Dark radiation from neutrino mixing after Big Bang Nucleosynthesis

Dark Radiation from Neutrino Mixing after Big Bang Nucleosynthesis

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A light ($m_{\nu d} \lesssim$ MeV) dark fermion mixing with the Standard Model neutrinos can naturally equilibrate with the neutrinos via oscillations and scattering. In the presence of dark sector interactions, production of dark fermions is generically suppressed above BBN, but then enhanced at later times. Over much of the parameter space, we find that the dark sector equilibrates, even for mixing angles θ_0 as small as 10^{-13}, and equilibration occurs at $T_{\rm equil} \simeq m_{\nu d}\left(\theta_0^2 M_{Pl}/m_{\nu d}\right)^{1/5}$ which is naturally at most a few orders of magnitude above the dark fermion mass. The implications of this are twofold: one, that light states are often only constrained by the CMB and LSS without leaving an imprint on BBN, and two, that sectors which equilibrate before recombination will typically have a mass threshold before recombination, as well. This can result in dark radiation abruptly transitioning from non-interacting to interacting, or vice-versa, a "step" in the amount of dark radiation, and dark matter with similar transitions in its interactions, all of which can leave important signals in the CMB and LSS, and may be relevant for cosmological tensions in observables such as H_0 or S_8. Minimal models leave an unambiguous imprint on the CMB above the sensitivity of upcoming experiments.

I. INTRODUCTION

The range of redshifts between $z \sim 10^9$ and $z \sim 10^3$ correspond to a great "desert" in ΛCDM. As the temperature cools below the MeV scale where Big Bang Nucleosynthesis (BBN), neutrino decoupling, and $e^+ e^-$ annihilation take place no new threshold is reached for almost six orders of magnitude until the eV scale where matter-radiation equality, CMB decoupling, and eventually the sum of the neutrino masses can be found. The ΛCDM desert originates from a coincidence of the large gap in the mass spectrum of the Standard Model between the electron mass and the scale of neutrino masses with the unrelated but perfectly overlapping gap between nuclear and atomic binding energies.

However, additional dark sectors can have new particles with masses in or below these scales, possibly leading to a rich phenomenology during this desert period. A simple extension of the standard model that can realize this involves the presence of an additional neutral fermion ν_d, with mass $m_{\nu d}$, which mixes with the SM neutrino via a small Dirac mass. A combination of oscillations and weak interaction scattering can easily populate this species for large enough mixing. The relevant rate of this process Γ/H peaks near $T \sim 100\,\mathrm{MeV}$ [1], yielding a fully thermalized fermion $\triangle N_{\rm eff} \approx 1$ for mixing angles $\sin\theta \gtrsim 10^{-3}$ for a keV-mass state. This additional radiation affects BBN and is excluded from the measurements of light element abundances which require $\triangle N_{\rm eff}|_{T\sim 1\mathrm{MeV}} \leq 0.407\,(95.45\%)$ [2]. Smaller mixing angles yield dark fermions which are unthermalized and cosmologically uninteresting as radiation (absent a population from pre-TeV processes), and highly constrained as dark matter.

Such a picture raises many questions, in particular regarding the origin of the new particle's mass. A natural expectation would be that the mass arises from some dynamics, and there would be other particles and interactions, such as self-interactions, connected to it. The consequences of such an interaction can be significant. Light fermions with large mixings can have their oscillations suppressed in the early universe [3–5], changing the cosmological constraints significantly. In the presence of a self-interaction regions of parameter space arise where a ∼ keV fermion with small mixings can be dark matter [6–9]. In contrast, absent self-interactions, direct production of such dark matter through weak interactions is excluded by a combination of X-ray data and the presence of small scale structure [10][1]. Thus it is clear that a dark fermion with interactions is qualitatively different from the "unnaturally minimal" scenario of an inert dark state.

Upcoming CMB and LSS observations will probe the ΛCDM desert, motivating a broader exploration of such models. In this letter we study the equilibration of dark sectors with the SM neutrinos after BBN due to the potential impact on these observations. For concreteness, we consider a single dark fermion ν_d which mixes with a SM neutrino by an amount $\sin\theta_0$ in vacuum. We further assume that ν_d has a self-interaction mediated by a force carrier ϕ with $m_\phi \ll m_{\nu d}$ and coupling strength α_d. In studying this simple setup we find two important results:

- The dark sector comes into equilibrium with the neutrinos over a very large parameter space roughly bounded only by $\theta_0^2 \alpha_d^2 M_{Pl} > m_{\nu d}$, allowing mixing angles ranging from 1 to 10^{-13}.

[1] A famous loophole exists when SM neutrinos have chemical potentials and a lepton asymmetry [11].

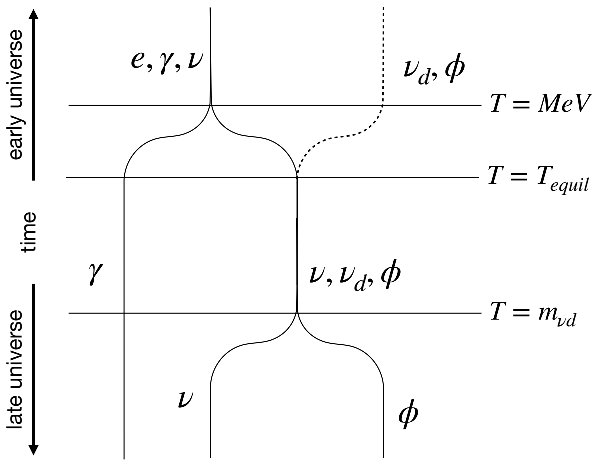

FIG. 1: Thermal history of a universe with dark sector thermalization from neutrino mixing after BBN. The dark sector intially has negligible energy density (dashed line). After neutrino decoupling and electron annihilation it equilibrates with the SM neutrinos at $T_{\rm equil}$. After ν_d annihilation at $T \sim m_{\nu d}$ the SM neutrinos re-decouple and free-stream. The remaining dark sector, ϕ, and any other DS particles may continue to interact or free-stream.

- Over most of the parameter space the temperature at which ν_d equilibrates is α_d-independent and given by

$$T_{\rm equil} \simeq m_{\nu d} \left(\theta_0^2 \frac{M_{Pl}}{m_{\nu d}}\right)^{1/5} \quad . \tag{1}$$

Thus even though the range of allowed values of θ_0 and α_d is huge, ν_d naturally equilibrates at temperatures near $m_{\nu d}$, and at most a few orders of magnitude higher, because of the $1/5$ power. Consequently, dark sectors with light ($<$ MeV) fermions often equilibrate after BBN and are therefore unconstrained by primordial light element abundances.

The simplest thermal history is sketched in Fig. 1. After neutrino decoupling and electron self-annihilation at $T \sim$ MeV, the dark sector ϕ and ν_d come into equilibrium with the SM neutrinos. At the lower temperature $T \sim m_{\nu d}$, the dark fermions ν_d annihilate away. This causes the SM neutrinos to decouple and become free-streaming again, and the entropy of ν_d is shared between ϕ and the SM neutrinos.

Importantly, dark sector equilibration with SM neutrinos *after* neutrino decoupling does not change the relativistic energy density because the total energy in neutrinos + dark sector is conserved in the equilibration process. Thus, constraints on $N_{\rm eff}$ from the CMB and LSS do not *a priori* constrain equilibration. However, in this thermal history, prior to $T \sim \nu_d$, the $\nu - \nu_d - \phi$ fluid is tightly coupled, and below $m_{\nu d}$ there is a"step" [12–14] in the total energy density as ν_d annihilates away. Thus, there is an inevitable imprint on the density perturbations of the universe.

Should other particles have couplings to ϕ and ν_d, they, too, will naturally come into equilbrium with the SM neutrinos below $T \sim$ MeV. As a result, there is a possibility for other interesting dynamics within a dark sector to affect cosmology, such as the thermalization and freeze-out of dark matter, the presence of a second "step" [12–14] in the energy density of the dark sector due to the annihilation of additional massive particles into lighter ones. Alternatively, in a minimal scenario with $m_{\nu d} \lesssim$ eV, self-interactions in (a portion of) the relativistic energy density may arise only at late times, near recombination. Neutrino-dark sector equilibration after BBN thus has very interesting and model-dependent impact on the CMB and structure formation with possible implications for H_0 and S_8, all of which will be probed by a wide range of upcoming experiments.

II. INTERACTIONS AND DARK SECTOR THERMALIZATION

A generic dark sector which contains a fermion ν_d that mixes with the SM neutrinos can equilibrate with the SM neutrinos very efficiently by the combined effect of $\nu - \nu_d$ oscillations and scattering. The relevant formalism is well-developed, see [15–17]. For simplicity we consider the case of one dark fermion oscillating with one SM neutrino. The rate of conversion of a SM neutrino into a dark fermion can be written as

$$\Gamma(E) = \frac{1}{2} \sin^2 2\theta_m \frac{\Gamma_{\rm int}}{2} \quad . \tag{2}$$

where we assume averaging over many oscillations, Γ_{int} is the rate of scattering, θ_m is the in-medium mixing angle between the SM neutrino and the dark fermion and both depend on the incoming neutrino energy E. The process of dark sector equilibration is the usual competition between the production rate in Eq. (2) and Hubble. The mixing angle is generally suppressed by the presence of large diagonal effective thermal masses and thus the overall conversion rate grows rapidly as T declines, both in an absolute sense and in comparison to Hubble, even as the scattering rate drops, as we see below.

The in-medium mixing angle is given by

$$\sin^2 2\theta_m = \frac{\sin^2 2\theta_0}{(\cos 2\theta_0 - 2E \triangle V_{\rm eff}/\triangle m^2)^2 + \sin^2 2\theta_0} \ , \tag{3}$$

where θ_0 is the in-vacuum angle that parameterizes the mixing between the SM neutrino and the dark fermion, $\triangle m^2 \simeq m_{\nu d}^2$ is the mass-squared difference between the two mass eigenstates and is dominated by the dark fermion mass, and $\triangle V_{\rm eff} = V_{\rm eff}^{SM} - V_{\rm eff}^{DS}$. The effective potential of ν from the SM weak interactions is well-known [1] and given by $V_{\rm eff}^{SM} \simeq -c_V G_F^2 T_\nu^4 E$ where $c_V \simeq 22$ (for mixing with ν_μ or ν_τ), and we assume vanishing lepton asymmetry [11]. The dark sector effective potential arises due to scattering with light particles and a light mediator in the dark thermal bath and can be parameterized as $2EV_{\rm eff}^{DS} \equiv \alpha_d T_d^2$ [4]. In what follows we take this as the definition of α_d. The expression for the effective potential (and interaction discussed below) assumes that the dark sector is self-equilibrated with temperature T_d and vanishing chemical potentials.[2] The exact expression can vary with Dirac/Majorana, internal symmetries and other model dependencies which amount to an overall O(few) rescaling of α_d. The precise mapping onto a specific model Lagrangian is straightforward and not important for our discussion.

The scattering rate is the sum of the SM weak interaction $\Gamma_{SM} = n_\nu \langle\sigma v\rangle_{SM} = c_\Gamma T_\nu^4 G_F^2 E$ with $c_\Gamma \simeq 0.92$ [1], and the scattering rate of the dark fermions which we parameterize as $\Gamma_{\rm DS} = n_{\rm DS}\langle\sigma v\rangle_{DS} \equiv \kappa\, \alpha_d^2 T_d^2/E$. This assumes that the cross section scales as $\langle\sigma v\rangle_{DS} \simeq \langle \kappa\alpha_d^2/E_{\rm CM}^2\rangle_{DS} \simeq \kappa\alpha_d^2/(ET_d)$ and $n_{\rm DS} \propto T_d^3$. Here κ is a number greater than one, which allows for the presence of additional dark states which scatter via ϕ exchange. For simplicity, we set $\kappa = 3$, and in general it would shift the precise region of parameter space but not make it much larger or smaller.

Finally, averaging the conversion rate Γ over the thermal distribution of the SM neutrinos approximately replaces $E \to 3T_\nu$ so that

$$\langle\Gamma\rangle = \frac{\frac{1}{4}\sin^2 2\theta_0 (3c_\Gamma T_\nu^5 G_F^2 + \alpha_d^2 \frac{T_d^2}{T_\nu})}{\left(\cos 2\theta_0 + \alpha_d \frac{T_d^2}{m_{\nu d}^2} + 18 c_V \frac{G_F^2 T_\nu^6}{m_{\nu d}^2}\right)^2 + \sin^2 2\theta_0} . \quad (4)$$

Armed with this expression for the conversion rate we can now determine if and when the dark sector equilibrates with the neutrinos by comparing Γ with the expansion rate, $H \simeq T_\nu^2/M_{Pl}$. There are two important limits to consider. First, in the Dodelson-Widrow [1] limit of vanishing dark sector interactions, $\alpha_d = 0$, the maximum conversion rate occurs when $G_F T_\nu^3/m_{\nu d} \sim 0.1$. This peak temperature is above an MeV so that full equilibration from DW would yield a thermalized dark sector before BBN which is excluded. The dark sector equilibrates if $\Gamma = H$ at the peak, therefore we obtain the constraint (in the DW limit) that $\theta_0^2 m_{\nu d} M_{Pl} G_F \lesssim 100$.

A qualitatively different solution is obtained when the dark sector interactions dominate over the weak interactions. Then $\langle\Gamma\rangle/H$ grows monotonically with decreasing temperature, and we can solve for the equilibration temperature (when $T_d = T_\nu$) by setting

$$1 \simeq \frac{\langle\Gamma\rangle}{H} \simeq \frac{\theta_0^2\alpha_d^2 T_\nu}{(1+\alpha_d\frac{T_\nu^2}{m_{\nu d}^2})^2}\frac{M_{Pl}}{T_\nu^2} \simeq \theta_0^2 \frac{M_{Pl}}{m_{\nu d}}\frac{m_{\nu d}^5}{T_\nu^5} , \quad (5)$$

giving $T_{\rm equil} = m_{\nu d}\,(\theta_0^2 M_{Pl}/m_{\nu d})^{1/5}$. It is remarkable both that this is independent of α_d and the dependence on $\theta_0^2 M_{pl}$ is mild because of the $1/5$ power. Thus for a very broad range in parameter space the dark sector equilibrates with the neutrinos, and it does so at a temperature which is at most a few orders of magnitude above the dark fermion mass. This yields the important qualitative result that in the presence of a light ($\ll$ MeV) fermion, the natural equilibration scale is *below* the BBN scale,[3] but also above recombination, as shown in the bulk regions of Fig. 3. Note that the small "fin" regions on the right of Fig. 3 correspond to parameter space in which $\alpha_d T_{\rm equil}^2/m_{\nu d}^2 < 1$.

This intuition is borne out by a numerical integration of the Boltzmann equation, as we now discuss. We assume that the dark sector starts out cold and initial dark abundances arise from Higgs decay and weak interactions. A small additional population below a thermal abundance does not significantly change the results.

One starts with the Boltzmann equation for the phase space distribution function of the dark fermions, which we can integrate and simplify under the assumption the dark sector is self-thermalized,

$$\frac{d}{d\log a}\left(a^4\rho_{\rm DS}\right) = \frac{\langle\Gamma\rangle}{H}a^4\left(\rho_\nu - \left.\frac{\rho_\nu}{\rho_{\rm DS}}\right|_{eq.}\rho_{\rm DS}\right) , \quad (6)$$

where $\rho_{\rm DS}$ is the total energy density in the DS including ν_d, ϕ and any other (relativistic) particles that the DS may have. Note that we added a simple back-reaction term on the right-hand side which accounts for conversions $\nu_d \to \nu$ and vanishes in equilibrium. This term is negligible until the dark sector and the SM neutrinos are close to equilibration and roughly models the correct approach to equilibrium. Any interactions of ν_d with dark sector particles can redistribute energy within the dark sector but because of energy conservation they cannot contribute to the evolution of the total energy density of the DS in (6).

The assumption that the DS is self-thermalized which allowed us to write $\langle\Gamma\rangle$ as a function of the dark sector temperature is not necessarily true. At early times, when DS particle number densities are still small the DS self-interactions could be smaller than Hubble. However, the DS always reaches kinetic equilibrium before it equilibrates with the neutrinos and in most of parameter space number-changing interactions in the DS also erase any chemical potentials before equilibration with the neutrinos. For simplicity, we assume that the DS self-thermalizes rapidly in Figs. 2 and 3. For small α_d this may require additional interactions which could be in $V(\phi)$ or involve additional dark sector particles.

The evolution of the DS energy density depends on the model parameters. Fig. 2 shows the DS temperature

[2] We discuss the assumption that the dark sector is self-thermalized later. We also ignore a possible shift of the scalar expectation value in the thermal background which would change the mass of ν_d.

[3] A similar phenomenology can be achieved in models of neutrinos which couple to a Majoron, and resonantly produce dark matter at late times [18].

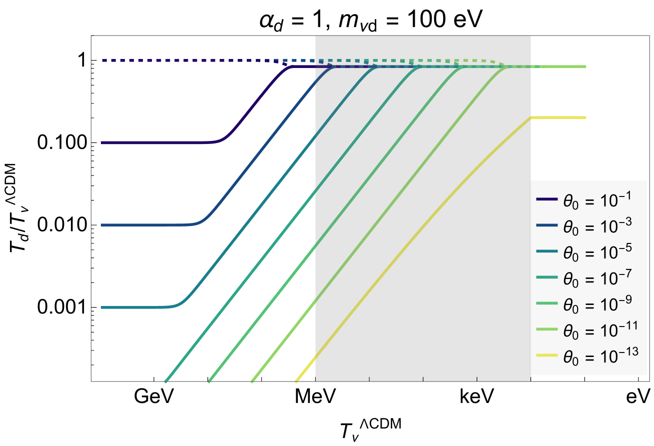

FIG. 2: The figure shows the ratio $T_d/T_\nu^{\Lambda\mathrm{CDM}}$ obtained from solving Eq. (6) as a function of $T_\nu^{\Lambda\mathrm{CDM}}$, the temperature of the active neutrinos in a reference ΛCDM (with no dark sector). For a large range of mixing angles the dark sector equilibrates after BBN and before $T_\nu^{\Lambda\mathrm{CDM}} = m_{\nu d}$, in the gray region. Dashed lines show the reduction of the neutrino temperature due to DS equilibration relative to the ΛCDM neutrino temperature $T_\nu/T_\nu^{\Lambda\mathrm{CDM}}$. In this figure we assumed equilibration with all three SM neutrinos and a minimal dark sector of one Majorana fermion and a real scalar, thus $\rho_\nu/\rho_{\mathrm{DS}}|_{eq.} = g_*^\nu/g_*^{\mathrm{DS}} = 21/11$.

evolution resulting from integrating Eq. (6) for different values of θ_0 with fixed $m_{\nu d} = 100$ eV and $\alpha_d = 1$. (We note that $\alpha_d \simeq 1$ may require higher orders in perturbation theory for precise predictions. Nevertheless, we use it as an example because it allows the largest range of angles θ_0 to equilibrate, see Fig. 3). The Figure shows $T_d/T_\nu^{\Lambda\mathrm{CDM}}$ as a function of a reference ΛCDM neutrino temperature in solid lines. At equilibration, T_d merges with the SM neutrino temperature which is shown with dashed lines. Equilibration in the gray region corresponds to equilibration between BBN ($T_\nu \simeq$ MeV) and $T_\nu = m_{\nu d}$. If equilibration is not reached before $T_\nu = m_{\nu d}$ it is never reached because the interaction $\langle\Gamma\rangle/H$ rapidly shuts off for $T_\nu < m_{\nu d}$.

Our primary result is contained in Fig. 3 which shows the large regions of parameter space where the dark sector comes into equilibrium with the SM neutrinos at some point before $T_\nu = m_{\nu d}$ (the lower boundary) and where equilibration is reached below $T_\nu =$ MeV, i.e. after neutrino decoupling and BBN. For the purposes of this Figure we define the equilibration temperature T_{equil} as the temperature at which ρ_{DS} crosses $\rho_\nu\, g_*^{\mathrm{DS}}/g_*^\nu$ with ρ_{DS} obtained from solving (6) with the backreaction term omitted.

It is worth noting that because of mixing of the SM neutrinos, for most of parameter space all three SM neutrinos equilibrate with the DS in rapid succession. That only a single SM neutrino equilibrates with the DS can occur for special regions in parameter space. Either the couplings of ν_d are tuned such that it only couples to a single SM neutrino mass eigenstate, or the dark parameters are such that equilibration with the first of the SM neutrinos occurs at a temperature just above $m_{\nu d}$ so that $\nu - \nu_d$ conversion shuts off because $m_{\nu d}$ is reached before another SM neutrino can equilibrate.

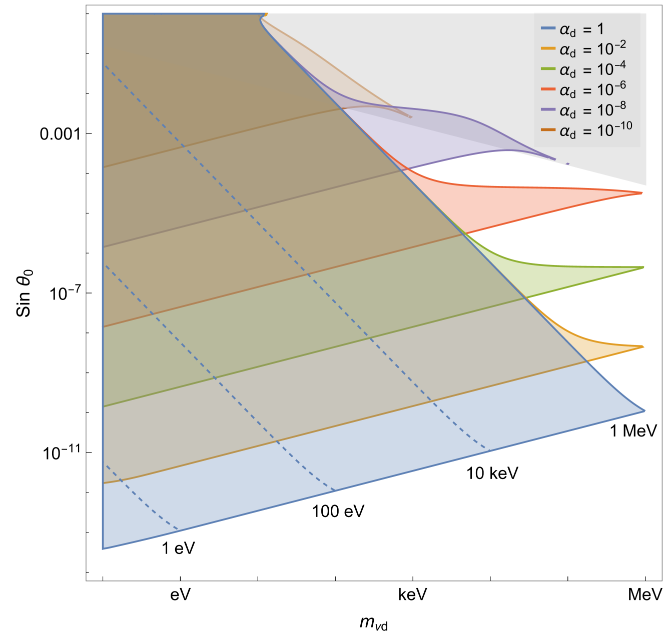

FIG. 3: Colored regions indicate the parameter space over which the dark sector comes into equilibrium with the SM neutrinos after BBN, for different values of α_d. The lower boundary of each region is determined by $T_{\mathrm{equil}} = m_{\nu d}$, while the upper (right) boundary comes corresponds to equilibration at BBN, i.e. $T_{\mathrm{equil}} =$ MeV. Also shown are contours of fixed equilibration temperatures T_{equil} (dashed contours labeled 1 eV, 100 eV, 10 keV, 1 MeV) for the $\alpha_d = 1$ case. The gray region in the upper right shows the parameter space over which equilibration would occur above BBN in absence of dark interactions ($\alpha_d = 0$) via Dodelson-Widrow production.

III. DISCUSSION

One of the simplest extensions of the standard model is to include a massive neutral fermion that mixes with the SM neutrino. It is natural - perhaps expected - that it should come with its own interaction, as well. In the presence of such an interaction, we find that even for very small couplings and mixings, a new eV— MeV mass fermion is equilibrated with the neutrino bath at a temperature within a few orders of magnitude of its mass, and often much less. Consequently, it typically equilibrates after BBN, leaving no imprint on light element abundances. Its implications for the CMB and LSS, however, can be significant. Once the dark fermion equilibrates at T_{equil}, a whole series of additional particles can come into equilbrium as well, including dark matter, which can have mass above T_{equil}, including above an MeV.

Although the equilibration of the dark sector does not immediately increase the energy density in radiation, it can transform some or all of the radiation into an interacting fluid. The associated mass threshold can change the relative amount of relativistic radiation, turn on or off interactions in a dark sector, and provide a basis for equilibrating a broader dark sector which may contain part or all of the dark matter.

- At high values of $100\,\mathrm{eV} \lesssim m_{\nu d} \lesssim$ MeV, the dark sector equilibrates with neutrinos and then goes through the mass threshold of the dark fermion before the CMB is directly sensitive to the transition. One consequence is the increase in N_{eff} by $\Delta N_{eff} = ((g_*^{UV}/g_*^{IR})^{1/3} - 1)N_{eq}$, where N_{eq} is the number of neutrinos that come into equilibrium with the dark sector, and $g_*^{UV}(g_*^{IR})$ is the total number of effective degrees of freedom above (below) the mass threshold, including the thermalizing neutrinos. The relativistic energy below this threshold could be interacting, non-interacting or a combination.

- At intermediate values of $O(1)\,\mathrm{eV} \lesssim m_{\nu d} \lesssim 100\,\mathrm{eV}$, equilibration typically happens before $100\,\mathrm{eV}$, but the mass threshold occurs in a period which is directly probed by the CMB and LSS. This can have important implications for many observables, including H_0 [12, 13] and S_8 [14].

- At very low values of $m_{\nu d}$, the equilibration can happen below $100\,\mathrm{eV}$, and the signal could appear as a transition of the relativistic energy from free-streaming to strongly interacting. This transition would occur sequentially for the three SM neutrino mass eigenstates and would lead to observable signals in the CMB if it occurred at times near recombination. These implications for the CMB are beyond our scope and warrant their own study.

It is interesting to consider what might be a minimal setup, where a single dark Majorana fermion comes into equilibrium with all three SM neutrinos after BBN, but then annihilates away into a real scalar ϕ before the CMB or LSS are directly sensitive. The late universe would have $N_{\rm eff} \simeq 3.30$ with $(1-f)N_{\rm eff} = 2.78$ free streaming neutrinos and $fN_{\rm eff} = 0.53$ interacting "neutrinos" (arising from ϕ). Even in this minimal model, the resulting radiation ($\triangle N_{\rm eff} \simeq 0.26$) is well above the sensitivity of Simons Observatory [19] and CMB-S4 [20]; and the fraction $f = 1/(1+3\cdot 7/4)$ of the "neutrinos" that is interacting can be measured from phase shifts of the CMB peaks [21–28].

If additional particles couple to ν_d or ϕ, they, too, will equilibrate at or after T_{equil} and the thermal history can be yet richer. If additional light particles are present, then the requirement that $m_\phi \ll m_{\nu d}$ is no longer necessary for a viable cosmology. Instead only $m_\phi \ll T_{equil}$ is needed for our calculations to hold, and in this case the neutrinos would become free-streaming again at m_ϕ rather than $m_{\nu d}$. With additional stable particles, dark matter could be produced through thermal processes. For freeze-out, in particular, the dark matter can have masses which are above T_{equil}, and dark matter would have naturally strong couplings to a radiation bath, at least for some period. In all of these cases, ΔN_{eff} can be found simply by an appropriate counting of degrees of freedom in the UV and IR (and intermediate steps, if needed).

In summary, we have considered the thermal history of dark fermions which mix with the SM neutrinos and have self-interactions through a light ($m_\phi \ll m_{\nu d}$) mediator. We find that such particles equilibrate at temperatures near their mass, and thus typically at late times. This implies that later universe observables, such as LSS and the CMB are independent probes when compared to BBN for such models. This can have important implications for models attempting to address cosmological tensions. As we look forward to upcoming results from CMB telescopes such as SPT, ACT, Simons Observatory, CMB-S4 as well as LSS studies well as from LSS measurements KiDS, DES, HSC and future galaxy surveys with Rubin, Roman and UNIONS, such models provide an example of natural late-universe phenomena which may have significant impact. Should such particles populate the ΛCDM desert, these upcoming studies may show striking deviations from ΛCDM expectations.

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