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Cosmological signatures of a mirror twin Higgs

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ABSTRACT: We explore the cosmological signatures associated with the twin baryons, electrons, photons and neutrinos in the Mirror Twin Higgs framework. We consider a scenario in which the twin baryons constitute a subcomponent of dark matter, and the contribution of the twin photon and neutrinos to dark radiation is suppressed due to late asymmetric reheating, but remains large enough to be detected in future cosmic microwave background (CMB) experiments. We show that this framework can lead to distinctive signals in large scale structure and in the cosmic microwave background. Baryon acoustic oscillations in the mirror sector prior to recombination lead to a suppression of structure on large scales, and leave a residual oscillatory pattern in the matter power spectrum. This pattern depends sensitively on the relative abundances and ionization energies of both twin hydrogen and helium, and is therefore characteristic of this class of models. Although both mirror photons and neutrinos constitute dark radiation in the early universe, their effects on the CMB are distinct. This is because prior to recombination the twin neutrinos free stream, while the twin photons are prevented from free streaming by scattering off twin electrons. In the Mirror Twin Higgs framework the relative contributions of these two species to the energy density in dark radiation is predicted, leading to testable effects in the CMB. These highly distinctive cosmological signatures may allow this class of models to be discovered, and distinguished from more general dark sectors.

KEYWORDS: Cosmology of Theories beyond the SM, Beyond Standard Model

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1 Introduction

Mirror Twin Higgs (MTH) models [1–3] provide an interesting and distinctive approach to the little hierarchy problem. In their original incarnation, these theories contain a mirror (“twin”) sector that has exactly the same particle content and gauge interactions as the Standard Model (SM). A discrete Z_2 twin symmetry interchanges the particles and interactions of the two sectors and ensures that the loop corrections to the Higgs mass from the mirror particles cancel the problematic quadratic divergences from the SM gauge and top loops. In contrast to most other symmetry-based solutions of the little hierarchy problem, the twin particles are not charged under the SM gauge groups, and are therefore not subject to the strong constraints from top partner searches at the Large Hadron Collider (LHC) [4–10].

The only coupling between the SM and the mirror sector that is required by the Twin Higgs mechanism is a Higgs portal interaction between the SM Higgs doublet and its counterpart in the twin sector. After electroweak symmetry breaking, the Higgs bosons in the two sectors mix. This mixing leads to a suppression of the couplings of the Higgs particle to SM states. In addition, the Higgs can now decay into invisible twin sector states. Both these effects result in a reduction of the number of Higgs events at the Large Hadron Collider (LHC) as compared to the SM prediction [11]. At present, this is the strongest collider constraint on the MTH model. In order to satisfy this constraint, we require a mild hierarchy between the scale of electroweak symmetry breaking in the twin sector, denoted

by v_B , and the corresponding scale in the SM sector v_A , so that $v_B/v_A \gtrsim 3$ [12]. This hierarchy can be realized by introducing a soft explicit breaking of the twin symmetry, albeit at the expense of mild tuning of order $2v_A^2/(v_A^2 + v_B^2) \simeq 20\%$. Then the elementary fermions and gauge bosons in the twin sector are heavier by a factor of v_B/v_A than their SM counterparts.

MTH models stabilize the hierarchy up to scales of order 5-10 TeV, beyond which an ultraviolet completion is required. Ultraviolet completions have been constructed based on supersymmetry [13–19] (for early work along the same lines, see [20]) and on the composite Higgs framework [21–23] that raise the cutoff to the Planck scale. In supersymmetric UV completions, the breaking of the global symmetry is realized linearly. Then the radial mode in the Higgs potential is present in the spectrum, and can be searched for at colliders [16, 24–26]. In general, composite Twin Higgs models predict new exotic states that carry both SM and mirror gauge charges, which can potentially be discovered at the LHC [27–29]. These theories have been shown to be consistent with precision electroweak constraints [30] and flavor bounds [31].

Cosmology places severe constraints on the MTH framework. The contribution of the light twin neutrinos and twin photons to the energy density of the universe speeds up the Hubble expansion, and this effect is tightly constrained by the existing CMB data. The Higgs portal interaction keeps the SM and the twin sector in thermal equilibrium until temperatures of order a GeV [2]. Below this temperature the twin sector continues to contribute almost half of the total energy density. Even after the other states in the twin sector have decoupled, the twin photon and neutrinos survive as thermal relics, resulting in a large contribution to the energy density in dark radiation during the CMB epoch, $\Delta N_{\text{eff}} = 5.7$ [32, 33]. A correction of this magnitude is ruled out by the current CMB constraints, which require $\Delta N_{\text{eff}} \lesssim 0.45 (2\sigma)$ [34–36].

Several ideas have been put forward to address this problem.¹ One approach is to admit hard breaking of the discrete Z_2 symmetry in the Yukawa couplings in the twin sector. This allows a large reduction in the number of degrees of freedom in the twin sector at the time when the two sectors decouple, leading to a suppression of ΔN_{eff} [38–41]. More radical proposals that produce the same result involve making the mirror sector vector-like [42], or even simply removing from the theory the first two generations of twin fermions, which do not play a role in solving the little hierarchy problem. This latter construction, known as the Fraternal Twin Higgs model [43], leads to distinctive collider signatures involving displaced vertices that can be seen at the LHC [44, 45], and contains several promising dark matter candidates [46–49].

An alternative approach to resolve this problem in mirror models is to incorporate an asymmetric reheating process that preferentially heats up the SM sector rather than the mirror sector [50, 51]. In the case of MTH models this reheating process must occur at late times, after the two sectors have decoupled, but before Big Bang nucleosynthesis (BBN). This has the effect of diluting the fraction of energy density contained in the twin sector,

¹This issue can be avoided if the reheat temperature after inflation lies at or below a GeV, with the inflaton decaying preferentially to the visible sector [37]. However, it is not simple to explain the origin of the baryon asymmetry within such a framework, and we do not consider it further.

allowing the bounds on ΔN_{eff} to be satisfied [32, 33]. In general, late asymmetric reheating can be realized without requiring further breaking of the discrete Z_2 symmetry that relates the two sectors. For example, in the ν MTH model [32], right-handed neutrinos with mass $\mathcal{O}(10)$ GeV decouple from the thermal bath while still relativistic, and come to dominate the energy density of the universe at temperatures of order a GeV. Their decays occur after SM-twin decoupling, with higher branching fractions into the visible sector because the SM W/Z bosons that mediate this process are lighter than their twin counterparts. This process has the effect of making the twin sector much colder than the SM, resulting in the suppression of $\Delta N_{\text{eff}} \sim 7.4(v_A/v_B)^2$. Late asymmetric reheating could also arise from the decay of a scalar field that is related to the spontaneous breaking of the Z_2 symmetry, or from Z_2 breaking in inflationary dynamics [33]. Although these scenarios suppress ΔN_{eff} , it is typically large enough to be seen in future CMB measurements. It is straightforward to accommodate baryogenesis within such a framework [52].

In this paper we explore in detail the cosmological signatures associated with the twin baryons, electrons, photons and neutrinos in the Mirror Twin Higgs scenario. We work in a framework in which the asymmetric twin baryon relic is assumed to constitute only a subcomponent of dark matter, rather than the primary component. In addition, the contribution of the twin photon and neutrinos to dark radiation is assumed to be suppressed due to late asymmetric reheating, but large enough to be detected in future cosmic microwave background (CMB) experiments. We primarily focus on the case in which the discrete Z_2 symmetry is only softly broken, so the mirror particles are heavier than their SM counterparts by a factor of v_B/v_A . Then the relative fractions of hydrogen and helium in the mirror sector can be determined as a function of v_B/v_A , ΔN_{eff} and the baryon asymmetry. We show that this class of theories gives rise to distinctive signals in large scale structure (LSS) and the CMB that can potentially be detected in future experiments.

To understand the origin of these signatures, we consider the thermal history of the MTH model after asymmetric reheating has occurred. Once the temperature of the universe falls below a few MeV, twin Big Bang nucleosynthesis (TBBN) begins in the mirror sector. This determines the relative abundances of twin hydrogen and twin helium at later times. Later, at temperatures below a keV, while density perturbations in cold dark matter (CDM) are growing logarithmically, the mirror baryons are scattering off the mirror electron and mirror photon, leading to twin baryon acoustic oscillations (TBAO). This prevents the mirror particles from contributing to structure growth, resulting in smaller inhomogeneities in the matter distribution (for modes that entered the horizon during the radiation dominated era) than would be expected from Λ CDM. At temperatures of order an eV, as the universe approaches the epoch of matter domination, recombination occurs in the twin sector, first for mirror helium and subsequently for mirror hydrogen. After this time, neutral mirror hydrogen and helium atoms behave as CDM and start to clump, contributing to structure growth. The oscillations in the mirror sector, apart from suppressing structure at short wavelengths, leave a residual oscillatory imprint in the matter power spectrum that can potentially be seen in future LSS measurements. In detail, this imprint depends sensitively on the relative abundances and ionization energies of both twin

hydrogen and helium, and is therefore highly characteristic of this class of models. In the absence of a signal, these measurements will be able to set an upper bound on the twin baryon density, and the results can be translated into an upper bound on the twin baryon asymmetry. As we discuss later, future high precision galaxy redshift surveys and cosmic shear surveys [53] are expected to be sensitive to both the suppression in the matter power spectrum and the oscillatory feature.

The twin neutrinos and twin photons have different effects on the CMB anisotropies. At early times, while the twin neutrinos free stream, the twin photons are prevented from free streaming by scattering off of the twin electrons. Consequently, the density perturbations in these two species evolve differently, with the result that their imprints in the CMB are distinct. In the MTH framework, although ΔN_{eff} itself is a free parameter, the relative energy densities in these two species are predicted. This leads to a testable prediction for the corrections to the heights and locations of the CMB peaks that can potentially be tested in future experiments.

All of these characteristic cosmological signals of the MTH framework, including dark BAO and contributions to ΔN_{eff} , are in fact features of the more general class of models in which the states in a mirror sector constitute some or all of the observed dark matter. Reviews of the status of mirror dark matter, which include many references, may be found in [54–56]. In detail, however, these signals depend sensitively on the masses of the mirror particles and the temperature in that sector. From this perspective, our paper therefore represents an updated, detailed study of the cosmological signals of mirror models, in the region of parameter space motivated by the hierarchy problem.

In this paper our primary focus is on the case in which the discrete Z_2 symmetry is only softly broken, which leads to a prediction for the relative abundances of twin hydrogen and helium as a function of v_B and the temperature of the mirror sector. However, for the purposes of comparison we also study the scenarios in which, for a given electron mass, the nuclei in the mirror sector are composed entirely of hydrogen, or entirely of helium. These studies therefore provide some insight into the cosmology of MTH models in which the Yukawa couplings of the light quarks exhibit hard breaking of the discrete Z_2 symmetry, so that the spectrum of mirror nuclei is composed of only a single species, either hydrogen or helium. In particular, this allows us to capture the cosmological signatures of the interesting scenario in which the mirror neutron is lighter than the mirror proton, and constitutes the primary component of the observed dark matter [2, 41, 57], while mirror helium represents an acoustic subcomponent that gives rise to the signals we discuss. This scenario can also arise in two Higgs doublet extensions of mirror models even in the absence of hard Z_2 breaking if $\tan\beta$, the ratio of the VEVs of the two doublets, is different in the two sectors [58]. As a dark matter candidate, mirror neutrons have the attractive feature that their self-interactions are parametrically of the right size to explain the observed small scale cosmological anomalies [59]. Interestingly, we find that the LSS of the framework in which both hydrogen and helium are present exhibits distinctive features that may allow it to be distinguished from the case of atomic dark matter with just a single type of nucleus.

The twin sector may already be playing a role in resolving some existing cosmological puzzles [60]. For almost two decades, the Λ CDM model has provided an excellent fit to

cosmological data on large scales. However, with the advent of precision measurements, the standard paradigm has come into tension with the data. In particular, there is a $\sim 3\sigma$ discrepancy between the value of the Hubble rate H_0 obtained from a fit to the CMB and baryon acoustic oscillation (BAO) data [34] and the results from local measurements [61]. In addition, the inferred value of the parameter σ_8 , which corresponds to the amplitude of matter density fluctuations at a scale of $8h^{-1}$ Mpc, is in $2\text{--}3\sigma$ tension with the direct measurements obtained from weak lensing surveys [62, 63]. Although the recently published results from the Dark Energy Survey (DES) data [64] exhibit smaller disagreement with the Planck results (less than 2σ), the fact that the low redshift measurements of σ_8 consistently give lower values is intriguing. Resolving these anomalies would require a framework that generically reduces the value of σ_8 as compared to the Λ CDM model, while enhancing H_0 . Several ideas have been proposed that make use of a non-minimal dark sector to address these problems [60, 65–70], and the MTH appears to possess all of the necessary features.

The outline of this paper is as follows. In section 2 we describe the early thermal history of the universe within the MTH framework. A set of model inputs that parameterizes the cosmology of the mirror sector is defined in section 2.1. These inputs also determine the mass splitting between the twin proton and neutron, which is important for TBBN. In section 2.2 we compute the neutron-proton density ratio right before TBBN, which can be translated into the relative fractions of twin hydrogen and twin helium. In section 2.3 we discuss the physics relevant for twin recombination, and derive the ionization fraction of the twin electron as a function of the mirror baryon density. In section 3, we study TBAO and estimate the suppression of the matter power spectrum resulting from oscillations in the mirror sector. We obtain a constraint on the twin baryon density using current results from LSS measurements, and discuss the future observation of oscillation patterns from TBAO. In section 4, we discuss the distinct effects of twin neutrinos and twin photons on the CMB spectrum, and show that this leads to a testable prediction. Our conclusions are in section 5.

2 Thermal history

In this section we describe the early thermal history of the universe within the MTH framework. This will set the stage for the computation of LSS and CMB observables.

2.1 Input parameters

Our focus is on the cosmological signatures of MTH models in which the twin baryons constitute a subcomponent of dark matter. We restrict our analysis to the case when the Yukawa couplings respect the discrete Z_2 symmetry that relates the two sectors, so that the twin fermions are heavier than their visible counterparts by a factor of v_B/v_A . The energy density in twin radiation is assumed to be diluted by late time asymmetric reheating after the two sectors have decoupled, allowing the current CMB and BBN constraints to be satisfied. For simplicity we neglect the masses of both the SM and twin neutrinos. In this framework, the effects on late time cosmology are determined by the following three

parameters,

$$\Delta N_{\text{eff}}, \quad v_B/v_A, \quad r_{\text{all}} = \Omega_{\text{all mirror baryons}}/\Omega_{\text{DM}}. \quad (2.1)$$

Here ΔN_{eff} represents the energy density in twin radiation parametrized in terms of the effective number of neutrinos, while r_{all} denotes the *total* asymmetric mirror baryon density relative to the total dark matter density today. Given these three parameters, we can determine from TBBN the fractional contributions of twin hydrogen and twin helium to the total dark matter energy density, $r_{\hat{\text{H}}}$ and $r_{\hat{\text{He}}}$,

$$r_{\hat{\text{H}}} = \Omega_{\hat{\text{H}}}/\Omega_{\text{DM}}, \quad r_{\hat{\text{He}}} = \Omega_{\hat{\text{He}}}/\Omega_{\text{DM}}. \quad (2.2)$$

The magnitude of ΔN_{eff} depends on the details of the asymmetric reheating process that occurs after the two sectors have decoupled, and so we simply treat it as an input parameter for our study. However, the relative contributions of the twin photons and twin neutrinos to ΔN_{eff} are independent of the nature of the reheating process. During the CMB era, prior to recombination, while the twin neutrinos free stream, the twin photons scatter off the ionized twin electrons. Consequently, as we discuss in section 4, the inhomogeneities associated with these two species do not evolve in the same way, and so their effects on the CMB are different. Then the fact that the relative energy densities in these two species are known leads to a prediction that can potentially be tested in future CMB experiments.

For late time cosmology, the masses of the twin electrons and twin baryons are especially important. These depend on the ratio v_B/v_A . While Higgs coupling measurements at the LHC constrain $v_B/v_A \gtrsim 3$, the requirement that the Higgs mass be only modestly tuned limits $v_B/v_A \lesssim 5$. The mass of the twin electron is simply v_B/v_A times the corresponding value in the SM. To determine the masses of the twin baryons, note that the quark masses are also v_B/v_A times larger than in the SM. This affects the running of the mirror QCD coupling, and leads to a larger confinement scale in the mirror sector than in the SM by 30-50% for $v_B/v_A = 3-5$. The proton and neutron masses, which are almost entirely dictated by Λ_{QCD} , scale the same way in the mirror sector,

$$\frac{m_{\hat{p}}}{m_p} \approx \frac{m_{\hat{n}}}{m_n} \approx \frac{\Lambda_{\text{QCD}_B}}{\Lambda_{\text{QCD}_A}} \approx 0.68 + 0.41 \log(1.32 + v_B/v_A) \quad (2.3)$$

The function of v_B/v_A in eq. (2.3) is a numerical fit valid in the range $v_B/v_A \in (2, 10)$ that gives excellent agreement with the solution from the $\overline{\text{MS}}$ 1-loop RGEs. Given the twin proton mass $m_{\hat{p}}$, we can relate r_{all} to the baryon asymmetry in the twin sector $\eta_{\hat{b}}$,

$$\frac{\eta_{\hat{b}}}{\eta_b} = \frac{r_{\text{all}} \Omega_{\text{DM}} m_p}{\Omega_b m_{\hat{p}}}. \quad (2.4)$$

In this expression, the baryon asymmetries are defined as the ratio of SM or twin baryon number density to the total entropy density of the universe. As in the SM, the contributions of the twin sector to the energy density in dark matter are almost entirely from mirror hydrogen and helium,

$$r_{\text{all}} = r_{\hat{\text{H}}} + r_{\hat{\text{He}}} \quad (2.5)$$

The LSS signals, as well as the twin baryon distribution in galaxies, depend on the relative abundances of mirror hydrogen and helium. Prior to recombination, the mirror ions, electrons and photons undergo dark acoustic oscillations, and do not contribute to the buildup of inhomogeneities. Only after recombination do the neutral mirror atoms contribute to structure growth. Since the ionization energies of mirror hydrogen and helium are different, recombination occurs at different times for these two species. Therefore the matter power spectrum in the MTH framework is very sensitive to the relative abundances of mirror hydrogen and helium.

The relative fractions of mirror hydrogen and helium in the early universe are determined by the dynamics of BBN in the twin sector. As in the SM, the result is very sensitive to the mass splitting between the proton and neutron, ΔM_{np} . This mass difference depends on the quark masses and Λ_{QCD} as,

$$\Delta M_{np} \approx C(m_d - m_u) - D\alpha_{EM}\Lambda_{QCD}. \quad (2.6)$$

We extract the (C, D) coefficients from figure 3 of the lattice QCD study [71], which gives $(C, D) \approx (0.86, 0.40)$. Then, using eq. (2.3) and the fact that the twin quark masses are larger by a factor of v_B/v_A than their SM counterparts, we can relate the neutron-proton mass differences in the two sectors,

$$\frac{\Delta M_{\hat{n}\hat{p}}}{\Delta M_{np}} \approx 1.68v_B/v_A - 0.68, \quad \Delta M_{np} = 1.29 \text{ MeV}. \quad (2.7)$$

For the range we are interested in, $3 \leq v_B/v_A \leq 5$, the separation between twin proton and neutron masses is $\approx 5\text{-}12$ MeV. We are now in a position to determine the relative fractions of mirror hydrogen and mirror helium in the twin sector as a function of ΔN_{eff} , v_B/v_A and r_{all} .

2.2 BBN

As in the SM, the most abundant elements in the mirror sector are expected to be twin hydrogen and twin helium. Their relative abundance is determined by the dynamics of TBBN. Under the assumption that the mirror sector contains a baryon asymmetry leading to stable mirror baryon relics, the mirror helium fraction affects LSS formation and TBAO. A precise calculation of the light element abundances, even in the SM, is extremely involved [72, 73]. Fortunately, the helium fraction can be estimated using the much simpler calculation of neutron-proton freeze-out, without considering the nuclear reactions in detail. This allows us to determine the mirror helium fraction with remarkable precision, despite large uncertainties due to mirror nuclear physics. Earlier studies of BBN in mirror models may be found in [56, 74, 75].

2.2.1 BBN in the SM

We first review a simple estimate of the helium fraction in the SM, following closely the analytical procedure of [76] (see also e.g. [77] and [72, 73]). The first step of the calculation involves computing the neutron-proton ratio “after freeze-out” but before the onset of

nuclear reactions and neutron decay. The $n \leftrightarrow p$ weak conversion rates are approximated by following integrals over thermal distributions,

$$\Gamma_{n\nu_e \rightarrow pe^-} = \frac{1+3g_a^2}{2\pi^3} G_F^2 Q^5 J(1; \infty), \quad \Gamma_{ne^+ \rightarrow p\bar{\nu}_e} = \frac{1+3g_a^2}{2\pi^3} G_F^2 Q^5 J\left(-\infty; -\frac{m_e}{Q}\right), \quad (2.8)$$

where

$$J(a, b) \equiv \int_a^b \sqrt{1 - \frac{(m_e/Q)^2}{q^2}} \frac{q^2(q-1)^2 dq}{\left(1 + e^{\frac{Q}{T\nu}}(q-1)\right) \left(1 + e^{-\frac{Q}{T}q}\right)}. \quad (2.9)$$

The inverse reaction rates are derived from detailed balance,

$$\Gamma_{pe^- \rightarrow n\nu_e} = e^{-Q/T} \Gamma_{n\nu_e \rightarrow pe^-}, \quad \Gamma_{p\bar{\nu}_e \rightarrow ne^+} = e^{-Q/T} \Gamma_{ne^+ \rightarrow p\bar{\nu}_e}. \quad (2.10)$$

Here $g_a \simeq 1.27$ is the standard nucleon axial-vector coupling, $Q = m_n - m_p \simeq 1.293$ MeV, G_F is the Fermi constant, and J is evaluated numerically. Electrons are assumed to annihilate away in a step-function approximation at $T \approx m_e/20$, and the neutrino temperature is $T_\nu = T$ ($T_\nu = (4/11)^{1/3}T$) before (after) electron annihilation. The differential equation for the neutron fraction $X_n \equiv n_n/(n_n + n_p)$ is

$$\frac{dX_n}{dT} = \frac{\Gamma_{n\nu_e \rightarrow pe^-} + \Gamma_{ne^+ \rightarrow p\bar{\nu}_e}}{TH(T)} \left(X_n - (1 - X_n)e^{-Q/T} \right). \quad (2.11)$$

Solving this differential equation numerically, we find that X_n reaches the freeze-out value

$$X_n^{\text{FO}} = 0.15 \quad (2.12)$$

around $T = T_n^{\text{FO}} \approx 0.2$ MeV, in agreement with [72, 73]. That temperature corresponds to $t \sim 20$ s, much less than the neutron lifetime $\tau_n \approx 880$ s, which is why we did not include a neutron decay term in eq. (2.11).

Following neutron-freeze-out one has to consider the onset of nuclear reactions, which eventually give rise to the light elemental abundances, as well as the competing process of neutron decay. The time scale of nucleosynthesis is dominated by the “deuterium bottleneck”, since the formation of helium and other elements proceeds via deuterium. The binding energy of deuterium is very small, $B_D \approx 2.2$ MeV, causing it to break apart at temperatures above ~ 0.1 MeV. Once the temperature has dropped below that threshold and deuterium is stable, the other elements form extremely rapidly at $t = t_{\text{ns}} \approx 180$ s, with almost all of the remaining neutrons being used up to form helium. The final helium fraction can therefore be computed from the remaining neutron fraction,

$$X_n(t_{\text{ns}}) \approx X_n^{\text{FO}} e^{-t_{\text{ns}}/\tau_n} \approx 0.122. \quad (2.13)$$

The primordial helium mass fraction is then

$$Y_p(^4\text{He}) \equiv \frac{\rho_{\text{He}}}{\rho_{\text{He}} + \rho_{\text{H}}} \approx \frac{4n_{\text{He}}}{n_{\text{H}} + 4n_{\text{He}}} = 2X_n(t_{\text{ns}}) \approx 0.245. \quad (2.14)$$

Note that only about 20% of neutrons decay before they are bound up in helium nuclei following freeze-out.

Figure 1. From left to right: the mirror (n/p) ratio, ratio of the mirror ${}^4\text{He}$ and H number density, and ratio of the mirror ${}^4\text{He}$ and H matter density in MTH model. We are interested in the region $3 \leq v_B/v_A \leq 5$. The corresponding value in the visible sector is indicated below each plot.

2.2.2 BBN in the Mirror Sector

The first step of this calculation can be easily repeated for the mirror sector to obtain $X_{\hat{n}}^{\text{FO}}$, by replacing $m_e \rightarrow m_{\hat{e}} = (v_B/v_A)m_e$, $G_F \rightarrow \hat{G}_F = (v_B/v_A)^{-2}G_F$ and $Q \rightarrow \hat{Q} = \Delta M_{\hat{n}\hat{p}} = \frac{\Delta M_{\hat{n}\hat{p}}}{\Delta M_{np}}Q$ using eq. (2.7). Since H is dominated by the visible sector, T still refers to the visible sector temperature, but in the integrated distribution functions J , the mirror sector temperature $\hat{T}$ must be used. This is related to the visible sector temperature by

$$r_T \equiv \frac{\hat{T}}{T} = \left(\frac{g_{\star A}}{g_{\star B}} \right)^{1/3} \left(\frac{\Delta N_{\text{eff}}}{7.4} \right)^{1/4} < 1. \quad (2.15)$$

The mirror sector neutrino temperature is given by the usual $\hat{T}_\nu = \hat{T}$, ($\hat{T}_\nu = (4/11)^{1/3}\hat{T}$) before (after) mirror electron annihilation. The mirror neutron-proton ratio, helium-hydrogen number density ratio, and mirror helium mass fraction derived from the resulting $X_{\hat{n}}^{\text{FO}}$ are shown in figure 1 as a function of v_B/v_A and ΔN_{eff} .

Before discussing the physical ramifications of these results, we have to justify our use of $X_{\hat{n}}^{\text{FO}}$ to directly derive the helium fraction. In obtaining figure 1, we assumed that $X_{\hat{n}} \approx X_{\hat{n}}^{\text{FO}}$ in the mirror version of eq. (2.13). This corresponds to the assumption that mirror neutron decay is much slower than the onset of TBBN and can be neglected. This requires that the mirror deuteron binding energy $B_{\hat{D}}$ is significantly larger than in the SM, even relative to the shorter mirror neutron lifetime.² We now argue that this is indeed the case.

Because of its unnaturally small binding energy, the deuteron is understood to be a fine tuned system [78–80]. While this prevents us from calculating the mirror deuteron binding energy analytically, we can reuse lattice calculations of the SM deuteron binding energy for different pion masses to obtain an estimate, see [81] for a review. The lattice studies find

²We also assume that the mirror baryon density is within a few orders of magnitude of the SM baryon density, so as not to prohibitively suppress the rate of nucleon collisions.

that the binding energy increases with pion mass, but only a handful of such calculations have been performed, and different methods appear to yield somewhat different results. Even so, we can bracket the range of possibilities for the binding energy as a function of pion mass with two linear parameterizations:

$$\begin{aligned} B_D^{\min} &= -(0.66 \text{ MeV}) + 0.021 m_\pi, \\ B_D^{\max} &= -(9.2 \text{ MeV}) + 0.084 m_\pi, \end{aligned} \quad (2.16)$$

both of which reduce to $B_D = 2.2 \text{ MeV}$ when $m_\pi = 135 \text{ MeV}$. To apply this parameterization to the mirror sector, we rescale the dimensionful constant by $\hat{\Lambda}_{\text{QCD}}/\Lambda_{\text{QCD}}$, see eq. (2.3), and replace m_π by $m_{\hat{\pi}}$, which is given by

$$m_{\hat{\pi}} = \sqrt{\frac{\hat{\Lambda}_{\text{QCD}}}{\Lambda_{\text{QCD}}} \frac{v_B}{v_A}} m_\pi \approx \sqrt{[0.68 + 0.41 \log(1.32 + v_B/v_A)] \frac{v_B}{v_A}} m_\pi. \quad (2.17)$$

Now define T_n ($T_{\hat{n}}$) as the *visible* sector temperature when the SM (mirror) neutron decay width equals the Hubble expansion: Γ_n ($\Gamma_{\hat{n}}$) = $H(T)$. This is related by a numerical factor (same in the visible and mirror sectors) to the temperature at which neutrons would typically decay. Note that $\Gamma_n \propto \Delta M_{np}^5/v^4$, $\Gamma_{\hat{n}} = \Gamma_n (\Delta M_{\hat{n}p}/\Delta M_{np})^5 (v_B/v_A)^{-4}$, where the ratio of $n - p$ mass splittings is given in eq. (2.7). We also define T_D ($T_{\hat{D}}$) as the *visible* sector temperature when the SM (mirror) deuterium bottleneck is resolved and SM (mirror) BBN starts. Up to a common prefactor, this is given by

$$T_D = B_D, \quad T_{\hat{D}} = \frac{B_{\hat{D}}}{r_T}.$$

Since $t \sim T^{-2}$,

$$\frac{t_{ns}}{\tau_n} \propto \left(\frac{T_n}{T_D} \right)^2, \quad (2.18)$$

and hence

$$\frac{\hat{t}_{ns}}{\tau_{\hat{n}}} = r_{nD} \frac{t_{ns}}{\tau_n}, \quad r_{nD} \equiv \left(\frac{T_{\hat{n}}/T_{\hat{D}}}{T_n/T_D} \right)^2. \quad (2.19)$$

Recall how in the SM calculation eq. (2.13), X_n^{FO} is scaled down by $e^{-t_{ns}/\tau_n} \approx 0.8$ to give the final neutron and hence helium yield. Therefore, r_{nD} parameterizes how important neutron decay is for TBBN. If $r_{nD} < 1$, neutron decay is *less* important in the mirror sector than in the SM.

We can compute r_{nD} for each of the two parameterizations of $B_{\hat{D}}$ in eq. (2.16). The dependence of r_{nD} on v_B/v_A and ΔN_{eff} is very modest, much weaker than the dependence on the parameterizations of the mirror deuterium binding energy. For $B_D = B_D^{\max(\min)}$, $r_{nD} \approx \frac{1}{16}$ (1). The final neutron yield $X_{\hat{n}}$ must then satisfy

$$0.8 \lesssim \frac{X_{\hat{n}}}{X_{\hat{n}}^{\text{FO}}} \lesssim 0.8^{1/16} \approx 1. \quad (2.20)$$

Unless the mirror deuterium binding energy is very close to our minimum estimate, $X_{\hat{n}}/X_{\hat{n}}^{\text{FO}}$ will be very close to 1. This justifies our use of $X_{\hat{n}}^{\text{FO}}$ to estimate the helium fraction in

figure 1. At worst, $n_{\hat{n}}/n_{\hat{p}}$, $n_{\hat{\text{He}}}/n_{\hat{\text{H}}}$ and $\hat{Y}_p(^4\hat{\text{He}})$ will be lower than the values shown in figure 1 by about 20%, 40% and 10% respectively, which will not significantly change our conclusions.

Figure 1 allows us to make a remarkable prediction for the MTH model. The primordial mirror neutron-to-proton ratio is ~ 0.6 - 0.7 , compared to the SM value of 0.14 . As a result, the mirror helium mass fraction is $\hat{Y}_p(^4\hat{\text{He}}) \approx 75\%$, much higher than in the SM. As we will show below, this has important consequences for LSS formation.

2.3 Recombination

When the temperature in the twin sector becomes much lower than the binding energy of mirror atoms, twin electrons $\hat{e}^-$ start to combine with twin hydrogen and helium nuclei into neutral bound states. This recombination process terminates the acoustic oscillations in the twin sector, and plays an important role in structure formation.

2.3.1 Recombination in the SM

Before considering recombination in the twin sector, it is helpful to first recall how this process occurs in the SM. Hydrogen provides the dominant contribution to the matter density in the SM. Therefore, in our analysis, we neglect the effects of helium, which are subdominant. The primary contribution to the recombination of hydrogen arises from the reaction $e^- + p \rightarrow \text{H}(n \geq 2) + \gamma$, followed by the decay from the excited state down to the $\text{H}(1s)$ state, rather than from direct capture to the ground state [82–84]. This is because the direct capture of an electron into the $\text{H}(1s)$ state results in the emission of a hard photon that quickly ionizes a neighboring atom in the ground state, and therefore gives no net contribution to the recombination process.³ We can simplify the process by considering just three electron states: ionized electrons, electrons in the $n = 2$ state, and electrons in the ground state. Recombination then arises from the capture of an ionized electron into the $n = 2$ state, followed by the de-excitation of the $n = 2$ electron down to $1s$, either through two photon emission, $2s \rightarrow 1s + 2\gamma$, or through Lyman- α decay, $2p \rightarrow 1s + \gamma$. In the case of de-excitation through Lyman- α decay, a net contribution to recombination only arises if the photon loses energy to redshift before colliding with another $1s$ electron.

We denote the ionized fraction of e as $\chi_e \equiv n_e/n_{\text{H,tot}} = n_p/n_{\text{H,tot}}$, where $n_{\text{H,tot}} = n_p + n_{\text{H}(1s)}$ is the sum of both neutral and ionized hydrogen, $n_{\text{H,tot}} = 8.6 \times 10^{-6} \Omega_b h^2 a^{-3} \text{cm}^{-3}$, and we have made use of the fact that helium has already recombined. The Boltzmann equation for χ_e takes the form [83–85],

$$\frac{d\chi_e}{dt} = -\alpha^{(2)} n_{\text{H,tot}} \chi_e^2 + \beta \chi_2, \quad (2.21)$$

$$\chi_i \equiv \frac{n_{\text{H}(n=i)}}{n_{\text{H,tot}}}, \quad \beta \equiv \frac{\alpha^{(2)}}{4} \left(\frac{m_e T}{2\pi} \right)^{3/2} e^{-\epsilon_0/4T}. \quad (2.22)$$

³The ionization cross section near threshold to the $1s$ atom is $\sim 10^7$ times larger than the Thomson cross section, which allows no opportunity to lose the photon energy before ionizing another $1s$ atom. Hubble expansion is also insufficient to redshift the photons arising from this transition to a low enough energy [82]. Since both the ionization cross section and Thomson cross section are proportional to m_e^{-2} , the relative sizes of the two scattering processes remain the same in the MTH case.

Here ϵ_0 denotes the ground state energy of hydrogen, 13.6 eV. In eq. (2.21), the first term on the right corresponds to the capture of ionized electrons into the $n = 2$ state. Since the excited states of hydrogen are in thermal equilibrium, and the energy splitting between the $2s$ and $2p$ states $\sim \alpha_{em}^2 \epsilon_0$ is much lower than the recombination temperature ($\sim \epsilon_0/10$), we do not distinguish between the two $n = 2$ states. Therefore χ_2 includes the contributions from both these states.

The second term on the right in eq. (2.21) corresponds to the ionization of the $n = 2$ state, which releases electrons back into the ionized state. Both terms on the right in eq. (2.21) depend on the recombination cross section to the $n = 2$ state, which can be approximated as [86, 87]

$$\alpha^{(2)} = 0.448 \frac{64\pi}{\sqrt{27}\pi} \frac{\alpha^2}{m_e^2} \left(\frac{\epsilon_0}{T}\right)^{1/2} \ln\left(\frac{\epsilon_0}{T}\right), \quad (2.23)$$

for a general mass and coupling of a hydrogen-like atom.⁴ The net rate of production of $n = 2$ hydrogen atoms is given by the equation

$$\frac{d\chi_2}{dt} = \alpha^{(2)} n_{H,tot} \chi_e^2 - \beta \chi_2 - \left(\Lambda_{2\gamma} + \frac{H \omega_{Ly\alpha}^3}{\pi^2 n_{H,tot} \chi_1} \right) \left(\frac{\chi_2}{4} - \chi_1 e^{-\omega_{Ly\alpha}/T} \right). \quad (2.24)$$

Here $\Lambda_{2\gamma} = 8.227 \text{ sec}^{-1}$ is the two photon decay rate, corresponding to the transition $2s \rightarrow 1s + 2\gamma$, where neither photon has enough energy to excite a ground state hydrogen atom. The de-excitation can also come from the Lyman- α decay $2p \rightarrow 1s + \gamma$, provided the redshift of the Lyman- α photon due to the expansion rate H is faster than the re-absorption from the $n = 1$ state determined by $(n_{H,tot} \chi_1 \omega_{Ly\alpha}^{-3})$, where $\omega_{Ly\alpha} = 3\epsilon_0/4$ is the energy of Lyman- α transition. For both these de-excitation processes, we have included the detailed balance correction corresponding to the reverse processes that arise from thermal excitation by background photons.

When the production and destruction of the $n = 2$ state is in equilibrium, $\frac{d\chi_2}{dt} = 0$, we have

$$\chi_2 = 4 \frac{\alpha^{(2)} n_{H,tot} \chi_e^2 + (\Lambda_{2\gamma} + \Lambda_\alpha) \chi_1 e^{-\omega_{Ly\alpha}/T}}{\Lambda_{2\gamma} + \Lambda_\alpha + 4\beta}, \quad \Lambda_\alpha = \frac{H(3\epsilon_0)^3}{(8\pi)^2 n_{H,tot} \chi_1}. \quad (2.25)$$

$\Lambda_\alpha \simeq 10 \text{ s}^{-1}$ in the SM, which relates to the decay rate of $2p$ state by Lyman- α emission. The net rate of electron ionization in eq. (2.21) is

$$\frac{d\chi_e}{dt} = -\frac{\Lambda_\alpha + \Lambda_{2\gamma}}{\Lambda_\alpha + \Lambda_{2\gamma} + 4\beta} \alpha^{(2)} \left[n_{H,tot} \chi_e^2 - (1 - \chi_e) \left(\frac{m_e T}{2\pi} \right)^{3/2} e^{-\epsilon_0/T} \right]. \quad (2.26)$$

⁴As pointed out in ref. [88], eq. (2.23) may not be a good approximation when the dark radiation temperature is higher or much lower than the binding energy, or when the dark proton is much colder than dark radiation. For the range that we are interested in, $3 \leq (v_B/v_A) \leq 5$, the relative sizes of the twin recombination temperature and twin particle masses is not very different from the SM, and so the twin electron remains in thermal equilibrium with twin photon. We therefore expect eq. (2.23) to provide a good approximation in our case.

Here we have used the fact that $\chi_2 \ll \chi_1$, which follows from detailed balance $\chi_2 \simeq 4\chi_1 \exp(-\omega_{\text{Ly}\alpha}/T)$, to write $\chi_1 \simeq 1 - \chi_e$. The ratio of the Λ terms in front is of eq. (2.26) is the Peebles correction, and the terms inside the square bracket correspond to those in the Saha equation that is derived from the thermal equilibrium of $p^+ + e^- \leftrightarrow \text{H}(1s) + \text{photons}$

$$\frac{n_p n_e}{n_{\text{H}}(1s)} = \frac{n_{\text{H,tot}} \chi_e^2}{1 - \chi_e} = \left(\frac{m_e T}{2\pi} \right)^{3/2} e^{-\epsilon_0/T}. \quad (2.27)$$

Although the Saha equation does not predict the correct relic abundance of electrons, it does approximate the starting point of recombination to a 15% level precision.⁵ We will use eq. (2.26) as the basis of our analysis for both SM and mirror hydrogen.

2.3.2 Recombination in the mirror sector

In order to generalize the analysis above to the case of twin hydrogen, we rescale the binding energy, mass, Lyman- α transition energy and temperature to their values in the mirror sector,

$$\epsilon_0 \rightarrow \hat{\epsilon}_0 = \frac{v_B}{v_A} \epsilon_0, \quad m_e \rightarrow m_{\hat{e}} = \frac{v_B}{v_A} m_e, \quad \omega_{\text{Ly}\alpha} \rightarrow \hat{\omega}_{\text{Ly}\alpha} = \frac{v_B}{v_A} \omega_{\text{Ly}\alpha}, \quad T \rightarrow \hat{T} = r_T T, \quad (2.28)$$

where r_T is defined in eq. (2.15). Since the Hubble expansion is mainly driven by the SM energy density, the expansion rate remains a function of SM temperature $H = H(T)$. Based on the fraction of twin hydrogen density $r_{\hat{\text{H}}}$ we use, the number density $n_{\text{H,tot}}$ is changed to the mirror density

$$n_{\hat{\text{H,tot}}} = \frac{\Omega_{\text{DM}} r_{\hat{\text{H}}} m_p}{(1 - Y_p(^4\text{He})) \Omega_b m_{\hat{p}}} n_{\text{H,tot}}. \quad (2.29)$$

The two photon transition rate of a twin atom is given by [88, 89]

$$\Lambda_{2\hat{\gamma}} = \left(\frac{\hat{\alpha}_{\text{em}}}{\alpha_{\text{em}}} \right)^6 \left(\frac{\hat{\epsilon}_0}{\epsilon_0} \right) \Lambda_{2\gamma} = \left(\frac{v_B}{v_A} \right) \Lambda_{2\gamma}. \quad (2.30)$$

We solve eq. (2.26) to obtain the ionization fraction of mirror hydrogen as a function of redshift for different values of $r_{\hat{\text{H}}}$ and ΔN_{eff} . The results are plotted in figure 2, with the SM ionization fraction provided for comparison. We see from the plot that larger ΔN_{eff} , which corresponds to a higher twin sector temperature, is associated with later recombination. A larger mirror hydrogen density and a lower twin temperature result in a smaller value of χ_e at freeze out.

The ionized states of mirror helium, $\hat{\text{He}}^{2+}$ and $\hat{\text{He}}^+$, have ionization energies $54.4(v_B/v_A)$ and $24.6(v_B/v_A)$ eV respectively. For $v_B/v_A \geq 3$, the universe is deep in the radiation dominated era when the temperature in the twin sector becomes comparable to the binding energy of mirror helium. However, due to its relatively large mass fraction shown in figure 1 (right), mirror helium still plays an important role in the formation of LSS, and cannot be neglected in the TBAO study. Mirror helium stops oscillating with

⁵For example, the Peebles equation predicts that half of the available hydrogen recombines by redshift $z \approx 1200$, while the Saha equation predicts $z \approx 1400$.

Figure 2. The SM (blue) and Twin electron ionization fraction as a function of the scale factor. The MTH results are shown with different values of $r_{\hat{H}} = \Omega_{\hat{H}}/\Omega_{\text{DM}}$.

the twin plasma after $\hat{\text{He}}^+$ recombination. As with helium in the SM, the recombination of $\hat{\text{He}}^+$ proceeds through complicated transitions between a network of excited levels. Instead of studying this process in detail, we estimate the time scale of recombination using the Saha equation. This approximation is justified because, in the SM, the Saha equation is known to reproduce the timescale of He^+ recombination to a precision of order 25%. From the Saha equation we obtain the $\hat{\text{He}}^+$ ionization fraction as

$$\frac{n_{\hat{\text{H}},\text{tot}} \chi_{\hat{e}} \chi_{\hat{\text{He}}^+}}{1 - \chi_{\hat{\text{He}}^+}} = 4 \left(\frac{m_{\hat{e}} r_T T}{2\pi} \right)^{3/2} e^{-24.6 \text{ eV} \left(\frac{v_B}{v_A} \right) / r_T T}. \quad (2.31)$$

For $v_B/v_A = 3$ and ΔN_{eff} ranging from 0.1 to 0.4, the ionization fraction $\chi_{\hat{\text{He}}^+}$ drops to 1% at scale factors ranging from $a = 5 \cdot 10^{-5}$ to $7 \cdot 10^{-5}$, corresponding to conformal times from 15 to 20 $h^{-1} \text{ Mpc}$. This means the $\hat{\text{He}}^+$ scattering in the mirror plasma is expected to modify matter density perturbations starting from wavenumbers $\gtrsim 0.05 h \text{ Mpc}^{-1}$. We will include the effects of twin helium oscillations on the matter power spectrum in the following section.

3 LSS signals

Prior to twin recombination, oscillations in the mirror baryon-photon fluid suppress the growth of structure in the twin sector. In contrast to PAcDM [67, 69, 70] and non-Abelian dark matter [90], which also exhibit dark matter-dark radiation scattering, the oscillations in the MTH stop at a much earlier time because of twin recombination. Consequently, neutral twin atoms still give a sizable contribution to the matter density perturbations, and TBAO leaves an interesting residual oscillation pattern in the matter power spectrum. The overall suppression of structure on scales that enter prior to recombination, and the oscillatory pattern in the matter power spectrum, are characteristic features, not just of the MTH framework, but of the larger class of mirror models. For earlier work on LSS in

the context of mirror models, see [91–95]. Detailed studies of LSS for the general case of atomic dark matter may be found in [88, 96].

To determine the size of the corrections to LSS, we solve a set of linearized Boltzmann equations for density perturbations and calculate the ratio of the matter power spectrum in the MTH *relative to* Λ CDM+DR (which is just Λ CDM with some extra dark radiation included to adjust the value of ΔN_{eff}). In the MTH model, we consider a simplified scenario that contains only CDM χ , ionized twin baryons $\hat{b} = \{\hat{H}, \hat{\text{He}}\}$, massless twin photons $\hat{\gamma}$, and the SM photons γ and protons p . Due to their small energy density, (twin) electrons are neglected, but their effects are implicitly included since they mediate the interactions between (twin) protons and (twin) photons. We work in the conformal Newtonian gauge

$$ds^2 = a^2(\tau) \left[-(1 + 2\psi)d\tau^2 + (1 - 2\phi)\delta_{ij}dx^i dx^j \right], \quad (3.1)$$

where the fields ψ and ϕ describe scalar perturbations on the background metric. They are determined by four scalar quantities associated with the perturbed energy-momentum tensor δT_μ^ν , namely, $\delta \equiv \delta\rho/\bar{\rho} = -\delta T_0^0/\bar{\rho}$, $\delta P = \delta T_i^i/3$, $\theta \equiv -\partial_i \delta T_0^i/(\bar{\rho} + \bar{P})$, and $\sigma \equiv -\hat{\partial}_i \hat{\partial}^j (\delta T_j^i - \delta P \delta_j^i)/(\bar{\rho} + \bar{P})$, where $\bar{\rho}$ and $\bar{P}$ are the unperturbed total energy density and pressure, and $\hat{\partial}^i \equiv \partial_i/\sqrt{\partial_j \partial_j}$. For each particle species s , we define $\delta_s \equiv \delta\rho_s/\bar{\rho}_s$, $\theta_s \equiv -\partial_i \delta T_{s0}^i/(\bar{\rho}_s + \bar{P}_s)$, etc. We can also re-express θ_s as the divergence of comoving 3-velocity, $\theta_s = \partial_i v_s^i$, where $v^i \equiv dx^i/d\tau$. For each s we assume an equation of state of the form $P_s = w_s \rho_s$ with constant w_s , so the pressures and energy densities are not independent quantities. The total δ and θ are given in terms of the individual δ_s and θ_s as $\delta = \sum_s \bar{\rho}_s \delta_s / \bar{\rho}$ and $\theta = \sum_s (\bar{\rho}_s + \bar{P}_s) \theta_s / (\bar{\rho} + \bar{P})$.

To linear order in the perturbations, the evolution of the dominant, collisionless component of dark matter, χ , is described by [87]

$$\dot{\delta}_\chi = -\theta_\chi + 3\dot{\phi}, \quad \dot{\theta}_\chi = -\frac{\dot{a}}{a}\theta_\chi + k^2\psi. \quad (3.2)$$

Here the derivatives are with respect to conformal time $\frac{d}{d\tau}$. For the oscillating component, the equations for the twin baryons are given by

$$\dot{\delta}_{\hat{b}} = -\theta_{\hat{b}} + 3\dot{\phi}, \quad (3.3)$$

$$\dot{\theta}_{\hat{b}} = -\frac{\dot{a}}{a}\theta_{\hat{b}} + \frac{4\rho_\gamma}{3\rho_{\hat{b}}} a n_{e^\pm}(a) \hat{\sigma}_T (\theta_{\hat{\gamma}} - \theta_{\hat{b}}) + k^2\psi. \quad (3.4)$$

The corresponding equations for the SM baryons take a similar form. We solve for the $\hat{\text{He}}^+$ recombination time using eq. (2.31) and approximate the process as a step function in scale factor. As in the SM, neutral twin helium remains tightly coupled to $\hat{H}^+$ even after $\hat{\text{He}}$ recombination. Therefore we continue to have $\hat{b} = \hat{H} + \hat{\text{He}}$ even after $\hat{\text{He}}$ recombination. However, the number density of free mirror electrons is reduced after twin helium recombination. This affects the time at which the twin photons decouple from the twin baryons, resulting in a reduction in the suppression of the matter power spectrum as compared to the case when only hydrogen is present. Therefore the matter power spectrum is sensitive to the relative abundances of $\hat{H}$ and $\hat{\text{He}}$. The term that contains the Thomson cross section $\hat{\sigma}_T = 6.7 \cdot 10^{-25} (v_A/v_B)^2 \text{ cm}^2$ captures the effect of $\hat{b}$ - $\hat{\gamma}$

scattering in the twin sector. The number density of the ionized twin electrons can be approximated as $n_{\hat{e}^\pm}(a) = [\Omega_{\hat{H}}(a) + \frac{1}{4}\Omega_{\hat{H}e}(a)]\rho_c/m_{\hat{H}}$ before twin helium recombination and $n_{\hat{e}^\pm}(a) = \chi_{\hat{e}}(a)\Omega_{\hat{H}}(a)\rho_c/m_{\hat{H}}$ afterwards, where the ionized twin baryon densities are given by $\Omega_{\hat{b}} = r_{\text{all}}\Omega_{\text{DM}}$ and $\Omega_{\hat{H}} = r_{\hat{H}}\Omega_{\text{DM}}$, and the ionization function of the twin electron $\chi_{\hat{e}}(a)$ is calculated numerically using the procedure outlined in section 2.3.

The twin photon perturbations, including higher modes in the Legendre polynomials, evolve as

$$\dot{\delta}_{\hat{\gamma}} = -\frac{4}{3}\theta_{\hat{\gamma}} + 4\dot{\phi}, \quad (3.5)$$

$$\dot{\theta}_{\hat{\gamma}} = k^2 \left(\frac{1}{4}\delta_{\hat{\gamma}} - \frac{1}{2}F_{\hat{\gamma}2} \right) + an_{\hat{e}}\hat{\sigma}_T(\theta_{\hat{b}} - \theta_{\hat{\gamma}}) + k^2\psi, \quad (3.6)$$

$$\dot{F}_{\hat{\gamma}2} = \frac{8}{15}\theta_{\hat{\gamma}} - \frac{3}{5}kF_{\hat{\gamma}3} - \frac{9}{10}an_{\hat{e}}\hat{\sigma}_TF_{\hat{\gamma}2}, \quad (3.7)$$

$$\dot{F}_{\hat{\gamma}l} = \frac{k}{2l+1} [lF_{\hat{\gamma}(l-1)} - (l+1)F_{\hat{\gamma}(l+1)}] - an_{\hat{e}}\hat{\sigma}_TF_{\hat{\gamma}l}, \quad l \geq 3 \quad (3.8)$$

$$\dot{F}_{\hat{\gamma}l_{\text{max}}} = kF_{\hat{\gamma}(l_{\text{max}}-1)} - \frac{l_{\text{max}}+1}{\tau}F_{\hat{\gamma}l_{\text{max}}} - an_{\hat{e}}\hat{\sigma}_TF_{\hat{\gamma}l_{\text{max}}}. \quad (3.9)$$

Here the $F_{\hat{\gamma}l}$ are related to the spatial variations in the density fluctuations in the twin photons, $\delta_{\hat{\gamma}} \equiv F_{\hat{\gamma}0}$, $\theta_{\hat{\gamma}} \equiv \frac{3}{4}kF_{\hat{\gamma}1}/4$, and the shear stress $\sigma \equiv \frac{1}{2}F_{\hat{\gamma}2}$. We truncate the Boltzmann hierarchy at order $l_{\text{max}} = 5$, making use of the approximation outlined in ref. [87]. As shown in the power spectrum ratio plot figure 3 (upper left), the result including only up to $l_{\text{max}} = 4$ (dotted) exhibits only a small deviation from $l_{\text{max}} = 5$ (solid). Since our focus is on describing the clustering of twin baryons in this framework, we do not solve for the twin photon polarization.

We take $\psi = -\phi$ in our analysis, ignoring a small correction arising from the presence of free streaming radiation.⁶ Gravity perturbations are sourced by the density fluctuations as described by the Einstein equation,

$$k^2\psi + 3\frac{\dot{a}}{a} \left(\dot{\psi} + \frac{\dot{a}}{a}\psi \right) = -\frac{a^2}{2M_{\text{pl}}^2} \sum_{i=\chi, \hat{b}, \hat{\gamma}, p, \gamma} \rho_i \delta_i. \quad (3.10)$$

For the initial conditions, the modes that enter before matter-radiation equality satisfy

$$\delta_{\gamma, \hat{\gamma}} = \frac{4}{3}\delta_{\chi, \hat{b}, p} = -2\psi, \quad \theta_{\gamma, \hat{\gamma}, \chi, \hat{b}, p} = \frac{k^2\eta}{2}\psi, \quad (3.11)$$

while for those come in during the era of matter domination,

$$\frac{3}{4}\delta_{\gamma, \hat{\gamma}} = \delta_{\chi, \hat{b}, p} = -2\psi, \quad \theta_{\gamma, \hat{\gamma}, \chi, \hat{b}, p} = \frac{k^2\eta}{3}\psi. \quad (3.12)$$

We set the initial values of the higher modes $F_{\hat{\gamma}\ell \geq 2} = 0$, since these higher angular modes quickly damp away when the Thompson scattering is large. We neglect the tilt in the

⁶In the presence of free streaming radiation, the superhorizon gravity perturbation can be written as $\psi = -(1 + \frac{2}{5}R_{FS})\phi$ and $R_{FS} = \left[1 + \frac{8}{7(N_\nu + \Delta N_{eff, \nu})(T_\nu/T_\gamma)^4}\right]^{-1}$. However, once the mode enters the horizon, the anisotropic stress associated with the free streaming radiation quickly decreases and ψ approaches $-\phi$.

primordial spectrum ($n_s = 1$) and take a k -independent value of $\psi = 10^{-4}$. The final results are independent of the precise value of ψ since we are interested in the ratio of the matter power spectra with and without the twin oscillations. In the numerical study, we choose the values $h = 0.68$, $\Omega_\gamma h^2 = 2.47 \times 10^{-5}$, $\Omega_\Lambda h^2 = 0.69$, $\Omega_b h^2 = 2.2 \times 10^{-2}$ and $\Omega_\nu = 0.69\Omega_\gamma$ [34].

We take as input the parameters of the MTH model (r_{all} , v_B/v_A and ΔN_{eff}). Once these numbers are fixed, the mirror hydrogen and helium density fractions $r_{\hat{\text{H}}}$ and $r_{\hat{\text{He}}}$ in figure 1, the ionization function in figure 2, and the rate of Thompson scattering in the twin sector are all determined. After solving for the density perturbations $\delta_{\chi, \hat{b}, p}$, we calculate the total matter perturbation δ_{tot} and determine the relative suppression of the matter power spectrum with respect to $\Lambda\text{CDM}+\text{DR}$ as,

$$\delta_{\text{tot}}(k) = \sum_{i=\chi, \hat{b}, p} (\Omega_i/\Omega_m) \delta_i(k), \quad \text{P.S. Ratio}(k) \equiv \frac{\delta_{\text{tot}}^2(k) \Big|_{\Lambda\text{CDM}+\text{MTH}}}{\delta_{\text{tot}}^2(k) \Big|_{\Lambda\text{CDM}+\text{DR}}}. \quad (3.13)$$

Before studying the full problem that includes both twin helium and hydrogen, let us first consider the $\hat{H}$ -only scenario to gain some physical intuition for the result. In figure 3 we plot the results for four combinations of (v_B/v_A and ΔN_{eff}), assuming that $\hat{\text{He}}$ is absent or has recombined much earlier so that $r_{\text{all}} = r_{\hat{\text{H}}}$. As compared to the $\text{P}\Lambda\text{CDM}$ model, we see that the MTH exhibits a relatively sudden drop in the ratio of the matter power spectra. This happens at $k \approx 0.04 h \text{ Mpc}^{-1}$, which corresponds to the inverse of the conformal time $\tau_{\text{rec}} \approx 25 h^{-1} \text{ Mpc}$ at twin recombination. For larger ΔN_{eff} , corresponding to a higher twin sector temperature, and smaller v_B/v_A , corresponding to a lower ionization energy, recombination happens later. The power spectrum oscillates around a constant suppression

$$\text{P.S. Ratio}(k \gg \tau_{\text{rec}}^{-1}) \simeq (1 - r_{\hat{\text{H}}})^2. \quad (3.14)$$

This scaling behavior is easy to understand. Twin recombination happens around the time of matter-radiation equality. Prior to this the density perturbations in cold dark matter grow logarithmically, $\delta_\chi(k) \simeq 6 \delta_{\chi, (i)} \ln \frac{k\tau}{\sqrt{3}}$ [97]. However, the twin protons and electrons undergo oscillations with the twin photons, leading to $\delta_{\hat{\text{H}}} \ll \delta_\chi$. It follows that the net matter power spectrum at the end of twin recombination is smaller than the ΛCDM result,

$$\text{P.S. Ratio}(k \gg \tau_{\text{rec}}^{-1}) \simeq \frac{\left[(1 - r_{\hat{\text{H}}}) 6 \delta_{\chi, (i)} \ln \left(\frac{k\tau_{\text{rec}}}{\sqrt{3}} \right) + r_{\hat{\text{H}}} \delta_{\hat{\text{H}}} \right]^2}{\left[6 \delta_{\chi, (i)} \ln \left(\frac{k\tau_{\text{rec}}}{\sqrt{3}} \right) \right]^2} \approx (1 - r_{\hat{\text{H}}})^2. \quad (3.15)$$

After twin recombination, the dark matter density perturbations in both χ and $\hat{\text{H}}$ grow in the same way as for a single species of cold dark matter in ΛCDM , with the result that the ratio above is preserved.

In addition to this overall suppression of the matter power spectrum, we see a residual oscillation pattern in the ratio of power spectra, with a period $\Delta k \approx 0.3 h \text{ Mpc}^{-1}$. There is also a subdominant oscillation with period $\Delta k \approx 0.06 h \text{ Mpc}^{-1}$ arising from interference with the SM BAO. In particular, the total density perturbation δ_{tot} in eq. (3.13) contains

Figure 3. Ratio of matter power spectrum between Λ CDM+MTH and Λ CDM+DR, both with $\Delta N_{\text{eff}} = 0.4$ (left) and 0.1 (right) for $v_B/v_a = 3$ (top) and 5 (bottom). Here we neglect the effect of twin Helium on the oscillation, assuming it is absent or has recombined much earlier. The orange dashed-dotted curve is from the PAcDM model [67], in which a sub-component dark matter never decouples from the dark radiation scattering. For the MTH results, the solid (dotted) curves come from solutions of linearized Boltzmann equations discussed in section 3 with the inclusion of Boltzmann hierarchy modes up to $\ell_{\text{max}} = 5$ (4). The vertical k_{rec} band corresponds to the values of k at the time of twin recombination over the range of r_{H} considered in the plot. The line k_{eq} corresponds to the k value of fluctuation mode that enters at matter-radiation equality.

contributions from the density fluctuations in both the SM and twin proton components. Since the SM protons contribute more to the energy density than twin baryons, it is the SM BAO that generates the dominant oscillation pattern in δ_{tot}^2 . However, when considering the ratio of the two power spectra, this contribution cancels out so that the leading oscillatory effect arises from TBAO. As in SM BAO, the twin sector perturbations carry a $\cos(k\hat{r}_s)$ dependence in δ_{H} , with the sound horizon at recombination defined as

$$\hat{r}_s \equiv \int_0^{\tau_{\text{twin rec}}} d\tau' c_s(\tau'), \quad c_s(\tau) \equiv \left[3 \left(1 + \frac{3 r_{\text{scatt}} \Omega_{\text{DM}}(\tau)}{4 \Omega_{\hat{\gamma}}(\tau)} \right) \right]^{-\frac{1}{2}}. \quad (3.16)$$

Here r_{scatt} represents the dark matter mass fraction of the scattering twin ions. For $v_B/v_A = 3$ and $\Delta N_{\text{eff}} = 0.4$, the sound horizon at time of last scattering in the twin sector can be estimated as $\hat{r}_s \approx 20 h^{-1} \text{Mpc}$. On the other hand, the oscillations in δ_p end at SM recombination, corresponding to a larger sound horizon $r_s \approx 100 h^{-1} \text{Mpc}$. In eq. (3.13), the dominant contribution from TBAO shows up linearly in $\cos(k\hat{r}_s)$, corresponding to oscillations with period $\Delta k = 2\pi/\hat{r}_s \approx 0.3 h \text{Mpc}^{-1}$. Interference between the SM and dark BAO also has an effect, but this is suppressed by an additional $r_{\text{H}} \Omega_b/\Omega_{\text{DM}}$

Figure 4. Ratio of matter power spectrum between Λ CDM+MTH and Λ CDM+DR, including only twin helium oscillations. $k_{\text{rec,He}}$ corresponds to the time of twin helium recombination as obtained from the Saha equation eq. (2.27).

Figure 5. Ratio of matter power spectrum between Λ CDM+MTH and Λ CDM+DR, including both the twin helium and hydrogen oscillations. The twin helium mass fractions are taken from figure 1. $k_{\text{rec,He}}$ corresponds to the time of twin helium recombination as obtained from the Saha equation eq. (2.27).

in the ratio of power spectra as compared to TBAO, and has a period $\approx 1/5$ times shorter than twin oscillations.

We now consider the effects of mirror helium on LSS. In contrast to the SM, TBBN produces a much larger mass density of twin helium than twin hydrogen. Therefore twin helium plays an important role in the TBAO process, and cannot be neglected. From the Saha equation discussed in section 2.3, $\hat{\text{He}}^+$ recombines at conformal time $\approx 20 h^{-1} \text{ Mpc}$, and its scattering in the twin plasma leads to additional suppression in the power spectrum for $k \gtrsim 0.05 h \text{ Mpc}^{-1}$. This means that $\hat{\text{He}}^+$, while extremely important for the nonlinear regime with $k \gtrsim 0.2 h \text{ Mpc}^{-1}$, also has a large impact on σ_8 . To illustrate these effects, in figure 4 we show two representative matter power spectra in which only the effects of mirror helium oscillations are included. The sound horizon at time of helium recombination has $\hat{r}_s \approx 17 h^{-1} \text{ Mpc}$, corresponding to oscillations with period $\Delta k \approx 0.6 h \text{ Mpc}^{-1}$, larger than in the case of hydrogen.

In figure 5 we present two examples of the power spectrum suppression, which take into account oscillations in both mirror hydrogen and helium. We take the twin helium

Figure 6. Ratio of the matter power spectra between Λ CDM+MTH and Λ CDM+DR. Results are shown for two different abundances of mirror H and He, but for the same value of v_B/v_A . The blue curve considers $r_{\text{all}} = 0.1$ with $Y_{\hat{\text{H}}\text{e}} = 0.75$, which is close to the value from the full MTH scenario studied in section 2.2. The red dashed curve corresponds to the hydrogen-only scenario, again with $r_{\text{all}} = 0.1$. The two curves begin to exhibit percent level differences for $k \gtrsim 0.1 h \text{ Mpc}^{-1}$. However, since the blue and red curves correspond to the same r_{all} for the blue and red curves the average suppression for large k -modes is the same, as discussed in eq. (3.15).

density from figure 1 (right) for different values of $(\frac{v_B}{v_A})$ and ΔN_{eff} and terminate the helium oscillations at a recombination time when $\chi_{\hat{\text{H}}\text{e}^+} = 1\%$ in eq. (2.31). At later times twin helium continues to behave as an oscillating component of dark matter, but the number of free electrons is reduced. As compared to the hydrogen-only scenario, the overall suppression in the power spectrum is now dominated by mirror helium. The magnitude of the suppression is still approximately given by eq. (3.15), but with the replacement $r_{\hat{\text{H}}} \rightarrow r_{\hat{\text{H}}} + r_{\hat{\text{H}}\text{e}}$. The sound horizon at the time of $\hat{\text{H}}^+$ recombination is given by $\hat{r}_s \approx 20 h^{-1} \text{ Mpc}$ as shown in figure 5 (left), so the oscillation pattern exhibits a period which is similar to the hydrogen-only case. The small distortions in the curves arise from interference between the mirror and SM baryon oscillations. For a smaller ΔN_{eff} , corresponding to a lower temperature in the twin sector, the earlier freeze out of the TBBN processes results in more mirror helium. Furthermore, the lower temperature makes the mirror helium recombine earlier. For a given r_{all} and v_B/v_A this results in the same overall suppression of the power spectrum deep in the nonlinear regime, but a smaller reduction in σ_8 .

In figure 5, the power spectrum suppression for $k \lesssim 0.1 h \text{ Mpc}^{-1}$ depends on the time scale of $\hat{\text{H}}$ recombination. For a given r_{all} , the relative densities of mirror hydrogen and helium determine the number of ionized electrons that survive after $\hat{\text{H}}\text{e}$ recombination. Consequently the time at which the acoustic oscillations in the mirror sector cease depends on the ratio $r_{\hat{\text{H}}\text{e}}/r_{\text{all}}$. Therefore the matter power spectrum is sensitive to the relative abundances of mirror hydrogen and helium, as shown in figure 6. This opens the door

Figure 7. Ratio of the matter power spectra between Λ CDM+MTH and Λ CDM+DR, for three different abundances of mirror H and He. The blue curve considers $r_{\text{all}} = 0.1$ with $Y_{\hat{\text{H}}\text{e}} = 0.75$, which is close to the value from the full MTH scenario studied in section 2.2. The red dashed curve corresponds to the hydrogen-only scenario, again with $r_{\text{all}} = 0.1$, but with a larger value of v_b/v_A corresponding to a heavier mirror electron. We see that this gives a good fit to the MTH result for $k \lesssim 0.12 h \text{ Mpc}^{-1}$, but the two curves begin to exhibit differences in the σ_8 region. However, since the blue and red curves correspond to the same r_{all} for the blue and red curves the average suppression for large k -modes is the same, as discussed in eq. (3.15). The green dotted curve corresponds to values of r_{all} and v_b/v_A that have been chosen to mimic as closely as possible the blue curve in the linear regime, $k \lesssim 0.2 h \text{ Mpc}^{-1}$. However, since r_{all} differs from 0.1, the average suppression for large k -modes deviates from the blue curve. The percent level differences between the blue and red curves at $k \approx 0.2 h \text{ Mpc}^{-1}$ and between the blue and green curves at $k > 0.2 h \text{ Mpc}^{-1}$ may allow future matter power spectrum measurements to distinguish the MTH from theories with a single species of dark atom.

to the possibility of distinguishing the MTH universe from scenarios with a single species of dark atom. As shown in figure 7, although it is possible to match part of the $\hat{\text{H}} + \hat{\text{He}}$ result (blue curve) using two different twin hydrogen-only (red dashed and green dotted) scenarios, the fit necessarily leaves residual differences with the MTH in either the linear regime, the nonlinear regime, or both. If we hold r_{all} the same as in the MTH case but allow v_B/v_A to float (the red curve), we can match the blue curve at low k , but find that percent level differences remain even in the linear regime near $k = 0.2 h \text{ Mpc}^{-1}$. In the nonlinear regime, the differences between the red and blue curves are much larger, but the average suppression of the power spectrum in the two cases remains the same. If, however, we allow both r_{all} and v_B/v_A to float (the green curve), it is possible to match the blue curve in the entire linear regime. However, there are still large differences between the shapes of the blue and green curves in the nonlinear regime, and even the average suppression of the power spectrum in the two cases is different. Future experiments from weak lensing data are expected to constrain the matter power spectrum for $k \lesssim 0.5 h \text{ Mpc}^{-1}$ to percent level

Figure 8. Estimation of projected constraints on the twin baryon mass fraction r_{all} from future LSS measurements, as a function of v_B/v_A . Next generation lensing measurements are expected to bound the power spectrum suppression at the few percent level [98], so we show the constraints arising from the 90, 95 and 98% lower limits on the ratio of power spectra defined in eq. (3.13) as the orange, green and red curves. The ratio of the densities of twin hydrogen and helium are obtained from figure 1. The region above the solid (dashed) blue line for $\Delta N_{\text{eff}} = 0.1$ (0.4) represents an approximation to the Lyman- α constraint that excludes scenarios with δA , the deviation of the integrated power spectrum defined in eq. (3.17), larger than 38%.

precision. Although the $k \gtrsim 0.5 h \text{ Mpc}^{-1}$ region is plagued by nonlinear effects, even here future measurements based on 21 cm tomography will probe the universe at much higher redshifts, $z \sim 100$, and therefore allow a detailed study of $\mathcal{O}(1) \text{ Mpc}$ k -modes before they enter the nonlinear regime. Therefore, future observations of the matter power spectrum may be able to reveal the existence of more than one species of dark atom, which is a striking prediction of the MTH framework.

If TBAO is to provide an explanation for the current $\approx 10\%$ discrepancy in σ_8 between the low redshift measurements and the Planck fit, the fractional contribution of twin hydrogen to the overall dark matter density is required to be $r_{\text{H}} \approx 1\%$, corresponding to $r_{\text{all}} \sim 3.6\%$, for $v_B/v_A = 3$ and $\Delta N_{\text{eff}} = 0.4$. For these values of r_{all} and v_B/v_A , the superimposed oscillation of the modes with $k > 0.2 h \text{ Mpc}^{-1}$ leads to an extra $\approx 4\%$ variation of the power spectrum as a function of k . This is particularly significant keeping in mind the expected $\mathcal{O}(1)\%$ sensitivity of future LSS observations [98], which will provide an observational channel to probe the MTH model. Nonlinear effects, however, become important for $k \gtrsim 0.2 h \text{ Mpc}^{-1}$, and a detailed study in the higher k -mode region is required to determine the full oscillation pattern.

In the absence of a signal, future measurements of LSS will be able to place stringent limits on the energy density in twin baryons. In figure 8 we show the upper bounds on the

fractional contribution of twin baryons to the dark matter density, r_{all} , corresponding to different current and projected lower bounds on the matter power spectrum suppression, as defined in eq. (3.13). The effects of both twin helium and hydrogen are included. We present the results for $k = 0.2 h \text{ Mpc}^{-1}$, which exhibits a significant suppression of the power spectrum, but for which nonlinear effects are still under theoretical control. Current measurements of the matter power spectrum allow a $\lesssim 10\%$ deviation from the ΛCDM prediction at this scale, which constrains the fractional contribution of twin baryons to the dark matter density to lie in the range $r_{\text{all}} \lesssim 4\text{--}10\%$, depending on the other parameters. Given the expected $\mathcal{O}(1)\%$ sensitivity of future LSS observations, if the observed result is fully consistent with the ΛCDM prediction, we expect to be able to constrain the density fraction to $r_{\text{all}} \lesssim 1\%$. Since cooler dark radiation (corresponding to a smaller ΔN_{eff}) and a heavier twin electron (larger v_B/v_A) lead to earlier recombination, the corrections to the power spectrum are smaller in this case. This explains the weaker bounds on r_{all} for the $\Delta N_{\text{eff}} = 0.1$ case with large v_B/v_A .

The TBAO suppression persists out to $k > 0.5 h \text{ Mpc}^{-1}$. Then the Lyman- α observations, which probe the matter power spectrum on scales $0.5 h^{-1} \text{ Mpc} < \lambda < 100 h^{-1} \text{ Mpc}$ [99] at $z \approx 3\text{--}5$ can also be used to constrain r_{all} . A precise Lyman- α bound requires detailed N-body simulations of the MTH plasma, which are beyond the scope of this work. Here we only give a rough estimate of the current bound following the strategy adopted in ref. [100] for the warm dark matter (WDM) study, which constrains the integrated matter power spectrum over a range of wave numbers $0.5 < k < 20 h \text{ Mpc}^{-1}$.⁷ We define the ratio factor

$$\delta A \equiv \frac{A_{\Lambda\text{CDM}+\text{DR}} - A_{\Lambda\text{CDM}+\text{MTH}}}{A_{\Lambda\text{CDM}+\text{DR}}}, \quad A_{\Lambda\text{CDM}+\text{MTH}} = \int_{k_{\text{min}}}^{k_{\text{max}}} dk \text{ P.S.Ratio}(k) \quad (3.17)$$

and determine the result for redshift $z = 3$. By construction, $A_{\Lambda\text{CDM}+\text{DR}} = k_{\text{max}} - k_{\text{min}}$. We take $\delta A < 0.38$ from ref. [100] as a 2σ bound on the suppression in power from the existing Lyman- α forest data from the MIKE/HIRES [99] and XQ-100 [102] datasets used in ref. [103]. Applied to the MTH, we obtain a bound $r_{\text{all}} \lesssim 16\text{ (12)\%}$ for $\Delta N_{\text{eff}} = 0.1\text{ (0.4)}$ that is quite insensitive to the v_B/v_A ratio. This agrees with the analytical estimate in eq. (3.15) for the average suppression at higher k -modes. Given the uncertainties involved, we should only consider this as a rough guide to the actual constraint.⁸ The result with $r_{\text{all}} \lesssim 16\text{ (12)\%}$ is shown as the region above the solid (dashed) blue line in figure 8.

⁷The study is based on the assumption that the bias function $b(k)$ between the flux power spectrum $P(k)_F$ and the linear power spectrum $P(k)$, written as $P(k)_F = b^2(k) P(k)$, only differs a little between ΛCDM and the new model. The integration over k -modes is justified by the fact that velocities in the Intergalactic Medium tend to redistribute the power spectrum within a range of wave numbers in the probed region [101].

⁸A study of the Lyman- α constraint using the SDSS [104, 105] data has been performed in ref. [106] on scenarios in which all dark matter particles couple to dark radiation. Although the setup is different, a similar study of dark matter scattering off dark radiation may be applied to the MTH model.

4 CMB signals

The MTH framework predicts a new contribution to the energy density of the universe in the form of dark radiation associated with the relativistic degrees of freedom in the mirror sector. In order to satisfy the current CMB bounds on dark radiation, $\Delta N_{\text{eff}} \lesssim 0.45$ [34–36], we assume asymmetric reheating of the SM bath after the two sectors have decoupled. However, while asymmetric reheating increases the energy density in the SM sector, it does not erase any pre-existing energy density in the mirror sector. Therefore, a small nonvanishing ΔN_{eff} is expected to be a general feature of this framework. This dark radiation could potentially be detected at future CMB Stage-IV experiments, which are expected to be sensitive to $\Delta N_{\text{eff}} \gtrsim 0.02$. We treat ΔN_{eff} as a free parameter, since its value depends on the precise details of the asymmetric reheating mechanism.

The dark radiation in the MTH is composed of the mirror photon and mirror neutrinos. Their relative contributions to ΔN_{eff} depend only on the number of degrees of freedom in the twin bath at the time when the mirror neutrinos decouple. For the parameter range of interest, $3 \lesssim v_B/v_A \lesssim 5$, the mirror electron goes out of the twin bath after the mirror neutrinos have decoupled, just as in the visible sector. Then, even though ΔN_{eff} itself is a free parameter, we have a prediction for the relative contributions of twin photons and twin neutrinos to the energy density in dark radiation

$$\frac{\Delta N_{\text{eff}}^\nu}{\Delta N_{\text{eff}}^\gamma} = \frac{3}{4.4} . \quad (4.1)$$

As we now explain, the fact that this ratio is known leads to a testable prediction for the CMB. This prediction is a consequence of the fact that although twin photons and twin neutrinos both constitute dark radiation and contribute to ΔN_{eff} , in detail their effects on the CMB are distinct, and can be distinguished. While the twin neutrinos free stream, the twin photons are prevented from free streaming by scattering off of twin electrons. Scattering and free streaming species have different effects on the CMB anisotropies [107–109]. Only after recombination has occurred in the mirror sector can twin photons free stream. The distinct effects of these two forms of dark radiation on the CMB anisotropies can be parametrized in terms of the free streaming fraction, defined as

$$f_\nu = \frac{\rho_{\text{free}}}{\rho_r} = \frac{\rho_{\text{free}}}{\rho_{\text{free}} + \rho_{\text{scatt}}} \quad (4.2)$$

Here ρ_{free} represents the energy density in free streaming radiation, ρ_{scatt} the energy density in scattering radiation, and ρ_r the total energy density in radiation. We use $N_{\text{eff}}^{\text{free}}$ and $N_{\text{eff}}^{\text{scatt}}$ to parametrize the energy densities in free streaming and scattering radiation in terms of the effective number of neutrinos. Then,

$$f_\nu = \frac{N_{\text{eff}}^{\text{free}}}{N_{\text{eff}}^{\text{free}} + N_{\text{eff}}^{\text{scatt}}} . \quad (4.3)$$

For small ΔN_{eff} , we have [110],

$$f_\nu - f_\nu^{\text{SM}} = \frac{0.41}{3} \left(0.59 \Delta N_{\text{eff}}^{\text{free}} - 0.41 \Delta N_{\text{eff}}^{\text{scatt}} \right) \quad (4.4)$$

For a given ΔN_{eff} , the amplitudes of the CMB modes depend on f_ν as [107, 108],

$$\frac{\delta C_\ell}{C_\ell} = -\frac{8}{15} f_\nu. \quad (4.5)$$

We see from this that the sign of the correction to the amplitude, relative to the SM, depends on whether the dark radiation free streams or scatters. For a given ΔN_{eff} the locations of the CMB peaks also depend on f_ν [109],

$$\delta \ell \approx -57 f_\nu \frac{\ell_A}{300}. \quad (4.6)$$

Here $\ell_A \approx 300$ represents the average angular spacing between the CMB peaks at large ℓ . Once again we see that the sign of the shift depends on whether the dark radiation free streams or scatters.

The MTH framework predicts the ratio of the energy densities in free streaming and scattering species. Prior to recombination in the twin sector, we have

$$\frac{\Delta N_{\text{eff}}^{\text{free}}}{\Delta N_{\text{eff}}^{\text{scatt}}} = \frac{\Delta N_{\text{eff}}^{\nu}}{\Delta N_{\text{eff}}^{\gamma}} = \frac{3}{4.4}, \quad f_\nu = f_\nu^{\text{SM}}. \quad (4.7)$$

Since ΔN_{eff} and f_ν can be independently determined from the CMB, this prediction for f_ν allows the MTH to be distinguished from other dark sector scenarios. The point is that while a non-zero contribution to ΔN_{eff} is a characteristic feature of any extension of the SM that contains some form of dark radiation, in general there is no reason to expect that $f_\nu = f_\nu^{\text{SM}}$. The fact that in the MTH the free streaming fraction is exactly the same as in the SM is because, in this framework, the dark radiation is composed of the twin counterparts of the SM photons and neutrinos, with the same relative energy densities. The prediction eq. (4.7) is not sensitive to the relative fractions of mirror hydrogen and mirror helium, but only requires that the energy density in the mirror component of dark matter be large enough to sustain acoustic oscillations during the CMB epoch, $r_{\text{all}} \gtrsim 0.1\%$.

A detailed calculation is needed to obtain a precise prediction for CMB observations, since the twin photon becomes free streaming after twin recombination. However, the point remains that the MTH makes a unique prediction for the ratio of energy densities, $\Delta N_{\text{eff}}^{\text{free}}/\Delta N_{\text{eff}}^{\text{scatt}}$, prior to twin recombination. Since CMB Stage-IV experiments are expected to be very sensitive to not just the total energy density in dark radiation, ΔN_{eff} , but also the free streaming fraction f_ν [35, 36], this provides a unique handle to test the MTH framework and distinguish it from other possible dark sectors.

5 Conclusions

In this paper, we have explored the cosmological signals associated with the mirror baryons, electrons, photons, and neutrinos in the MTH model. We have worked in a framework in which the cosmological problems of the original MTH proposal are assumed to be solved by late time asymmetric reheating after the two sectors have decoupled. We have primarily focused on the case in which the discrete Z_2 symmetry that relates the two sectors is only

softly broken. Then, in order for the little hierarchy problem to be addressed, the masses of particles in the mirror sector are restricted to lie in a limited range. This means that, although many of the late time thermal processes, such as BBN and recombination, are sensitive to the particle masses and couplings, we have still been able to draw a clear picture of the cosmology of the MTH framework, which exhibits several characteristic features. Therefore cosmological observations may allow this class of models, which are difficult to test in collider experiments, to be discovered.

From a study of TBBN, we have found that, in contrast to the SM, $\approx 75\%$ of the mass density in mirror particles density is contained in helium. Mirror hydrogen and helium remain ionized until close to the time of matter-radiation equality and, as in the SM, scattering off twin electrons and photons generates TBAO that suppress matter density perturbations. Although mirror helium recombines into neutral atoms before mirror hydrogen, its relatively larger abundance means that its impact on LSS, especially on shorter wavelengths, cannot be neglected. TBAO is ended by recombination in the mirror sector, leading to an oscillatory feature in the matter power spectrum in k -space with a frequency lower than that of SM BAO. This may allow an observation of TBAO in the near future in LSS observations. Current observations allow up to 5 (10)% of the matter density to come from twin hydrogen and helium for $\Delta N_{\text{eff}} = 0.4(0.1)$. In the absence of a signal, future LSS measurements will be able to tighten this bound to the 1% level.

In the MTH framework, we expect observable contributions to ΔN_{eff} from mirror photons and neutrinos. Although ΔN_{eff} itself is a free parameter, the relative energy densities in these two species are known. The mirror photons and mirror neutrinos have distinct effects on the CMB because, until recombination, the twin photons scatter off the twin electrons, which prevents them from free streaming. For any given ΔN_{eff} , this leads to a prediction for the corrections to the heights and locations of the CMB peaks that can be potentially be tested in future experiments.

Our analysis has primarily focused on the case in which the discrete Z_2 symmetry is only softly broken, so that the relative abundances of twin hydrogen and helium are predicted. However, for the purposes of comparison we have also explored scenarios in which the nuclei in the mirror sector are composed entirely of hydrogen, or entirely of helium. In particular, our studies capture the features of the important case in which, because of hard breaking of the discrete Z_2 symmetry in the Yukawa couplings of the light quarks, the mirror neutron is lighter than the mirror proton, and constitutes the dominant component of the observed dark matter, while mirror helium represents an acoustic subcomponent that gives rise to the signals we consider. Interestingly, we find that because hydrogen and helium recombine at different times, the matter power spectrum of the framework in which both nuclei are present exhibits distinctive features that may allow it to be distinguished from the case of atomic dark matter with just a single type of nucleus.

While we have focused on cosmological signals of the MTH framework, it is worth keeping in mind that this scenario can also give rise to striking astrophysical signals. During the formation of the galactic halo, the mirror baryonic component of dark matter in the Milky Way may have become re-ionized. If the ionized mirror baryons subsequently dissipated enough energy, they may have collapsed into a dark disk, as in the scenario

discussed in [111]. The details of this process depend sensitively on the MTH parameters, and on the initial distribution of dark matter in the galaxy. If the mirror baryons do form a second disk aligned with our own, current Gaia observations [112] already constrain the mirror baryons to constitute less than $\sim 1\%$ of the dark matter density. However, a careful study is needed to draw robust conclusions about the allowed parameter space. The distribution of mirror baryons and electrons within our galaxy will also determine how they might be discovered at current or future dark matter direct detection experiments. In the MTH framework, the SM and mirror sectors interact through the Higgs portal. Unfortunately, the resulting signal is far too small for direct detection of mirror hydrogen or mirror helium nuclei in any current or proposed experiment. However, in the event of a small kinetic mixing between the SM photon and its twin counterpart, the mirror baryons and electrons will acquire a tiny electric charge. The mirror sector can then interact with the SM through this portal, opening a new pathway for direct detection. Avoiding recoupling of the mirror sector after asymmetric reheating via $e\hat{e}$ scattering for $T \gtrsim \text{MeV}$ leads to an upper bound on the kinetic mixing parameter, $\epsilon \lesssim 10^{-9}$. In the MTH model, no kinetic mixing is generated through 3-loop order [1], and therefore even such small values of ϵ are radiatively stable. It follows that this bound can naturally be satisfied provided that the contributions to ϵ from UV physics are sufficiently small. Therefore ϵ of order 10^{-9} , corresponding to *nano-charged dark matter*, constitutes a natural sensitivity goal for future direct detection experiments. Such small kinetic mixings are also compatible with supernova bounds [113]. Matter direct detection have been previously studied in the context of mirror models (see [56] and references contained therein), the signals again depend sensitively on the details of the hidden sector. These striking astrophysical signals of the MTH framework, which clearly warrant further study, will be the subject of a companion paper [114].

The MTH scenario with minimal Z_2 breaking avoids almost all collider bounds, since it addresses the hierarchy problem with top partners that are neutral under all the SM forces. Our study demonstrates that in this framework the mirror baryons, electrons, photons and neutrinos can give rise to a rich plethora of distinctive cosmological and astrophysical signatures. Taken together, these may allow the nature of the hidden sector, and its possible relation to the SM and the hierarchy problem, to be discovered and probed in considerable detail. This opens the door to the tantalizing possibility that the first hints of naturalness may come from the sky, rather than from colliders.

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