Decaying dark matter, the H_0 tension, and the lithium problem

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It has long been known that the sharpened tension between the observed and inferred values of the Hubble constant H_0 can be alleviated if a fraction of dark-matter particles of type χ were produced nonthermally in association with photons γ through the decays of a heavy and relatively long-lived state, viz., $X \to \chi\gamma$. It was recently proposed that this model can also resolve the long-standing lithium (also known as ^{7}Li) problem if $M = 4$ MeV and $m = 0.04$ keV, where M and m are, respectively, the masses of X and χ. We confront this proposal with experiment and demonstrate that cold dark matter decaying before recombination cannot resolve the H_0 problem. Moreover, we show that the best-case scenario for alleviating the H_0 tension within the context of cold dark matter decaying before recombination arises when the particles decay exclusively into dark radiation, while leaving completely unmodified the production of light elements. To this end, we calculate the general functional form describing the number of equivalent light neutrino species $\Delta N_{\rm eff}$ carried by χ. We show that to resolve the H_0 tension at the 1σ level, a 55% correction in m is needed and that the required $\Delta N_{\rm eff}$ is excluded at 95% C.L. by Planck data. We argue in favor of a more complex model of dynamical dark matter to relax the H_0 tension.

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I. INTRODUCTION

Over the past decade, cosmological parameters have been measured to unprecedented precision. The most reliable measurement of the Hubble constant $H_0 = 74.03 \pm 1.42$ km/s/Mpc comes from Hubble Space Telescope (HST) observations of Cepheid variables in the host of recent, nearby type Ia supernova to build a three-rung distance ladder [1]. In addition, a prediction of H_0 can be obtained from the sound horizon observed from the cosmic microwave background (CMB). A fit to the data from the Planck mission, under the assumption of a flat Λ cold dark-matter (ΛCDM) cosmological model leads to $H_0 = 67.27 \pm 0.60$ km/s/Mpc [2]. These two H_0 values are discrepant by about 4.4σ, which gives rise to the so-called H_0 tension [3,4].

Big bang nucleosynthesis (BBN) has played a central role in the development of precision cosmology. Accurate knowledge of the baryon-to-photon ratio from observations of CMB anisotropies has made BBN a parameter-free theory. BBN predictions match the observed primordial deuterium (D) and helium (^{3}He, ^{4}He) abundances. However, the abundance of primordial lithium (^{7}Li) predicted by BBN is about a factor of 3 larger than the amount of cosmic ^{7}Li measured in the atmospheres of metal-poor stars in the halo of our Galaxy. On the basis of the rather strong assumptions that ^{7}Li has not been depleted at the surface of these stars and that ^{7}Li is independent of Galactic nucleosynthesis and is primordial, one would infer that there is a mismatch with the BBN prediction. This $4-5\sigma$ mismatch constitutes the so-called "cosmic ^{7}Li problem" [5]

Both short-lived ($\tau \ll t_{\rm LS}$) and long-lived ($\tau \gg t_{\rm LS}$) dark-matter particles decaying into dark radiation provide promising scenarios to tackle the tension on the expansion rate (where τ is the particle's lifetime and $t_{\rm LS}$ denotes the time of last scattering) [6–10]. To understand why this is so, we begin by noting that the CMB anisotropy power spectrum tightly constrains the angular size of the sound horizon at recombination θ_*, which in a flat universe is given by the ratio of the comoving sound horizon to the comoving angular diameter distance to last-scattering surface: $\theta_* = r_s(z_{\rm LS})/D_M(z_{\rm LS})$. The comoving linear size of the sound horizon and the comoving angular diameter distance are linked to the expansion history of the Universe via $r_s(z) = \int_z^\infty c_s(z')dz'/H(z')$ and $D_M(z) = \int_0^z dz'/H(z')$, respectively, with c_s the speed of sound and $H(z)$ the Hubble parameter at redshift z [2].

For $\tau \ll t_{\rm LS}$, matter is depleted into radiation at redshifts $z > z_{\rm LS}$. The main effect of adding radiation density to the early Universe is to increase the expansion rate $H(z)$, which

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in turn reduces the sound horizon during the era leading up to recombination. Since the location of the acoustic peak is accurately measured, θ_* must be kept fixed. This can be accomplished, e.g., by simultaneously increasing H_0. However, adding radiation density to the early Universe also alters the damping scale θ_D of the CMB power spectrum, with $\theta_D/\theta_* \propto \sqrt{H(z_{\rm LS})}$ [11]. Increasing $H(z_{\rm LS})$ at fixed θ_* leads to an increase of θ_D, and so the damping kicks in at larger scales reducing the power in the damping tail. An accurate measurement of the small-scale CMB anisotropies thus constrains the fraction of short-lived CDM decaying into radiation.

For $\tau \gg t_{\rm LS}$, matter is depleted into radiation at redshifts $z < z_{\rm LS}$. Therefore, the sound horizon is regulated by the state of the Universe prior to last scattering, and so the value of $r_s(z_{\rm LS})$ does not differ appreciably from that obtained assuming ΛCDM for the same choice of cosmological parameters. However, as a result of the cosmic expansion, the radiation density decreases more rapidly than the matter density, and therefore, the local expansion rate of the Universe is lower in the late-time decaying dark-matter scenario than it would have been in ΛCDM with the same value of $H(z_{\rm LS})$. Now, a consistently lower value of $H(z)$ at low redshifts leads to a larger value of $D_M(z_{\rm LS})$, which in turn would result in a smaller value of θ_*. Again, the location of the acoustic peak is accurately measured, and so θ_* must be kept fixed. This can be accomplished, e.g., by increasing the dark energy density Ω_Λ. As a consequence of the larger Ω_Λ, the matter–dark-energy equality is shifted to earlier times than it otherwise would be in ΛCDM, yielding a larger value of H_0. Remarkably, such late-time decays can also resolve the growing tension between the cosmological and local determination of $S_8 \equiv \sigma_8\sqrt{\Omega_m(\text{today})/0.3}$, which quantifies the rms density fluctuations when smoothed with a top-hat filter of radius $8h^{-1}/\text{Mpc}\ (\equiv\sigma_8)$ as a function of the present day value of the nonrelativistic matter density parameter $\Omega_m(\text{today})$, where h is the dimensionless Hubble constant [12–14].

It was recently proposed that if a fraction of dark-matter particles χ were produced nonthermally in association with photons through the decays of a heavy and relatively long-lived state, both the H_0 and ^{7}Li problems can be simultaneously resolved [15] (see [16] for a similar proposal). In this paper, we confront this proposal with experiment and demonstrate that short-lived CDM cannot resolve the H_0 problem. Moreover, we show that the best-case scenario for alleviating the H_0 tension within the context of short-lived CDM arises when the particles decay exclusively into dark radiation, while leaving completely unmodified the production of light elements. The paper is structured as follows: In Sec. II, we calculate the general functional form describing the number of equivalent light neutrino species $\Delta N_{\rm eff}$ carried by χ. We show that to resolve the H_0 tension at the 1σ level, a 55% correction in the mass of χ is needed and that the required $\Delta N_{\rm eff}$ is excluded at 95% C.L. by Planck data. In Sec. III, we argue in favor of a more complex hidden sector that could combine "early-time" and "late-time" decaying dark-matter solutions of the H_0 tension. The paper wraps up with some conclusions in Sec. IV.

II. CONSTRAINTS ON SHORT-LIVED CDM

Following [15], we assume that of the total dark-matter (DM) density around today $\Omega_{\rm DM}(\text{today})$, a small fraction $f = \Omega_\chi(\text{today})/\Omega_{\rm DM}(\text{today})$ is of particles of type χ produced via decay of a heavy relic X with mass M and lifetime τ: $X \to \chi\gamma$. At any time after the decay of X, the total DM energy density is found to be

$$\rho_{\rm DM}(t) = \frac{mn_\chi(\text{today})}{a^3(t)}\gamma(t) + (1-f)\rho_c\frac{\Omega_{\rm DM}(\text{today})}{a^3(t)}, \quad (1)$$

where m is the mass of the χ, $\gamma(t)$ its Lorentz boost, $n_\chi(t)$ its number density, ρ_c the critical density, and $a(t)$ is the expansion scale factor normalized by $a(\text{today}) = 1$. In the center-of-mass frame of X (this should also be a good approximation of any frame as we assume that X is nonrelativistic so its mass energy dominates), we have

$$M = E_0 + p_0 = \sqrt{p_0^2 + m^2} + p_0 \Rightarrow p_0 = \frac{M^2 - m^2}{2M}, \quad (2)$$

where $E_0 = \gamma(\tau)m$ is the initial energy of the particle χ at $z_{\rm decay}$ and p_0 its momentum, with

$$\gamma(\tau) = \frac{E_0}{m} = \frac{M}{2m} + \frac{m}{2M}. \quad (3)$$

Because of the cosmic expansion, any particle with momentum p becomes redshifted at a rate $p(t) = p_0 a(\tau)/a(t)$. Using the relation $E^2 - p^2 = m^2$, which holds in the Robertson-Walker metric [17], we find that the particle's energy gets redshifted according to

$$E^2(t) = p_0^2\frac{a(\tau)^2}{a(t)^2} + m^2. \quad (4)$$

The scale factor dependence on the Lorentz boost is found to be

$$\begin{aligned}\gamma(t) &= \sqrt{1 + \left[\frac{a(\tau)}{a(t)}\right]^2\left(\frac{p_0}{m}\right)^2} \\ &= \sqrt{1 + \left[\frac{a(\tau)}{a(t)}\right]^2\left(\frac{E_0^2 - m^2}{m^2}\right)} \\ &= \sqrt{1 + \left[\frac{a(\tau)}{a(t)}\right]^2[\gamma^2(\tau) - 1]}. \end{aligned}\quad (5)$$

Expansion of the square root in (5) leads to

$$\gamma(t) \approx 1 + \frac{1}{2}\left[\frac{a(\tau)}{a(t)}\right]^2 [\gamma^2(\tau) - 1] - \frac{1}{8}\left[\frac{a(\tau)}{a(t)}\right]^4 [\gamma^2(\tau) - 1]^2 + \cdots. \tag{6}$$

Now, $\Omega_\chi(\text{today}) = m\, n_\chi(\text{today})/\rho_c$, because the χ is nonrelativistic. To obtain such a nonrelativistic limit, we demand the magnitude of the second term in the expansion of (6) to be greater than the third term, which results in $[a(\tau)/a(t)]^2[\gamma^2(\tau) - 1] < 4$. Contrariwise, by this criteria the particle X is relativistic if $\gamma(t) > \sqrt{5}$. A point worth noting at this juncture is that the general functional form of the Lorentz boost given in (5) and its expansion given in (6) are substantially different from the approximate expression of $\gamma(t)$ given in Eq. (3) of [6]. However, for $\gamma^2(\tau) \gg 1$ and $\gamma^2(\tau)[a(\tau)/a(t)]^2 \gg 1$, both expressions give similar Lorentz factors.

The total "dark" relativistic energy density (including the three left-handed neutrinos of the Standard Model) is usually characterized by the number of "equivalent" light neutrino species $N_{\text{eff}} \equiv (\rho_R - \rho_\gamma)/\rho_{\nu_L}$ in units of the density of a single Weyl neutrino ρ_{ν_L}, where ρ_γ is the energy density of photons, and ρ_R is the total energy density in relativistic particles [18]. Following [15], we obtain the χ contribution to N_{eff} at the time of matter-radiation equality assuming that the χ decouples from the plasma prior to ν_L decoupling, conserving the temperature ratio $T_\gamma/T_{\nu_L} = 11/4$ from ΛCDM cosmology,

$$\begin{aligned} \Delta N_{\text{eff}} &= \frac{8}{7}\left(\frac{11}{4}\right)^{4/3} \frac{\rho_\chi(t_{\text{EQ}})}{\rho_\gamma(t_{\text{EQ}})} \\ &= \frac{8}{7}\left(\frac{11}{4}\right)^{4/3} \frac{\Omega_{\text{DM}}(\text{today})}{\Omega_\gamma(\text{today})} a(t_{\text{EQ}}) f\, \gamma(t_{\text{EQ}}) \\ &= \frac{8}{7}\left(\frac{11}{4}\right)^{4/3} \frac{\Omega_{\text{DM}}(\text{today})}{\Omega_\gamma(\text{today})} a(t_{\text{EQ}}) f \\ &\quad \times \sqrt{1 + \left[\frac{a(\tau)}{a(t_{\text{EQ}})}\right]^2 [\gamma^2(\tau) - 1]}, \end{aligned} \tag{7}$$

where $\rho_\gamma(t_{\text{EQ}}) = \rho_c \Omega_\gamma(\text{today})/a^4(t_{\text{EQ}})$, and the factor of 8/7 is due to the difference between the Fermi and Bose integrals [19]. Note that the functional form of (7) is different from that in Eq. (5) of [6]; the latter was adopted in the recent study of [15]. For $\gamma^2(\tau) \gg 1$ and $\gamma^2(\tau)[a(\tau)/a(t)]^2 \gg 1$, both (7) and Eq. (5) of [6] give similar contributions to ΔN_{eff}.

Model considerations set bounds on free parameters. On the one hand, consistency with large-scale structure observations implies $f \lesssim 0.01$ [6]. On the other hand, the decay of the X's could significantly alter the light element abundances synthesized during BBN. Of particular interest here, the threshold energy of the photon for the process $^7\text{Be}(\gamma, {}^3\text{He})^4\text{He}$ is $E^{\text{th}}_{^7\text{Be}} \simeq 1.59$ MeV, which is lower than that of the photodissociation of D ($E^{\text{th}}_{\text{D}} \simeq 2.22$ MeV) and ^{4}He ($E^{\text{th}}_{^4\text{He}} \sim 20$ MeV). Hence, if the energy of the injected photons is in the range $E^{\text{th}}_{^7\text{Be}} < E_\gamma < E^{\text{th}}_{\text{D}}$, the photodissociation of ^{7}Be could take place to solve the ^{7}Li problem without significantly affecting the abundances of other light elements [20]. With this in mind, to destroy enough ^{7}Li without affecting the abundance of other elements, we set $f \simeq 0.01$ and must fine-tune simultaneously the X lifetime $\tau \simeq 2 \times 10^{-4}$ s [15] and the initial electromagnetic energy release in each X decay $E^{\text{th}}_{^7\text{Be}} < E_\gamma = (M^2 - m^2)/(2M) < E^{\text{th}}_{\text{D}}$. For $M \gg m$, the latter leads to $M \simeq 4$ MeV.

The correlation between H_0 and N_{eff} has been estimated numerically

$$\Delta H_0 = H_0 - H_0|_{\Lambda\text{CDM}} \approx 6.2 \Delta N_{\text{eff}}, \tag{8}$$

where $H_0|_{\Lambda\text{CDM}}$ is the value of H_0 inferred within ΛCDM [21].[1] The rescaled posterior distributions of H_0 from the parameter fit with different choices of ΔN_{eff} are displayed in Fig. 1. To accommodate the H_0 tension at 1σ level, we require $\Delta H_0 \approx 4.7$, which implies $\Delta N_{\text{eff}} \approx 0.76$. The 95% C.L. bound on the extra equivalent neutrino species derived from a combination of CMB, BAO, and BBN observations is $\Delta N_{\text{eff}} < 0.214$ [2]. This limit combines the helium measurements of [25,26] with the latest deuterium abundance measurements of [27] using the PArthENoPE code [28] considering $D(p,\gamma)^3$He reaction rates from [29]. Note that even considering the most optimistic helium abundance measurement of [30] in place of [25,26], the 95% C.L. bound $\Delta N_{\text{eff}} < 0.544$ [2] still precludes accommodating the HST observations within 1σ.

Following [6], we take a radiationlike scale factor evolution $a(t) \propto t^{1/2}$, $\Omega_{\text{DM}}(\text{today}) \approx 0.227$, $\Omega_\gamma(\text{today}) \approx 0.0000484$, $a(t_{\text{EQ}}) \approx 3 \times 10^{-4}$, and $a(\tau)/a(t_{\text{EQ}}) = 7.8 \times 10^{-4}\sqrt{10^{-6}\tau/\text{s}}$. Substituting these figures into (7) while demanding the constraint $\Delta N_{\text{eff}} \approx 0.76$ via (8) leads to $m \simeq 0.018$ keV.

If $m = 0.07$ keV, the contribution to $\Delta N_{\text{eff}} \simeq 0.2$ is consistent with the upper limit reported by the Planck Collaboration [2], and the $X \to \chi + \gamma$ decays still produce enough electromagnetic energy [6]

$$\varepsilon_\gamma \approx 1.5 \times 10^{-6} f \left(\frac{M}{m} - \frac{m}{M}\right) \text{MeV} \tag{9}$$

to dilute the ^{7}Li [20]. However, a recent study of the photodissociation of light elements in the early Universe

[1] We note in passing that the relation (8) has been derived on the basis of the Planck 2015 TT + lowP + BAO + Pantheon dataset combination [22,23]. Should we instead have used Planck 2018 TTTEEE + lowE + BAO + Pantheon [2], the proportionality constant would turn out to be 5.9 rather 6.2 [24]. This is because small-scale polarization data come into play, making it even harder to accommodate high H_0 values with ΔN_{eff}. Herein, we remain conservative in our calculations and adopt the value given in [21].

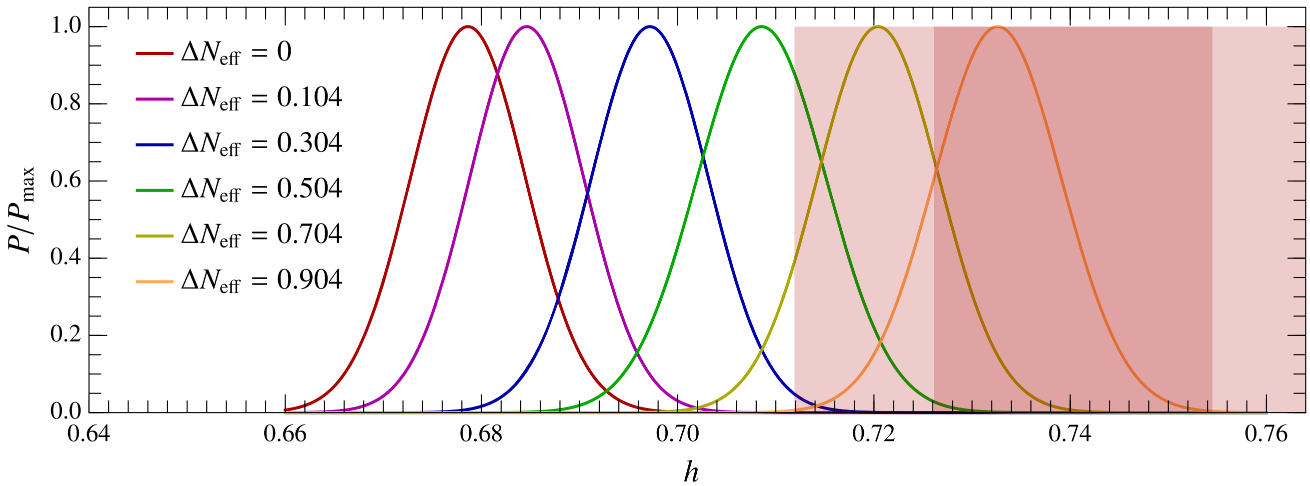

FIG. 1. Rescaled posterior distributions of $H_0 = 100h$ km/s/Mpc (due to marginalization over additional free parameters) with different choices of $\Delta N_{\rm eff}$ from the seven parameter fit of [21]. The shaded areas indicate the 1σ and 2σ regions as determined with HST observations [1].

indicates that the bound on the electromagnetic energy in 2 MeV photons that can be injected during BBN is $\epsilon_\gamma < 10^{-4}$ MeV [31]. This further constrains the contribution to $\Delta N_{\rm eff}$ via $X \to \chi\gamma$. On the other hand, if the X particle decays exclusively into dark radiation, the contribution of this process to $\Delta N_{\rm eff}$ is only limited by $\Delta N_{\rm eff} < 0.214$ [2], allowing for $\Delta H_0 \sim 1.3$.

III. H_0 TENSION AND DYNAMICAL DARK MATTER

Dynamical dark matter (DDM) [32] provides a self-sustaining framework to unify short-lived and long-lived decaying dark-matter models. In the DDM framework, the requirement of dark-matter stability is swapped by a balancing of lifetimes against cosmological abundances across an ensemble of individual dark-matter components with different masses, lifetimes, and abundances. This DDM ensemble collectively plays the role of the dark-matter "candidate" while collectively describing the observed dark-matter abundance.

Following [33], we consider an ensemble comprising a large number N of individual constituent particle species X_n, which decay via $X_n \to \psi\bar{\psi}$, where ψ is a massless dark-sector particle which behaves as dark radiation and the index $n = 1, 2, \ldots, N$ labels the particles in order of increasing mass m_n. The total decay widths Γ_n scale across the ensemble satisfying $\Gamma_1 < \Gamma_2 < \cdots < \Gamma_N$, where Γ_1 is the decay width of the lightest particle in the ensemble. The initial abundances $\Omega_n(z_{\rm prod})$ are regulated by early Universe processes and fixed at $z_{\rm prod} \gg z_{\rm LS}$, with $t_{\rm prod} \ll \tau_N$.

For simplicity, we first assume that all particles in the ensemble are cold, in the sense that their equation-of-state parameter may be taken to be $w_n \approx 0$ for all $t > t_{\rm prod}$. If we further assume that X_{N-1} is almost stable ($\Gamma_{N-1} \ll H_0$), then the DDM framework can accommodate short-lived or long-lived CDM models decaying exclusively into dark radiation.

On the one hand, if $\Gamma_N \gtrsim 10^6$ Gyr^{-1}, the presence of additional energy at around the time of matter-radiation equality contributing to the value of $\Delta N_{\rm eff}$ constrains the abundance of the dark radiation field $\Omega_\psi(z_{\rm LS}) \lesssim 0.1\Omega_\gamma(z_{\rm LS})$, where we have used the 95% C.L. upper limit $\Delta N_{\rm eff} < 0.214$ [2].

On the other hand, if $\tau_N \gtrsim t_{\rm LS}$, the decay of X_N gives late-time solutions to the H_0 problem. These long-lived decaying CDM models can be further subclassified by means of the particle's decay width:

(i) $\Gamma_N \gtrsim H_0 \sim 0.7$ Gyr^{-1}, in which most of the X_N particles have disappeared by $z = 3$.

(ii) $\Gamma_N \lesssim H_0$, in which only a fraction of the X_N particles had time to disappear.

The total initial CDM abundance of the DDM ensemble $\sum_n \Omega_n(z_{\rm prod})$ is essentially fixed by the requirement that at recombination $\sum_n \Omega_n(z_{\rm LS})$ accommodates the dark-matter abundance $\Omega_{\rm DM}(z_{\rm LS})$ derived from Planck data. The initial fraction $F \equiv \Omega_N(z_{\rm prod})/\sum_n \Omega_n(z_{\rm prod})$ is a free parameter of the ensemble.

A full resolution of the Hubble tension within the first subclass requires $F \sim 0.1$ [8], with possible implications for the flux of cosmic neutrinos detected by IceCube [34]. However, a large amount of X_N decay significantly suppresses the amount of DM at low redshifts and thus reduces the power of the CMB lensing effect, which is at odds with Planck observations [35]. Indeed, CMB and BAO data constrain the fraction of decaying DM $F < 0.01$ [36–38], and hence disfavor a full resolution of the H_0 tension if $0.7 \lesssim \Gamma_N/{\rm Gyr}^{-1} \lesssim 10^6$.

An explanation of the Hubble tension within the second subclass allows for smaller values of F but requires a more complex structure of the ensemble with two additional assumptions: (i) allow for intraensemble decays, viz., decays of X_N into final states that include other, lighter X_n; (ii) not all the X_n in the ensemble are cold particle species [9]. If this were the case, then the massive daughter particle could be born relativistic at $z_{\rm decay}$ when the expansion rate is given by $H(z_{\rm decay})$ but behave like CDM as the Universe evolves, yielding a dynamic equation of state $w_n(z)$. Therefore, at any z the collective behavior of any X_n which were born relativistic needs to be averaged over all particles that were born at higher redshifts. This implies that the evolution of the energy density $\Omega_{X_n}(z)$ of these warm particles depends on the sum of all contributions of particles born during the interval $0 \leq z \leq z_{\rm decay}$, some of which were born relativistic and redshifted away by $z = 0$ and some that are born at late times but had no time to be redshifted. A related model requires a DDM ensemble in which the decaying dark matter has an appreciable free-streaming length, and therefore does not cluster on small scales as CDM does, viz., X_N is relativistic at production [10]. However, if we look underneath the seats, we find devils carved: Low multipoles amplitude of the CMB anisotropy power spectrum severely constrains the feasibility of late-time decay models (including those with intraensemble decays) to fully resolve to the H_0 tension [39,40].

The solution of this conundrum lies perhaps in a large ensemble with particle interactions in the hidden sector (see, e.g., [41–44]) and at least $\Gamma_{N-1} \gtrsim H_0$ to combine all of these models. The decays of X_n would also leave observable imprints on the matter power spectrum [45]. This spectrum could play an archaeological role in reconstructing the properties of the underlying dark sector. A comprehensive study of the parameter space of such a complex hidden sector is beyond the scope of this paper and will be presented elsewhere.

IV. CONCLUSIONS

We have studied a set of models endowed with a fraction of CDM decaying in the early Universe to see whether they are able to address the H_0 tension. We have shown that short-lived CDM cannot fully resolve the H_0 tension. The largest contribution to $\Delta N_{\rm eff}$ in agreement with Planck data corresponds to $\Delta H_0 \sim 1.3$ and comes from dark-matter particles which decay exclusively into the hidden sector (e.g., $X \to \psi\bar{\psi}$), while leaving completely unmodified the production of light elements. CDM decaying into photons plus dark radiation is severely constrained by BBN. It would be interesting to study more complex models of dynamical dark matter to see whether they can resolve the H_0 tension by combining short-lived and long-lived CDM.

In closing, we comment on relevant aspects of the ^{7}Li discrepancy. Primordial ^{7}Li abundance is inferred from observations of absorption lines in the photospheres of primitive, low-metallicity stars in the Galactic halo. These objects are warm ($5700 \leq T/{\rm K} \leq 6250$) metal-poor (with small Fe/H abundances relative to the Sun) dwarf stars. ^{7}Li is destroyed in red giants with core temperatures $T \gtrsim 10^6$ K via the reaction $^7{\rm Li}(p,\alpha)^4{\rm He}$, and this is why white dwarfs at moderate temperatures have been used to deduce the ^{7}Li abundance. For each star, the lithium line strength is used to deduce the Li/H abundance. Observations from the early 1980s appear to indicate that in low-metallicity stars the ^{7}Li abundance is roughly constant [46,47]. This constant ^{7}Li abundance, usually referred to as the "Spite plateau," has been interpreted as corresponding to the BBN ^{7}Li yield. This interpretation, of course, assumes that lithium has not been depleted at the surface of these stars so that the presently observed abundance is supposed to be equal to the primordial one. However, more recent observations of low-metallicity stars seem to contradict the conclusions drawn from the Spite plateau. Indeed, the existence of a metallicity trend in the abundance of ^{7}Li in very metal-poor stars, as well as a large dispersion of data have been now observed by many groups (see, e.g., [48,49]) suggesting that the observed values ^{7}Li may not be representative of the cosmological production mechanism [50,51]. All in all, it may well be that the solution of the ^{7}Li discrepancy lies inside the stars and not in cosmology.

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