

Cosmological bounds on heavy neutral leptons with dark decay channels

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Heavy neutral leptons (HNLs) are well-motivated new physics candidates. The mixing of sub-GeV HNLs with active neutrinos is severely constrained by cosmology. In particular, the success of big bang nucleosynthesis (BBN) requires the HNL lifetime to be shorter than about 0.02 sec if they were in thermal equilibrium, thus excluding a wide range of mixing angles accessible to terrestrial experiments. In order to justify the laboratory searches in this cosmologically forbidden region, it is often argued that adding new dark sector decay modes of HNLs can evade the stringent BBN constraint. Here we rule out this possibility and show that, contrary to the naive expectation, HNLs with significant dark decay modes actually lead to stronger cosmological bounds. This is mainly because of the increase in the extra radiation energy density in the Universe around the BBN epoch, which causes observable effects in the primordial helium fraction and ΔN_{eff} . Our result has major implications for laboratory searches of HNLs.

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I. INTRODUCTION

The origin of neutrino mass is a fundamental open question. Among the many theoretical possibilities, perhaps the simplest is to extend the Standard Model (SM) particle content by adding right-handed neutrinos, popularly known as heavy neutral leptons (HNLs), thus generating the Dirac mass (m_D) for neutrinos after electroweak symmetry breaking. Moreover, the HNLs, being SM singlets, can also acquire a Majorana mass (m_N), which leads to the famous seesaw relation for explaining the light neutrino masses: $m_\nu \simeq -m_D m_N^{-1} m_D^T$ [1–4]. Irrespective of the origin of the Majorana mass term and its ultraviolet completion, even the minimal ν SM scenario with just SM + HNLs can not only explain neutrino masses, but also account for the observed dark matter and baryon asymmetry of the Universe [5,6], thus making it an elegant candidate for beyond-the-SM (BSM) physics.

In the ν SM, the HNLs interact with the SM sector only via their mixing with the active neutrinos ν_α (with $\alpha = e, \mu, \tau$) given by $U_\alpha \simeq m_D m_N^{-1}$. The seesaw relation predicts

$$U_\alpha^2 \simeq \frac{m_\nu}{m_N} \quad (\text{seesaw line}), \quad (1)$$

where $m_\nu \sim 0.01\text{--}0.1$ eV is the neutrino mass scale needed to explain the solar and atmospheric mass-squared differences, respectively. Note that this relation is strictly valid only in the single-HNL limit, while the general case with two or more HNLs [7–10], or specific model constructions [11–16], can yield significantly larger mixing, thus allowing m_N and U_α to be essentially treated as free parameters. Nevertheless, Eq. (1) offers a theoretically well-motivated target for experimental searches.

A summary of the existing constraints in the HNL mass-mixing plane can be found in Refs. [17–19]. For $m_N \lesssim \mathcal{O}(\text{GeV})$, the most stringent constraint on the HNL mixing comes from big bang nucleosynthesis (BBN), which requires that the HNL lifetime must be shorter than $\sim 0.02\text{--}0.1$ sec [20–25]. In this paper, we explore the possibility of whether this lifetime constraint can be evaded or at least relaxed in the presence of additional HNL decays to dark sector particles. We show that, contrary to the naive expectation, just making the HNL more short lived by dominantly decaying to the dark sector does not actually

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relax the cosmological bound, but makes it stronger. This comes from a combination of the BBN constraints on the primordial abundance of light elements and the cosmic microwave background (CMB) constraint on the effective number of relativistic degrees of freedom. This conclusion is qualitatively expected from energy conservation, since the shortened lifetime generally implies that the energy of HNLs is dominantly transferred to the dark sector, so its impact on the Hubble expansion cannot be hidden. Nonetheless, a quantitative analysis is still required to evaluate this impact, as we carry out in this paper.

II. HNLs IN THE EARLY UNIVERSE

In this work, we are mainly interested in the MeV-GeV scale HNLs and experimentally relevant mixing above the seesaw line. In this regime, the mixing yields a sufficiently high interaction rate of HNLs with the SM thermal bath and thus brings the HNLs into thermal equilibrium at an early epoch. As the Universe cools down to the temperature

$$T_i \simeq \frac{1 \text{ MeV}}{U_\alpha^{2/3}}, \quad (2)$$

HNLs freeze out from the thermal bath while being relativistic (for the allowed values of mixing), similar to neutrino decoupling in the SM. Throughout this paper, we use the subscript “ i ” to indicate the moment of $T = T_i$. After decoupling, HNLs start free-streaming in the Universe and eventually decay when the age of the Universe is comparable to their lifetime τ_N .

The age of the Universe during the radiation-dominated era can be estimated using $t \approx 1/(2H)$, where H is the Hubble parameter. Using $H \approx 1.66 g_\star^{1/2} T^2 / m_{\text{pl}}$, where $g_\star$ is the effective number of degrees of freedom, T is the temperature of the Universe, and $m_{\text{pl}} \approx 1.22 \times 10^{19} \text{ GeV}$ is the Planck mass, one obtains¹

$$t \approx 0.74 \text{ sec} \cdot \left(\frac{\text{MeV}}{T} \right)^2 \left(\frac{10.75}{g_\star} \right)^{1/2}. \quad (3)$$

At $T \approx 2.7 \text{ MeV}$, the age of the Universe is 0.1 sec. As the Universe evolves beyond this point, the SM thermal bath undergoes several crucial processes, including (i) neutrino decoupling, (b) freeze-out of the neutron-to-proton ratio, and (iii) e^+e^- annihilation. All of these processes occur at around the MeV scale and, if influenced by new physics, may alter standard cosmological predictions of BBN and CMB observables such as the helium abundance (Y_P) and the effective number of relativistic species (N_{eff}). In order to keep the successful predictions of the standard

¹Note that in the presence of HNLs, Eq. (3) could be modified significantly, in particular when HNLs dominate the total energy. In this work, Eq. (3) is only used for qualitative discussions and not used in quantitative calculations.

cosmology unaffected, HNLs must decay before these crucial processes start to take effect, i.e., before the Universe cools down to the MeV scale. Consequently, the cosmological bound on HNLs in the freeze-out regime roughly corresponds to an upper bound on their lifetime τ_N . For instance, $\tau_N \lesssim 0.1 \text{ sec}$ was used in earlier studies [6,26,27] as a conservative estimate of the BBN bound. More recent studies performed a dedicated analysis of the HNL’s influence on the MeV-scale Universe and obtained lower values of τ_N varying between 0.02–0.1 sec [20,21]; see Appendix A.

III. CAN DARK DECAY RELAX THE BBN BOUND?

Now let us consider the possibility that HNLs may also decay via additional nonstandard interactions to some dark sector particles. For fixed U_α and m_N , a dark decay mode of N would enlarge its total decay width, resulting in a shorter lifetime. Therefore, one may naively expect that adding new dark sector decay modes can evade or relax the BBN constraint (see, e.g., the discussion in Ref. [18] and references therein). The general reasoning is that decays into truly dark states avoid direct electromagnetic/charged injections and can reduce visible damage to light-element abundances or CMB signals; see, e.g., Refs. [28–34]. However, in what follows, we show that this, in general, is not true for HNL decays, especially in the region of interest to laboratory HNL searches.

Let us first consider a relativistic dark decay mode, $N \rightarrow X$, where X denotes dark sector final states which we assume to consist of relativistic species only. Hence, once produced, X would behave as dark radiation. We denote the total decay width of N by $\Gamma_N \equiv 1/\tau_N$, the SM contribution to it by $\Gamma_N^{(\text{SM})}$, and the decay width of $N \rightarrow X$ by $\Gamma_{N \rightarrow X}$. The dark branching ratio thus reads

$$\text{Br}_X \equiv \frac{\Gamma_{N \rightarrow X}}{\Gamma_N}. \quad (4)$$

After N has decoupled from the thermal bath [see Eq. (2)], it starts free-streaming relativistically until its momentum is redshifted to a low value comparable to m_N . Due to the smallness of the mixing, N decays nonrelativistically in most cases. From relativistic free-streaming to nonrelativistic decay, the energy density of N is approximately given by (see Appendix B for derivation):

$$\rho_N \approx \begin{cases} \frac{7\pi^2}{120} \left(\frac{a_i}{a} \right)^4 T_i^4 & \left(a \lesssim 3.15 \frac{a_i T_i}{m_N} \right) \\ \frac{3\zeta(3)}{2\pi^2} \left(\frac{a_i}{a} \right)^3 m_N T_i^3 e^{-\Gamma_N t} & \left(a \gtrsim 3.15 \frac{a_i T_i}{m_N} \right) \end{cases}. \quad (5)$$

Here a is the scale factor and $\zeta(3) \approx 1.202$ is the Riemann zeta function.

The first case in Eq. (5) is applicable to relativistic free-streaming, during which the comoving energy density $\rho_N a^4$

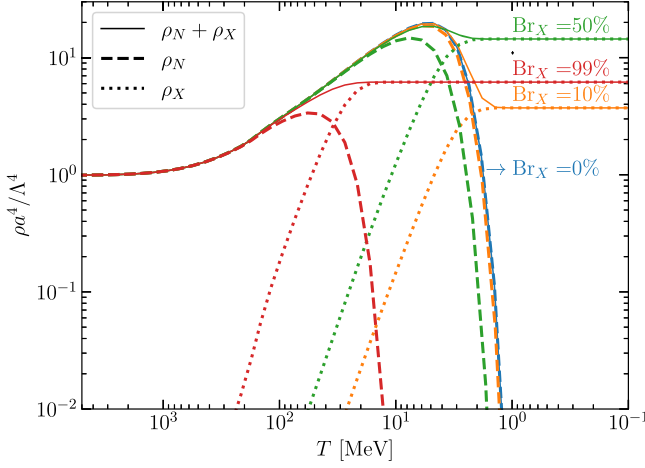


FIG. 1. The evolution of energy density of N (dashed lines), its decay product X (dotted lines), and the total energy density of N and X (solid lines). The y axis has been normalized to a dimensionless form using $\Lambda^4 \equiv \rho_{N,i} a_i^4$. For shown examples, we take $m_N = 1.08$ GeV, $U_e^2 = 10^{-10}$, and $\text{Br}_X \in \{0\%, 10\%, 50\%, 99\%\}$.

remains constant. When N becomes non-relativistic, corresponding to the second case in Eq. (5), ρ_N scales as a^{-3} if $\Gamma_N t \ll 1$, implying that $\rho_N a^4$ increases during this period. At $t \sim 1/\Gamma_N$, the exponential decay factor $e^{-\Gamma_N t}$ in Eq. (5) starts to suppress the energy density. The dashed curves in Fig. 1 show the exact evolution of $\rho_N a^4$ obtained by numerically solving the Boltzmann equation of N . They agree well with our expectation of the behavior of ρ_N described above.

The decay of N leads to the production of X , whose energy density can also be calculated analytically (see Appendix B). After a sufficiently long period of evolution, ρ_X scales as a^{-4} , corresponding to the flat part of the dotted lines in Fig. 1.

Although introducing a dark decay mode makes N decay faster, the total energy stored in the BSM sector at the BBN epoch is not necessarily reduced. This is illustrated by the solid lines in Fig. 1. For instance, when Br_X increases from 10% (blue solid line) to 50% (orange solid line), the total energy at $T \lesssim 1$ MeV increases significantly. This significant amount of dark radiation would modify the Hubble expansion rate and thereby alter BBN predictions, as we discuss below.

It is noteworthy that for the shown examples, the final value of $(\rho_N + \rho_X) a^4$ is higher than its initial. This has an important implication: the energy in the dark sector obtained in this way is even higher than what can be carried by a hypothetical neutrino species decoupled at T_i . If there was no variation of $g_\star$ during the evolution, these examples would lead to $\Delta N_{\text{eff}} > 1$, which would be excluded by BBN and CMB observations [35]. Including the variation of $g_\star$ leads to a smaller ΔN_{eff} , but still often large enough to be ruled out. *Therefore, introducing the*

dark decay mode $N \rightarrow X$ with sizable Br_X would generally cause more deviations of BBN predictions from their standard values.

For N decaying nonrelativistically to SM particles and X , we find that the resulting ΔN_{eff} can be approximated by (see Appendix B)

$$\Delta N_{\text{eff}} \approx \frac{\text{Br}_X \sqrt{1 - \text{Br}_X}}{0.16(1 - \text{Br}_X)^{3/2} + 0.08 g_{\star i}^{13/12} \xi_N}, \quad (6)$$

with $\xi_N \equiv \sqrt{\Gamma_N^{(\text{SM})} m_{\text{pl}}/m_N}$. Equation (6) implies that the dependence of ΔN_{eff} on Br_X is not monotonic. Indeed, as shown in Fig. 1, if one increases Br_X from 50% to 99%, the final value of $\rho_X a^4$ decreases. One may wonder what would happen if Br_X is increased to a value very close to 100%. In the limit of $\text{Br}_X \rightarrow 1$ with fixed U_α and m_N , Eq. (6) loses validity because it is derived based on the assumption of nonrelativistic decay. The limit of $\text{Br}_X \rightarrow 1$ implies $\Gamma_{N \rightarrow X} \gg \Gamma_N^{(\text{SM})}$, drastically enhancing the total decay width Γ_N . The enhanced Γ_N causes N to decay before it evolves into the nonrelativistic regime, eliminating the rising part of the red solid curve in Fig. 1. Consequently, the curve of $(\rho_N + \rho_X) a^4$ becomes almost flat. The relativistic decay of N with an almost 100% branching ratio to X simply converts ρ_N to ρ_X without gain or loss.

Before ending this section, we would like to discuss two possible variants:

- (i) X being a massive but stable species; and
- (ii) X being an unstable species.

Compared to the case of X being dark radiation, the first variant (see, e.g., Refs. [36,37]) would result in more dark-sector energy at the BBN epoch. This is due to the dilution-resistant effect of the mass—the energy stored in the form of mass cannot be diluted by the cosmological expansion, as opposed to that stored in the form of radiation—see Ref. [38] for a detailed discussion. So in this case, the solid lines in Fig. 1 would exhibit a second rise caused by the mass of X , implying that the scenario could be more easily excluded by BBN observations, though in this case its effect cannot be fully accounted for by ΔN_{eff} .

If X is unstable, the result crucially depends on whether it dominantly decays to SM particles or other dark species. For the latter, the conclusion is similar: the dark-sector energy at the BBN epoch cannot be effectively reduced, rendering the BBN constraints more restrictive. If it is the former (as in the case of HNL decays to explain the short-baseline anomalies [39–41]), one can indeed alleviate the BBN constraint, provided that most of the energy in the dark sector is released before the BBN epoch. But this essentially corresponds to opening new channels for N to release its energy into the SM sector. We will briefly comment on it later.

IV. BBN AND CMB OBSERVABLES

The effect of HNLs and their decay products on cosmological observables is twofold: (a) they contribute to the total energy density and thereby affect cosmological expansion; (b) they may directly participate in some particle physics processes that can affect cosmological observables.

For effect (a), the new physics contribution could be quantified by N_{eff} if the new particles stay relativistic during all relevant epochs from neutrino decoupling to recombination. For decaying nonrelativistic HNLs, this requirement is often not satisfied. However, within the parameter space considered in this work, HNLs must have exhaustively decayed before recombination. Assuming X is dark radiation, the impact on CMB observations can be fully accounted for by the N_{eff} used in CMB data analyses. To avoid potential confusion, we denote this N_{eff} by $N_{\text{eff}}^{\text{CMB}}$. From Planck + BAO data [35]:

$$N_{\text{eff}}^{\text{CMB}} = 2.99 \pm 0.17. \quad (7)$$

During the BBN epoch, altered cosmological expansion also influences the evolution of the fractional neutron abundance $X_n \equiv n_n/(n_n + n_p)$, where n_p and n_n denote the proton and neutron number densities, respectively. Typically, the freeze-out of neutron-proton conversion processes ($n + \nu_e \leftrightarrow p + e^-$, $n + e^+ \leftrightarrow p + \bar{\nu}_e$) is advanced due to the faster expansion caused by the new physics contribution to the total energy, giving rise to a higher freeze-out value of X_n ; see Fig. 2. Consequently, the helium abundance Y_p , which is approximately proportional to the freeze-out value of X_n , increases.

For effect (b), the most prominent example is meson-driven $n - p$ conversion. Among the decay products of

HNLs, mesons including pions and kaons have the largest cross sections with baryons. Therefore, processes like $\pi^- + p \rightarrow n + \pi^0/\gamma$ and $\pi^+ + n \rightarrow p + \pi^0$ can efficiently maintain a balance between the neutron and proton abundance, tending to pull X_n towards 1/2.

To quantitatively evaluate the new physics effects, we solve the Boltzmann equations of ρ_N , ρ_X , and X_n numerically. They are formulated as follows:

$$\frac{d\rho_N}{dt} = -3H(1+w)\rho_N - \rho_N\Gamma_{N-} + \rho_N^{(\text{eq})}\Gamma_{N+}, \quad (8)$$

$$\frac{d\rho_X}{dt} = -4H\rho_X + \rho_N\Gamma_{N-}\text{Br}_X, \quad (9)$$

$$\frac{dX_n}{dt} = -\Gamma_{n \rightarrow p}X_n + \Gamma_{p \rightarrow n}(1 - X_n), \quad (10)$$

where w is the equation of state of N , $\Gamma_{N\pm}$ are production and depletion rates of N , $\rho_N^{(\text{eq})}$ is the thermal equilibrium value of ρ_N , and $\Gamma_{n \rightarrow p(p \rightarrow n)}$ are neutron-to-proton (proton-to-neutron) conversion rates. The SM contributions to $\Gamma_{n \rightarrow p(p \rightarrow n)}$ can be found in Refs. [42,43]. Additionally, there are also meson-driven contributions from pions and kaons generated by HNL decays. We detail the calculation of these contributions in Appendix E. When solving the Boltzmann equations, we replace $dt \rightarrow da/(aH)$ with H including both SM and BSM contributions. Here we do not include backreaction terms responsible for energy transfer from X to N . The effect of such terms is expected to be insignificant as long as the temperature of the dark sector is low, which is guaranteed by N_{eff} constraints. By solving Eqs. (8)–(10), we obtain the evolutions of $\rho_{N,X}$ and X_n shown in Figs. 1 and 2, respectively.

In the SM, when neutrons and protons are in chemical equilibrium, the neutron-to-proton ratio is proportional to $e^{-Q/T}$ where $Q = m_n - m_p \approx 1.3$ MeV. The chemical equilibrium is no longer maintained when the exponential suppression $e^{-Q/T}$ becomes significant. Hence after $T \lesssim Q$, X_n approximately approaches a constant until neutron decay becomes significant. Since the neutron lifetime is $\tau_n \approx 877$ sec [44], most neutrons decay roughly at $t \sim \tau_n$, corresponding to $T \sim 0.03$ MeV according to Eq. (3). Shortly before the decay, nucleosynthesis produces deuterium (^2H) via $n - p$ fusion and then helium (^4He) from deuterium. This happens at $t \sim 200\text{--}300$ sec, represented by the vertical gray band in Fig. 2 with $T \in [0.06, 0.08]$ MeV.

The relic abundance of helium after nucleosynthesis is approximately determined by Ref. [45]

$$Y_p \approx 2X_n(t_{\text{nuc}}), \quad (11)$$

where $t_{\text{nuc}} \approx 250$ sec is the typical nucleosynthesis time.

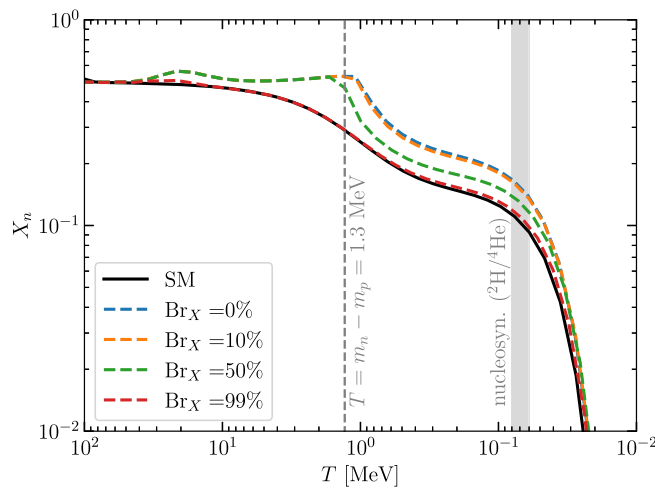


FIG. 2. The fractional neutron abundance $X_n \equiv n_n/(n_n + n_p)$ in the standard (solid line) and HNL decay scenarios (dashed lines). For the latter, we assume $m_N = 1.08$ GeV, $U_e^2 = 10^{-10}$, and $\text{Br}_X \in \{0\%, 10\%, 50\%, 99\%\}$.

In the presence of HNLs, the meson-driven processes pull X_n towards a value slightly higher than $1/2$, as can be seen from the colored dashed lines in Fig. 2. This effect persists until all HNLs and the produced mesons have decayed. The blue, orange, and green lines, corresponding to $\text{Br}_X = 0\%/10\%/50\%$, have sufficiently long-lived HNLs such that the pulling effect lasts until $T \lesssim Q$. Consequently, their X_n at t_{nuc} is significantly enhanced compared to the SM value (black line). Increasing Br_X to 99%, which implies a much shorter HNL lifetime, results in the red line, which exhibits a small deviation from the SM curve at $T \sim 20$ MeV. However, this deviation is rapidly washed out in the subsequent evolution by SM thermal processes. Nevertheless, this does not necessarily imply that all new physics effects in this case are eliminated. If the decay width of $N \rightarrow X$ is large, the background receives considerable corrections from the dark sector—see effect (a). Therefore, at $T \lesssim Q$, the red curve starts to deviate from the SM curve again. This deviation, though small (about 6.5%), is still sufficient to be observed, since current observations have measured the primordial helium abundance to a percent level [46–49]:

$$Y_P = 0.245 \pm 0.003. \quad (12)$$

We have used the PDG recommended value instead of the latest measurement of Y_P with higher precision [50] (which could have strengthened our bounds) because the latter has mild tension with the SM and Planck results.

To quantitatively evaluate how much the resulting $\Delta N_{\text{eff}}^{\text{CMB}}$ and Y_P are disfavored by cosmological observations, we employ the following χ^2 function:

$$\chi^2 = \left(\frac{\Delta N_{\text{eff}}^{\text{CMB}}}{\sigma_N} \right)^2 + \left(\frac{\Delta Y_P}{\sigma_Y} \right)^2, \quad (13)$$

where $\Delta N_{\text{eff}}^{\text{CMB}}$ and ΔY_P denote the new physics contributions to $N_{\text{eff}}^{\text{CMB}}$ and Y_P , respectively. Their uncertainties, denoted by σ_N and σ_Y , are indicated by Eqs. (7) and (12), respectively. We perform a numerical scan of the parameter space with $m_N \in [0.1, 10]$ GeV and $U_e^2 \in [10^{-13}, 10^{-6}]$ (similar results can be obtained for U_μ^2 and U_τ^2), and present the result in Fig. 3. In this plot, the yellow region is strongly disfavored by cosmological observation at a confidence limit (C.L.) above 5σ ($\chi^2 \geq 25$). A less conservative constraint is represented by the white dashed contour, corresponding to 2σ C.L. For comparison, we also add the standard 2σ BBN bound for $\text{Br}_X = 0$ on each panel, presented as the blue dashed line, which is always found to be less restrictive than the white dashed contour. For very small mixing, HNLs may never have reached thermal equilibrium. In Fig. 3, we add a thermalization line below which HNLs do not thermalize (see Appendix C for an analytical derivation). The part below this line corresponds to the freeze-in regime.

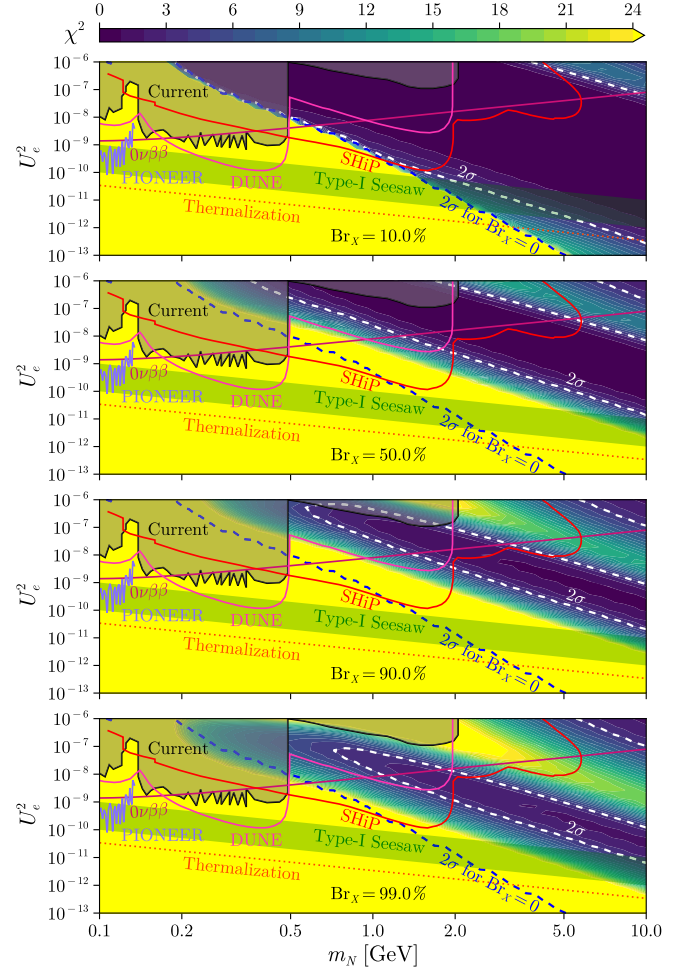


FIG. 3. The combined χ^2 values to reveal the cosmologically allowed (blue) and disfavored (yellow) parameter space for different Br_X values. The current laboratory constraints and future sensitivities, as well as the seesaw band and thermalization line, are shown for comparison.

Also shown in Fig. 3 is the Type-I seesaw band, plotted using Eq. (1) with m_ν varying from 0.01 to 0.1 eV. Targeting this theoretical benchmark, current laboratory bounds and future sensitivity reaches are also shown. The current laboratory bounds are taken from Refs. [17,51,52] presented as gray shaded regions. We also show a bound from $0\nu\beta\beta$ experiments, which is however not shaded because it can be significantly weakened in the presence of more than one HNLs [17]. For future experiments, we select three representative ones, namely, SHiP [53,54], DUNE [55–57], and PIONEER [58], as shown by the solid lines, which could potentially reach the seesaw band. However, we find that the laboratory-accessible parts of the seesaw band are strongly disfavored by cosmological data. Our results suggest that the possible loophole of introducing dark decay channels to evade the cosmological bounds on HNLs is firmly ruled out. Instead, dark decay modes strengthen the cosmological constraints.

V. DISCUSSION AND SUMMARY

We now comment on some possible scenarios that might evade the cosmological bounds on HNLs: (i) The most straightforward scenario that one can imagine is that N and X can rapidly release their energy back into the SM thermal bath before the BBN epoch. However, this scenario would imply additional interactions of the dark sector with the SM. And the interaction strength cannot be too small for it to play a significant role. This might cause N to interact more with the SM via the dark sector than via the mixing with active neutrinos. (ii) Another possible scenario is that X contains both SM and dark-sector particles. For instance, in Ref. [59], the dominant HNL decay channel is $N \rightarrow \nu a$ where a is an axionlike particle. This channel efficiently depletes the energy of N when it becomes non-relativistic, leaving only the SM thermal plasma and a in the BBN epoch, with the number density of a lower than the neutrino one. (iii) In addition to modifying HNL decay, it is also possible to relax the BBN constraint via entropy dilution [60]. As one can see from Eq. (6), if the thermal bath contains a very large number of species when N decouples (i.e., very large $g_{\star,i}$), its effect at the BBN epoch would be significantly diluted. (iv) HNLs with very small mixing may have never been in thermal equilibrium and can also evade the BBN bound. This is the freeze-in regime (see Ref. [22] for a recent study), in contrast to the freeze-out regime considered here. Note however that the photodisintegration effect on formed nuclei in this regime can be significant and may set strong constraints [61]. The very small mixing renders HNLs almost stable, allowing keV-scale HNLs to serve as dark matter [62–66]. Besides, the mixing could be suppressed by new interactions contributing to the thermal effective potential of HNLs, which has been employed to evade the cosmological bounds on eV-scale HNLs [67–72]. (v) Another way to avoid the BBN bound is to have a very low reheating temperature such that the HNL contribution to the radiation energy density could be reduced [73–76]. However, the impact of low reheating temperature on cosmological observables has to be kept in mind [77].

In conclusion, we emphasize that modern precision cosmology imposes strong bounds on HNLs in the parameter space relevant to laboratory searches. These bounds cannot be alleviated simply by introducing new dark decay channels. If future laboratory searches detect a positive HNL signal in the cosmologically disfavored regime, it would imply some nonstandard cosmology or a nontrivial HNL scenario.

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DATA AVAILABILITY

The data that support the findings of this article are openly available [78].

APPENDIX A: HNL DECAY TO SM PARTICLES

For $m_N < m_W$, HNLs can decay into various leptonic and hadronic modes, among which the most relevant ones are $N \rightarrow 3\nu$, $N \rightarrow \nu + \ell^\pm + \ell^\mp$, $N \rightarrow \ell^\mp + \mathbf{m}^\pm$, and $N \rightarrow \nu + \mathbf{m}^0$, where $\mathbf{m}^{\pm,0}$ denotes charged and neutral mesons (e.g., $\pi^\pm$ and π^0). The invisible decay mode $N \rightarrow 3\nu$ has a compact and simple expression for its decay width [79]:

$$\Gamma_{N \rightarrow 3\nu} = \frac{G_F^2 m_N^5}{192\pi^3} \sum_\alpha U_\alpha^2, \quad (\text{A1})$$

where G_F is the Fermi constant. Other decay channels are more complicated due to the involved mass thresholds and hadronic matrix elements. We refer to Ref. [79] for a comprehensive calculation of all HNL decay modes. Here we adopt a simple way to calculate the total HNL decay width for SM final states:

$$\Gamma_N^{(\text{SM})} = \Gamma_{N \rightarrow 3\nu} / \text{Br}_{N \rightarrow 3\nu}, \quad (\text{A2})$$

where $\text{Br}_{N \rightarrow 3\nu}$ is the branching ratio of the invisible mode presented in Fig. 18 of Ref. [79]. Figure 4 shows the lifetime $\tau_N = 1/\Gamma_N^{(\text{SM})}$ obtained using Eqs. (A1) and (A2), as a function of U_α and m_N . We present τ_N in both linear (left panel) and log scales (right panel) because the former is more relevant to the BBN epoch and the latter offers a broader scope, allowing us to relate the lifetime with other potentially relevant epochs in the Universe. The black dot-dashed and the white dashed lines correspond to $\tau_N = 0.1$ and 0.02 sec, respectively. Imposing the BBN bounds obtained in Refs. [20,21] on this plot, we can see that they can be well approximated by the lifetime contours of $\tau_N = 0.1$ and 0.02 sec.

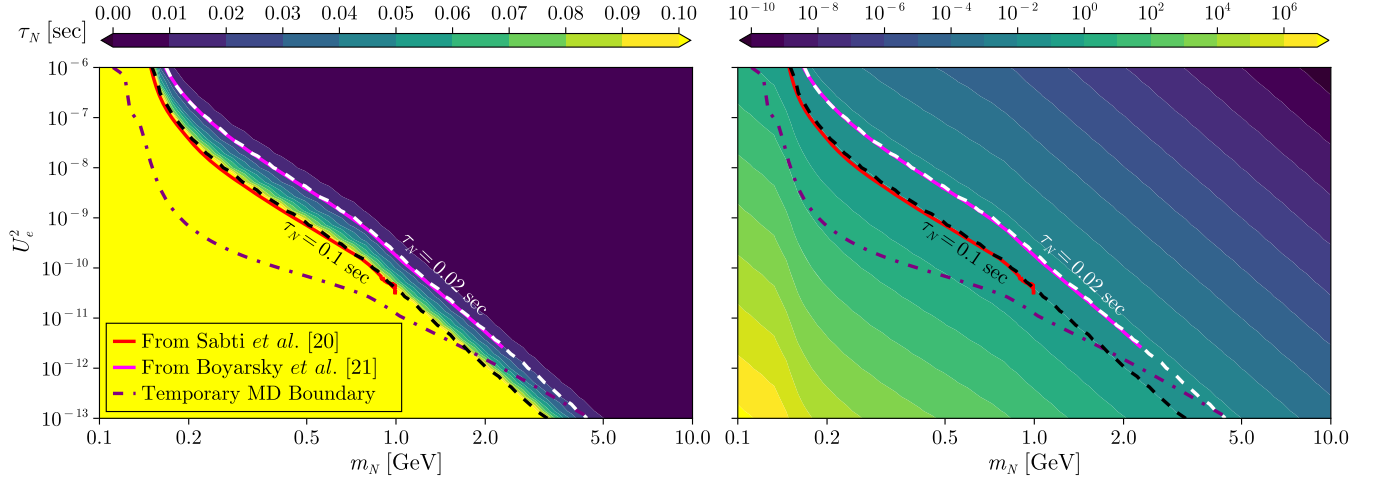


FIG. 4. The lifetime τ_N as a function of the mixing U_e , and the mass, m_N , assuming the HNL only mixes with ν_e . The left panel presents τ_N in linear scales, and the right panel in log scales. The white and black dashed lines correspond to $\tau_N = 0.02$ and 0.1 sec, respectively. The purple dash-dotted line represents the boundary of temporary matter domination caused by HNLs, obtained using Eq. (B9). The cosmological bounds from Sabti *et al.* [20] and Boyarsky *et al.* [21] can be well approximated by the two τ_N contours.

APPENDIX B: ANALYTICAL ESTIMATE IN THE FREEZE-OUT REGIME

In the freeze-out regime, HNLs decouple from the thermal bath at a certain freeze-out temperature. Below this temperature, HNLs start to free-stream from relativistic states to nonrelativistic states, and eventually decay to the SM particles and the dark sector final states X , if any.

For decoupled HNLs, when the decay is insignificant, the number density is given by

$$n_N = n_{N,i} \left(\frac{a_i}{a} \right)^3, \quad (\text{B1})$$

where the subscript i denotes an initial point which can be set at decoupling. The energy density is approximately given by

$$\rho_N \approx \begin{cases} 2 \times \frac{7}{8} \times \frac{\pi^2}{30} T_N^4 & (\text{UR}) \\ m_N n_N & (\text{NR}) \end{cases}. \quad (\text{B2})$$

Here UR and NR indicate ultrarelativistic and nonrelativistic limits, respectively. The temperature of HNLs, T_N , is determined by

$$T_N = T_i \frac{a_i}{a}, \quad (\text{B3})$$

where T_i is the temperature of HNLs at decoupling (also the temperature of the thermal bath). Note that after decoupling, T_N may be lower than the thermal bath temperature T , which is affected by the variation of $g_\star$:

$$T = \left(\frac{g_{\star,i}}{g_\star} \right)^{\frac{1}{3}} T_i \frac{a_i}{a}. \quad (\text{B4})$$

Strictly speaking, one should use the effective number of entropy degrees of freedom ($g_{\star,s}$) in the above expression. However, since the difference between $g_{\star,s}$ and $g_\star$ is small above the MeV scale, we neglect it in the analytical calculation.

To estimate the transition point from the UR to the NR regime, we equate the UR and NR expressions in Eq. (B2) and obtain the following solution:

$$a_{\text{tr}} \approx \frac{7\pi^4}{180\zeta(3)} \frac{a_i T_i}{m_N} \approx 3.15 \frac{a_i T_i}{m_N}, \quad (\text{B5})$$

where ζ is the Riemann zeta function and $\zeta(3) \approx 1.202$ arises from computing the relativistic fermion number density $n_{N,i}$. For $a \lesssim a_{\text{tr}}$ or $a \gtrsim a_{\text{tr}}$, ρ_N can be approximately calculated using the UR or NR limit, respectively.

In most cases, N decays in the nonrelativistic regime, allowing us to take the decay effect into account via

$$\rho_N \rightarrow \rho_N e^{-\Gamma_N t}, \quad (\text{B6})$$

where Γ_N is the total decay rate of N at rest. Equation (B6) only applies to nonrelativistic decay. For relativistic decay, the energy spectrum of N would be distorted, requiring a much more complex treatment—see, e.g., Appendix C in Ref. [80].

Assembling the above pieces together, we obtain the following analytical expression for ρ_N :

$$\rho_N \approx \begin{cases} \frac{7\pi^2}{120} \left(\frac{a_i}{a} \right)^4 T_i^4 & (a \lesssim a_{\text{tr}}) \\ \frac{3\zeta(3)}{2\pi^2} \left(\frac{a_i}{a} \right)^3 T_i^3 e^{-\Gamma_N t} & (a \gtrsim a_{\text{tr}}) \end{cases}. \quad (\text{B7})$$

Since ρ_N scales as a^{-3} in the nonrelativistic regime before it decays while ρ_{SM} roughly scales as a^{-4} , it is

possible that at some point ρ_N may exceed ρ_{SM} , causing a temporary period of matter domination (MD). To estimate the condition for MD, we compute the ratio ρ_N/ρ_{SM} using Eq. (B7) and find its maximum, which is given by

$$\left(\frac{\rho_N}{\rho_{\text{SM}}}\right)_{\text{max}} \approx 0.43 \frac{m_N}{g_{\star,s,i} \sqrt{\Gamma_N m_{\text{pl}}}} \left(\frac{g_{\star,s}}{g_{\star}^{3/4}}\right)_{\text{decay}}, \quad (\text{B8})$$

where the factor $(g_{\star,s}/g_{\star}^{3/4})_{\text{decay}}$ should be evaluated at $t \sim 1/\Gamma_N$.

Substituting Eq. (A2) into Eq. (B8) and solving $(\rho_N/\rho_{\text{SM}})_{\text{max}} = 1$ with respect to U_α^2 , we obtain

$$U_\alpha^2 \approx 6.8 \times 10^{-7} \text{Br}_{N \rightarrow 3\nu} \left(\frac{\text{GeV}}{m_N}\right)^3 \frac{1}{g_{\star,s,i}^2} \left(\frac{g_{\star,s}}{g_{\star}^{3/4}}\right)_{\text{decay}}^2. \quad (\text{B9})$$

This is the critical value of U_α^2 to reach a temporary MD period. In Fig. 4, we plot it as a purple dash-dotted line. Below this line, the universe undergoes a short period of MD before HNLs decay.

The energy density of the dark radiation X is governed by the following Boltzmann equation,

$$\frac{d\rho_X}{dt} + 4H\rho_X = \rho_N \Gamma_N \text{Br}_X, \quad (\text{B10})$$

where Br_X is the branching ratio of the dark decay mode. Since the left-hand side can be written as $\frac{d(\rho_X a^4)}{a^4 dt}$ and $dt = H^{-1} d \ln a$, Eq. (B10) implies

$$\rho_X a^4 = \Gamma_N \text{Br}_X \int \frac{a^4}{H} \rho_N d \ln a. \quad (\text{B11})$$

Using the NR expression in Eq. (B7) and $H \approx H_i a_i^2/a^2$, one can calculate the above integral analytically. The result reads

$$\rho_X \approx \frac{3\zeta(3)}{4\pi^2} \left(\frac{a_i}{a}\right)^4 \sqrt{\frac{\pi}{\gamma}} \text{Br}_X m_N T_i^3 \mathcal{F}\left(\frac{a}{a_i} \sqrt{\gamma}\right), \quad (\text{B12})$$

with

$$\gamma \equiv \frac{\Gamma_N}{2H_i}. \quad (\text{B13})$$

The $\mathcal{F}$ function is defined as

$$\mathcal{F}(x) \equiv \text{erf}(x) - \frac{2}{\sqrt{\pi}} x e^{-x^2}, \quad (\text{B14})$$

which varies from 0 to 1 for $x \in [0, \infty)$.

The above calculation can be straightforwardly adapted for computing the energy injection into the SM thermal plasma. Since the branching ratio of $N \rightarrow \text{SM}$ is simply

$1 - \text{Br}_X$, the SM energy density receives the following contribution:

$$\Delta\rho_{\text{SM}} \approx \frac{3\zeta(3)}{4\pi^2} \left(\frac{a_i}{a}\right)^4 \sqrt{\frac{\pi}{\gamma}} (1 - \text{Br}_X) m_N T_i^3 \mathcal{F}\left(\frac{a}{a_i} \sqrt{\gamma}\right). \quad (\text{B15})$$

In particular, the neutrino energy density of a single flavor is given by

$$\rho_\nu = \frac{7\pi^2}{120} \left(\frac{g_{\star,i}}{g_{\star}}\right)^{\frac{1}{3}} T_i^4 \left(\frac{a_i}{a}\right)^4 + \frac{2 \times \frac{7}{8}}{g_{\star}} \Delta\rho_{\text{SM}}. \quad (\text{B16})$$

Equation (B16) is valid only before neutrino decoupling. After neutrino decoupling, ρ_ν scales as a^{-4} , provided that the decay has effectively stopped. If after neutrino decoupling there is still a significant amount of HNLs decaying, then one needs to further calculate the fraction of $\Delta\rho_{\text{SM}}$ that is injected into the neutrino sector. This is complicated for an analytical estimate, but straightforward for numerical implementation.

The ratio ρ_X/ρ_ν after neutrino decoupling can be directly interpreted as the new physics contribution to N_{eff} . By combining the above analytical results, we also get an estimate of the contribution to N_{eff} :

$$\Delta N_{\text{eff}} \approx \frac{\text{Br}_X}{0.16(1 - \text{Br}_X) + 0.15 g_{\star,i}^{4/3} \gamma^{1/2} \frac{T_i}{m_N}}. \quad (\text{B17})$$

Using $\Gamma_N = \Gamma_N^{(\text{SM})}/(1 - \text{Br}_X)$, and γ from Eq. (B13) with $H_i \approx 1.66 g_{\star}^{1/2} T_i^2/m_{\text{pl}}$, it is straightforward to obtain Eq. (6).

APPENDIX C: THE FREEZE-IN REGIME AND THERMALIZATION

When the mixing is sufficiently small, HNLs cannot thermalize, implying that they are in the freeze-in regime. While both freeze-out and freeze-in are included in our numerical code that directly solves the relevant Boltzmann equations, here we would like to provide an analytical estimate of the condition for thermalization.

The production rate of HNLs is approximately given by (see, e.g., Ref. [37])

$$\Gamma_{N,\text{prod}} \approx \theta_{\text{eff}}^2 \Gamma_{\nu_L,\text{prod}}, \quad (\text{C1})$$

with

$$\Gamma_{\nu_L,\text{prod}} \approx 3c_1 G_F^2 T^5, \quad (\text{C2})$$

$$\sin^2 2\theta_{\text{eff}} \approx \frac{\sin^2 2\theta}{\sin^2 2\theta + (\cos 2\theta - 6TV/m_N^2)^2}, \quad (\text{C3})$$

$$V \approx -c_2 G_F^2 T^5. \quad (\text{C4})$$

Here θ_{eff} is the effective mixing angle of N with SM neutrinos (ν_L). It is different from the in-vacuum mixing angle $\theta \approx U_\alpha^{1/2}$ due to the thermal MSW potential V in Eq. (C3). The two coefficients c_1 and c_2 have a mild temperature dependence; however, for a simple estimate we treat them as constants, taking $c_1 \approx 1.27$ and $c_2 \approx 2.5 \times 10^2$. The production rate is to be compared with H . Hence we define

$$\gamma_{N,\text{prod}} \equiv \frac{\Gamma_{N,\text{prod}}}{H} \approx \frac{3c_1 G_F^2 T^3 m_{\text{pl}} \sin^2 2\theta}{4g_H((6c_2 G_F^2 m_N^{-2} T^6 + \cos 2\theta)^2 + \sin^2 2\theta)}, \quad (\text{C5})$$

where $g_H \equiv m_{\text{pl}} H / T^2 \approx 1.66 g_\star^{1/2}$. The condition for thermalization corresponds to $\gamma_{N,\text{prod}} \gtrsim 1$. Note that $\gamma_{N,\text{prod}}$ asymptotically approaches T^{-9} and T^3 at very high and low T , respectively, implying that it peaks at an intermediate temperature. Taking $\cos 2\theta \rightarrow 1$ and $\sin 2\theta \rightarrow 0$ in the denominator, the peak temperature and the corresponding peak value of $\gamma_{N,\text{prod}}$ can be calculated analytically:

$$T^{(\text{peak})} \approx \left(\frac{m_N}{3\sqrt{2}c_2 G_F} \right)^{\frac{1}{3}} \approx 11 \text{ GeV} \times \left(\frac{m_N}{\text{GeV}} \right)^{\frac{1}{3}}, \quad (\text{C6})$$

$$\gamma_{N,\text{prod}}^{(\text{peak})} \approx \frac{9c_1 G_F m_N m_{\text{pl}} U_\alpha^2}{16g_H \sqrt{2}c_2}. \quad (\text{C7})$$

From Eq. (C7), one can see that the condition for thermalization corresponds to

$$U_\alpha^2 \gtrsim \frac{16g_H \sqrt{2}c_2}{9c_1 G_F m_N m_{\text{pl}}} \approx 3 \times 10^{-12} \times \left(\frac{g_\star}{83.5} \right)^{\frac{1}{2}} \times \frac{\text{GeV}}{m_N}. \quad (\text{C8})$$

Here, $g_\star$ is evaluated at the peak temperature. For $m_N = 1 \text{ GeV}$, it is 83.5, which is the value of $g_\star$ at peak temperature of 11 GeV, as suggested by Eq. (C6). For smaller or larger m_N , $g_\star$ needs to be adjusted accordingly but the variation is not significant (varying from 82 to 87 for $0.1 < m_N/\text{GeV} < 10$) within the parameter space considered in this work.

APPENDIX D: THE ABUNDANCE OF HNLS AT $T = 1 \text{ MeV}$

The abundance of HNLS at $T = 1 \text{ MeV}$ is useful for understanding their impact on BBN, since this temperature is close to neutron freeze-out. If the abundance of HNLS at this point is high, the neutron-to-proton ratio is expected to be significantly affected. In Fig. 5, we present the ratio $R_{N/\nu} \equiv \rho_N/\rho_\nu$ at $T = 1 \text{ MeV}$, obtained from our numerical calculation assuming $\text{Br}_X = 0$. This ratio is exponentially suppressed in the large-mixing regime (above the upper branches of the orange curves) due to fast decay. For very small mixing, it enters the freeze-in regime, in which this ratio is suppressed by the production of HNLS.

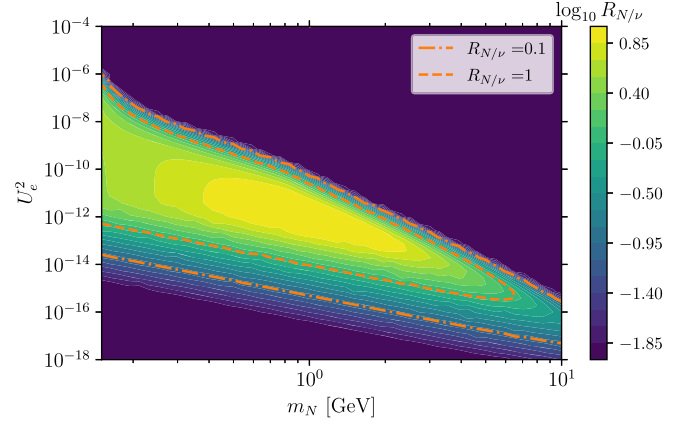


FIG. 5. The ratio $R_{N/\nu} \equiv \rho_N/\rho_\nu$ at $T = 1 \text{ MeV}$ as an indicator of the abundance of HNLS around neutron freeze-out.

APPENDIX E: MESON-DRIVEN $p - n$ CONVERSION

One of the most significant effects of HNLS on BBN is that mesons produced from N decay may cause meson-driven $p - n$ conversion. Taking $\pi^\pm$, for example, the following processes are possible:

$$\pi^- + p \rightarrow n + \pi^0/\gamma, \quad \pi^+ + n \rightarrow p + \pi^0. \quad (\text{E1})$$

Due to the Sommerfeld enhancement, meson-driven $p \rightarrow n$ conversion processes are more efficient than meson-driven $n \rightarrow p$ conversion processes.

Note that the lifetime of $\pi^\pm$, $\tau_{\pi^\pm} \approx 2.6 \times 10^{-8} \text{ sec}$, is much shorter than the timescale of the early Universe around the BBN epoch. Due to the short lifetime, if all processes for $\pi^\pm$ production have effectively ceased, the abundance of $\pi^\pm$ would decrease rapidly via decay. In the presence of HNLS, which serve as a steady source of $\pi^\pm$, the number density $n_{\pi^\pm}$ can maintain a balanced value [21]:

$$n_{\pi^\pm} \Gamma_{\pi^\pm} = n_N \Gamma_{N \rightarrow \ell^\mp + \pi^\pm}. \quad (\text{E2})$$

The depletion of $\pi^\pm$ via Eq. (E1) is negligible compared to that via decay. Equation (E2) implies that the balanced value of $n_{\pi^\pm}$ is proportional to the lifetime of $\pi^\pm$. This further implies that the inverse processes of Eq. (E1), i.e., $n + \pi^0 \rightarrow \pi^- + p$ and $p + \pi^0 \rightarrow \pi^+ + n$, are unimportant because π^0 has a much shorter lifetime ($\tau_{\pi^0} \approx 8.5 \times 10^{-17} \text{ sec}$).

The effect of the processes in Eq. (E1) on BBN can be taken into account by adding the following contributions to $\Gamma_{n \rightarrow p}$ and $\Gamma_{p \rightarrow n}$ in the Boltzmann equation of X_n [cf. Eq. (10)]:

$$\Gamma_{n \rightarrow p}^{(\pi^+)} = n_{\pi^+} \langle \sigma_{n \rightarrow p}^{\pi^+} v \rangle, \quad (\text{E3})$$

$$\Gamma_{p \rightarrow n}^{(\pi^-)} = n_{\pi^-} \langle \sigma_{p \rightarrow n}^{\pi^-} v \rangle, \quad (\text{E4})$$

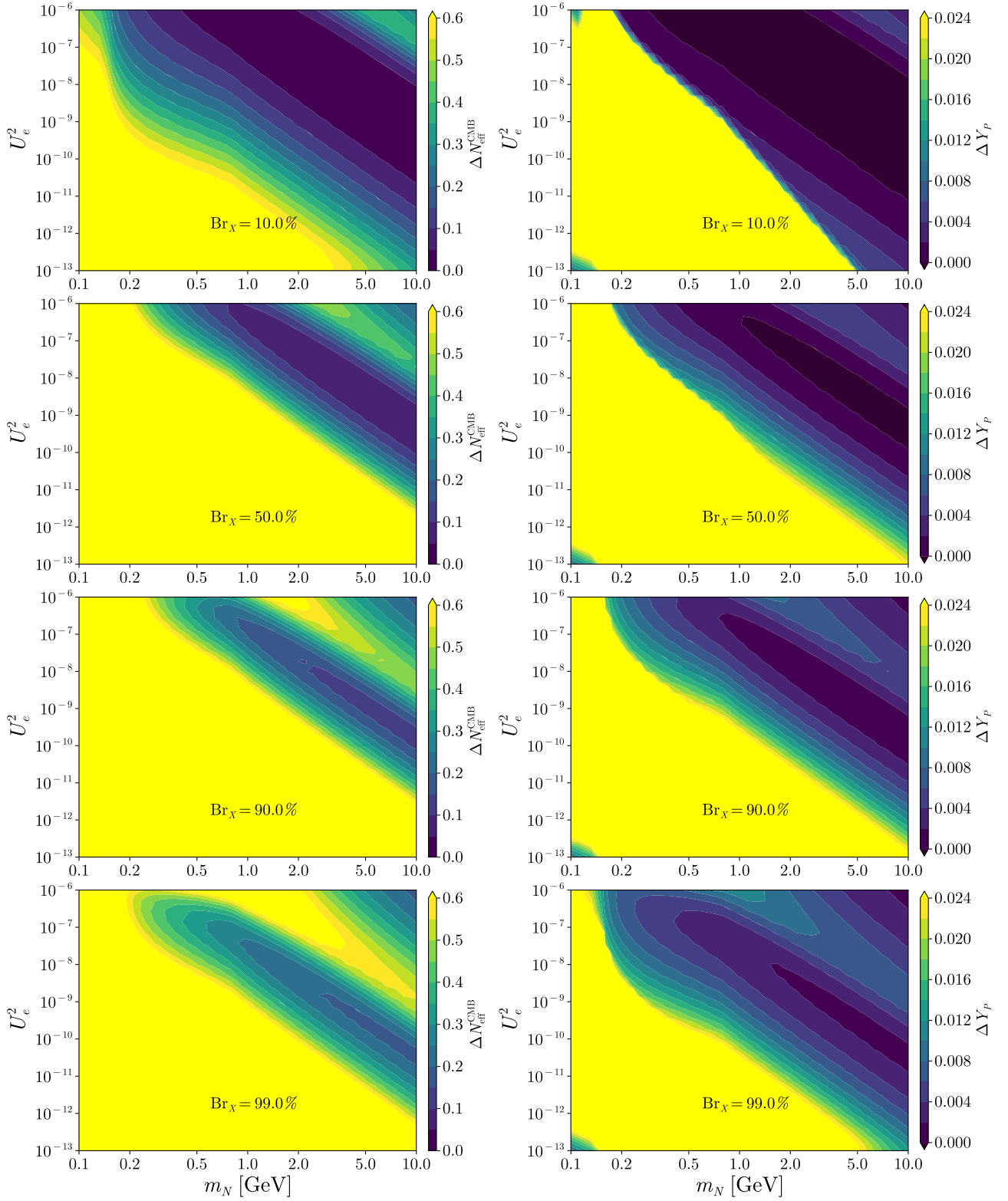


FIG. 6. The CMB and BBN observables, $\Delta N_{\text{eff}}^{\text{CMB}}$ and Y_p , affected by HNLs in the presence of a dark-decay channel with the branching ratio Br_X .

where the thermally averaged cross sections $\langle\sigma_{n\rightarrow p}^{\pi^+}v\rangle$ and $\langle\sigma_{p\rightarrow n}^{\pi^-}v\rangle$ can be found in Ref. [81]:

$$\langle\sigma_{n\rightarrow p}^{\pi^+}v\rangle = 1.7 \text{ mb} = 4.4 \times 10^{-18} \text{ eV}^{-2}, \quad (\text{E5})$$

$$\langle\sigma_{p\rightarrow n}^{\pi^-}v\rangle/C_\pi = 1.5 \text{ mb} = 3.9 \times 10^{-18} \text{ eV}^{-2}. \quad (\text{E6})$$

Due to the opposite charges of the initial particles in $\pi^- + p \rightarrow n + \pi^0/\gamma$, Equation (E6) contains a Sommerfeld enhancement factor: $C_\pi = \epsilon/(1 - e^{-\epsilon})$ with $\epsilon = 2\pi\alpha\sqrt{\frac{m_\pi m_p}{2T(m_\pi + m_p)}}$, where α is the fine-structure constant. We note here that similar results can also be found in Refs. [21,82], despite a minus sign typo in Eq. (D3) of Ref. [21].

APPENDIX F: CONSTRAINTS FROM $N_{\text{eff}}^{\text{CMB}}$ AND Y_P SEPARATELY

As discussed in the main text, HNL decays to dark sector particles affect both $N_{\text{eff}}^{\text{CMB}}$ and Y_P , and therefore, can be separately constrained by these observables. This is presented in Fig. 6, where we plot the values of $\Delta N_{\text{eff}}^{\text{CMB}}$ and

ΔY_P in left and right columns, respectively. The figure shows that Y_P is more restrictive for small Br_X , while $N_{\text{eff}}^{\text{CMB}}$ is more restrictive for large Br_X . A remarkable feature of the $N_{\text{eff}}^{\text{CMB}}$ constraint is that as Br_X increases, it even starts ruling out larger U_e^2 regions (e.g., $m_N \sim 2 \text{ GeV}$, $|U_e|^2 \sim 10^{-6}$), in addition to the small U_e^2 region as in the $N \rightarrow \text{SM}$ case. This is because a large U_e together with a large $\Gamma_{N\rightarrow X}$ allows N to produce X while still being relativistic and in thermal equilibrium. When $\Gamma_{N\rightarrow X}$ is sufficiently large, X reaches thermal equilibrium before N decouples. Consequently, the contribution of X to $N_{\text{eff}}^{\text{CMB}}$ becomes significant, even for short-lived N . Similarly, the right panels show that this region is also disfavored by Y_P since the effect of X on the background evolution during BBN becomes non-negligible. Another notable feature is that the bottom left corners of Y_P plots show relatively small values of ΔY_P , implying that it is in the freeze-in regime. Therefore, the BBN bound becomes weaker as U_e decreases in this region. This roughly corresponds to the lower edge of the BBN bound shown in Ref. [22]. The combined constraint from $N_{\text{eff}}^{\text{CMB}}$ and Y_P measurements is shown in Fig. 3.

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