Toward a realistic model of dark atoms to resolve the Hubble tension

Nikita Blinov[1,*] Gordan Krnjaic,[2,3,4,†] and Shirley Weishi Li[2,‡]
[1]*Department of Physics and Astronomy, University of Victoria, Victoria, British Columbia V8P 5C2, Canada*
[2]*Fermi National Accelerator Laboratory, Batavia, IL 60510, USA*
[3]*University of Chicago, Kavli Institute for Cosmological Physics, Chicago, Illinois 60637, USA*
[4]*University of Chicago, Department of Astronomy and Astrophysics, Chicago, Illinois 60637, USA*

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It has recently been shown that a subdominant hidden sector of atomic dark matter in the early Universe provides a novel avenue toward resolving the Hubble (H_0) tension while maintaining good agreement with cosmic microwave background era observables. However, such a mechanism requires a hidden sector whose energy density ratios are the same as in our sector and whose recombination also takes place at redshift $z \approx 1100$, which presents an apparent fine-tuning. We introduce a realistic model of this scenario that dynamically enforces these coincidences without fine tuning. In our setup, the hidden sector contains an identical copy of Standard Model (SM) fields but has a smaller Higgs vacuum expectation value (VEV) and a lower temperature. The baryon asymmetries and reheating temperatures in both sectors arise from the decays of an Affleck-Dine scalar field, whose branching ratios automatically ensure that the reheating temperature in each sector is proportional to the corresponding Higgs VEV. The same setup also naturally ensures that the hydrogen binding energy in each sector is proportional to the corresponding VEV, so the ratios of binding energy to temperature are approximately equal in the two sectors. Furthermore, our scenario predicts a correlation between the SM/hidden temperature ratio and the atomic dark matter abundance and automatically yields values for these quantities favored by concordant early- and late-Universe measurements of H_0.

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I. INTRODUCTION

The long-standing tension between the early Universe [1] and local [2] extractions of the Hubble constant H_0 may signal the breakdown of the standard Λ cold dark matter (ΛCDM) paradigm (see Ref. [3] for reviews). Although recent local measurements using the tip of the red giant branch method suggest that late-time measurements might be more compatible with early Universe extractions [4], it remains to be seen whether these results will ultimately converge without the need for new physics [5]. While the tension may still be an artifact of systematic error in at least one measurement technique, many models of new physics have been proposed to resolve the discrepancy (see Refs. [6–8] for recent reviews).

It has recently been shown that a subdominant component of atomic dark matter (ADM) [9–16] can viably increase the early Universe value of H_0 while preserving a good fit to other cosmic microwave background (CMB) and baryon acoustic oscillation (BAO) observables [17]. In this scenario, the usual ΛCDM model is supplemented with a hidden sector of dark atoms, photons, and neutrinos; the former accounts for a few percent of the dark matter, and the latter contribute to the radiation density during the CMB era. Unlike other models that propose modifications to the early Universe, this scenario mimics the behavior of visible matter by maintaining the same matter/radiation ratio and undergoing recombination at $z \approx 1100$. The addition of this sector approximately mimics the scaling symmetry

$$\rho_i \to f^2 \rho_i, \qquad \sigma_T n_e \to f \sigma_T n_e, \qquad A_s \to A_s f^{1-n_s}, \qquad (1)$$

where ρ_i is the ith energy density component, n_e is the electron density, $\sigma_T = 8\pi\alpha^2/(3m_e^2)$ is the Thomson cross section, A_s is the amplitude of scalar fluctuations, n_s is the spectral tilt, and f is an arbitrary constant. The transformation in Eq. (1) preserves the form of all cosmological

*nblinov@uvic.ca
†krnjaicg@fnal.gov
‡shirleyl@fnal.gov

perturbation equations in linear theory, thereby retaining the good agreement of CMB/BAO predictions with data [17].

As noted in Ref. [17], this framework presents two main observational challenges: (1) the model favors a small CMB value of the cosmological helium fraction, $Y_p \approx 0.17$, to within a few percent, which is in clear tension with the consensus value from big bang nucleosynthesis (BBN) $Y_p = 0.245 \pm 0.003$ [18], and (2) the best fit hidden/visible temperature ratio satisfies $T'/T \approx 0.7$, corresponding to a large value of $\Delta N_{\rm eff} \approx 1.6$ during BBN, assuming identical Standard Model (SM) field content in the hidden sector. Thus, the ADM mechanism shifts the tension into parameters that affect a different era of cosmological history.

Furthermore, at face value, mimicking this approximate symmetry in Eq. (1) requires an *ad hoc* coincidence to ensure that dark and visible recombination *both* occur at $z \approx 1100$. Since the reheating temperature in each sector is an initial condition and the hydrogen binding energy depends on strong, electromagnetic, and Higgs couplings, such a coincidence across sectors with different masses and thermal histories seems extremely unlikely at first glance; this situation calls for a dynamical explanation.

In this paper, we show how such a coincidence can arise in a realistic model that accounts for the full cosmological history of the atomic hidden sector. We model the hidden sector as an identical copy of the Standard Model with lighter elementary particles and derive the baryon asymmetry and initial temperature in each sector through its coupling to an Affleck-Dine scalar field that dominates the early, postinflationary Universe.

Our model preserves all of the beneficial features identified in Ref. [17] while eliminating fine-tuning needed to time hidden sector recombination. Furthermore, this scenario correlates the interacting DM fraction $f_{\rm adm} \equiv \Omega_{\rm adm}/\Omega_{\rm cdm}$ and the T'/T ratio, thereby removing one free parameter to make our scenario more predictive. However, since we recover the same hidden sector studied in Ref. [17], it inherits the tension in $\Delta N_{\rm eff}$ and Y_p.

II. MODEL OVERVIEW

Inspired by twin Higgs models [19], we postulate a mirror hidden sector which contains an identical copy of all SM fields, coupling constants, and gauge interactions. As in twin Higgs models, the two sectors here have different values for the Higgs vacuum expectation values (VEVs), but in our scenario, we demand that $v'/v < 1$ where $v^{(\prime)}$ is the SM (hidden) Higgs VEV and we use primed symbols to refer to hidden sector quantities throughout this work. Since all other couplings are identical, the QCD confinement scales in both sectors satisfy $\Lambda_{\rm qcd} \approx \Lambda'_{\rm qcd} \approx 200$ MeV [20], which also yields a similar proton mass for both sectors; all other elementary particle masses in the hidden sector are scaled down by an overall factor of v'/v relative to the SM. Unlike in Twin Higgs models, our setup does not invoke any direct couplings between the two sectors and does not address the electroweak hierarchy problem, so there are no collider constraints on the VEV ratio.

To ensure viable bound state formation to explain the Hubble tension in this framework, we must satisfy the following conditions:

(1) There must be a hidden baryon asymmetry to prevent hidden sector particles from completely annihilating into radiation.
(2) Any interaction between sectors must be sufficiently feeble to prevent them from reaching thermal equilibrium.
(3) The reheating temperature in each sector must be directly proportional to the corresponding Higgs VEV. Since we have exhausted all the freedom in choosing field content and coupling constants, the VEV dependence in the hydrogen binding energy $B \propto v$ is compensated by $T \propto v$, so $B/T \approx B'/T'$ and both sectors undergo recombination at the same time.

In what follows, we realize all of these requirements by coupling both sectors to an Affleck-Dine field whose decays simultaneously yield the requisite particle asymmetries and temperature relations to reconcile local and CMB measurements of H_0 with a subdominant atomic dark sector; we assume the remaining CDM in this scenario arises from a different source.

While our model and the phenomenological scenario in Ref. [17] both approximately realize the symmetry in Eq. (1), we add one additional source of scaling violation in the hidden sector Thomson cross section $\sigma'_T/\sigma_T = (v/v')^2 > 1$. The hidden radiation is more tightly coupled to its matter content as a result. Nonetheless, we find that this additional symmetry violating detail does not spoil the good CMB/BAO fit from Ref. [17].

III. COSMOLOGICAL EVOLUTION

We assume the postinflationary early Universe is dominated by a complex scalar field $\phi = re^{i\theta}/\sqrt{2}$ that carries baryon number B_ϕ and realizes Affleck-Dine baryogenesis [21] as depicted schematically in Fig. 1. In polar coordinates, the scalar potential is

$$V(r,\theta) = \frac{m_\phi^2}{2} r^2 + \frac{\lambda}{8} r^4 - \frac{\kappa}{8} r^4 \cos 4\theta, \tag{2}$$

where m is the ϕ mass, $|\kappa| \ll |\lambda|$ are dimensionless couplings, and the explicit θ dependence in the last term provides a source of baryon number violation. After inflation,[1] ϕ starts rolling at $t = t_i$, and the ϕ baryon asymmetry $n_\phi(t) = B_\phi r^2 \dot{\theta}$ evolves according to

[1]In principle, the scalar ϕ could itself be the inflaton field as in Ref. [22], but this is not required for our scenario. Exploring this connection further is beyond the scope of this work.

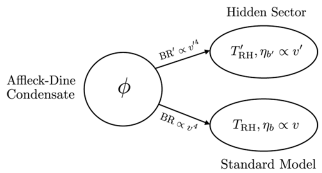

FIG. 1. Schematic illustration of our setup. A nonrelativistic condensate of Affleck-Dine scalar field ϕ dominates the energy density of the postinflationary early Universe and carries net baryon number. Upon decay, ϕ transfers its asymmetry to SM and hidden sector fields with branching ratios proportional to the Higgs VEV in each sector.

$$\frac{1}{a^3}\frac{\partial}{\partial t}(a^3 n_\phi) = -B_\phi \frac{\partial V}{\partial \theta}, \tag{3}$$

where a is the Friedman-Robertson-Walker scale factor and we assume a symmetric initial condition, $n_\phi(t_i) = 0$. Integrating Eq. (3) approximately yields [23]

$$n_\phi(t) \approx -\frac{B_\phi}{\mathcal{H}(t_i)}\frac{a(t_i)^3}{a(t)^3}\frac{\partial V}{\partial \theta}(t_i), \tag{4}$$

where $\mathcal{H} = \dot{a}/a$ is the Hubble rate and this expression gives the baryon number stored in ϕ until it decays through baryon conserving interactions to transfer the asymmetry to the two sectors. Since the baryon density in each sector is set by the corresponding ϕ branching fraction, the baryon-to-entropy ratios $\eta_b^{(\prime)} \equiv n_b^{(\prime)}/s^{(\prime)}$ satisfy

$$\frac{\eta_{b'}}{\eta_b} \approx \frac{\mathrm{BR}_{\phi\to\mathrm{SM}'}}{\mathrm{BR}_{\phi\to\mathrm{SM}}}\frac{g_{\star,s}(T_{\mathrm{RH}})}{g'_{\star,s}(T'_{\mathrm{RH}})}\left(\frac{T_{\mathrm{RH}}}{T'_{\mathrm{RH}}}\right)^3, \tag{5}$$

where $g_{\star,s}^{(\prime)}$ is the number entropic degrees of freedom in each sector and $T_{\mathrm{RH}}^{(\prime)}$ is the visible (hidden) reheating temperature. Thus, the baryon asymmetry transferred to each sector will differ based on model parameters and initial conditions in this framework.

To calculate the branching ratios, we postulate baryon conserving interactions of the form

$$\mathcal{L}_{\mathrm{int}} = \frac{\phi}{\Lambda^n}(H^2\hat{\mathcal{O}} + H'^2\hat{\mathcal{O}}') + \mathrm{H.c.}, \tag{6}$$

where $H^{(\prime)}$ is the visible (hidden) Higgs doublet, $\hat{\mathcal{O}}^{(\prime)}$ is an operator with compensating baryon number $-B_\phi$, Λ is the cutoff scale of the effective interaction, and n is an integer chosen to ensure that the full expression has mass dimension 4. The form of the operator in Eq. (6) is schematic, and a nontrivial contraction of $SU(2)_L$ indices may be required to ensure the leading ϕ coupling is proportional to v^2 or v'^2; for our purposes, any operator will suffice as long as the coefficient preserves this proportionality. The ϕ branching ratios to each sector satisfy

$$\frac{\mathrm{BR}_{\phi\to\mathrm{SM}'}}{\mathrm{BR}_{\phi\to\mathrm{SM}}} \approx \left(\frac{v'}{v}\right)^4. \tag{7}$$

This approximate expression neglects the contributions from loop level decays through virtual Higgs propagators that need not be proportional to the same powers of v; however, such process are suppressed by loop factors of order $(16\pi^2)^{-2} \approx 4\times 10^{-5}$ and can be safely neglected.

Assuming instantaneous reheating through ϕ decays, the energy density of each sector is proportional to the corresponding branching fraction, so the reheating temperatures satisfy

$$\frac{T'_{\mathrm{RH}}}{T_{\mathrm{RH}}} \approx \frac{v'}{v}\left[\frac{g_\star(T_{\mathrm{RH}})}{g'_\star(T'_{\mathrm{RH}})}\right]^{1/4}, \tag{8}$$

where $g_\star^{(\prime)}$ is the number of relativistic degrees of freedom in each sector. Thus, using Eqs. (7) and (8), Eq. (5) becomes

$$\frac{\eta_{b'}}{\eta_b} \approx \frac{v'}{v}\left[\frac{g'_\star(T'_{\mathrm{RH}})}{g_\star(T_{\mathrm{RH}})}\right]^{1/4}, \tag{9}$$

yielding a simple relationship between the asymmetries of our sectors. Note that in the $g_\star = g'_\star$ limit, Eqs. (8) and (9) imply that all energy density ratios are equal in the two sectors (e.g., $\rho_b/\rho_\gamma = \rho_{b'}/\rho_{\gamma'}$), as required to approximate the symmetry in Eq. (1).

Since the hidden sector satisfies $m_{e'} \ll m_{p'}$, the binding energy of hydrogen in both sectors obeys

$$B^{(\prime)} = \frac{\alpha^2}{2}\mu^{(\prime)} \propto v^{(\prime)}, \tag{10}$$

where $\mu^{(\prime)}$ is the electron-proton reduced mass. Therefore, for $g_\star = g'_\star$ in Eq. (8), we predict $B'/T' \approx B/T$, which suffices to trigger hidden recombination around $z \approx 1100$. In order to ensure that $T'_{\mathrm{RH}}/T_{\mathrm{RH}} < 1$, we require $v' < v$, so massive elementary particles are uniformly lighter in the hidden sector. Note that the expression in Eq. (10) is only an approximate equality because the redshift of recombination is logarithmically sensitive to η_b and our scenario predicts $\eta_b \neq \eta_{b'}$ from Eq. (5). However, since the temperature ratio of the two sectors is only of order 1 [17], this mild deviation is negligible for our purposes.

Finally, we note that our model is more predictive than the phenomenological study in Ref. [17] because the ADM fraction of the total dark matter density is

$$f_{\mathrm{adm}} \equiv \frac{\rho_{\mathrm{adm}}}{\rho_{\mathrm{cdm}}} = \frac{\rho_b}{\rho_{\mathrm{cdm}}}\left(\frac{v'}{v}\right)^4 \approx 0.05\left(\frac{T'/T}{0.7}\right)^4, \tag{11}$$

where we have assumed $\rho_b/\rho_{\rm cdm} \approx 1/5$ and used Eq. (8) with $g_\star = g'_\star$.[2] Thus, $f_{\rm adm}$ is correlated with the temperature ratio and lies naturally in the range favored to reconcile early- and late-time measurements of H_0 for the best fit value $T'/T \approx 0.7$ in Ref. [17].

IV. CAVEATS AND COMMENTS

A. Affleck-Dine mass scale

Assuming ϕ decays take place during a cold, matter-dominated phase, both sectors are in the broken electroweak phase throughout reheating; if this were not the case, the branching fractions would not necessarily scale according to the relation in Eq. (7). Furthermore, in the broken vacuum, the H^2 proportionality in Eq. (6) can be expanded using $H = [0, (v+h)/\sqrt{2}]^T$ to generate interactions of the form $vh\phi\hat{\mathcal{O}}/\Lambda^n$, whose branching fraction scales as $\propto v^2$, not v^4, as desired in Eq. (7). Such decays can be kinematically forbidden if $m_\phi < m_{h'}$.

B. Reheat temperature

To ensure that the branching ratios in Eq. (7) are satisfied throughout the early Universe, we demand that both sectors reheat to temperatures below the scale of electroweak symmetry breaking, $T^{(\prime)}_{\rm RH} \lesssim v^{(\prime)}$. Since this requirement necessarily implies that some fields in each sector will not be produced, we must nonetheless ensure that $g_\star(T_{\rm RH}) = g'_\star(T'_{\rm RH})$. However, since the field content in the two sectors is identical, this can be achieved across a wide range of temperature ratios. For example, with $v'/v = 1/2$ and $T_{\rm RH} = 50$ GeV, all leptons, light quarks, and massless gauge bosons are produced in both sectors, but neither thermalizes its $W^\pm, Z^0, h$, or t particles.

Assuming a nonrelativistic ϕ-dominated universe, instantaneous ϕ decays, and rapid equilibration (compared to Hubble expansion) in each sector, the reheating temperature of each sector can be approximated as

$$T^{(\prime)}_{\rm RH} \approx 1\ {\rm GeV}\left(\frac{{\rm BR}_{\phi\to{\rm SM}^{(\prime)}}}{0.33}\right)^{1/4}\left(\frac{0.5\ \mu{\rm s}}{\tau_\phi}\right)^{1/2}, \quad (12)$$

where τ_ϕ is the scalar lifetime. However, the scaling in Eq. (12) is highly model dependent and can be modified, for example, with additional decay channels for ϕ or by parametric resonance effects [24]. Our scenario is compatible with any of these reheating variations as long as ϕ decays in the broken electroweak vacuum and the relation in Eq. (7) is preserved to good approximation.

[2]Note that the precise value of $f_{\rm adm}$ needs to be determined self-consistently by varying v'/v, ρ_b, and $\rho_{\rm cdm}$ simultaneously while fitting to cosmological data.

C. Choosing the decay operator

To realize the branching ratio relation in Eq. (7), the operators in Eq. (6) must be chosen with care. Since the early Universe is always in the broken electroweak phase, ϕ decays must directly generate a net baryon asymmetry; a purely lepton number asymmetry would not yield a baryon asymmetry here since sphalerons are always out of equilibrium in our scenario.

Furthermore, since SM operators with net baryon number involve many insertions of quark fields [e.g., $\hat{\mathcal{O}} = u^c d^c d^c$ where u^c and d^c are respectively up- and down-type $SU(2)_L$ singlet quarks], the exponent n in Eq. (6) is a large integer. This suppression makes it generically difficult to reheat the Universe above the MeV scale while keeping $m_\phi < m_h$, required to maintain the relation in Eq. (7). However, in the presence of additional baryon-charged fields in each sector, this problem can be avoided.

As a toy example, we can add to each sector a gauge singlet Weyl fermion $\chi^{(\prime)}$ and its Dirac partner $\chi^{c(\prime)}$ with baryon number $\mp B_\phi/2$, respectively. This enables us to posit ϕ decay interactions of the form $\phi|H|^2\chi\chi/\Lambda^2$, which can ensure the relations in Eq. (7) while giving ϕ sufficiently prompt $\phi \to \chi^{(\prime)}\chi^{(\prime)}$ decay channels. Once the Universe is populated with χ particles at $T_{\rm RH} >$ MeV, prompt $\chi \to 3q$ decays can proceed through a $\chi^c u^c d^c d^c/M^2$ operator[3] to transfer the baryon asymmetry to SM particles, where M is the mass of a heavy particle that has been integrated out. While this realization satisfies all of our requirements, the gauge and baryon charge assignments for χ also allow direct $\phi\chi\chi$ couplings, which would induce decays that violate the VEV scaling in Eq. (7). This issue can be avoided if ϕ carries baryon minus lepton number, which allows the operator $\phi(LH)^2\chi\chi/\Lambda^5$ and forbids the renormalizable $\phi\chi\chi$ interaction[4]; as noted below, this operator dimension can still yield a sufficiently high reheating temperature for BBN.

For all candidate decay operators, it is important to forbid mixed terms that involve fields from both sectors. From the example above, a complete ultraviolet model must forbid operator variations of the form $\phi(LH)\chi(L'H')\chi'/\Lambda^5$, which would spoil the required relation in Eq. (7). Such a sequestration can be achieved, for example, in extra-dimensional models in which the two sectors live on different four-dimensional branes, but ϕ lives in a higher-dimensional bulk, so hybrid operators involving both sectors are forbidden by locality (see Refs. [25,26] for examples). The mixed operators can also be suppressed if $\chi^{(\prime)}$ is

[3]It is generically easy to ensure $\tau_\phi \gg \tau_\chi$, where $\tau_\chi \sim M^4/m_\chi^5 \sim 10^{-9}\ {\rm s}\,(M/100\ {\rm TeV})^4(10\ {\rm GeV}/m_\chi)^5$ is the χ lifetime. Thus, M can be sufficiently large to evade empirical bounds.

[4]Forbidding the $\phi\chi\chi$ operator but allowing the $\phi(LH)^2\chi\chi/\Lambda^5$ interaction may require baryon minus lepton number to be gauged and spontaneously broken at low energies.

endowed with a parity symmetry under which $\chi\chi'$ transforms nontrivially; this symmetry is then spontaneously broken to allow for $\chi^{(\prime)}$ decays to baryons.

D. Avoiding thermalization

Since the two sectors in our setup must not thermalize with each other, we conservatively demand that ϕ never thermalize with the SM whose energy density is always greater. For an interaction rate based on Eq. (6) in the broken electroweak phase, $\Gamma_{\phi\text{-SM}} \sim v^4 T_{\rm RH}^{2n-3}/\Lambda^{2n} < \mathcal{H}$ at reheating, so the suppression scale must satisfy

$$\Lambda \gtrsim \left(\frac{m_{\rm Pl} v^4 T_{\rm RH}^{2n-5}}{\sqrt{g_\star(T_{\rm RH})}}\right)^{1/2n} \approx 2\ \text{TeV}\left(\frac{T_{\rm RH}}{10\ \text{GeV}}\right)^{1/2}, \qquad (13)$$

where $m_{\rm Pl} = 1.22\times10^{19}$ GeV is the Planck mass and in the last step we took $n=5$. Thus, one can ensure that $T_{\rm RH} \gtrsim$ MeV as required by the success of standard BBN [27–30] without thermalizing ϕ with either sector. We expect that a careful treatment of the phase space in the above thermalization rate would only change the bound on Λ by $\mathcal{O}(1)$ because of the $2n$ root.

In addition to avoiding thermalizing the two sectors through the ϕ interactions, we must also ensure that mixed operators of the form $\lambda|H|^2|H'|^2$ are either forbidden or sufficiently suppressed such that $hh \leftrightarrow h'h'$ reactions are always slower than Hubble expansion. As noted above, such hybrid interactions can be naturally suppressed by locality in extra dimensional theories that require the two sectors to live on different branes.

V. DISCUSSION

We have introduced a realistic model of atomic dark matter that brings local and CMB measurements of H_0 into concordance following the phenomenological study in Ref. [17]. Our approach is inspired by twin Higgs models in which a hidden sector contains an identical copy of the SM, but with a slightly different Higgs VEV, v', and temperature, T'. The baryon asymmetry and reheating temperature in each sector is set by the VEV-dependent branching ratio of an Affleck-Dine scalar field ϕ, and all differences between sectors are governed by v'/v. Since the ϕ branching ratio to the visible and hidden sectors scales as $\propto v^4$ and v'^4, respectively, the hydrogen binding energy-to-temperature ratio, B/T, is the same in both sectors, and recombination occurs for all atomic species at $z\approx 1100$. As a result, there is no fine-tuning required to ensure this coincidence.

Our model features several notable differences with respect to the phenomenological treatment of Ref. [17]:

(1) Because our hidden sector has a smaller electron mass, the Thomson cross section satisfies $\sigma_T' = (v/v')^2\sigma_T = 2.1\sigma_T$. So, even if recombination still occurs at $z\approx 1100$ on account of $B/T = B'/T'$, this larger σ_T' can change the acoustic oscillations of the hidden sector baryons and thus modify the CMB/BAO observables indirectly. To test the importance of this effect, we implemented a mirror sector model in the Boltzmann code CLASS3.1.1 [31] by modifying the existing interacting DM module [32]; we simply rescaled the DM-dark radiation interaction terms by the free-electron fraction of the SM (this is a reasonable approximation since fast variations in this quantity, i.e., recombination, occur simultaneously in the visible and dark sectors by construction) and scaled the interaction coefficient, A_IDM_DR, to match σ_T. The redshift dependence of this interaction matches that of Thomson scattering for $n=2$ in the notation of Ref. [32].[5] In Fig. 2, we show that a different scattering cross section has very little impact on the quality of the fit to CMB and BAO observables. There, we take cosmological parameters motivated by the detailed Markov Chain Monte Carlo analysis in Ref. [17]: $v'/v = 0.68$, $Y_p = 0.17$, $f_{\rm adm} = 0.027$. Note that other parameters were not explicitly provided in that paper, so we picked values that give a reasonable fit by eye: $\Omega_{\rm cdm}h^2 = 0.142$, $\Omega_b h^2 = 0.0224$, $100\theta_s = 1.0425$, $\ln 10^{10}A_s = 3.03$, $n_s = 0.958$, $\tau_{\rm reio} = 0.0539$, and $N_{\rm fs} = 3.30$, where $N_{\rm fs}$ is the number of free-streaming degrees of freedom (i.e., SM and mirror neutrinos). The total number of relativistic degrees of freedom, $N_{\rm eff}\approx 4.24$, includes the contribution of the non-free-streaming dark photons. A similar Boltzmann code implementation of a twin mirror sector was recently studied in Ref. [34], but the authors' focus was on models with $v'/v > 1$. It would be interesting to perform a detailed cosmological analysis for our specific realization of the mirror sector in a future work—we expect this would give rise to only small shifts in the preferred values of the cosmological parameters.

(2) From Eq. (11), our scenario features a one-to-one correspondence between $f_{\rm adm}$ and the hidden/visible temperature ratio, so our model has fewer free parameters than Ref. [17]. It is also notable that the relation in Eq. (11) is automatically consistent with the approximate best fit values $f_{\rm adm}\approx 0.05$ and $T'/T \approx 0.7$ (see also Footnote 2).

(3) Unlike Ref. [17], which assumed $\eta_{b'} = \eta_b \approx 9\times10^{-11}$ [18], from Eqs. (5) and (9), our framework predicts a different baryon asymmetry in the hidden sector.

[5] The interacting DM module contains additional parameters: α_ℓ encodes the damping of higher multipole moments of the dark photon distribution—we take $\alpha_2 = 9/10$ and $\alpha_{\ell>2} = 1$ by matching the notation of Ref. [32] to ΛCDM photons [33]. STAT_F_IDR is set to 1 for bosonic dark radiation; the interacting DM mass M_IDM is taken to be m_p.

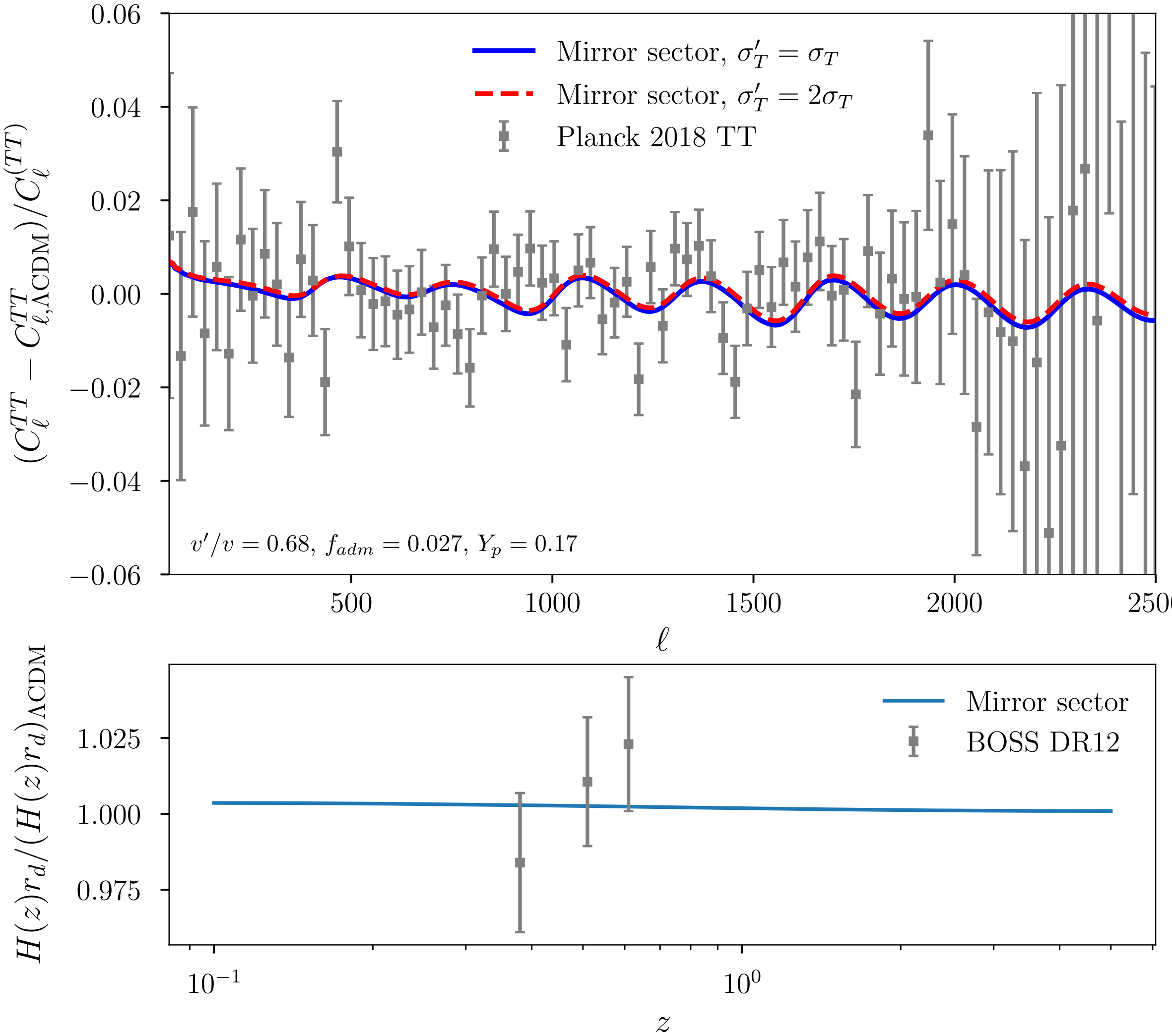

FIG. 2. Impact of a mirror sector on representative CMB (upper panel) and BAO (lower panel) observables. The upper panel shows the relative shift of the TT power spectrum for two mirror sector models: one with the dark Thomson cross section σ'_T equal to the SM one, σ_T (as assumed in Ref. [17]), and one where $\sigma'_T \approx 2\sigma_T$ as predicted in our setup. The differences are negligible, and both models are compatible with the observed spectra, whose binned residuals relative to ΛCDM are shown by the points with error bars [40]. The mirror sector models have $H_0 \approx 73.4$ km/s/Mpc, $f_{\mathrm{adm}} = 0.027$, $T'/T = 0.68$, and $Y_p = 0.17$ following Ref. [17] (the other cosmological parameters are given in the text). In the lower panel, we show $H(z)r_d$ relative to its ΛCDM value, where r_d is the sound horizon at the drag epoch, along with BAO measurements of this quantity from BOSS [41]. Note that ΛCDM and the mirror sector models are nearly indistinguishable in this observable because of the scaling symmetry $H \to fH$, $r \to r/f$ that the mirror sector implements [17].

Since we, nonetheless, recover $\rho_{b'}/\rho_{\gamma'} = \rho_b/\rho_\gamma$, this difference does not affect any CMB observables but can play an important role in hidden sector BBN. While hidden BBN does not affect visible sector CMB observables, it may have interesting observational consequences that warrant further study [35–39].

As noted in Ref. [17], this framework is in generic tension with both the direct measurement of the primordial helium fraction and the theoretical prediction of big bang nucleosynthesis accounting for the large value of $\Delta N_{\mathrm{eff}} \approx 1.6$ (assuming the full SM-like field content in the hidden sector). BBN with such a large enhancement to the expansion rate predicts $Y_p \approx 0.266 \pm 0.0053$, higher than in the standard cosmology with $Y_p \approx 0.247 \pm 0.0046$ (with $n_b/n_\gamma = 6.13 \times 10^{-10}$ [1])[6]; thus, the theoretical prediction of Y_p within this model appears to be even more inconsistent with the CMB inference of the same quantity. Even if the hidden sector is populated after BBN, through, e.g., the decay of a nonrelativistic particle, the CMB-preferred value still disagrees both with the direct late-Universe measurement and the (now) standard Y_p yield from BBN.

[6]These predictions depend on the specific BBN code and reaction network used, but the discrepancy is unaltered. We used AlterBBN2.1 [42] to estimate the theoretical uncertainties using a Monte Carlo procedure.

These points of tension persist in our scenario, but because we approximately realize the scaling relation in Eq. (1), we also inherit a good fit to CMB observables and a larger value of $H_0 \approx 73\ \mathrm{km\,s^{-1}\,Mpc^{-1}}$, provided that Y_p and ΔN_{eff} are allowed to float.

It would be interesting to study whether hidden sector model variations can overcome these observational limitations. For example, it may be possible to realize our scenario with only a subset of SM generations in the hidden sector, in analogy with fraternal twin Higgs models [43], but such studies are beyond the scope of this work.

Finally, we note that our scenario may imply several interesting consequences that are worth exploring in detail. For example, an atomic dark sector with identical field content may yield dark nuclei [35–37], galactic disks [44,45], and stars [46–51], but exploring the observational implications of these structures is beyond the scope of this work. Moreover, UV completions of the various effective operators required to realize our proposal might lead to additional signatures. For example, it may be possible for neutral hidden-visible oscillations between each sector's neutrinos, neutrons, and photons, as long as the mixing interactions that enable these processes do not thermalize the two sectors at early times.

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[1] N. Aghanim *et al.* (Planck Collaboration), Astron. Astrophys. **641**, A6 (2020).
[2] A. G. Riess, S. Casertano, W. Yuan, J. B. Bowers, L. Macri, J. C. Zinn, and D. Scolnic, Astrophys. J. Lett. **908**, L6 (2021).
[3] L. Verde, T. Treu, and A. G. Riess, Nat. Astron. **3**, 891 (2019).
[4] W. L. Freedman, Astrophys. J. **919**, 16 (2021).
[5] G. S. Anand, R. B. Tully, L. Rizzi, A. G. Riess, and W. Yuan, arXiv:2108.00007.
[6] L. Knox and M. Millea, Phys. Rev. D **101**, 043533 (2020).
[7] E. Di Valentino, O. Mena, S. Pan, L. Visinelli, W. Yang, A. Melchiorri, D. F. Mota, A. G. Riess, and J. Silk, Classical Quantum Gravity **38**, 153001 (2021).
[8] N. Schöneberg, G. F. Abellán, A. P. Sánchez, S. J. Witte, c. V. Poulin, and J. Lesgourgues, arXiv:2107.10291.
[9] R. Foot and S. Mitra, Phys. Rev. D **66**, 061301 (2002).
[10] R. Foot and R. R. Volkas, Phys. Rev. D **68**, 021304 (2003).
[11] D. E. Kaplan, G. Z. Krnjaic, K. R. Rehermann, and C. M. Wells, J. Cosmol. Astropart. Phys. 05 (2010) 021.
[12] D. E. Kaplan, G. Z. Krnjaic, K. R. Rehermann, and C. M. Wells, J. Cosmol. Astropart. Phys. 10 (2011) 011.
[13] J. M. Cline, Z. Liu, and W. Xue, Phys. Rev. D **85**, 101302 (2012).
[14] F.-Y. Cyr-Racine and K. Sigurdson, Phys. Rev. D **87**, 103515 (2013).
[15] J. M. Cline, Z. Liu, G. Moore, and W. Xue, Phys. Rev. D **89**, 043514 (2014).
[16] J. M. Cline, arXiv:2108.10314.
[17] F.-Y. Cyr-Racine, F. Ge, and L. Knox, arXiv:2107.13000.
[18] P. A. Zyla *et al.* (Particle Data Group), Prog. Theor. Exp. Phys. **2020**, 083C01 (2020).
[19] Z. Chacko, H.-S. Goh, and R. Harnik, Phys. Rev. Lett. **96**, 231802 (2006).
[20] Z. Chacko, D. Curtin, M. Geller, and Y. Tsai, J. High Energy Phys. 09 (2018) 163.
[21] I. Affleck and M. Dine, Nucl. Phys. **B249**, 361 (1985).
[22] J. M. Cline, M. Puel, and T. Toma, Phys. Rev. D **101**, 043014 (2020).
[23] V. A. Rubakov and D. S. Gorbunov, *Introduction to the Theory of the Early Universe: Hot Big Bang Theory* (World Scientific, Singapore, 2017).
[24] L. Kofman, A. Linde, and A. A. Starobinsky, Phys. Rev. D **56**, 3258 (1997).
[25] C. Csaki, in *Theoretical Advanced Study Institute in Elementary Particle Physics (TASI 2002): Particle Physics and Cosmology: The Quest for Physics Beyond the Standard Model(s)* (World Scientific, 2004), 10.1142/5565.
[26] R. Sundrum, in *Theoretical Advanced Study Institute in Elementary Particle Physics: Physics in $D \geqq 4$* (World Scientific, 2005), 10.1142/6110.
[27] M. Kawasaki, K. Kohri, and N. Sugiyama, Phys. Rev. D **62**, 023506 (2000).
[28] S. Hannestad, Phys. Rev. D **70**, 043506 (2004).
[29] P. de Salas, M. Lattanzi, G. Mangano, G. Miele, S. Pastor, and O. Pisanti, Phys. Rev. D **92**, 123534 (2015).
[30] T. Hasegawa, N. Hiroshima, K. Kohri, R. S. Hansen, T. Tram, and S. Hannestad, J. Cosmol. Astropart. Phys. 12 (2019) 012.
[31] J. Lesgourgues and T. Tram, J. Cosmol. Astropart. Phys. 09 (2011) 032.
[32] M. Archidiacono, D. C. Hooper, R. Murgia, S. Bohr, J. Lesgourgues, and M. Viel, J. Cosmol. Astropart. Phys. 10 (2019) 055.
[33] C.-P. Ma and E. Bertschinger, Astrophys. J. **455**, 7 (1995).

[34] S. Bansal, J. H. Kim, C. Kolda, M. Low, and Y. Tsai, arXiv:2110.04317.
[35] G. Krnjaic and K. Sigurdson, Phys. Lett. B **751**, 464 (2015).
[36] E. Hardy, R. Lasenby, J. March-Russell, and S. M. West, J. High Energy Phys. 06 (2015) 011.
[37] W. Detmold, M. McCullough, and A. Pochinsky, Phys. Rev. D **90**, 115013 (2014).
[38] M. Redi and A. Tesi, J. High Energy Phys. 04 (2019) 108.
[39] A. Mathur, S. Rajendran, and H. Ramani, arXiv:2102.11284.
[40] N. Aghanim *et al.* (Planck Collaboration), Astron. Astrophys. **641**, A5 (2020).
[41] S. Alam *et al.* (BOSS Collaboration), Mon. Not. R. Astron. Soc. **470**, 2617 (2017).
[42] A. Arbey, J. Auffinger, K. P. Hickerson, and E. S. Jenssen, Comput. Phys. Commun. **248**, 106982 (2020).
[43] N. Craig, A. Katz, M. Strassler, and R. Sundrum, J. High Energy Phys. 07 (2015) 105.
[44] J. Fan, A. Katz, L. Randall, and M. Reece, Phys. Dark Universe **2**, 139 (2013).
[45] M. McCullough and L. Randall, J. Cosmol. Astropart. Phys. 10 (2013) 058.
[46] M. Hippert, J. Setford, H. Tan, D. Curtin, J. Noronha-Hostler, and N. Yunes, arXiv:2103.01965.
[47] F. Sandin and P. Ciarcelluti, Astropart. Phys. **32**, 278 (2009).
[48] D. Curtin and J. Setford, Phys. Lett. B **804**, 135391 (2020).
[49] D. Curtin and J. Setford, J. High Energy Phys. 03 (2020) 041.
[50] D. Curtin and J. Setford, J. High Energy Phys. 03 (2021) 166.
[51] H. Winch, J. Setford, J. Bovy, and D. Curtin, arXiv:2012.07136.

Correction: An erroneous modification to the title during the review process has been rectified.