Cogenesish from neutron-dark matter oscillations

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It was recently suggested that dark matter consists of $\sim$GeV particles that carry baryon number and mix with the neutron. We demonstrate that this could allow for resonant dark matter-neutron oscillations in the early universe, at finite temperature, leading to low-scale baryogenesis starting from a primordial dark matter asymmetry. In this scenario, the asymmetry transfer happens around 30 MeV, just before big bang nucleosynthesis. We illustrate the idea using a model with a dark $U(1)'$ gauge interaction, which has recently been suggested as a way of addressing the neutron lifetime anomaly. The asymmetric dark matter component of this model is both strongly self-interacting and leads to a suppression of matter density perturbations at small scales, allowing to mitigate the small-scale problems of cold dark matter cosmology. Future CMB experiments will be able to consistently probe, or firmly exclude, this scenario.

I. INTRODUCTION

One of the curious coincidences of the $\Lambda$CDM cosmological model is the rough similarity of the contributions from baryons and dark matter (DM) to the energy density, $\Omega_b \simeq 0.049$ versus $\Omega_{CDM} \simeq 0.26$ [1], given that in typical scenarios of the early universe they have a completely different origin. An appealing feature of many models of asymmetric dark matter is that this coincidence could be due to DM having a mass at the GeV scale, like baryons, and a shared mechanism for the generation of the two asymmetries. Then instead of needing separate mechanisms for baryogenesis and “darkogenesis,” there could be cogenesis through a single event, followed by some means of sharing the asymmetry between the standard model (SM) and DM sectors (for a review see [2]). Ref. [3] analyzes this scenario in terms of the lowest-dimension, gauge-invariant effective operators that would allow such a redistribution of a primordial asymmetry starting in the dark sector. For example, the operator $\chi u d d / A^2$ would induce roughly equal asymmetries between DM particles $\chi$ and quarks if it was in equilibrium at high temperatures [4].

In this work, we consider a novel situation where the $\chi u d d / A^2$ coupling only arises below the QCD phase transition. Then it is replaced by the unique relevant operator that could connect Dirac dark matter to the SM, namely

$$\mathcal{L}_{\text{mix}} = -\delta m \bar{n}\chi + \text{h.c.}$$

Our purpose is to show that it can produce the baryon asymmetry at low temperatures $\sim 30\text{ MeV}$ starting from a $\chi$ asymmetry (whose origin we do not try to specify here). This comes about by oscillations between $\chi$ and $n$, in analogy to neutrino oscillations, that are resonantly enhanced by finite-temperature effects. Oscillations of neutrons to a mirror sector partner have been studied extensively in [4-7], however, to our knowledge these oscillations have never been considered as a means for cogenesis. Similar ideas for producing the baryon asymmetry at low temperatures via oscillations of SM baryons have been recently suggested in Refs. [8, 9].

Recently motivation for the operator in Eq. (1) came from a quite different direction. It was suggested [10] that mixing between the neutron and DM could resolve a long-standing discrepancy between determinations of the neutron lifetime from decay-in-flight versus bottle measurements, by having a small dark decay channel [11]. The proposed decays $n \rightarrow \chi \gamma$ and $n \rightarrow \chi e^+ e^-$ were quickly ruled out by experimental searches [12, 13], suggesting a completely hidden channel, like $n \rightarrow \chi' \gamma'$ where $\gamma'$ is a dark photon. Even this model is ruled out by neutron star properties [14, 17], unless repulsive $\chi$ self-interactions are strong enough, requiring

$$m_{\chi'} \lesssim (45 - 60)\text{ MeV}$$

$^1$ This example requires $m_\chi < m_n$ since it would otherwise allow the DM to decay into quarks.

$^2$ We note that the residual symmetric baryon component is negligible by this time; see e.g. Fig. 5 of Ref. [8].

$^3$ It has also been proposed that neutron-mirror neutron oscillations could explain this puzzle [1].
to be satisfied \[18\]. (The uncertainty in the bound is due to the unknown nuclear equation of state.) This scenario is highly constrained by a number of observables sensitive to \(\gamma'\), including big bang nucleosynthesis (BBN), the cosmic microwave background (CMB), DM direct detection, supernovae, and structure formation effects from DM self-interactions.

In the present work we will show that it is possible to explain both the baryon asymmetry and the neutron decay anomaly within the same model, while at the same time evading all cosmological bounds, if \(m_{\gamma'} \sim 60\text{keV}\) and an additional species of dark radiation \(\nu'\) is introduced to enable the decay \(\gamma' \rightarrow \nu'\bar{\nu}'\). This entails moderate tension with BBN and CMB limits on extra radiation species, implying that it may be testable already in the near future. Alternatively, larger \(m_{\gamma'}\) consistent with \(\gamma' \rightarrow e^+e^-\), and hence without the need of introducing the additional degrees of freedom contributed by \(\nu'\), can be accommodated if we disregard the neutron lifetime anomaly.

II. THEORETICAL FRAMEWORK

We consider the same model as in Ref. \[18\], with an elementary Dirac dark matter particle \(\chi\) that carries baryon number and mixes with the neutron through the mass term \(|\tilde{\chi}\tilde{\chi}\rangle\), and a dark photon coupling to \(\chi\) with strength \(g'\). The \(n-\chi\) mass splitting must satisfy

\[ \Delta m \equiv m_n - m_\chi < 1.665\text{MeV} + m_{\gamma'} \]  

for stability of \(^9\text{Be}\), which is the most constraining nuclear decay \[11\].

Another relevant parameter is the ratio of the temperatures in the two sectors,

\[ \xi = \frac{T_{\gamma'}}{T} \]  

(or \(T_{\gamma'}/T\) at later times where \(\gamma' \rightarrow \nu'\bar{\nu}'\) has already occurred). We will show how \(\xi < 1\) is determined by the decoupling temperature between the two sectors, which turns out to be independent of details of the UV physics. The baryon asymmetry depends mainly on \(\Delta m\) and the mixing mass \(\delta m\), with weak dependence on the combination \(\gamma'\xi\).

The dark photon mass must satisfy \(m_{\gamma'} < \Delta m\) to explain the neutron lifetime anomaly. We will further show that \(\Delta m < 0.2\text{ MeV}\) is required to simultaneously respect the limit from \(n \rightarrow \chi\gamma\) and to obtain the observed baryon asymmetry. This bound kinematically forbids \(\gamma' \rightarrow e^+e^-\) decays, which in turn makes the dark photon longlived, such that it overcloses the universe \[18\] \[19\]. To avoid this, we supplement the model with an inert light particle, dark radiation, taken to be a massless Dirac neutrino \(\nu'\) that couples to \(\gamma'\) with strength \(Q_{\nu'\gamma'}\) and charge \(Q_{\nu'} \neq 1\). This is anomaly-free and allows for \(\gamma' \rightarrow \nu'\bar{\nu}'\), while the charge difference between \(\nu'\) and \(\chi\) forbids mass mixing between the states, that would lead to proton decay.

A further requirement is that the symmetric component of the dark matter must be sufficiently diluted by annihilations. The channels \(\tilde{\chi}\chi \rightarrow \gamma'\gamma', \nu'\bar{\nu}'\) turn out to be too inefficient since we need a small gauge coupling \(g' \sim 10^{-3}\) to satisfy previous requirements. This can be overcome by introducing a heavy dark fermion \(\psi\) with coupling \(\tilde{\chi}\phi\psi\) to the DM and the dark Higgs \(\phi\), allowing for \(\tilde{\chi}\chi \rightarrow \phi\phi^*\) through exchange of \(\psi\).

Finally, one should avoid leakage of the primordial \(\chi\) asymmetry into the SM sector at high temperatures, since this would generally produce too large a baryon asymmetry relative to the DM abundance. We point out a simple mechanism for enforcing this requirement in the present model, by adding a very weakly coupled interaction \(\tilde{\chi}\phi\psi\) involving a heavy Majorana dark sector fermion \(\psi'^*\).

III. \(\chi-n\) OSCILLATIONS

The oscillations of \(\chi\) into neutrons are determined by a \(2 \times 2\) Hamiltonian that includes the mass terms and thermal self-energy corrections from the forward scattering of the fermions on particles in the plasma. The most important such interactions are the elastic scattering of neutrons on pions and DM on dark photons.

The Hamiltonian is given by

\[ \mathcal{H} = \begin{pmatrix} \Delta E_n + m_n & \delta m \\ \delta m & \Delta E_\chi + m_\chi \end{pmatrix} \]  

The thermal energy correction \(\Delta E_n\) to the neutron has been calculated in Refs. \[20\] \[24\] using various techniques. Some of these results are plotted in Fig. 1 (left panel). The most reliable ones use dispersion relations, supplemented by experimentally measured scattering cross sections \[20\] \[21\]. We have digitized the results of Ref. \[21\] as our estimate of \(\Delta E_n\). Using thermal field theory techniques similar to Ref. \[25\] (see appendix \[A\] for details), we find for the DM energy shift

\[ \Delta E_\chi = g'^2 \frac{T^2_{\gamma'}}{8m_\chi}. \]  

This result is also plotted, for \(g'\xi = 1\), in Fig. 1(right panel).

The combination \(\Delta E_n - \Delta E_\chi\) is relevant for getting resonantly enhanced oscillations, which occur when \(\delta E = \Delta m + \Delta E_n - \Delta E_\chi\) vanishes and the eigenvalues of \(\mathcal{H}\) become degenerate. For illustration, we plot in Fig. 1(right panel) the difference in thermal self-energies, with \(g'\xi = 0, 0.5, 1\), along with a fiducial value of \(\Delta m = 0.2\text{ MeV}\). Clearly there are two resonant temperatures, where the curves cross the horizontal line. We will show that baryon production is dominated by the resonance at the lower temperature.
III.1. Heuristic approach

To find the efficiency of \( \chi \to n \) oscillations for producing the baryon asymmetry we first use an approximate formalism \[26\] that has the advantage of being simple and intuitive. Diagonalizing \( \mathcal{H} \) leads to the mixing angle \( \theta \),

\[
\tan(2\theta) = \frac{2\frac{\delta m}{\Delta m + \Delta E_n - \Delta E_\chi}}{\delta E} = \frac{2\frac{\delta m}{\Delta E}}{\Delta E}.
\]

The difference between the eigenvalues is given by

\[
|\delta\omega| = \sqrt{(\delta E)^2 + 4\frac{\delta m^2}{\Delta m + \Delta E_n - \Delta E_\chi}}.
\]

Imagine starting with a state that is purely \(|\chi\rangle\) at \( t = 0 \). Evolving it with the Hamiltonian \( \mathcal{H} \) in Eq. (5) leads to

\[
|\psi(t)\rangle = e^{-i(\omega_1 + \omega_2)t/2} \left[ \cos \left( \frac{\delta \omega}{2} t - i \cos 2\theta \sin \frac{\delta \omega}{2} t \right) |\chi\rangle - i \left( \sin 2\theta \sin \frac{\delta \omega}{2} t \right) |n\rangle \right].
\]

Therefore the probability to oscillate into a neutron is given by

\[
P_n(t) = \sin^2(2\theta) \sin^2(\delta \omega t/2).
\]

Because of the large rate of interactions \( \Gamma_n \) of neutrons on heat bath pions, however, there may not be time for a full oscillation. In general one should carry out a time average, over the short time scale \( 1/\Gamma_n \),

\[
\bar{P}_n = \Gamma_n \int_0^\infty dt e^{-\Gamma_n t} P_n(t) = \frac{2\frac{\delta m^2}{\delta E^2 + \Gamma_n^2}}{\delta \omega^2 + \Gamma_n^2},
\]

where we used \( \sin^2 2\theta = \frac{\delta m^2}{\delta m^2 + \delta E^2/4} \).

The rate of production of neutrons via oscillations is then \( \Gamma_{\text{osc}} = \bar{P}_n \Gamma_n \) per \( \chi \) particle (ignoring the much smaller elastic scattering rate \( \Gamma_\chi \) of DM on dark photons and scalars\[3\]). The inverse process of \( \chi \) production must proceed with the same rate, per neutron, so the number density of DM overall decreases as \( n_\chi = -\dot{\Gamma}_{\text{osc}} (n_\chi - n_n) \).

Since baryon number, and hence the total number of \( \chi \)'s plus neutrons is conserved, the equation determining the fraction \( f = n_n / (n_\chi + n_n) \) of DM converting to neutrons is thus

\[
f = \Gamma_{\text{osc}} (1 - 2f),
\]

which has the solution

\[
f = \frac{1}{2} \left( 1 - \exp \left( -2 \int \frac{dt}{\Gamma_{\text{osc}}} \right) \right) \cong \frac{1}{2} \left( 1 - \exp \left( -2 \int \frac{dT}{HT} \Gamma_n \bar{P}_n \right) \right)
\]

(13)

(with \( H \approx 1.66 \sqrt{g_s} T^2 / M_p, M_p = 1.22 \times 10^{19} \text{GeV} \) and \( g_s \approx 10.75 \) so long as \( P_{\text{ osc}} < 1 \). We need

\[
f \approx 0.16 \]

(14)

to get the desired abundance ratio \( \Omega_{\chi} / \Omega_B = (1 - f) / f \approx 5.3 \), dictated by the coincidence that \( m_\chi \approx m_n \).

It remains to determine the rate \( \Gamma_n \). The cross section for neutrons to scatter on pions at low energy is \[27, 28\]

\[
\sigma_{n\pi} = 4\pi a_0^2 \approx \frac{0.1}{m_n^2} \approx 2 \text{ mb},
\]

(15)

\[4\] The cross section for \( \phi \chi \) scattering from the interaction \[46\] is \( 10^{-7} \text{ mb} \), saturating the bound \( \Phi \). For \( m_\phi = 60 \text{ MeV} \), for example, \( n_\phi / n_{\chi} \) is only ~7 at the resonance temperature. See also Appendix A.

\[5\] We average the contributions from the \( I = 1/2 \) and \( 3/2 \) isospin scattering lengths as \( \sigma = 4\pi(a_1^2/2 + a_3^2/3) \).

Figure 1. Left panel. Thermal shifts \( \Delta E_n \) to neutron mass from dispersion relations \[20, 21\] (solid) and chiral perturbation theory \[22\] (dotted). Right panel. Effect of including the DM thermal mass. The dotted curve is \( \Delta E_n \) for \( g'_\xi = 1 \) from Eq. (6). Lower curves show the difference in the thermal self-energies, \( \Delta E_n - \Delta E_\chi \), for \( g'_\xi = 0, 0.5, 1 \). For comparison we also indicate the maximal mass difference \( \Delta m \equiv m_n - m_\chi \lesssim 0.2 \text{ MeV} \) compatible with baryon asymmetry generation and stability of \( ^7\text{Be} \) through \( n \to \chi + \gamma \).
The integrand $\Gamma_n P_n/HT$ of Eq. (13) versus the photon temperature $T$, showing how baryon number generation is dominated by the resonant enhancement of $\chi$-$n$ oscillations, using the fiducial parameters of Eq. (21). Figure 2.

\[
\Gamma_n = n_\pi \langle \sigma v \rangle
\]

with $\langle v \rangle = \sqrt{8T/(\pi m_\pi)} \approx 1.6 \sqrt{T/m_\pi}$, where $n_\pi$ is the thermal density of pions, including a factor of 3 for isospin multiplicity.

This formalism makes it transparent that baryon production is dominated by resonant enhancement of the oscillations by the finite-T contributions to $\delta E$, which allow for $\delta E(Tr) = 0$ at the resonance temperature $Tr$. At $T = 0$, tight constraints on $\Delta m$ and $\delta m$ (to be discussed below) imply a vacuum mixing angle $\theta = 2\delta m/\Delta m$ that is by far too small to allow oscillations efficient enough for baryogenesis, but at $Tr$ the mixing is maximal, although the oscillation probability is reduced by the damping from $\Gamma_n$ in Eq. (11). This is illustrated in Fig. 2 for a benchmark model, Eq. (21) below, for which $Tr \approx 31$ MeV. Because of the damping, the final efficiency for baryon production scales as $\delta m^2/\Gamma_n$.

### III.2. Boltzmann equations

A more rigorous approach is to solve Boltzmann equations for the relative abundances, including off-diagonal terms arising from the density matrix, that take into account the oscillations. The formalism is described for oscillations of dark matter and its antiparticle in Ref. [29, 30], and requires only small modifications for neutron-dark matter oscillations. In terms of the independent variable $x = m_n/T$, the equation for the matrix of abundances $Y = n/s$ is

\[
dY = -\frac{i}{Hx} \text{[H, Y]} - \frac{\Gamma_n}{2Hx} [P_n, [P_n, Y]] - \frac{\Gamma_\chi}{2Hx} [P_\chi, [P_\chi, Y]]
\]

\[
= -\frac{i}{Hx} \begin{pmatrix} -\delta m Y_+ & -\delta m \delta Y - \delta E Y_{12} \\ \delta m \delta Y - \delta E Y_{21} & \delta m Y_- \end{pmatrix} - \frac{\Gamma_n + \Gamma_\chi}{2Hx} \begin{pmatrix} 0 & 12 \\ 21 & 0 \end{pmatrix},
\]

where we defined $\delta Y \equiv Y_{11} - Y_{22} = Y_n - Y_\chi$, $Y_+ = Y_{12} \pm Y_{21}$, and $P_{n,\chi}$ projects onto the state 1 (the neutron) or 2 (the $\chi$), respectively. As before, we will ignore $\Gamma_\chi$ since it is much smaller than $\Gamma_n$.

Since $\chi$ can be assigned a baryon number, which is conserved, the combination $Y_{11} + Y_{22} = Y_n + Y_\chi$ does not evolve. It is then convenient to recast the equations in terms of $\delta Y$, $Y_+$ and $\eta_- \equiv i \delta m Y_-,

\[
\delta Y' = \frac{2}{Hx} \eta_-,
\]

\[
\eta_- = \frac{\delta m}{Hx} (-2\delta m \delta Y + \delta E Y_+) - \frac{\Gamma_n}{2Hx} \eta_-,
\]

\[
Y_+ = -\frac{\delta E}{\delta m Hx} \eta_- - \frac{\Gamma_n}{2Hx} Y_+.
\]

This form has the advantage that all the functions are now real-valued, and the coefficients are of order unity, for cases of interest. The rescaling of $\eta_-$ by $\delta m$ alleviates the problem of stiffness in the numerical integration. We integrate the system (18) with initial conditions $Y_{22} = Y_\chi^\infty$ and $Y_{ij} = 0$ for the other components. Since $Y_{11} + Y_{22}$ is conserved, the fraction of DM that converts to neutrons is $f = Y_{11}/Y_\chi^\infty = 1 - Y_{22}/Y_\chi^\infty$. We note that this fraction is necessarily independent of the initial DM abundance $Y_\chi^\infty$, so for any value of $f$ we can simply rescale $Y_\chi^\infty$ such that the observed DM density $\rho_{CDM} = sm_\chi Y_\chi^\infty$ is recovered at late times (in practice, we use $Y_\chi^\infty = 1$ as initial condition). Given that the DM and neutron are nearly degenerate, we can identify $f/(1 - f) = \Omega_b/\Omega_{CDM}$ at late times.

### III.3. Results

Fig. 3 shows the correlation between $\delta m$ and $\Delta m$ needed to get the observed baryon asymmetry, by fixing $f$ as stated in Eq. (14), for several values of $g' \xi$. Each curve is doubled, with a small splitting, showing the good

6. Here $s$ is taken to be the entropy in the visible sector. Since baryon number $Y_1 + Y_2$ is conserved, we are only concerned with the fraction $f$ of DM that converts to baryons, from which $s$ divides out.

7. It is worth pointing out that the elastic scatterings on $n$ and $\chi$ enter equally into the damping of the oscillations, regardless of whether one considers $n \to \chi$ or $\chi \to n$ transitions. This is because either interaction can measure the state of the system, regardless of the outcome of the measurement.
which the mixing mass should satisfy \[10\] the baryon asymmetry can be explained only if \[\delta m\] is larger due to the higher density of pions, explaining why \[\delta m\] must increase with \[\Delta m\] to keep the baryon asymmetry constant.

On the same figure we plot the upper limit on the vacuum mixing angle from null searches for the decay \[n \rightarrow \chi\gamma\] (blue curve), whose branching ratio must be \[\lesssim 0.17\] times the 0.9\% needed to resolve the lifetime discrepancy \[12\]. This gives

\[
\delta m \lesssim 5 \times 10^{-11} \text{MeV} \left( \frac{1 \text{MeV}}{\Delta m} \right) \frac{1}{2} \tag{19}
\]

Fig. 3 thus implies that for reasonable values of \[g'\xi\], the baryon asymmetry can be explained only if \[\Delta m \lesssim 0.2 \text{MeV}\].

We next consider the neutron lifetime anomaly, for which the mixing mass should satisfy \[10\]

\[
\delta m = \frac{4.2 \times 10^{-13}}{(1 - (m_{\gamma'}/\Delta m)^2)^{3/4}} \left( \frac{m_{\gamma'}}{g'} \right) \left( \frac{1 \text{MeV}}{\Delta m} \right)^{1/2} \tag{20}
\]

Combining this with the limit on \[m_{\gamma'}/g'\] from neutron stars in Eq. \[2\], this becomes a constraint on \[\delta m\] as a function of \[\Delta m\] and \[g'\] (with a slight dependence upon the assumed nuclear equation of state). Using \[m_{\gamma'}/g' = 60 \text{MeV}\], we indicate in Fig. 3 the values of \[\delta m\] that are required to resolve the lifetime anomaly for two representative values of \[g' = 0.001, 0.002\] (solid red and orange curves, respectively). It is possible to choose a smaller value of \[m_{\gamma'}/g'\] (which is consistent with the neutron star constraint) to plot these constraints on resolving the lifetime anomaly (and still satisfy the neutron star constraint \[2\]), leading to smaller values of this would lead to \[\delta m\]. The choice \[m_{\gamma'}/g' = 60 \text{MeV}\] is plotted for illustration (dashed red and orange curves), for the same two values of \[g'\].

The dark photon mass in these scenarios is significantly below the \[2m_e\] threshold. In the absence of additional channels the dominant decay mode would then be \[\gamma' \rightarrow 3\gamma\], which is so slow that the initially thermally distributed dark photons would be metastable and overclose the universe \[19\] (see Ref. \[18\] for additional constraints). This motivates us to introduce the additional massless Dirac neutrino \[\nu'\] already mentioned in Section II enabling the much more efficient decay \[\gamma' \rightarrow \nu'\nu'\].

Alternatively, if we ignore the neutron decay anomaly (assuming, for example, that it is due to experimental error) and insist only on low-scale baryogenesis, it is not necessary for \[g'\] to be small. This in turn allows heavier dark photons and the possibility for fast \[\gamma' \rightarrow e^+e^-\] decays through kinetic mixing, via the Lagrangian term \[-(e/2) F_{\mu\nu}F^{\mu\nu}\], without the need to introduce additional light degrees of freedom like \[\nu'\].

In this simpler scenario, the additional heavy fermion \[\psi\] described in Sec. II is also no longer required, since with 

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8 Some extra field carrying U(1) charge will be needed to insure U(1) charge neutrality of the universe at early times, as we will discuss in Section IV.D. However it need not be light, and so will not necessarily contribute extra radiation degrees of freedom that contribute to \[N_{\text{eff}}\]. See appendix E for more details.
sizeable values of $g'$ the process $\chi \bar{\chi} \to \gamma'\gamma'$ can efficiently annihilate away the symmetric DM component. Possible benchmark values for such a minimal cogenesis scenario are

$$\begin{align*}
\text{B2} & \quad \Delta m = 0.15 \text{ MeV}, \quad \delta m = 1.2 \times 10^{-10} \text{ MeV} \\
g'\xi = 0.5, \quad m_{\gamma'} = 50 \text{ MeV}.
\end{align*}$$

(22)

In the B2 model, an upper bound on the kinetic mixing parameter $\epsilon$ comes from the scattering interactions of $\chi$ off protons, with a cross section

$$\sigma_{\chi p} = \frac{4\alpha(g')^2 \mu_{px}^2}{m_A^2},$$

(23)

where $\mu_{px}$ is the reduced mass of $\chi$ and the proton. Results from CRESST-III limit $g'\epsilon \lesssim 1.2 \times 10^{-7}$ [31]. Even more constraining are the beam dump experiments at Orsay [32] and SLAC (E137) [33], which limit $\epsilon \gtrsim 1.5 \times 10^{-5}$ or $\epsilon \lesssim 4.2 \times 10^{-8}$ [34], and observations of the flux of neutrinos from supernova 1987A, which constrain $\epsilon \gtrsim 4.7 \times 10^{-8}$ or $\epsilon \lesssim 3.2 \times 10^{-10}$ [35].

The decay of the dark photon to electrons and positrons proceeds with a width and lifetime given by

$$\Gamma = \frac{\alpha\epsilon^2}{3m_{\gamma'}} \left( m_{\gamma'}^2 + 2m_e^2 \right) \left( 1 - 4m_e^2/m_{\gamma'}^2 \right)^{1/2},$$

$$\tau \cong 0.5 \left( \frac{10^{-10}}{\epsilon} \right)^2 \text{s}.$$

(24)

The above constraints combine to limit $\tau \gtrsim 0.05 \text{s}$, which is sufficiently early to avoid disturbing the production of light elements during BBN [36, 37]. As will be discussed in section IV.3, interactions other than the kinetic mixing can thermalize the dark and visible sectors, so there is no constraint on $\epsilon$ from this requirement.

IV. THERMAL HISTORY

To have a fully consistent scenario, we need to address several issues: (1) redistribution of the presumed initial $\chi$ asymmetry into baryons should not have taken place at a higher temperature through the interactions present in our model, since this might be more important than the $\chi$-$n$ oscillations; (2) the $U(1)'$-breaking transition where $\phi$ gets its VEV must occur before the oscillations go through resonance; (3) the additional light degrees of freedom in the dark sector must be compatible with constraints from BBN and CMB; (4) the symmetric component of the $\chi$ relic density must be sufficiently suppressed, to avoid having more dark matter than is observed; (5) the self-interaction cross section of $\chi$ should respect constraints from structure formation. We consider these in turn. The constraints (2–5) are automatically satisfied for the benchmark point B2 if they are fulfilled for B1, so we adopt the B1 parameters in this section.

IV.1. Sequestering of $B$ at high $T$

Our assumption is that initial asymmetries in $\chi$, $\phi$, and $\nu'$ were somehow created at a high temperature, while the standard model baryon and lepton asymmetries were initially vanishing. Asymmetries in $\phi$ or $\nu'$ are needed in addition to that in $\chi$ to maintain $U(1)'$ charge neutrality of the universe, since $U(1)'$ remains unbroken until $T \sim 100 \text{ MeV}$ (see section IV.2). However the UV model of reference [18], augmented for the present purposes (see section IV.4), includes interactions which at high temperatures will redistribute these initial asymmetries into the other dark sector particles as well as into the visible sector. The relevant interactions are

$$\begin{align*}
\mathcal{L}_{\text{UV}} &= \lambda_1 d^a P_L \chi_{1,a} + \lambda_2 \epsilon^{abc} \bar{u}^a_c P_R d_b \Phi_2, \\
&+ \lambda_3 \bar{\chi} \phi \psi + \mu \phi \Phi_{1,a} \Phi^*_{2,a} + \text{H.c.}.
\end{align*}$$

(25)

$\lambda_{1,2,3}$ are dimensionless coupling constants, while $\mu$ is a constant with units of mass. $\Phi_1$ and $\Phi_2$ are both TeV scale scalar fields that are colored triplets with baryon number $-2/3$, and $\Phi_1$ also carries $U(1)'$ charge. The resulting decays and inverse decays will be in equilibrium at high temperatures, redistributing any initial asymmetry amongst all particles in the plasma. Generically, this will induce an early baryon asymmetry of the same order as that in dark matter, contrary to our hypothesis that only neutron-DM oscillations are important.

To see this, it is sufficient to consider temperatures just below the electroweak phase transition (EWPT), $T_c \cong 160 \text{ GeV}$, at which the relevant SM interactions are sphalerons (still in equilibrium since the EWPT is a cross-over transition) and $W^- \leftrightarrow d \bar{u}$, $W^- \leftrightarrow e \bar{v}$. Following Ref. [39] we take the up- and down-type quarks of all generations to have common chemical potentials $\mu_u$ and $\mu_d$, and denote the sum of neutrino chemical potentials by $\mu_\nu$. (The charged lepton potentials $\mu_\ell$ are eliminated in favor of $\mu_W$ and $\mu_\nu$.) At this scale the effective operator from integrating out $\Phi_i$ is

$$\frac{\lambda_1 \lambda_2 \mu}{m_{\phi_1}^2 m_{\phi_2}^2} \epsilon^{abc} (\bar{u}^c \gamma_i P_R d_b) \left( \bar{\chi} \gamma_i P_R d_c \right),$$

(26)

which is in equilibrium down to $T \cong 11 \text{ GeV}$ for a light dark scalar $\phi$ (see appendix B). The coefficient of (26) is fixed, since it gives rise to the $n$-$\chi$ mass mixing term when $\phi$ gets its VEV,

$$\delta m = \frac{\lambda_1 \lambda_2 \beta(\phi) \mu}{m_{\phi_1}^2 m_{\phi_2}^2} \sim 10^{-10} \text{ MeV}.$$

(27)

where $\beta = 0.014 \text{ GeV}^3$ [35] from the lattice matrix element $\langle n | u d d | 0 \rangle$.

9 After this interaction no longer keeps $\chi$ in equilibrium, it will lead to $\chi$ decays before the QCD phase transition. However, the lifetime for such a decay is $\sim 10$ seconds for our benchmark points, so the effect of these decays on the abundance of $\chi$ particles will be minimal.
The equilibrium constraints from Eq. (26) and the SM interactions [39] are then
\[ \mu_\chi = \mu_\phi + \mu_u + 2\mu_d \] (28)
\[ \mu_W = \mu_d - \mu_u \] (29)
\[ 0 = 9\mu_u + 6\mu_W + \mu_\nu \] (30)
These are supplemented by the vanishing of the conserved charges, electric and U(1)',
\[ Q \propto 3\mu_u - \mu_\nu - 9\mu_W = 0 \] (31)
\[ Q' \propto \mu_\chi + \mu_\phi + Q'_\nu \mu_\nu = 0 \] (32)
where we have used that \( Q'_\nu \) is the \( \nu' \) charge relative to that of \( \chi \) or \( \phi \). These equations give quark asymmetries
\[ \mu_u = \frac{1}{11}(\mu_\chi - \mu_\phi), \quad \mu_d = 5\mu_u \] (33)
unless \( \mu_\chi = \mu_\phi \) at this temperature. Such a cancellation does not come about without tuning the relative initial asymmetries of \( \chi \) and \( \nu' \). How these chemical potentials are related to the initial asymmetries, possibly generated at a much higher scale, is discussed in appendix C. The \( \nu' \) asymmetry is determined by Eq. (32) alone.

However, there is a simple low-energy mechanism for enforcing \( \mu_\phi = \mu_\chi \), supposing there is a massive Majorana fermion \( \psi' \) in the dark sector with coupling
\[ \lambda' \phi \psi' + H.c. \] (34)
If in equilibrium, this gives the desired relation \( \mu_\phi = \mu_\chi \), since \( \mu_\nu' = 0 \) by the Majorana nature of \( \psi' \). The decays \( \psi' \rightarrow \chi \phi \) are in equilibrium for \( T \geq T_c = 160 \text{ GeV} \) so long as
\[ \frac{\lambda'^2}{m_{\psi'}} \geq 0.9 \times 10^{-8} \left( \frac{100 \text{ GeV}}{m_{\psi'}} \right)^{1/2} \] (35)
This supplies the extra condition needed to ensure that no baryon asymmetry is induced prior to the onset of \( n-\chi \) oscillations.

The new interaction (34) is potentially dangerous since it violates baryon number. The dominant process is the \( \Delta B = 2 \) oscillations of \( \chi-\bar{\chi} \) due to the mass term
\[ \frac{1}{2} \delta m_\chi \bar{\chi} \chi^C + H.c. = \frac{\lambda'^2 v'^2}{2m_{\psi'}} \bar{\chi} \chi^C + H.c. \] (36)
arising after \( \phi \) gets its VEV \( v' / \sqrt{2} \) and integrating out \( \psi' \). In our scenario, this can be very small, \( \delta m_\chi \equiv 3 \times 10^{-18} \text{ MeV} \) for \( v' = 60 \text{ MeV} \) as we have assumed, and \( \lambda' \) saturating the bound (35) with \( m_\phi = 100 \text{ GeV} \). Such oscillations have been studied in detail in Ref. [30]. They are damped by the scattering of \( \chi \) on dark photons as long as the scattering rate \( \Gamma_\chi \propto g^4 T^3 / m_{\psi'}^2 \) exceeds the oscillation rate \( 2\delta m_\chi \), true for temperatures
\[ T \gtrsim \left( \frac{2\delta m_\chi m_{\psi'}^2}{g'^4} \right)^{1/3} \sim 3 \text{ keV} \] (37)
which is far below the scales of interest for co genesis through \( \chi-n \) oscillations. The estimate (37) applies for the B1 benchmark model, but is even lower for B2.

IV.2. U(1)'-breaking phase transition

The dark scalar \( \phi \) must get its VEV before the \( n-\chi \) oscillations begin, since the mass mixing term \( \delta m \) coming from (26) is proportional to \( \langle \phi \rangle \) in our model. Because the gauge coupling \( g' \) is very small, the phase transition in the dark sector is controlled by the dark scalar \( \phi \) alone. Supposing its potential is
\[ V_0 = \lambda_\phi (|\phi|^2 - v'^2 / 2)^2 \] (38)
the field-dependent masses for the real and imaginary parts of \( \phi = (\varphi_r + i\varphi_i) / \sqrt{2} \) are
\[ m_r^2 = \frac{\partial^2 V}{\partial \varphi_r^2} = \lambda_\phi (6|\phi|^2 - v'^2), \]
\[ m_i^2 = \frac{\partial^2 V}{\partial \varphi_i^2} = \lambda_\phi (2|\phi|^2 - v'^2) \] (39)
(assuming that \( \varphi_r \) is the component that gets the VEV). At mean-field level, the finite-temperature contribution to the potential is [40]
\[ V_T = \frac{T'^2}{24} (m_r^2 + m_i^2) \] (40)
where \( T' \) is the temperature of the dark sector. Combining (38) and (40) one finds that the critical temperature, where the curvature of the full potential at \( \phi = 0 \) vanishes, is given by
\[ T_c = \sqrt{3} v' = \sqrt{3} \frac{m_{\psi'}}{g'} \approx 104 \text{ MeV} \] (41)
independently of \( \lambda_\phi \), so long as \( \lambda_\phi \gg g'^2 \). Since \( T' = \xi T \) with \( \xi < 1 \), we see that the critical temperature as measured in the visible sector is even larger, \( T_c > 104 \text{ MeV} \). This is comfortably above the resonance temperature \( T_3 \sim 30 \text{ MeV} \) for our benchmark model.

IV.3. Constraints on extra degrees of freedom

The dark sector in our model has new light degrees of freedom at the scale of BBN, the ~60 keV dark photon and two massless Weyl neutrinos \( \nu' \), the latter being also present during temperatures relevant for the CMB. We show in appendix B that the effective operator (26) keeps the two sectors in equilibrium until \( T \approx 11 \text{ GeV} \), irrespective of the details of the UV model parameters. This relatively high decoupling temperature helps to reduce the number of extra radiation degrees of freedom, conventionally parametrized as the number of additional effective neutrino species, \( \Delta N_{\text{eff}} \).

Following decoupling, the entropies in the two sectors are separately conserved. During the QCD transition, the relative temperature \( \xi = T'/T \) goes down because the number of SM degrees of freedom, \( g_{\text{SM}}(11 \text{ GeV}) \approx 86.25 \),
decreases by a greater factor than the corresponding decrease in the dark sector \((\chi + \nu' + \phi + \gamma' \rightarrow \nu' + \gamma')\). We thus expect

\[
\xi_{\text{BBN}} \cong \left(\frac{86.25}{10.75}\right)^{-1/3} \left(\frac{6.5}{11}\right)^{-1/3} \cong 0.59
\]

at \(T > m_e\) when electrons are still in equilibrium. At this temperature, the dark photons and neutrinos are both present in the dark sector, and they contribute

\[
\Delta N_{\text{eff}} = \left(\frac{4}{7} \times 3 + 2\right) \xi_{\text{BBN}} \sim 0.47
\]

to the effective number of neutrino species. This is below the 3\(\sigma\) bound \(\Delta N_{\text{eff}} < 0.54\), but larger than the 2\(\sigma\) bound \(\Delta N_{\text{eff}} < 0.31\) allowed by BBN \[41\] (similar BBN limits are reported by Cyburt et al. \[42\]).

At lower temperatures, the number of relativistic degrees of freedom again changes in both sectors, as electrons and dark photons disappear. At the time of the CMB, we have a temperature ratio of

\[
\xi_{\text{CMB}} \cong \left(\frac{10.75}{3.91} \times \frac{3.5}{6.5}\right)^{-1/3} \xi_{\text{BBN}} \cong 0.52.
\]

Taking into account that the temperature ratio of SM neutrinos and photons has also increased by that time, we hence get

\[
\Delta N_{\text{eff}} = 2\xi_{\text{CMB}}^4 \left(11/4\right)^{4/3} \sim 0.56.
\]

This is above the 3\(\sigma\) limit of \(\Delta N_{\text{eff}} < 0.45\) from Planck \[43\], combining CMB and BAO measurements. But we note that this limit weakens to \(\Delta N_{\text{eff}} < 0.55\) (0.66) at 2\(\sigma\) (3\(\sigma\)) \[44\] when adding the direct measurement of the Hubble rate, \(H_0 = (73.45 \pm 1.66)\) km s\(^{-1}\) Mpc\(^{-1}\) \[43\], to these datasets.

For a more quantitative treatment we implemented the full temperature dependence of the energy density for each dark sector d.o.f. in DarkSUSY \[46\], rather than the relativistic limit used in the above estimates, and plot the expected \(\Delta N_{\text{eff}}\) as a function of decoupling temperature \(T_{\text{dec}}\) in Fig. 4. The red line shows \(\Delta N_{\text{eff}}\) at CMB times, while the blue line shows this quantity for \(T = 1\) MeV; the blue band corresponds to \(\Delta N_{\text{eff}}\) for the whole range of temperatures \(0.1\) MeV < \(T < 5\) MeV most strongly probed by BBN. We also indicate in this plot the already mentioned constraints from BBN and CMB. For the benchmark point B1, Fig. 4 confirms the above estimates for \(\Delta N_{\text{eff}}\). At the same time, this figure illustrates the effect of changing \(T_{\text{dec}}\). For example, for a smaller \(\langle \phi \rangle = m_{\gamma'}/g'\) the coefficient in front of Eq. \[26\] would increase, leading to a decrease in \(T_{\text{dec}}\) that sharpens the tension with BBN and CMB. A larger value of \(\langle \phi \rangle\), on the other hand, cannot be achieved in view of the neutron star constraint \[2\]. To decrease the tension with \(\Delta N_{\text{eff}}\) would thus require to modify the low-energy operator given in Eq. \[26\] – or to make the dark photon and Higgs non-relativistic already during BBN, as in benchmark point B2. Finally, if decoupling happens while some of the heavy degrees of freedom are still in equilibrium, this will increase the entropy ultimately dumped into \(\nu'\); we illustrate this in Fig. 4 by showing the effect of an additional fermion with mass \(m_\psi = 50\) GeV (dotted lines).

IV.4. Dark matter annihilation and kinetic decoupling

Although we are assuming that the asymmetric component of the DM abundance is generated by an unknown mechanism, the symmetric (thermal) component is determined by the annihilations of \(\chi\bar{\chi}\) into lighter particles. Because of the small value of \(g' = 10^{-5}\), the cross sections for \(\chi\bar{\chi} \rightarrow \gamma'\gamma'\) and \(\chi\bar{\chi} \rightarrow \nu'\nu'\) are inadequate for avoiding overclosure. However the channel \(\chi\bar{\chi} \rightarrow \phi\phi^*\) becomes available if we introduce a neutral Dirac fermion \(\psi\) with coupling

\[
\lambda_3 \chi\bar{\phi}\psi + \text{H.c.}
\]

This is consistent with stability of \(\chi\) as long as \(m_\psi > m_\chi\). The cross section for \(\chi\bar{\chi} \rightarrow \phi\phi^*\) by \(\psi\) exchange, to leading order in the center-of-mass velocity \(v_{\text{cm}}\) of the DM and for \(m_\psi \gg m_\chi\), is

\[
\sigma v_{\text{rel}} = \frac{\lambda_3^3 v_{\text{cm}}^2}{16 \pi m_\psi^2}.
\]

After therally averaging, we find that to avoid overclosure one needs \(m_\psi \lambda_3^2 \lesssim 1\) TeV \[45\].

For a more quantitative estimate one needs to take into account the non-standard temperature evolution in
the dark sector. To this end, we implemented the exact $\xi(T)$ dependence as well as the full expressions for $\sigma_{\text{rel}}$ in DarkSUSY. Solving the Boltzmann equation numerically for input parameters in the range of interest, we find that the following provides a reliable fit to the relic density of the symmetric component in the absence of an asymmetric component:

$$\left(\Omega h^2\right)_{\chi+\bar{\chi}} \approx 0.19 \left(\frac{m_{\psi}/\chi^2}{500 \text{ GeV}}\right)^{1.82} \ (48)$$

The additional presence of the asymmetric component, however, makes the annihilation of $\chi$ more efficient. From Fig. 4 in Ref. [45], we can read off that already an annihilation rate $\langle \sigma v \rangle$ a factor of 2 larger than the cross section $\langle \sigma v \rangle_0$ that would give the correct relic density in the purely symmetric case, i.e. in accordance with Eq. (18), results in a suppression of the symmetric component by a factor of almost 50. Choosing this reference relic density as $\left(\Omega h^2\right)_{\chi+\bar{\chi}} = 0.145$ (rather than the standard value of $(\Omega h^2)_{\text{DM}} \approx 0.12$ because 17% of the DM will convert to neutrons), this translates to the requirement

$$\frac{m_{\psi}}{\chi^2} \lesssim 300 \text{ GeV} \ . \ (49)$$

This condition is easily satisfied as long as $\lambda_3$ is not too small (note that decreasing $m_{\psi}$ below about 100 GeV would increase the tension with $\Delta N_{\text{eff}}$, c.f. Fig. 4).

Even after freeze-out, both the symmetric and the asymmetric DM component are kept in local (kinetic) equilibrium with $\nu'$. This is exactly one of the scenarios studied in Ref. [47] where, due to the presence of a light mediator, elastic scattering between $\chi$ and $\nu'$ is highly efficient and kinetic decoupling therefore happens significantly later than for standard WIMP candidates. The dark acoustic oscillations in the $\chi$-$\nu'$ fluid during and after decoupling lead to a cutoff in the power spectrum of matter density perturbations with an associated mass scale of

$$M_{\text{cut}} \simeq 2 \cdot 10^9 M_\odot \left(\frac{0.5}{\xi}\right)^2 Q_{\nu'}^2 \left(\frac{g'}{0.001}\right)^3 \left(\frac{m_{\nu'}}{60 \text{ keV}}\right)^{-3} \ . \ (50)$$

A warm DM candidate with mass $[10^{11} h^{-1} M_\odot/M_{\text{cut}}]^{1/4}$ keV generates an almost identical small-scale suppression of the nonlinear power spectrum [48]. This implies that Lyman-$\alpha$ constraints on our model can be easily evaded as long as the charge of the dark sector neutrino is not too large, $Q_{\nu'} \lesssim 1$ [49, 51]. For $Q_{\nu'} \sim 1$, in fact, the resulting mild suppression of the power spectrum might help to alleviate the missing satellites problem [47, 52].

IV.5. Dark matter self interactions

The cross section for $\chi\chi \rightarrow \chi\chi$ scattering is nonperturbatively enhanced by multiple $\gamma'$ exchange, and hence cannot be calculated in the Born limit. We use DarkSUSY to calculate the transfer cross section for a repulsive Yukawa potential in the classical limit $m_{\chi_{\text{rel}}} \gg m_{\nu'}$, based on parameterizations from Ref. [53], averaged over a Maxwellian velocity distribution with a most probable velocity $v_0$. For our benchmark model (21) and $v_0 = 30 \text{ km/s}$, this leads to $\langle \sigma v \rangle/m_{\chi} \sim 0.15 \text{ cm}^2/\text{g}$, which may be strong enough to visibly affect the structure of dwarf galaxies [54–56]. The cross sections required to fully address the small-scale structure problems of cold DM are typically quoted to be a bit larger. This can be achieved by choosing a slightly smaller dark photon mass: for $m_{\nu'} = 30 \text{ keV}$ for example, still respecting the constraint on $M_{\text{cut}}$ implied by Eq. (50), we find $\langle \sigma v \rangle/m_{\chi} \sim 0.6 \text{ cm}^2/\text{g}$. The dependence of the transfer cross section on $v_0$ is shown in Fig. 5. The self-interaction rate drops very sharply with the typical velocity, indicating that constraints on cluster scales [57–59] are not relevant for our model.

V. CONCLUSIONS

We have demonstrated that cogenesis of baryons, starting from an initial dark matter asymmetry, can be successfully implemented through dark matter-neutron oscillations. It is intriguing that the same model invented to address the neutron lifetime anomaly in Ref. [18] is capable, with small modifications, of producing the observed baryon asymmetry in this way. A key requirement is that a somewhat smaller mass splitting $\Delta m = m_n - m_{\chi} \sim 0.1 \text{ MeV}$ is now needed, relative to resolving only the lifetime anomaly. This arises because the baryogenesis curves in Fig. 5 fall in the region excluded by $n \rightarrow \chi\gamma$ limits at larger $\Delta m$. Simultaneously

![Figure 5. Phase space averaged self-interaction cross section over mass, $\langle \sigma v \rangle/m_{\chi}$, versus characteristic DM velocity $v_0$, for $g' = 0.001$ and $m_{\nu'} = 60 \text{ keV}$ (solid) or $m_{\nu'} = 30 \text{ keV}$ (dashed). The vertical bands indicate typical velocities encountered in dwarf spheroidal galaxies, Milky Way-sized galaxies, and clusters. Mitigating the ACDM small-scale problems requires $\langle \sigma v \rangle/m_{\chi} \sim 0.1 - 1 \text{ cm}^2/\text{g}$ at dwarf scales.](image)
satisfying the neutron lifetime anomaly requires small $g' \sim 10^{-3}$ and a consequently lower dark photon mass, $m'_\gamma \sim 60$ keV, mandating an additional dark radiation species $\nu'$ so that dark photons can decay fast enough to avoid problems with BBN and the CMB. Intriguingly, the combination of dark radiation and a very light dark photon may help to alleviate known issues of $\Lambda$CDM cosmology at small scales.

The most ambitious version of our scenario, where both baryogenesis and the neutron lifetime anomaly are treated, is challenged by extra radiation contributing to $N_{\text{eff}}$, being near the edge of current cosmological constraints. Updated BBN constraints and future CMB observations, like the planned Simons observatory [60], are likely to provide the most sensitive experimental test of the proposal. On the other hand, no dark radiation is needed in a simpler version, exemplified by the parameter choices (22), that provides only baryogenesis but no significant dark decay channel for the neutron. In this version, new physics signals could arise from kinetic mixing of the photon with the dark photon. In both cases it is possible that the heavy scalar triplets required by the UV complete version of the model may be accessible at the LHC.

Let us stress in closing that the main phenomenon, of asymmetry-sharing through $n$-$\chi$ oscillations, depends only upon the existence of $n$-$\chi$ mixing at low energies. It is remarkable that purely standard model physics, giving the neutron a negative self-energy correction at finite temperature, makes it possible for the oscillations to be resonant at temperatures not far above BBN. To make this part of a coherent picture including dark matter generation in the early universe requires lifting the effective theory to its UV completion. It could be interesting to find other examples of microscopic models leading to these phenomena.

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Appendix A: Thermal self-energy of $\chi$

The thermal contributions to the $\chi$ self-energy are shown in Fig. 6, in which the propagators of the intermediate light bosons are replaced by their thermal corrections (see for example ref. [25]),

$$\frac{1}{p^2 - m^2} \rightarrow \frac{2\pi\delta(p^2 - m^2)}{e^{E/T} - 1}.$$  \hspace{1cm} (A1)

Figure 6. Thermal self-energy corrections to $\chi$, where the intermediate boson has a finite-temperature propagator.

Evaluating the contribution from the hidden photon $\gamma'$ gives

$$\Sigma_\chi = g^2 \int \frac{d^4p}{(2\pi)^3} \gamma'^\mu \delta(p^2 - m'^2) \frac{-\eta_{\mu\nu} + p_\mu p_\nu/m'^2}{e^{p_\nu/T} - 1}.$$

(A2)

which simplifies in the limit $\tilde{k} \rightarrow 0$ of nonrelativistic $\chi$, and neglecting the dark photon mass, $m_\gamma' \ll T_{\gamma'}$. We get

$$\Delta E_\chi = \Sigma_\chi \approx g^2 T_{\gamma'}^2 \frac{m_\gamma'^2}{8m_\chi}.$$ \hspace{1cm} (A3)

in this limit.

A similar contribution to $\Delta E_\chi \sim \lambda'^2/8m_\psi$ arises when introducing the additional fermion $\psi$, with a coupling to $\chi$ given in Eq. (46). For large values of $g'$, this contribution is parametrically suppressed by a factor of $m_\chi/m_\psi$ with respect to Eq. (A3). If $g' \ll \lambda_3$, on the other hand, this can in fact be the dominant contribution to the thermal self-energy of $\chi$ — but would still be so small that it does not affect any of our numerical results (in Fig. 4, in particular, we would still remain on the curve labelled ‘$\xi g' = 0$’). In our analysis, we therefore neglect the contribution from $\psi$ to $\Delta E_\chi$.

Appendix B: Decoupling of dark sector

We estimate the freezeout temperature for the dimension-7 interaction (26) from the scattering process $\phi \chi \rightarrow udd$. At $T \gg m_\chi$, we can approximate all particles as being massless, with momenta $p_1 = p_\chi, p_2 = p_{\phi}$, and $p_{3,4,5}$ for the final-state quarks. The matrix element is

$$\langle |M|^2 \rangle = \frac{6}{\Lambda^6} p_1 \cdot p_3 p_4 \cdot p_5.$$

where $1/\Lambda^3 \equiv 1.7 \times 10^{-10}$ GeV$^{-3}$ is the coefficient of the operator (26). We can evaluate the cross section in the center of mass frame, choosing the 3-momenta to be

$$\vec{p}_1 = E_1(s_\alpha c_\phi, s_\alpha s_\phi, c_\alpha)$$  \hspace{1cm} (B1)

$$\vec{p}_2 = -\vec{p}_1$$  \hspace{1cm} (B2)

$$\vec{p}_3 = E_3(0, 0, 1)$$  \hspace{1cm} (B3)

$$\vec{p}_4 = E_4(s_\gamma, 0, c_\gamma)$$  \hspace{1cm} (B4)

$$\vec{p}_5 = -(\vec{p}_3 + \vec{p}_4)$$  \hspace{1cm} (B5)

where the angle between $p_3$ and $p_4$ is fixed by energy-momentum conservation,

$$c_\gamma = 1 + \frac{s - 2\sqrt{s}(E_3 + E_4)}{2E_3E_4}.$$  \hspace{1cm} (B6)
and \( E_5 = \sqrt{s} - E_3 - E_4 \). Then \( p_1 \cdot p_3 p_4 \cdot p_5 = s(1 - c_\alpha)E_3(\sqrt{s} - 2E_3)/4 \), and the cross section is given by

\[
\sigma = \frac{1}{16(2\pi)^4} s \int dE_3 dE_4 \int_0^{\sqrt{s}/2} dE_5 \sqrt{\frac{s}{E_5}} dE_4 ([M]^2) = \frac{s^2}{1024 \pi^3 \Lambda^6} \tag{B7}
\]

The thermal average is

\[
\langle \sigma v \rangle = \frac{1}{32 T^5} \int_0^\infty ds \sigma s^{3/2} K_1(\sqrt{s}/T) = \frac{9 T^4}{256 \pi^3 \Lambda^6} \tag{B8}
\]

Multiplying by the thermal density of a complex scalar and equating the scattering rate to the Hubble rate yields the freezeout temperature \( T_f = 11 \text{ GeV} \).

**Appendix C: Relations between chemical potentials and initial asymmetries**

In this appendix we present a calculation of the chemical potentials of both standard model and dark sector species when there is no baryon number violating interaction with \( \psi' \) which sets \( \mu_\chi = \mu_\phi \). After the electroweak phase transition, the chemical potentials of the particle species in equilibrium are given by the equations in Sec. [IV.1] with an additional relation derived from the fact the difference between total baryon number and lepton number is conserved when \( \psi' \) is not present:

\[ 2\mu_\chi + B - L = 2\mu_0. \tag{C1} \]

Here \( B = 6(\mu_\mu + \mu_d) \) and \( L = \nu_\mu + 2\mu_\ell \) where \( \mu_\ell \) is the sum of the charged lepton chemical potentials. \( \mu_0 \) is the initial \( \chi \) chemical potential.

From these equations, the chemical potentials of the various species can be determined in terms of the initial asymmetries:

\[
B = \frac{18}{61}(2\mu_0 + Q'_\nu \mu_\nu) \tag{C2a}
\]

\[
L = -\frac{25}{12}B - 6\mu_W = \frac{2}{3}B \tag{C2b}
\]

\[
2\mu_\chi = \frac{1}{122}(-2\mu_0 - 111Q'_\nu \mu_\nu) \tag{C2c}
\]

\[
2\mu_\phi = \frac{1}{122}(-22\mu_0 - 133Q'_\nu \mu_\nu) \tag{C2d}
\]

The coefficients of the chemical potentials in the above expressions are chosen so that their ratios correspond to the ratios of the asymmetries, e.g. \( 2\mu_\chi/B = (n_\chi - n_\bar{\chi})/(n_B - n_\bar{B}) \). If \( Q'_\nu \mu_\nu = -2\mu_0 \), then \( B = L = \mu_W = 0 \), \( 2\mu_\chi = 2\mu_\phi = 2\mu_0 \), and the analysis of Sec. [III] directly applies. However, other values for the initial asymmetries can lead to a baryon asymmetry when these interactions freeze out that is greater than what is observed today. Such scenarios are not viable – the subsequent oscillations will only drive the ratio of baryon to \( \chi \) asymmetries to 1, rather than reducing the baryon asymmetry to the observed value.

**Appendix D: Limit on \( \nu' \) mass**

Generically, one needs to introduce dark sector particles \( \nu' \) (though not necessarily massless fermions) with a non-vanishing chemical potential, which are stable because they carry a charge \( Q' \neq 1 \). This extra field is needed to ensure that the universe is neutral in \( U(1)' \) charge at early times. These new states \( \nu' \) should not be too heavy in order not to overclose the universe. For the case of benchmark scenario B2, we can estimate this constraint by noting that the annihilation into \( \gamma' \gamma' \) keeps both \( \chi \) and \( \nu' \) in equilibrium, respectively, until they are non-relativistic. This allows to relate \( \mu_\nu \) to \( n_\nu - n_\bar{\nu} \) in the non-relativistic limit until freeze-out at \( T_\nu, t_0 \) after which the comoving number density stays constant. The same applies, correspondingly, to the \( \chi \) particles. For chemical potentials of roughly the same size, \( |\mu_\nu/\mu_\chi| \sim O(1) \), we then find that \( m_{\nu'} \sim 100 \text{ MeV} \) is already sufficient to suppress \( \rho_{\nu'}/\rho_\chi \ll 1 \) at late times. While the existence of the additional states \( \nu' \) is necessary for the cogeneration scenario to work, they can thus be heavy enough to leave the phenomenology of B2 unaffected.

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