

Hunting for heavy Z' with IceCube neutrinos and gravitational waves

Basabendu Barman^{1,*}, Arindam Das^{2,3,†}, Suruj Jyoti Das^{4,‡} and Marco Merchand^{5,6,§}

¹*Department of Physics, School of Engineering and Sciences,
SRM University-AP, Amaravati 522240, India*

²*Institute for the Advancement of Higher Education, Hokkaido University, Sapporo 060-0817, Japan*

³*Department of Physics, Hokkaido University, Sapporo 060-0810, Japan*

⁴*Particle Theory and Cosmology Group, Center for Theoretical Physics of the Universe,
Institute for Basic Science (IBS), Daejeon, 34126, Korea*

⁵*KTH Royal Institute of Technology, Department of Physics, SE-10691 Stockholm, Sweden*

⁶*The Oskar Klein Centre for Cosmoparticle Physics, AlbaNova University Centre,
SE-10691 Stockholm, Sweden*



(Received 11 March 2025; accepted 1 August 2025; published 25 August 2025)

In the minimal gauged $B - L$ extension of the Standard Model, we demonstrate that PeV-scale dark matter (DM) and the baryon asymmetry of the Universe (BAU) can be simultaneously explained through the three right-handed neutrinos (RHNs) present in the theory. The DM candidate undergoes rare decay into light neutrinos, providing an explanation for the observed IceCube events, while the other two RHNs generate the BAU via leptogenesis. The breaking of gauge symmetry gives rise to detectable gravitational waves (GWs) from decaying cosmic strings (CS), making this framework testable at several future GW detectors—despite being beyond the reach of conventional collider experiments due to the extremely weak gauge coupling. The symmetry-breaking scale establishes a connection between particle masses, couplings, and the GW spectrum, offering a unified and predictive scenario.

DOI: [10.1103/n6jd-ltd5](https://doi.org/10.1103/n6jd-ltd5)

I. INTRODUCTION

Neutrino oscillation experiments have provided evidence for tiny neutrino masses and flavor mixing over time [1]. The PLANCK [2] observations estimate the baryon asymmetry of the universe (BAU), defined as the ratio of the baryon-antibaryon number density difference ($n_B - n_{\bar{B}}$) to the entropy density (s), to be $Y_B^o \simeq 8.75 \times 10^{-11}$ [2]. Furthermore, astrophysical and cosmological observations strongly support the existence of dark matter (DM), with its relic abundance measured as $\Omega h^2 \simeq 0.12$ [2–4]. These phenomena remain unexplained within the Standard Model (SM), necessitating exploration of beyond Standard Model (BSM) frameworks to uncover their fundamental origins.

In the recent past, the IceCube neutrino observatory reported the detection of three PeV neutrinos, roughly 3σ excess above the expected background rates [5–10]. These highest-energy events correspond to deposited energies of 1.04 PeV, 1.14 PeV and 2.0 PeV, respectively. Although the origin of these very high energy events is still unclear, it has been shown that such events could be originated from decays of superheavy DM [11–17]. The neutrino energy spectrum presents a high-energy cutoff at half of the DM mass [12,13] if two body decays including one neutrino are present. Moreover, the IceCube spectrum sets a lower bound on the DM lifetime $\tau_{\text{DM}} \simeq \mathcal{O}(10^{28})$ s [13,18], which is largely model-independent and significantly exceeds the age of the Universe.

On the other hand, current observations of gravitational waves (GWs) [19–25], have opened a complementary avenue to test BSM physics. For example, if symmetries are broken spontaneously at very high temperature then topological defects, such as cosmic strings and domain walls, may appear in the early stages of the Universe [26,27] and the system of these defects can be considered as a prominent source of a GW background, while the scale of symmetry breaking can be associated with the scale of new physics.

To explore these aspects within a simple yet elegant framework, we consider an extension of the SM featuring

*Contact author: basabendu.b@srmmap.edu.in

†Contact author: arindamdas@oia.hokudai.ac.jp

‡Contact author: surujjd@gmail.com

§Contact author: marcomm@kth.se

Published by the American Physical Society under the terms of the [Creative Commons Attribution 4.0 International](https://creativecommons.org/licenses/by/4.0/) license. Further distribution of this work must maintain attribution to the author(s) and the published article's title, journal citation, and DOI. Funded by SCOAP³.

an anomaly-free $U(1)_{B-L}$ gauge group [28,29], incorporating three generations of SM-singlet right-handed neutrinos (RHNs) and an SM-singlet scalar which acquires a nonzero vacuum expectation value (VEV) resulting in the breaking of the $U(1)_{B-L}$ symmetry. As a result, the Majorana masses for the RHNs are generated, which in turn induce tiny masses and flavor mixing for the observed left-handed neutrinos via the seesaw mechanism [30–34]. This setup naturally explains the BAU via vanilla leptogenesis if at least two RHN generations contribute to the asymmetry generation, while the third, lightest RHN can serve as a PeV-scale decaying DM candidate, addressing IceCube high-energy neutrino events through the freeze-in mechanism. Additionally, the $U(1)_{B-L}$ breaking gives rise to one-dimensional topological defects—cosmic strings (CS)—characterized by a string tension $G\mu \sim BGv_\phi^2$, where v_ϕ is the VEV of the $U(1)_{B-L}$ symmetry breaking singlet scalar and $B \sim 0.1$ [35–37]. This minimal scenario, therefore, offers a unified explanation for (i) PeV-scale decaying DM linked to IceCube events, (ii) baryogenesis via leptogenesis and (iii) GW from CS which could be tested at future GW detectors. We thus constrain heavy neutral gauge boson mass beyond 1 TeV, with tiny gauge coupling ($\lesssim \mathcal{O}(10^{-5})$), which lies beyond the reach of high energy collider experiments, that typically provide constraints on gauge couplings around $\mathcal{O}(10^{-2})$ for the B – L scenario at the LHC [38]. At this stage it is worth pointing out that the IceCube events may also be accounted for by the high-energy neutrino flux originating from evaporating primordial black holes (PBHs) [39–44]. The central focus of the present work, however, is to establish a connection with the mechanism responsible for generating light neutrino masses—an aspect that unequivocally requires physics beyond the Standard Model (BSM). Rather than direct neutrino emission from PBHs, one could also consider scenarios where PBHs emit heavy BSM particles (such as heavy leptons), which subsequently decay into neutrinos. This indirect production channel has been widely explored in the literature and necessitates the inclusion of BSM fields such as right handed neutrinos (see, for example, Refs. [45–48]). In contrast, our approach does not rely on PBHs as the source, but considers the thermal SM radiation bath to be responsible for producing heavy neutral leptons (RHNs in present case), which ultimately lead to the observed neutrino signatures.

II. THE FRAMEWORK

Under the $SM \otimes U(1)_{B-L}$ gauge symmetry, the SM quark fields transform as $q_L^i = \{3, 2, \frac{1}{6}, \frac{1}{3}\}$, $u_R^i = \{3, 1, \frac{2}{3}, \frac{1}{3}\}$, $d_R^i = \{3, 1, -\frac{1}{3}, \frac{1}{3}\}$, respectively. The SM lepton fields transform as $\ell_L^i = \{1, 2, -\frac{1}{2}, -1\}$, $e_R^i = \{1, 1, -1, -1\}$, respectively, while the SM Higgs field $H = \{1, 2, \frac{1}{2}, 0\}$. We introduce three SM-singlet RHNs to cancel gauge and mixed gauge-gravity anomalies which transform as

$N_R^i = \{1, 1, 0, -1\}$ with $i = 1, 2, 3$ and one SM-singlet $U(1)_{B-L}$ scalar which transforms as $\Phi = \{1, 1, 0, 2\}$. The relevant Yukawa interactions read,

$$\mathcal{L} \supset -Y_{\nu_{\alpha\beta}} \bar{\ell}_L^\alpha \tilde{H} N_R^\beta - \frac{1}{2} Y_{N_\alpha} \Phi (\overline{N_R^\alpha})^c N_R^\alpha + \text{H.c.}, \quad (1)$$

where we start-off in a basis where the Y_{N_α} matrix is diagonal, and $\tilde{H} = i\tau^2 H^*$ with τ^2 being the second Pauli matrix. The scalar potential involving two scalar fields is given by

$$V = \sum_{\mathcal{I}=H,\Phi} [m_{\mathcal{I}}^2 (\mathcal{I}^\dagger \mathcal{I}) + \lambda_{\mathcal{I}} (\mathcal{I}^\dagger \mathcal{I})^2] + \lambda_{\text{mix}} (H^\dagger H) (\Phi^\dagger \Phi). \quad (2)$$

After the breaking of B – L and electroweak gauge symmetries, the scalar fields H and Φ develop their VEVs as

$$\langle H \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} v + h \\ 0 \end{pmatrix}, \quad \text{and} \quad \langle \Phi \rangle = \frac{v_\Phi + \phi}{\sqrt{2}}, \quad (3)$$

where electroweak scale is $v = 246$ GeV at the potential minimum. For $v_\Phi \gg v$, the mass of the B – L gauge boson can be written as $M_{Z'} = 2g_X v_\Phi$. The breaking of B – L symmetry induces the Majorana mass term for the RHNs while the electroweak symmetry breaking generates the Dirac mass term for the light left-handed neutrinos from Eq. (1) as

$$M_\alpha = \frac{Y_{N_\alpha}}{\sqrt{2}} v_\Phi, \quad m_{D_{\alpha\beta}} = \frac{Y_{\nu_{\alpha\beta}}}{\sqrt{2}} v. \quad (4)$$

From the mass matrices above, the light active neutrino masses can be derived using the standard seesaw formula $-m_D M_\alpha^{-1} m_D^T$ [32,49,50]. This mechanism successfully explains the tiny neutrino masses and their flavor mixing. Out of three generations of RHNs, we identify N_1 as a long-lived decaying DM. Its only decay channel, $N_1 \rightarrow \ell H$, produces boosted high-energy neutrinos. The DM lifetime, as required by the IceCube observations, in terms of its mass and Yukawa coupling reads [12,13,18,51]

$$\tau \simeq 10^{28} \text{ s} \left(\frac{(Y_\nu)_{11}}{2 \times 10^{-29}} \right)^2 \left(\frac{M_1}{4 \text{ PeV}} \right). \quad (5)$$

Due to extremely small Yukawa coupling strength, it is not possible to address right DM abundance via the inverse decay channel $\ell H \rightarrow N_1$. Thus, we consider its production via freeze-in [52,53] from the thermal bath. Consequently, production and decay of N_1 is disentangled, and hence right DM abundance can be satisfied even with a DM decay lifetime of $\tau_{N_1} \sim 10^{28}$ sec, complying with the IceCube data. Here onward we will consider the following set of parameters as independent in our analysis: $\{g_X, M_{Z'}, Y_{N_\alpha}\}$, while the DM mass ($M_1 = Y_{N_1} v_\Phi / \sqrt{2}$) is always fixed at

4 PeV. Note that, the long-lived RHN N_1 , which is a potential DM candidate, contributes negligibly to the active neutrino masses because of the tiny Yukawa coupling $(Y_\nu)_{11} \simeq \mathcal{O}(10^{-30})$ [cf. Eq. (5)]. As a result, the lightest active neutrino remains (almost) massless. This scenario is fully compatible with current neutrino oscillation data, which constrain only the mass-squared differences and not the absolute mass scale. This also provides a complementary prospect in experiments sensitive to the absolute neutrino mass scale. While such a tiny neutrino mass is out of reach of ongoing tritium beta decay experiments like KATRIN [54], future observation of neutrinoless double beta decay can falsify our scenario [55], particularly for normal ordering.

III. DM PRODUCTION VIA FREEZE-IN

The PeV scale N_1 DM can be produced via (i) on-shell decay of Z' , provided $M_{Z'} > 2M_1$, (ii) on-shell decay of ϕ , if $m_\phi > 2M_1$ and (iii) 2-to-2 scattering off of the bath particles, mediated by Z' . Notice that we consider the limit of zero mixing¹ between the extra scalar and the SM Higgs thus we do not include any ϕ mediated scattering in our analysis.² The universal interaction strength g_X must be feeble to ensure nonthermal DM production via freeze-in. As a result, the Z' cannot reach thermal equilibrium, and its comoving number density must be determined by solving a set of coupled Boltzmann equations (BEQs) are given in Appendix B. The first line of Eq. (B1) corresponds to the yield of ϕ , with Γ_ϕ being its decay rate as reported in Appendix A. The second line of Eq. (B1) takes care of the Z' yield, sourced from the decay of $\phi \rightarrow Z'Z'$. Finally, in the third line, we provide the evolution equation for DM yield, sourced from Z' -decay, ϕ -decay and Z' -mediated scattering channels. Since the mixing between ϕ and the SM Higgs is considered to be negligibly small, its decay into SM final states are extremely suppressed. Consequently ϕ decays dominantly into RHNs and Z' pairs. To fit the observed DM relic density, it is required that $Y_0 M_1 = \Omega h^2 \frac{1}{s_0} \frac{\rho_c}{h^2} \simeq 4.3 \times 10^{-10}$ GeV, where $Y_0 \equiv y_{N_1}(z \rightarrow \infty)$ is the present DM yield. We use the critical energy density $\rho_c \simeq 1.05 \times 10^{-5} h^2$ GeV/cm³, present entropy density $s_0 \simeq 2.69 \times 10^3$ cm⁻³ [65] and DM relic abundance $\Omega h^2 \simeq 0.12$, with $h \simeq H_0/100$ (km/s/Mpc) being the reduced Hubble rate, where $H_0 \simeq 67.4 \pm 0.5$ km/s/Mpc is the current Hubble rate [2].

Before moving on, let us emphasize that for the freeze-in to be valid, the DM has to be out of chemical equilibrium

¹Current bounds from LHC [56–58], LEP [59] prospective colliders like ILC [60] and CLIC [61] suggests that the scalar mixing requires to be < 0.001 , for ϕ mass up to 1 TeV. Since in our case $m_\phi \gg 1$ TeV, therefore we consider the mixing to be vanishingly small.

²Some dedicated studies of scalar effects on the DM production via freeze-in include [62–64].

with the SM bath. Now, the DM production rate from Z' decay is given by the production rate densities as

$$\gamma_D = \frac{g_a}{2\pi^2} m_a^2 \Gamma_{Z' \rightarrow N_1 N_1} T K_1 \left(\frac{m_a}{T} \right), \quad (6)$$

while production from 2-to-2 scattering off of the bath particles reads

$$\begin{aligned} \gamma_s = & \frac{T}{32\pi^4} g_a g_b \times \int_{\max[(m_a+m_b)^2, 4M_1^2]}^{\infty} ds \\ & \times \frac{[(s - m_a^2 - m_b^2)^2 - 4m_a^2 m_b^2]}{\sqrt{s}} \sigma(s)_{a,b \rightarrow N_1 N_1} K_1 \left(\frac{\sqrt{s}}{T} \right). \end{aligned} \quad (7)$$

Here, the subindices a, b correspond to the SM states, $g_{a,b}$ are the corresponding degrees of freedom and K_i denotes the modified Bessel functions of i th kind. In order to ensure out of equilibrium DM production, the following condition needs to be satisfied

$$\mathcal{R}_{D(s)} \equiv \frac{\gamma_{D(s)}}{n_{\text{eq}}^{\frac{1}{N_1}} H} < 1, \quad (8)$$

where H is the Hubble parameter during radiation domination and $n_{\text{eq}} = (T/(2\pi^2))m^2 K_2(m/T)$ is the equilibrium number density. Note that, here we use the equilibrium number density for the DM itself in order to obtain the conservative bound on the masses and couplings. In Fig. 1 we show the evolution of DM production rate, as a function of the temperature of the thermal bath. Due to g_X^2 dependence, it is not possible to have $g_X \gtrsim 10^{-4}$ in order to ensure $\mathcal{R} < 1$ till $T \simeq M_{Z'}$ in the case of DM genesis via decay. However, for scattering it is still possible to have g_X as large as about 10^{-3} , while ensuring $\mathcal{R} < 1$ because of the g_X^4 dependence of the scattering cross section. Note that, in this case the DM production starts at $T \sim M_1$, as prior to that the thermal bath is not energetic enough to produce such a massive DM.

The evolution of DM and Z' yields ($y_{Z'}$) are shown in the left panel of Fig. 2, as a function of $z = M_1/T$, with $M_1 = 4$ PeV. For the red curves, we fixed $m_\phi/2 > M_{Z'} > 2M_{N_1}$, such that the DM can be produced from the on-shell decay of Z' , as well as via Z' mediated scattering. Now, in presence of both channels, DM production via decay dominates since the production rate is proportional to g_X^2 , compared to scattering, which is proportional to g_X^4 . The DM yield y_{N_1} gradually grows with time, finally reaching the observed abundance in the asymptotic limit. In the first plateau region of the red solid curve, DM is dominantly produced from ϕ -decay. Due to large decay width, ϕ decay is completed earlier and the DM (as well as Z') production from ϕ -decay stops. However, Z' decay still goes on. Thus, the final DM abundance is set by $Z' \rightarrow N_1 N_1$ channel. On the other hand, the blue solid curve represents a scenario where the DM production is

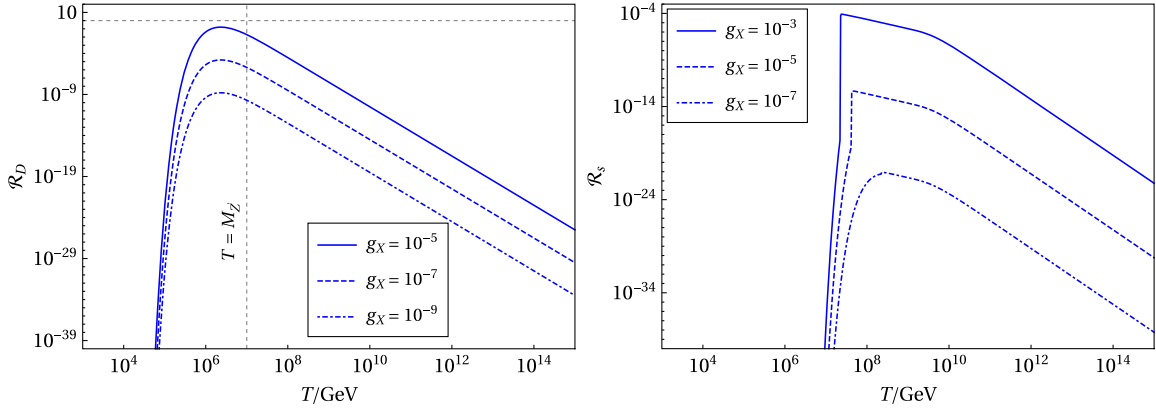


FIG. 1. Left: DM production rate from Z' decay, as a function of the bath temperature T , for different choices of g_X . Here we have fixed $M_{Z'} = 10^7$ GeV. Right: same as left, but considering DM production from 2-to-2 scattering, mediated by Z' , where $M_{Z'} = 10^5$ GeV. In all cases the DM mass is fixed at 4 PeV.

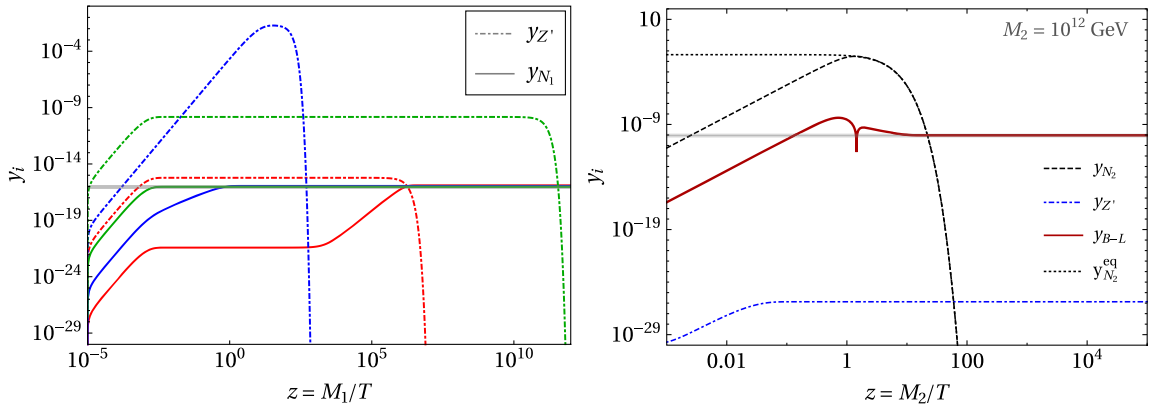


FIG. 2. In all cases DM mass is fixed to $M_{N_1} = 4$ PeV. Left: evolution of DM and Z' yields in solid and dashed curves, respectively, as function of $z = M_1/T$. The red, blue and green colors correspond to $M_{Z'} = \{10^8, 10^5, 10^3\}$ GeV with the corresponding $g_X = \{2 \times 10^{-12}, 1.2 \times 10^{-6}, 10^{-14}\}$, respectively, to produce the right DM abundance (shown via gray thick straight line). For the red and green curves we fix $m_\phi = 10^{10}$ GeV, while for the blue curve $m_\phi = 10^6$ GeV. Right: evolution of B – L yield (red solid) with M_2/T , for N_2 leptogenesis. The dark red curve produces the observed baryon asymmetry (shown via gray thick straight line). We show for comparison the equilibrium N_2 yield $y_{N_2}^{\text{eq}}$ in black dotted curve.

instead controlled by Z' mediated scattering, as the decay is kinematically forbidden when $M_{Z'} < M_1$, also $m_\phi < 2 M_1$, but $m_\phi > 2 M_{Z'}$ such that Z' 's are produced from on-shell ϕ -decay. For $M_{Z'} = 10^8(10^5)$ GeV and for DM mass of 4 PeV, the freeze-in occurs at the temperature $T \sim M_{Z'}(M_1)$. This is the typical IR feature [53], where the DM yield saturates when $T = \max[M_{Z'}, M_1]$. Finally, for $M_{Z'}/2 < M_1 \ll m_\phi$, the DM is dominantly produced from ϕ -decay (green solid curve), due to which, once again, a very tiny g_X is required since this decay rate is also proportional to g_X^2 .

IV. BARYON ASYMMETRY FROM N_2 LEPTOGENESIS

The observed BAU is generated via vanilla leptogenesis from the CP -violating, out of equilibrium decays of N_2 .

Due to very long lifetime, the lepton asymmetry generated through N_1 decay does not contribute to the generation of the BAU as its decay takes place far below the sphaleron equilibration temperature. The CP asymmetry from N_2 -decay can be expressed as [66,67]

$$\epsilon_{\Delta L} \simeq \frac{3\delta_{\text{eff}}}{16\pi} \frac{M_2 m_{\nu, \text{max}}}{v^2}, \quad (9)$$

where δ_{eff} is the effective CP violating phase in the neutrino mass matrix with $0 \leq \delta_{\text{eff}} \leq 1$ (see Appendix C) and we take $m_{\nu, \text{max}} = 0.05$ eV as the heaviest light neutrino mass following the normal hierarchy. The final B – L asymmetry is obtained by solving a set of coupled BEQs given in Appendix D. Once again, the first line of the set of equations above takes into account the production of Z' from the thermal bath, via the decay of ϕ and the decay of

Z' into several final states. The second line corresponds to the evolution of N_2 yield. The produced N_2 's then undergo CP violating out of equilibrium decay producing a net $B - L$ asymmetry. The evolution of this asymmetry is tracked by the last line of the equation, where we also include the contribution of inverse decays to the wash-out term.

The sphaleron interactions [66] are in equilibrium within the temperature range 100 GeV to 10^{12} GeV, and they convert a fraction of a nonzero $B - L$ asymmetry into a baryon asymmetry via

$$Y_B \simeq a_{\text{sph}} y_{B-L} = \frac{8N_F + 4N_H}{22N_F + 13N_H} y_{B-L}, \quad (10)$$

where $N_F (= 3)$ and $N_H (= 1)$ are numbers of fermion generations and Higgs doublets, $a_{\text{sph}} \simeq 28/79$ and $Y_B^o \simeq 8.75 \times 10^{-11}$. We show the evolution of yields for RHN (y_{N_2}) and $B - L$ (y_{B-L}) as a function of $z = M_2/T$ in the right panel of Fig. 2. The population of N_2 receives contributions from two sources: (i) the thermal bath, through the inverse decay process $\ell H \rightarrow N_2$ and (ii) decay of Z' and ϕ . Notably, the bath contribution is the dominant one, as all other production channels are suppressed due to tiny g_X to satisfy freeze-in conditions. Following the Davidson-Ibarra (DI) bound [68], $M_2 \gtrsim 10^9$ GeV is required for successful thermal leptogenesis. Since the RHN mass depends on v_ϕ from Eq. (4), in order to satisfy DI bound, $v_\phi \gtrsim 10^9$ GeV for $Y_{N_2} \sim \mathcal{O}(1)$.

V. RUNNING OF THE COUPLINGS

To solve the renormalization group equations (RGEs) and track the running of the couplings, we set initial conditions at the top quark mass scale $\mu = m_t = 172.4$ GeV. The SM

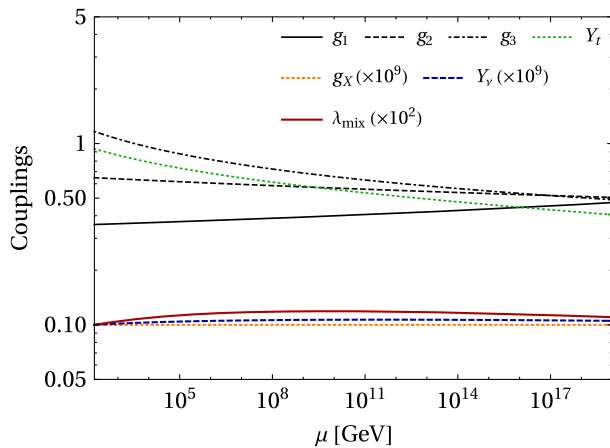


FIG. 3. Evolution of different couplings appearing in the present framework with the running scale $\mu: [m_t, M_P]$, with the top quark mass $m_t = 173.2$ GeV. Note that, g_X (shown in orange dot), Y_ν (shown in blue dashed), λ_{mix} (shown in red solid) are magnified.

parameters, including the top Yukawa coupling Y_t , are taken from [69], while we fix the new physics couplings as $g_X = Y_\nu = 10^{-10}$. For the scalar quartic couplings, we choose $\lambda_H = 0.12$, $\lambda_\phi = 10^{-2}$, and $\lambda_{\text{mix}} = 10^{-3}$. With these initial values, we evolve the couplings up to the Planck scale, $\mu = M_P$, considering the set of equations in Appendix E. Given that $\beta_X \sim g_X^3$, the new gauge coupling g_X remains nearly constant due to its tiny initial value, as shown by the orange dashed curve in Fig. 3. The same applies to the Yukawa coupling Y_ν , illustrated by the blue dashed curve. The scalar mixing coupling λ_{mix} also stays approximately unchanged (red solid curve), primarily because of its small initial value constrained by collider bounds. Its slight downward bend near $\mu \rightarrow M_P$ arises from the negative contribution of the term $-\frac{9}{2}g_2^2\lambda_{\text{mix}}$ in its β -function.

VI. GW SPECTRUM FROM COSMIC STRINGS

Numerical simulations based on the Nambu-Goto action [70,71] have found that for gauged symmetry, the dominant channel of energy loss from CS is through GW radiation from oscillating loops.³ The rate of energy loss or the power of GW emission is given by [80], $P_{\text{GW}} = \frac{G}{5} \langle \ddot{Q} \rangle^2$, where Q is the quadrupole moment of the oscillating loop, and the triple time derivative $\ddot{Q} \propto \mu$. Thus, the rate of energy loss can be written as $\frac{dE}{dt} = -\Gamma G \mu^2$, where $\Gamma \approx 50$ [81]. Due to GW emission, the loop starts to shrink from its initial length $l_i = \alpha t_i$ at the time of formation (t_i) as $l(t) = \alpha t_i - \Gamma G \mu (t - t_i)$, where α is the loop size parameter and considered to be $\alpha \approx 0.1$ from simulation studies [82,83]. The total energy loss from a loop constitutes a set of normal mode oscillations with frequencies $f_k = 2k/l(t)$, where k represents number of modes ($k = 1, 2, 3 \dots \infty$). The GW density parameter is defined as

$$\Omega_{\text{GW}}(t_0, f) = \frac{f}{\rho_c} \frac{d\rho_{\text{GW}}(t_0, f)}{df} = \sum_k \Omega_{\text{GW}}^{(k)}(t_0, f), \quad (11)$$

where f and t_0 represent the current frequency and the present time, while $\rho_c = 3M_P^2 H_0^2$ is the critical energy density where M_P is reduced Planck mass and Hubble parameter. Since the GW energy density redshifts as a^{-4} , we have [82]

³The present framework might also give rise to primordial GW signal via graviton Bremsstrahlung through decay [72–75], considering *minimal* gravitational interactions. In the context of leptogenesis this has been explored in [76,77]. However, such decays are suppressed by the Planck mass. The amplitude of Bremsstrahlung GW signal is also suppressed for the same reason, and their spectral peak typically lies in the GHz range—accessible only to specialized detectors, such as resonant cavities. If one introduces a nonminimal coupling between the DM and the Ricci scalar, then the DM can also decay via the nonminimal gravity portal [78,79], however we do not introduce such coupling in this analysis.

$$\frac{d\rho_{\text{GW}}^{(k)}}{df} = \int_{t_F}^{t_0} \left[\frac{a(t_E)}{a(t_0)} \right]^4 P_{\text{GW}}(t_E, f_k) \frac{dF}{df} dt_E, \quad (12)$$

where f_k denotes emitted frequency (f_E) at the time t_E , t_F is the loop formation time, $\frac{dF}{df} = f \left[\frac{a(t_0)}{a(t_E)} \right]$ accounts for the redshift of the frequency, and

$$P_{\text{GW}}(t_E, f_k) = \frac{2kG\mu^2\Gamma_k}{f \left[\frac{a(t_0)}{a(t_E)} \right]^2} n \left(t_E, \frac{2k}{f} \left[\frac{a(t_E)}{a(t_0)} \right] \right), \quad (13)$$

is the power emitted by the loops.

The GW spectrum depends on the nature of small-scale structure in the loops which can appear in the form of cusps or kinks [84,85]. For our study, we consider cusplike structures to dominate the GW spectra. Here we have $\Gamma_k = \sum_{m=1}^{\infty} \frac{\Gamma_k^{-4/3}}{m^{-4/3}}$, with $\sum_k \Gamma_k = \Gamma$ and n denotes the number density of loops which, in a cosmological background with scale factor $a \propto t^\beta$, is found from the velocity dependent one scale (VOS) model [86–88] and numerical simulations to be [82]

$$n(t_E, l_k(t_E)) = \frac{A_\beta}{\alpha} \frac{(\alpha + \Gamma G\mu)^{3(1-\beta)}}{[l_k(t_E) + \Gamma G\mu t_E]^{4-3\beta} t_E^{3\beta}}, \quad (14)$$

where A_β is a constant depending on the cosmological background. Using Eqs. (11)–(14), we obtain the current GW energy density for the mode k as

$$\Omega_{\text{GW}}^{(k)}(t_0, f) = \frac{2kG\mu^2\Gamma_k}{f\rho_c} \int_{t_{\text{osc}}}^{t_0} dt \left[\frac{a(t)}{a(t_0)} \right]^5 n(t, l_k), \quad (15)$$

where the integration over time is from the moment t_{osc} when loops start oscillating after damping due to thermal friction [89] and hence producing subdominant effect. For the case of loops formed and radiated during radiation domination, the GW spectrum has a typical flat plateau, with the amplitude given by

$$\Omega_{\text{GW}}^{(k=1), \text{plateau}}(f) = \frac{128\pi G\mu A_r}{9\zeta(4/3) \epsilon_r} \Omega_r [(1 + \epsilon_r)^{3/2} - 1], \quad (16)$$

where $\epsilon_r = \alpha/\Gamma G\mu$, and $A_r = 0.54$ [88] for radiation domination. CMB measurements require $G\mu \lesssim 10^{-7}$ [90] and hence we have $\alpha \gg \Gamma G\mu$. This gives $\Omega_{\text{GW}}^{(k=1)}(f) \propto v_\phi$, and hence a higher symmetry breaking scale is more likely to be probed by the GW detectors. Recent results from NANOGrav [22] puts a strong upper bound on $G\mu \lesssim 10^{-10}$. However, this bound is found to be relaxed, considering long-lived loops in a wider class of models [91], with the upper bound close to the CMB limit $G\mu \lesssim 10^{-7}$ [90,92].

The viable parameter space is furnished in Fig. 4, where the thick-solid red and blue contours correspond to the observed DM abundance for $M_{N_1} = 4$ PeV. The diagonal dashed lines denote different v_ϕ providing the correct value

of $M_{Z'}$ for the corresponding g_X . A few of such v_ϕ 's are within the reach of proposed GW detectors: Big Bang Observer (BBO) [93,94], ultimate DECIGO (uDECIGO) [95,96], LISA [97], the cosmic explorer (CE) [98] and the Einstein Telescope (ET) [99–102]. The gray shaded region in top left corner is disallowed by the DI bound on the RHN mass for $Y_{N_2} \lesssim \sqrt{4\pi}$. The light gray shaded region in the bottom left corner corresponds to the combination of g_X and $M_{Z'}$ that gives rise to $\tau_{Z'} = 1/\Gamma_{Z'} > 1$ sec, thereby jeopardizing the big bang nucleosynthesis (BBN) predictions. The dark gray shaded region in the bottom right corner demands (super-)Planckian v_ϕ . As the $Z' \rightarrow N_1 N_1$ decay channel opens up, one needs $g_X \lesssim 10^{-10}$ to satisfy the nonthermal freeze-in condition in order to produce the correct relic abundance. As a result, the thick-solid red and blue lines take a sharp bend at the mass threshold. Along the horizontal part of the red contour, we choose $m_\phi = 10^6$ GeV, such that DM is produced solely from Z' -mediated scattering, while Z' is produced from ϕ decay. In this case, the final DM yield approximately reads $Y(T_{\text{FI}}) \sim \sigma(T_{\text{FI}}) M_P T_{\text{FI}} \sim g_X^4 M_P / M_1$, where $\sigma(T_{\text{FI}}) \sim g_X^4 / T_{\text{FI}}^2$, with freeze-in temperature $T_{\text{FI}} \sim M_1$. Consequently, the DM abundance becomes almost independent of the mediator mass. For the thick-solid blue contour, $m_\phi = 10^{10}$ GeV $> 2 M_1$, as well as $m_\phi > 2 M_{Z'}$ (and $M_{Z'} > 2 M_1$). As a result, the DM production is decay-dominated, demanding g_X to be extremely small, following Fig. 2.

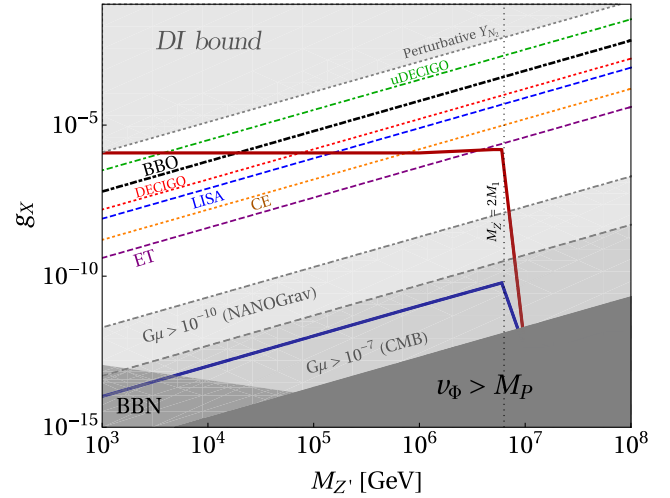


FIG. 4. Summary of parameter space: The red and blue thick contours correspond to right DM abundance by fixing $M_1 = 4$ PeV. Along the red contour we consider DM production both via scattering and decay, while along the blue contour, only production via decay is considered (see text). The diagonal contours correspond to different v_ϕ 's, some of which fall within the sensitivity of a few GW detectors, as denoted. The shaded regions are disallowed from Davidson-Ibarra (DI) bound ($M_2 < 10^9$ GeV), BBN bound on Z' -lifetime and by having super-Planckian VEV-values.

VII. CONCLUSIONS

We have proposed a scenario that simultaneously accounts for both DM and the BAU while remaining testable at IceCube as well as at several proposed GW detectors. Notably, our approach is broadly applicable to the general class of $U(1)_X$ models. A key element of our framework is the symmetry breaking scale, which unifies all BSM masses and couplings, offering a coherent and compact picture, as illustrated in Fig. 4. This provides bounds on extremely small g_X along with $M_{Z'} \gtrsim 1$ TeV, both lying beyond the reach of standard collider experiments. This work underscores the essential role of multimessenger astronomy in probing feebly coupled new physics.

ACKNOWLEDGMENTS

The work of S.J.D. was supported by IBS under the project code, IBS-R018-D1. We thank Osamu Seto and Rinku Maji for useful communications.

DATA AVAILABILITY

The data are not publicly available. The data are available from the authors upon reasonable request.

APPENDIX A: RELEVANT DECAY WIDTHS AND CROSS SECTIONS

The partial decay width of Z' into a pair of single generation SM fermions and DM reads,

$$\begin{aligned}\Gamma_{Z' \rightarrow f_i \bar{f}_i} &= N_c \frac{g_X^2 M_{Z'}}{24\pi} [(Q_L^2 + Q_R^2)], \\ \Gamma_{Z' \rightarrow N_1 N_1} &= \frac{g_X^2 M_{Z'}}{24\pi} \left(1 - \frac{4M_1^2}{M_{Z'}^2}\right)^{\frac{3}{2}},\end{aligned}\quad (\text{A1})$$

where $N_c = 1(3)$ is the color factor for leptons (quarks) and $Q_{L(R)}$ is the B – L charge for the left (right handed) SM fermions, and we have taken the limit in which $M_{Z'}$ is much heavier compared to the SM particles. The dominant decay rates of ϕ are given by,

$$\begin{aligned}\Gamma_{\phi \rightarrow Z' Z'} &= \frac{g_X^2}{8\pi r^4} \frac{M_{Z'}^2}{m_\phi} \sqrt{1 - 4r^2} (12r^4 - 4r^2 + 1), \\ \Gamma_{\phi \rightarrow N_i N_i} &= \frac{g_X^2}{2\pi} \left(\frac{M_i}{M_{Z'}}\right)^2 m_\phi \left(1 - \frac{4M_i^2}{m_\phi^2}\right)^{3/2},\end{aligned}\quad (\text{A2})$$

where $r = M_{Z'}/m_\phi$. For very heavy Z' and DM, the total annihilation cross section mediated by Z' reads,

$$\begin{aligned}\sigma(s)_{\text{SM SM} \rightarrow N_1 N_1} &= \frac{170 g_X^4}{2592 \pi} \frac{s}{[(s - M_{Z'}^2)^2 + \Gamma_{Z'}^2 M_{Z'}^2]} \left(1 - \frac{4M_1^2}{s}\right)^{3/2},\end{aligned}\quad (\text{A3})$$

where $\Gamma_{Z'}$ is the total decay width of Z' into all possible final states.

APPENDIX B: BOLTZMANN EQUATIONS FOR DM ANALYSIS

The coupled BEQs for freeze-in production of the DM, N_1 , read

$$\begin{aligned}\frac{dy_\phi}{dz} &= -\frac{z}{\mathcal{H}} \langle \Gamma_\phi \rangle y_{\text{eq}}^\phi + \frac{s}{\mathcal{H} z^2} \langle \sigma v \rangle_{\text{SM SM} \rightarrow \phi \phi} (y_{\text{eq}}^\phi)^2, \\ \frac{dy_{Z'}}{dz} &= -\frac{z}{\mathcal{H}} \langle \Gamma_{Z'} \rangle y_{Z'} + \frac{z}{\mathcal{H}} \langle \Gamma_{\phi \rightarrow Z' Z'} \rangle (y_{\text{eq}}^\phi - y_{Z'}), \\ \frac{dy_{N_1}}{dz} &= \frac{z}{\mathcal{H}} \langle \Gamma_{Z' \rightarrow N_1 N_1} \rangle y_{Z'} + \frac{z}{\mathcal{H}} \langle \Gamma_{\phi \rightarrow N_1 N_1} \rangle y_{\text{eq}}^\phi \\ &\quad + \frac{s}{\mathcal{H} z^2} \langle \sigma v \rangle_{\text{SM SM} \rightarrow N_1 N_1} y_{\text{eq}}^2,\end{aligned}\quad (\text{B1})$$

where $y_i \equiv n_i/s$ is the yield of a certain species i , with

$$y_j^{\text{eq}} = \frac{45}{4\pi^4} \frac{g_j}{g_{\star s}} z^2 K_2[z], \quad (\text{B2})$$

is the equilibrium yield, with g_j being the degrees of freedom for the corresponding j particle and $z = M_1/T$ is a dimensionless variable. The thermally averaged decay rate $\langle \Gamma_\phi \rangle$ reads,

$$\langle \Gamma_{\phi \rightarrow jj} \rangle = \frac{K_1(z)}{K_2(z)} \times \Gamma_{\phi \rightarrow jj}, \quad (\text{B3})$$

with $z = M_1/T$, and j represents the final state particle. Here, $\mathcal{H} = (\pi/3) \sqrt{g_\star/10} (T^2/M_P)$ is the Hubble parameter for a standard radiation dominated (RD) Universe and $s = (2\pi^2/45) g_{\star s}(T) T^3$ is the entropy density. The number of relativistic degrees of freedom in the bath corresponding to energy density and entropy density are tracked by $g_\star(T)$ and $g_{\star s}(T)$, respectively. At temperatures well above the QCD phase transition we have $g_\star(T) \simeq g_{\star s}(T) \approx 106$.

APPENDIX C: CP-ASYMMETRY FROM N_2 DECAY

Since N_1 is a long-lived DM, the CP asymmetry generated from N_2 decay is given by [103]

$$\begin{aligned}\epsilon_{\Delta L} &\equiv \frac{\Gamma_{N_2 \rightarrow \ell_i H} - \Gamma_{N_2 \rightarrow \bar{\ell}_i \bar{H}}}{\Gamma_{N_2 \rightarrow \ell_i H} + \Gamma_{N_2 \rightarrow \bar{\ell}_i \bar{H}}} \\ &\simeq \frac{1}{8\pi} \frac{1}{(Y_\nu^\dagger Y_\nu)_{22}} \text{Im} (Y_\nu^\dagger Y_\nu)_{23}^2 \times \mathcal{F} \left(\frac{M_3^2}{M_2^2} \right),\end{aligned}\quad (\text{C1})$$

where

$$\mathcal{F}(x) \equiv \sqrt{x} \left[\frac{1}{1-x} + 1 - (1+x) \log \left(\frac{1+x}{x} \right) \right]. \quad (\text{C2})$$

For $x \gg 1$, $\mathcal{F} \simeq -3/(2\sqrt{x})$, and Eq. (C1) becomes

$$\epsilon_{\Delta L} \simeq -\frac{3}{16\pi} \frac{1}{(Y_\nu^\dagger Y_\nu)_{22}} \times [\text{Im} (Y_\nu^\dagger Y_\nu)_{23}^2] \frac{M_2}{M_3}. \quad (\text{C3})$$

$$m_{\nu,3} = \frac{|(Y_\nu)_{23}|^2 v^2}{M_1}, \quad (\text{C6})$$

Then

$$\epsilon_{\Delta L} \simeq -\frac{3\delta_{\text{eff}} |(Y_\nu)_{23}|^2 M_2}{16\pi M_3}, \quad (\text{C4})$$

which corresponds to the heaviest left-handed neutrino in the normal hierarchy. This leads to

$$\epsilon_{\Delta L} \simeq -\frac{3\delta_{\text{eff}} M_2 m_{\nu,3}}{16\pi v^2}. \quad (\text{C7})$$

while the effective CP violating phase is given by

$$\delta_{\text{eff}} = \frac{1}{(Y_\nu)_{23}^2} \frac{\text{Im}(Y_\nu^\dagger Y_\nu)_{23}^2}{(Y_\nu^\dagger Y_\nu)_{22}}. \quad (\text{C5})$$

To connect with the light neutrino mass, we impose the seesaw relation

APPENDIX D: BOLTZMANN EQUATIONS FOR N_2 LEPTOGENESIS

To track the evolution of RHN number density and the corresponding $B-L$ asymmetry with time, we solve the following set of BEQs [along with the first one in Eq. (B1)],

$$\begin{aligned} \frac{dy_{Z'}}{dz} &= -\frac{z}{\mathcal{H}} \langle \Gamma_{Z'} \rangle y_{Z'} + \frac{z}{\mathcal{H}} \langle \Gamma_{\phi \rightarrow Z' Z'} \rangle (y_{\text{eq}}^\phi - y_{Z'}), \\ \frac{dy_{N_2}}{dz} &= \frac{z}{\mathcal{H}} \langle \Gamma_{Z' \rightarrow N_2 N_2} \rangle y_{Z'} + \frac{z}{\mathcal{H}} \langle \Gamma_{\Phi \rightarrow N_2 N_2} \rangle y_{\text{eq}}^\phi + \frac{s}{\mathcal{H}} \frac{1}{z^2} \langle \sigma v \rangle_{\text{SM SM} \rightarrow N_2 N_2} (y_{N_2}^{\text{eq}} - y_{N_2}^2) - \frac{z}{\mathcal{H}} \langle \Gamma_{N_2} \rangle (y_{N_2} - y_{N_2}^{\text{eq}}), \\ \frac{dy_{B-L}}{dz} &= \frac{\langle \Gamma_{N_2} \rangle}{\mathcal{H}} z \epsilon_{\Delta L} (y_{N_2} - y_{N_2}^{\text{eq}}) - \frac{\langle \Gamma_{N_2} \rangle}{\mathcal{H}} \frac{z^3}{4} \frac{\tilde{m}_2}{m_\star} K_1(z) y_{B-L}, \end{aligned} \quad (\text{D1})$$

where $z = M_2/T$. Here, $\tilde{m}_2 = (m_D^\dagger m_D)_{22}/M_2 \approx m_{\nu,2}^2/M_2$ and $m_\star \simeq 10^{-3}$ eV is the equilibrium neutrino mass [66]. In our scenario, we find, $\tilde{m}_2 > m_\star$, as a result we are always in the strong washout regime. The thermally averaged N_2 decay width is given by

$$\langle \Gamma_{N_2} \rangle = \frac{K_1(M_2/T)}{K_2(M_2/T)} \frac{M_2 m_{\nu,3} M_2}{4\pi v^2}, \quad (\text{D2})$$

where we have utilized the seesaw relations following normal hierarchy, $m_{\nu,1} \simeq 0$, $m_{\nu,2} \simeq 0.0086$ eV and $m_{\nu,3} \simeq 0.0506$ eV.

APPENDIX E: RENORMALIZATION GROUP EQUATIONS (RGE)

We extract 1-loop β -function using PyR@TE [104–106]. These are given by,

$$\begin{aligned} (4\pi)^2 \beta(g_X) &= 12g_X^3 \\ (4\pi)^2 \beta(Y_\nu) &= \frac{3}{2} Y_\nu Y_\nu^\dagger Y_\nu + 3\text{Tr}(Y_t^\dagger Y_t) Y_\nu + \text{Tr}(Y_\nu^\dagger Y_\nu) Y_\nu - \frac{3}{4} g^2 Y_\nu - 6g_X^2 Y_\nu - \frac{9}{4} g_2^2 Y_\nu \\ (4\pi)^2 \beta(\lambda_H) &= 24\lambda_H^2 + \lambda_{\text{mix}}^2 - 3g_1^2 \lambda_H - 9g_2^2 \lambda_H + \frac{3}{8} g_1^4 + \frac{3}{4} g_1^2 g_2^2 + \frac{9}{8} g_2^4 + 12\lambda_H Y_t^2 - 6Y_t^4 \\ (4\pi)^2 \beta(\lambda_\Phi) &= 20\lambda_H^2 + 2\lambda_\Phi^2 \\ (4\pi)^2 \beta(\lambda_{\text{mix}}) &= 12\lambda_H \lambda_{\text{mix}} + 8\lambda_\Phi \lambda_{\text{mix}} + 4\lambda_{\text{mix}}^2 - \frac{3}{2} g_1^2 \lambda_{\text{mix}} - 24g_X^2 \lambda_{\text{mix}} - \frac{9}{2} g_2^2 \lambda_{\text{mix}} + 6\lambda_{\text{mix}} \text{Tr}(Y_t^\dagger Y_t), \end{aligned} \quad (\text{E1})$$

together with

$$\begin{aligned} (4\pi)^2 \beta(g_1) &= \frac{41}{6} g_1^3, \\ (4\pi)^2 \beta(g_2) &= -\frac{19}{6} g_2^3, \\ (4\pi)^2 \beta(g_3) &= -7g_3^3, \\ (4\pi)^2 \beta(Y_t) &= \frac{3}{2} Y_t Y_t^\dagger Y_t + 3\text{Tr}(Y_t^\dagger Y_t) Y_t - \frac{17}{12} g_1^2 Y_t - \frac{9}{4} g_2^2 Y_t - 8g_3^2 Y_t. \end{aligned} \quad (\text{E2})$$

- [1] S. Navas *et al.* (Particle Data Group), Review of particle physics, *Phys. Rev. D* **110**, 030001 (2024).
- [2] N. Aghanim *et al.* (Planck Collaboration), Planck 2018 results. VI. Cosmological parameters, *Astron. Astrophys.* **641**, A6 (2020); *Astron. Astrophys.* **652**, C4(E) (2021).
- [3] G. Bertone and D. Hooper, History of dark matter, *Rev. Mod. Phys.* **90**, 045002 (2018).
- [4] J. de Swart, G. Bertone, and J. van Dongen, How dark matter came to matter, *Nat. Astron.* **1**, 0059 (2017).
- [5] M. G. Aartsen *et al.* (IceCube Collaboration), First observation of PeV-energy neutrinos with IceCube, *Phys. Rev. Lett.* **111**, 021103 (2013).
- [6] M. G. Aartsen *et al.* (IceCube Collaboration), Evidence for high-energy extraterrestrial neutrinos at the IceCube detector, *Science* **342**, 1242856 (2013).
- [7] M. G. Aartsen *et al.* (IceCube Collaboration), Observation of high-energy astrophysical neutrinos in three years of IceCube data, *Phys. Rev. Lett.* **113**, 101101 (2014).
- [8] M. G. Aartsen *et al.* (IceCube Collaboration), A combined maximum-likelihood analysis of the high-energy astrophysical neutrino flux measured with IceCube, *Astrophys. J.* **809**, 98 (2015).
- [9] M. G. Aartsen *et al.* (IceCube Collaboration), Evidence for astrophysical muon neutrinos from the northern sky with IceCube, *Phys. Rev. Lett.* **115**, 081102 (2015).
- [10] S. W. Li, P. Machado, D. Naredo-Tuero, and T. Schwemberger, Clash of the Titans: Ultra-high energy KM3NeT event versus IceCube data, *arXiv:2502.04508*.
- [11] Y. Bai, R. Lu, and J. Salvado, Geometric compatibility of IceCube TeV-PeV neutrino excess and its galactic dark matter origin, *J. High Energy Phys.* **01** (2016) 161.
- [12] T. Higaki, R. Kitano, and R. Sato, Neutrino Universe, *J. High Energy Phys.* **07** (2014) 044.
- [13] A. Esmaili, S. K. Kang, and P. D. Serpico, IceCube events and decaying dark matter: Hints and constraints, *J. Cosmol. Astropart. Phys.* **12** (2014) 054.
- [14] K. Murase, R. Laha, S. Ando, and M. Ahlers, Testing the dark matter scenario for PeV neutrinos observed in IceCube, *Phys. Rev. Lett.* **115**, 071301 (2015).
- [15] E. Dudas, T. Gherghetta, K. Kaneta, Y. Mambrini, and K. A. Olive, Gravitino decay in high scale supersymmetry with R-parity violation, *Phys. Rev. D* **98**, 015030 (2018).
- [16] T. Cohen, K. Murase, N. L. Rodd, B. R. Safdi, and Y. Soreq, γ -ray constraints on decaying dark matter and implications for IceCube, *Phys. Rev. Lett.* **119**, 021102 (2017).
- [17] C. Rott, K. Kohri, and S. C. Park, Superheavy dark matter and IceCube neutrino signals: Bounds on decaying dark matter, *Phys. Rev. D* **92**, 023529 (2015).
- [18] C. A. Argüelles, D. Delgado, A. Friedlander, A. Kheirandish, I. Safa, A. C. Vincent, and H. White, Dark matter decay to neutrinos, *Phys. Rev. D* **108**, 123021 (2023).
- [19] R. Abbott *et al.* (LIGO Scientific Collaboration, Virgo Collaboration, and KAGRA Collaboration), Constraints on cosmic strings using data from the third Advanced LIGO–Virgo observing run, *Phys. Rev. Lett.* **126**, 241102 (2021).
- [20] R. Caldwell *et al.*, Detection of early-universe gravitational-wave signatures and fundamental physics, *Gen. Relativ. Gravit.* **54**, 156 (2022).
- [21] G. Agazie *et al.* (NANOGrav Collaboration), The NANOGrav 15 yr data set: Evidence for a gravitational-wave background, *Astrophys. J. Lett.* **951**, L8 (2023).
- [22] A. Afzal *et al.* (NANOGrav Collaboration), The NANOGrav 15 yr data set: Search for signals from new physics, *Astrophys. J. Lett.* **951**, L11 (2023); *Astrophys. J. Lett.* **971**, L27(E) (2024).
- [23] H. Xu *et al.*, Searching for the nano-hertz stochastic gravitational wave background with the Chinese Pulsar Timing Array data release I, *Res. Astron. Astrophys.* **23**, 075024 (2023).
- [24] J. Antoniadis *et al.* (EPTA Collaboration and InPTA Collaboration), The second data release from the European Pulsar Timing Array—III. Search for gravitational wave signals, *Astron. Astrophys.* **678**, A50 (2023).
- [25] P. Auclair *et al.* (LISA Cosmology Working Group), Cosmology with the laser interferometer space antenna, *Living Rev. Relativity* **26**, 5 (2023).
- [26] H. B. Nielsen and P. Olesen, Vortex line models for dual strings, *Nucl. Phys. B* **61**, 45 (1973).
- [27] T. W. B. Kibble, Topology of cosmic domains and strings, *J. Phys. A* **9**, 1387 (1976).
- [28] A. Davidson, $B-L$ as the fourth color within an $SU(2)_L \times U(1)_R \times U(1)$ model, *Phys. Rev. D* **20**, 776 (1979).
- [29] R. E. Marshak and R. N. Mohapatra, Quark–lepton symmetry and $B-L$ as the $U(1)$ generator of the electroweak symmetry group, *Phys. Lett.* **91B**, 222 (1980).
- [30] P. Minkowski, $\mu \rightarrow e\gamma$ at a Rate of one out of 10^9 muon decays?, *Phys. Lett.* **67B**, 421 (1977).
- [31] T. Yanagida, Horizontal gauge symmetry and masses of neutrinos, *Conf. Proc. C* **7902131**, 95 (1979).
- [32] M. Gell-Mann, P. Ramond, and R. Slansky, Complex spinors and unified theories, *Conf. Proc. C* **790927**, 315 (1979).
- [33] R. N. Mohapatra and G. Senjanovic, Neutrino mass and spontaneous parity nonconservation, *Phys. Rev. Lett.* **44**, 912 (1980).
- [34] J. Schechter and J. W. F. Valle, Neutrino masses in $SU(2) \times U(1)$ theories, *Phys. Rev. D* **22**, 2227 (1980).
- [35] A. Babul, T. Piran, and D. N. Spergel, Bosonic superconducting cosmic strings. I. Classical field theory solutions, *Phys. Lett. B* **202**, 307 (1988).
- [36] A. Vilenkin and E. P. S. Shellard, *Cosmic Strings and Other Topological Defects* (Cambridge University Press, Cambridge, England, 2000).
- [37] J. A. Dror, T. Hiramatsu, K. Kohri, H. Murayama, and G. White, Testing the Seesaw mechanism and leptogenesis with gravitational waves, *Phys. Rev. Lett.* **124**, 041804 (2020).
- [38] A. Das, P. S. B. Dev, Y. Hosotani, and S. Mandal, Probing the minimal $U(1)_X$ model at future electron-positron colliders via the fermion pair-production channel, *Phys. Rev. D* **105**, 115030 (2022).
- [39] C. Lunardini and Y. F. Perez-Gonzalez, Dirac and Majorana neutrino signatures of primordial black holes, *J. Cosmol. Astropart. Phys.* **08** (2020) 014.
- [40] N. Bernal, V. Muñoz Alborno, S. Palomares-Ruiz, and P. Villanueva-Domingo, Current and future neutrino limits

- on the abundance of primordial black holes, *J. Cosmol. Astropart. Phys.* **10** (2022) 068.
- [41] Q.-f. Wu and X.-J. Xu, High-energy and ultra-high-energy neutrinos from primordial black holes, *J. Cosmol. Astropart. Phys.* **02** (2025) 059.
- [42] M. Chianese, A. Boccia, F. Iocco, G. Miele, and N. Saviano, Light burden of memory: Constraining primordial black holes with high-energy neutrinos, *Phys. Rev. D* **111**, 063036 (2025).
- [43] M. Zantedeschi and L. Visinelli, Ultralight black holes as sources of high-energy particles, *Phys. Dark Universe* **49**, 102034 (2025).
- [44] M. J. Baker, J. Iguaz Juan, A. Symons, and A. Thamm, Explaining the PeV neutrino fluxes at KM3NeT and IceCube with quasi-extremal primordial black holes, [arXiv:2505.22722](#).
- [45] T. Fujita, K. Harigaya, and T. Matsuda, Baryon asymmetry, dark matter, and density perturbation from primordial black holes, *Phys. Rev. D* **89**, 103501 (2014).
- [46] S. Datta, A. Ghosal, and R. Samanta, Baryogenesis from ultralight primordial black holes and strong gravitational waves from cosmic strings, *J. Cosmol. Astropart. Phys.* **08** (2021) 021.
- [47] B. Barman, D. Borah, S. J. Das, and R. Roshan, Non-thermal origin of asymmetric dark matter from inflaton and primordial black holes, *J. Cosmol. Astropart. Phys.* **03** (2022) 031.
- [48] N. Bernal, C. S. Fong, Y. F. Perez-Gonzalez, and J. Turner, Rescuing high-scale leptogenesis using primordial black holes, *Phys. Rev. D* **106**, 035019 (2022).
- [49] O. Sawada and A. Sugamoto, *Proceedings: Workshop on the Unified Theories and the Baryon Number in the Universe: Tsukuba, Japan, 1979* (Natl. Lab. High Energy Phys., Tsukuba, Japan, 1979).
- [50] R. N. Mohapatra and G. Senjanovic, Neutrino masses and mixings in gauge models with spontaneous parity violation, *Phys. Rev. D* **23**, 165 (1981).
- [51] M. Chianese and A. Merle, A consistent theory of decaying dark matter connecting IceCube to the sesame street, *J. Cosmol. Astropart. Phys.* **04** (2017) 017.
- [52] L. J. Hall, K. Jedamzik, J. March-Russell, and S. M. West, Freeze-in production of FIMP dark matter, *J. High Energy Phys.* **03** (2010) 080.
- [53] N. Bernal, M. Heikinheimo, T. Tenkanen, K. Tuominen, and V. Vaskonen, The dawn of FIMP dark matter: A review of models and constraints, *Int. J. Mod. Phys. A* **32**, 1730023 (2017).
- [54] M. Aker *et al.* (KATRIN Collaboration), Improved upper limit on the neutrino mass from a direct kinematic method by KATRIN, *Phys. Rev. Lett.* **123**, 221802 (2019).
- [55] M. J. Dolinski, A. W. P. Poon, and W. Rodejohann, Neutrinoless double-beta decay: Status and prospects, *Annu. Rev. Nucl. Part. Sci.* **69**, 219 (2019).
- [56] T. Robens and T. Stefaniak, Status of the Higgs singlet extension of the standard model after LHC run 1, *Eur. Phys. J. C* **75**, 104 (2015).
- [57] G. Chalons, D. Lopez-Val, T. Robens, and T. Stefaniak, The Higgs singlet extension at LHC Run 2, *Proc. Sci., ICHEP2016* (2016) 1180 [[arXiv:1611.03007](#)].
- [58] A. Das, S. Gola, S. Mandal, and N. Sinha, Two-component scalar and fermionic dark matter candidates in a generic $U(1)_X$ model, *Phys. Lett. B* **829**, 137117 (2022).
- [59] R. Barate *et al.* (LEP Working Group for Higgs boson searches, ALEPH, DELPHI, L3, OPAL Collaborations), Search for the standard model Higgs boson at LEP, *Phys. Lett. B* **565**, 61 (2003).
- [60] Y. Wang, M. Berggren, and J. List, ILD benchmark: Search for extra scalars produced in association with a Z boson at $\sqrt{s} = 500$ GeV, [arXiv:2005.06265](#).
- [61] J. de Blas *et al.* (CLIC Collaboration), The CLIC potential for new physics, *CERN Yellow Rep. Monogr.* **3**, 1 (2018).
- [62] K. Kaneta, Z. Kang, and H.-S. Lee, Right-handed neutrino dark matter under the $B - L$ gauge interaction, *J. High Energy Phys.* **02** (2017) 031.
- [63] S. Eijima, O. Seto, and T. Shimomura, Revisiting sterile neutrino dark matter in gauged $U(1)_{B-L}$ model, *Phys. Rev. D* **106**, 103513 (2022).
- [64] O. Seto, T. Shimomura, and Y. Uchida, Freeze-in sterile neutrino dark matter in feeble gauged $B - L$ model, *J. High Energy Phys.* **05** (2025) 147.
- [65] R. L. Workman *et al.* (Particle Data Group), Review of particle physics, *Prog. Theor. Exp. Phys.* **2022**, 083C01 (2022).
- [66] W. Buchmuller, P. Di Bari, and M. Plumacher, Leptogenesis for pedestrians, *Ann. Phys. (N.Y.)* **315**, 305 (2005).
- [67] K. Kaneta, Y. Mambrini, K. A. Olive, and S. Verner, Inflation and leptogenesis in high-scale supersymmetry, *Phys. Rev. D* **101**, 015002 (2020).
- [68] S. Davidson and A. Ibarra, A lower bound on the right-handed neutrino mass from leptogenesis, *Phys. Lett. B* **535**, 25 (2002).
- [69] D. Buttazzo, G. Degrandi, P. P. Giardino, G. F. Giudice, F. Sala, A. Salvio, and A. Strumia, Investigating the near-criticality of the Higgs boson, *J. High Energy Phys.* **12** (2013) 089.
- [70] C. Ringeval, M. Sakellariadou, and F. Bouchet, Cosmological evolution of cosmic string loops, *J. Cosmol. Astropart. Phys.* **02** (2007) 023.
- [71] J. J. Blanco-Pillado, K. D. Olum, and B. Shlaer, Large parallel cosmic string simulations: New results on loop production, *Phys. Rev. D* **83**, 083514 (2011).
- [72] K. Nakayama and Y. Tang, Stochastic gravitational waves from particle origin, *Phys. Lett. B* **788**, 341 (2019).
- [73] B. Barman, N. Bernal, Y. Xu, and O. Zapata, Gravitational wave from graviton Bremsstrahlung during reheating, *J. Cosmol. Astropart. Phys.* **05** (2023) 019.
- [74] B. Barman, N. Bernal, Y. Xu, and O. Zapata, Bremsstrahlung-induced gravitational waves in monomial potentials during reheating, *Phys. Rev. D* **108**, 083524 (2023).
- [75] K.-Y. Choi, E. Lkhagvadorj, and S. Mahapatra, Gravitational wave sourced by decay of massive particle from primordial black hole evaporation, *J. Cosmol. Astropart. Phys.* **07** (2024) 064.
- [76] A. Ghoshal, R. Samanta, and G. White, Bremsstrahlung high-frequency gravitational wave signatures of high-scale nonthermal leptogenesis, *Phys. Rev. D* **108**, 035019 (2023).

- [77] A. Datta and A. Sil, Probing Leptogenesis through gravitational waves, [arXiv:2410.01900](#).
- [78] O. Catà, A. Ibarra, and S. Ingenhütt, Dark matter decay through gravity portals, *Phys. Rev. D* **95**, 035011 (2017).
- [79] O. Catà, A. Ibarra, and S. Ingenhütt, Sharp spectral features from light dark matter decay via gravity portals, *J. Cosmol. Astropart. Phys.* **11** (2017) 044.
- [80] A. Vilenkin, Gravitational radiation from cosmic strings, *Phys. Lett. B* **107**, 47 (1981).
- [81] T. Vachaspati and A. Vilenkin, Gravitational radiation from cosmic strings, *Phys. Rev. D* **31**, 3052 (1985).
- [82] J.J. Blanco-Pillado, K.D. Olum, and B. Shlaer, The number of cosmic string loops, *Phys. Rev. D* **89**, 023512 (2014).
- [83] J.J. Blanco-Pillado and K.D. Olum, Stochastic gravitational wave background from smoothed cosmic string loops, *Phys. Rev. D* **96**, 104046 (2017).
- [84] T. Damour and A. Vilenkin, Gravitational wave bursts from cusps and kinks on cosmic strings, *Phys. Rev. D* **64**, 064008 (2001).
- [85] Y. Gouttenoire, G. Servant, and P. Simakachorn, Beyond the standard models with cosmic strings, *J. Cosmol. Astropart. Phys.* **07** (2020) 032.
- [86] C. J. A. P. Martins and E. P. S. Shellard, Quantitative string evolution, *Phys. Rev. D* **54**, 2535 (1996).
- [87] C. J. A. P. Martins and E. P. S. Shellard, Extending the velocity dependent one scale string evolution model, *Phys. Rev. D* **65**, 043514 (2002).
- [88] P. Auclair *et al.*, Probing the gravitational wave background from cosmic strings with LISA, *J. Cosmol. Astropart. Phys.* **04** (2020) 034.
- [89] A. Vilenkin, Cosmic string dynamics with friction, *Phys. Rev. D* **43**, 1060 (1991).
- [90] T. Charnock, A. Avgoustidis, E. J. Copeland, and A. Moss, CMB constraints on cosmic strings and superstrings, *Phys. Rev. D* **93**, 123503 (2016).
- [91] J. Kume and M. Hindmarsh, Revised bounds on local cosmic strings from NANOGrav observations, *J. Cosmol. Astropart. Phys.* **12** (2024) 001.
- [92] J. Lizarraga, J. Urrestilla, D. Daverio, M. Hindmarsh, and M. Kunz, New CMB constraints for Abelian Higgs cosmic strings, *J. Cosmol. Astropart. Phys.* **10** (2016) 042.
- [93] J. Crowder and N. J. Cornish, Beyond LISA: Exploring future gravitational wave missions, *Phys. Rev. D* **72**, 083005 (2005).
- [94] V. Corbin and N. J. Cornish, Detecting the cosmic gravitational wave background with the big bang observer, *Classical Quantum Gravity* **23**, 2435 (2006).
- [95] N. Seto, S. Kawamura, and T. Nakamura, Possibility of direct measurement of the acceleration of the universe using 0.1-Hz band laser interferometer gravitational wave antenna in space, *Phys. Rev. Lett.* **87**, 221103 (2001).
- [96] H. Kudoh, A. Taruya, T. Hiramatsu, and Y. Himemoto, Detecting a gravitational-wave background with next-generation space interferometers, *Phys. Rev. D* **73**, 064006 (2006).
- [97] P. Amaro-Seoane *et al.* (LISA Collaboration), Laser Interferometer Space Antenna, [arXiv:1702.00786](#).
- [98] D. Reitze *et al.*, Cosmic explorer: The U.S. contribution to gravitational-wave astronomy beyond LIGO, *Bull. Am. Astron. Soc.* **51**, 035 (2019).
- [99] S. Hild *et al.*, Sensitivity studies for third-generation gravitational wave observatories, *Classical Quantum Gravity* **28**, 094013 (2011).
- [100] M. Punturo *et al.*, The Einstein Telescope: A third-generation gravitational wave observatory, *Classical Quantum Gravity* **27**, 194002 (2010).
- [101] B. Sathyaprakash *et al.*, Scientific objectives of Einstein Telescope, *Classical Quantum Gravity* **29**, 124013 (2012); , *Classical Quantum Gravity* **30**, 079501(E) (2013).
- [102] M. Maggiore *et al.*, Science case for the Einstein Telescope, *J. Cosmol. Astropart. Phys.* **03** (2020) 050.
- [103] S. Davidson, E. Nardi, and Y. Nir, Leptogenesis, *Phys. Rep.* **466**, 105 (2008).
- [104] F. Lyonnet, I. Schienbein, F. Staub, and A. Wingerter, PyR@TE: Renormalization group equations for general gauge theories, *Comput. Phys. Commun.* **185**, 1130 (2014).
- [105] F. Lyonnet and I. Schienbein, PyR@TE 2: A Python tool for computing RGEs at two-loop, *Comput. Phys. Commun.* **213**, 181 (2017).
- [106] L. Sartore and I. Schienbein, PyR@TE 3, *Comput. Phys. Commun.* **261**, 107819 (2021).