of asymmetry from the dark sector to the visible sector is most

FIG. 9. Dark-sector asymmetry $Y_{\Delta\psi}$: (a)–(f) as described in the text, for the benchmark point given in Table VI.

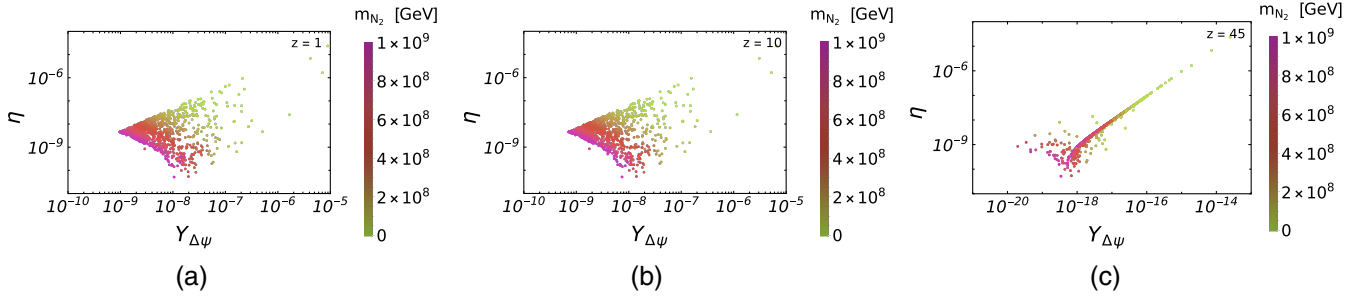


FIG. 10. Scan of the dark sector asymmetry as a function of $z = m_{N_1}/T$ for an alternative benchmark point given in Table VI, for three different values of z : (a) $z = 1$, (b) $z = 10$, and (c) $z = 45$.

efficient at early times, namely at $z = 1$ and $z = 10$. At later times ($z = 45$), the transfer becomes highly suppressed and has a negligible impact on the dark matter relic density and visible sector asymmetry. This is because dark matter decouples at a later epoch, by which time the number densities of ψ and $\bar{\psi}$ are approximately equal.

VI. CONCLUSION

In this work, we present an analysis of a simple extension of the SM that accommodates both a dark matter candidate and a viable mechanism for leptogenesis in the early Universe. The model introduces three Majorana right-handed neutrinos, a Z_2 -odd real scalar field ϕ , and a vectorlike neutral fermion ψ . In this setup, both ϕ and ψ are stable and serve as stable dark matter particles. To distinguish the scotino (ψ) from its antiparticle, we introduce a global $U(1)_{L'}$ corresponding to a generalized lepton number. The fermionic dark matter candidate ψ interacts with other fields only through its Yukawa couplings. In contrast, the scalar field ϕ couples through quartic scalar couplings as well. However, the quartic coupling involved must be extremely small in order to satisfy dark matter direct detection constraints. As a consequence, ϕ cannot serve as a viable single-component dark matter candidate in this scenario, when considered as a WIMP. Since ψ pair annihilation is prohibited when ϕ and N_1 are heavier, the only phenomenologically viable dark matter configuration in this model is a two-component fermion-scalar dark matter scenario, where the correct relic abundance is achieved predominantly through coannihilation channels. For simplicity, we have considered an exact mass degeneracy of ψ and ϕ in our study. The Yukawa coupling of ϕ , ψ interaction and the masses of a dark matter candidate, for the given masses of right handed neutrinos is constrained from the observed dark matter relic density. Minimum masses of the dark matter candidate depend on the masses of right handed neutrinos. For the lightest right-handed neutrino mass of 2000 GeV, the lightest dark matter mass can naturally lie around 200 GeV, while for $M_{N_1} = 1000$ GeV, this is slightly reduced to 187 GeV.

Canonical thermal leptogenesis with hierarchical right-handed neutrinos typically requires a very heavy right handed

neutrino in order to generate the observed baryon asymmetry of the Universe, unless at least two of them are nearly degenerate in mass. In our model, however, additional possibilities arise due to the Yukawa interactions between the right-handed neutrinos and the dark matter candidate, which contribute to the leptogenesis mechanism. In addition, the decay of N_1 into the dark matter candidate produces a primordial dark sector asymmetry. We have shown that this asymmetry can be efficiently transferred to the visible sector, thereby contributing to the observed baryon asymmetry. In the early Universe, the dark sector asymmetry is significantly larger and capable of sourcing the visible sector asymmetry through the transfer mechanism. However, in the present epoch, the residual dark sector asymmetry becomes negligible, ensuring that the dark matter relic density remains unaffected.

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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.

APPENDIX A: THE RELEVANT DECAY AND SCATTERING PROCESSES

The relevant processes to the Boltzmann equations (31), (32), and (33) of the yields, Y_{N_1} , $Y_L = Y_\ell - Y_{\bar{\ell}}$, $Y_{\Delta\psi} = (Y_\psi - Y_{\bar{\psi}})$, are listed as follows:

- (1) Decay and inverse decay: $N_1 \leftrightarrow \ell H$, $N_1 \leftrightarrow \bar{\ell} H^\dagger$,
 $N_1 \leftrightarrow \phi\psi$, $N_1 \leftrightarrow \phi\bar{\psi}$
- (2) $\Delta\ell = 1$ scattering processes:
 standard s -channel processes:
 $\ell N_1 \leftrightarrow d\bar{u}$, $\bar{\ell} N_1 \leftrightarrow \bar{d}u$,
 standard t -channel processes:
 $N_1 u \leftrightarrow d\bar{\ell}$, $N_1 \bar{u} \leftrightarrow \bar{d}\ell$,
 $N_1 d \leftrightarrow u\bar{\ell}$, $N_1 \bar{d} \leftrightarrow \bar{u}\ell$
 standard processes involving gauge boson

($V_\mu = W_\mu$ and B_μ):

$$N_1 H \xleftrightarrow{s} V_\mu \ell, N_1 H^\dagger \xleftrightarrow{s} V_\mu \bar{\ell},$$

$$N_1 \ell \xleftrightarrow{s} V_\mu H, N_1 \bar{\ell} \xleftrightarrow{s} V_\mu H^\dagger,$$

$$N_1 V_\mu \xleftrightarrow{t} H \ell, N_1 V_\mu \xleftrightarrow{t} H^\dagger \bar{\ell},$$

$$N_1 V_\mu \xleftrightarrow{u} H \ell, N_1 V_\mu \xleftrightarrow{u} H^\dagger \bar{\ell},$$

$$N_1 H^\dagger \xleftrightarrow{t} V_\mu \ell, N_1 H \xleftrightarrow{t} V_\mu \bar{\ell},$$

$$N_1 \ell \xleftrightarrow{t} V_\mu H^\dagger, N_1 \bar{\ell} \xleftrightarrow{t} V_\mu H,$$

Nonstandard scattering processes:

$$\ell H \xleftrightarrow{s} \phi \psi, \ell H \xleftrightarrow{s} \phi \bar{\psi},$$

$$\ell \phi \xleftrightarrow{t} H \psi, \bar{\ell} \psi \xleftrightarrow{t} H^\dagger \phi,$$

$$\bar{\ell} \phi \xleftrightarrow{t} H^\dagger \psi, H \phi \xleftrightarrow{t} \ell \psi,$$

$$\phi \psi \xleftrightarrow{s+t} \phi \bar{\psi}, \phi \bar{\psi} \xleftrightarrow{s+t} \phi \psi, \phi \phi \xleftrightarrow{t} \psi \bar{\psi}, \phi \phi \xleftrightarrow{t} \psi \psi,$$

$$\phi \phi \xleftrightarrow{t} \bar{\psi} \bar{\psi}; \text{ and}$$

$$(3) \Delta \ell = 2 \text{ scattering processes: } \ell \ell \xleftrightarrow{t+u} H^\dagger H^\dagger,$$

$$H \ell \xleftrightarrow{s+t} H^\dagger \bar{\ell}, \text{ and all the reverse reactions.}$$

As it is clear, some of the decay and scattering processes mentioned above are new, while others are studied in standard thermal leptogenesis, [65,67,68].

APPENDIX B: THERMAL AVERAGE OF DECAY WIDTH OF $X \rightarrow ab$

Consider a particle X that decays into particles a and b . Let Γ_X denote the decay width of the particle X . The thermally averaged decay width of X is given by

$$\begin{aligned} \gamma_X &= \int d\Pi_X e^{-E_X/T} \\ &\times \int d\Pi_a d\Pi_b (2\pi)^4 \delta^4(p_X - p_a - p_b) |M_{X \rightarrow ab}|^2 \\ &= n_X^{\text{eq}} \frac{K_1(z)}{K_2(z)} \Gamma_{X \rightarrow ab}, \end{aligned} \quad (\text{B1})$$

where n_X^{eq} is the equilibrium number density of X , and $\Gamma_{X \rightarrow ab}$ is the decay width of $X \rightarrow ab$ and $K_1(z)$ and $K_2(z)$ are the modified Bessel functions of the second kind of order one and two, respectively. Besides that the remaining details are the same as usual. The modified Bessel functions of the second kind of order one $K_1(z)$ and two $K_2(z)$ are as follows:

$$K_1(z) = \frac{1}{z} \int_z^\infty e^{-y} \sqrt{y^2 - z^2} dy, \quad \text{and} \quad (\text{B2})$$

$$K_2(z) = \frac{1}{z^2} \int_z^\infty y e^{-y} \sqrt{y^2 - z^2} dy, \quad (\text{B3})$$

respectively.

APPENDIX C: THERMAL AVERAGED CROSS SECTION OF $ab \rightarrow cd$ SCATTERING PROCESSES

The thermal average cross section γ for a $2 \rightarrow 2$ scattering process with a and b being incoming particles is given by

$$\gamma^{ab} = \frac{T}{64\pi^4} \int_{s_{\min}}^\infty ds \sqrt{s} \hat{\sigma}^{ab}(s) K_1\left(\frac{\sqrt{s}}{T}\right), \quad (\text{C1})$$

where,

(i) T is the temperature,

(ii) $\sqrt{s}$ is the center-of-mass energy,

(iii) K_1 is the modified Bessel function of the second kind of order one,

(iv) $\hat{\sigma}(s)$ is the *reduced cross-section*, defined as

$$\hat{\sigma}^{ab}(s) = 2s \sigma^{ab}(s) \lambda\left(1, \frac{m_a^2}{s}, \frac{m_b^2}{s}\right), \quad (\text{C2})$$

where $\sigma^{ab}(s)$ is the cross-section for the process, and

$$\lambda(x, y, z) = x^2 + y^2 + z^2 - 2(xy + yz + zx), \text{ and}$$

(v) m_a, m_b are the masses of the incoming particles.

The lower limit of integration, $s_{\min} = \max[(m_a + m_b)^2, (m_c + m_d)^2]$, ensures that the process is kinematically allowed.

Note. For scattering processes that violate CP asymmetry, the definition of the thermally averaged cross-section of the scattering process in Eq. (C1) slightly modifies to Eq. (C3). In this case, the reduced cross section $\hat{\sigma}$ is replaced with $\epsilon \hat{\sigma}$ (or $\hat{\sigma} \rightarrow \epsilon \hat{\sigma}$), where ϵ denotes the CP -asymmetry parameter. Finally, the thermally averaged cross section of scattering processes that violate CP asymmetry $\Theta_{\epsilon_a}^{ab}$ is given by

$$\Theta_{\epsilon_a}^{ab} = \frac{T}{64\pi^4} \int_{s_{\min}}^\infty ds \sqrt{s} \epsilon_a \hat{\sigma}^{ab}(s) K_1\left(\frac{\sqrt{s}}{T}\right), \quad (\text{C3})$$

where other details remain the same as in Eq. (C1).

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