Nuclear Fusion Inside Dark Matter

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A new dynamic is identified between dark matter and nuclei. Nuclei accelerated to MeV energies by the internal potential of composite dark matter can undergo nuclear fusion. This effect arises in simple models of composite dark matter made of heavy fermions bound by a light scalar field. Cosmologies and detection prospects are explored for composites that catalyze nuclear reactions in underground detectors and stars, including bremsstrahlung radiation from nuclei scattering against electrons in hot plasma formed in the composite interior. If discovered and collected, this kind of composite dark matter could in principle serve as a ready-made, compact nuclear fusion generator.

INTRODUCTION

The presence of dark matter has become manifest through galactic dynamics, the lensing of light, and temperature fluctuations in the cosmic microwave background. But setting aside these gravitational signifiers, little is known about dark matter despite extensive laboratory and astrophysical efforts. It is a high priority of modern science to uncover dark matter, identify its mass and couplings, and determine what influence it may have on other particles that compose the known universe.

In the past decade theorists have enunciated how a certain variety of dark matter could bear a striking resemblance to known matter. Atoms, nuclei, and nucleons, which comprise the bulk of known particles, are all built from fundamental fermions – electrons, protons, quarks – bound together by photons and gluons into composite states. Similarly, dark matter could also be comprised of many particles bound together in a composite state [1–14]. One simple composite dark matter model consists of fermions ($X$) bound by a new attractive force provided by a massive scalar field ($\phi$) [15–21]. If this force is strong enough, then in the early universe large dark matter states would be built from successive fusion of $X$ particles into increasingly massive states, in a process similar to big bang nucleosynthesis (BBN). In the absence of the repulsive Coulomb force between protons in Standard Model nuclei, these dark composites can become extremely massive after accumulating oodles of $X$ particles. As we will see in this work, if $X$ has a TeV-EeV mass, this can imply dark matter composite masses ranging from a few micrograms to thousands of tons.

We have found that such large dark matter (DM) composites imply novel dynamical interactions with Standard Model nuclei. In this paper we present these newly identified dynamics. Large composite dark matter can cause Standard Model (SM) nuclei to accelerate, radiate, and fuse in the composite interior, as shown in Fig. 1. These dynamics occur because the scalar field binding $X$ particles together can have an extremely high potential $\langle \phi \rangle$ inside the dark matter composite. Under the influence of this potential, SM nuclei are accelerated to energies $\Delta E \sim g_n \langle \phi \rangle \sim \text{MeV}$, sufficient to initiate nuclear fusion and radiation from high energy collisions, even for a miniscule Yukawa coupling $g_n$ between $\phi$ and nucleons. This implies new signatures and even potential uses for composite dark matter, including nuclear fusion and bremsstrahlung radiation as unique signatures in particle detectors, nuclear reactions in stars and planets, and speculatively the use of composites as compact fusion reactors. In addition, we find that the largest fusion-capable composites would make white dwarfs explode.

HEAVY COMPOSITE COSMOLOGY

We begin with the cosmology of very large composites, formed from heavy asymmetric fermions with masses ranging up to an EeV. As we will see, dark composites formed from such heavy fermions can have large internal potentials that accelerate Standard Model nuclei to fusion temperatures. The cosmology of up to $10^{10}$ GeV mass asymmetric dark matter, motivated by high scale baryogenesis mechanisms like Affleck-Dine [22, 23], has
been detailed in [24]. In asymmetric dark matter models [25, 26], an initial dark sector particle asymmetry sets the dark matter relic abundance, and typically dark matter freeze-out annihilation eliminates most of the symmetric dark matter abundance, i.e. \( X + \bar{X} \rightarrow \text{SM} \), leaving behind a residual asymmetric abundance of \( X \) particles. For a heavy asymmetric dark matter scenario [24], the abundance of dark fermions is subsequently depleted (along with the baryon abundance) by e.g. the decay of a field some time after freeze-out. The amount the asymmetric dark matter (and baryon) abundance is depleted by this decay is given by \( \Omega_{DM}^{dep} = \Omega_{DM} \zeta \), where \( \zeta = s_{before}/s_{after} \) is the ratio of entropy density in the universe before and after the field decays [24]. In the models that follow, we will assume that dark fermions freeze-out to an initial abundance that is later diluted through the decay of a metastable field [23, 27–31], or a phase transition/second phase of inflation [32–37]. This means that right after its freeze-out, dark matter’s abundance will be larger by a factor of \( \zeta^{-1} \), relative to a cosmology without subsequent depletion. We will see that this relative overabundance of \( X \) after freeze-out leads to the formation of rather large DM composites.

Asymmetric DM composites made of sub-TeV mass fermions have been studied at length in [15–21]. Here we consider heavier fermions. The Lagrangian,

\[
\mathcal{L} = \frac{1}{2} (\partial \varphi)^2 + \bar{X} (i \gamma^\mu \partial_\mu - m_X) X + g_X \bar{X} \varphi X - \frac{1}{2} m_\varphi^2 \varphi^2 + g_n \bar{n} n + \mathcal{L}_{SM},
\]

includes the scalar \( \varphi \) which provides an attractive force that binds together \( X \) fermions. The second to last term couples \( \varphi \) to SM nucleons \( n \), where this is the simplest renormalizable coupling to SM particles. Once enough \( X \) particles are bound together, fermionic composites will reach a saturation point after the composite radius exceeds \( R_X \gtrsim m_\varphi^{-1} \), at which point the interior density becomes approximately constant, \( \rho_c = \bar{m}_X^2/3 \pi^2 \), where \( \bar{m}_X \) is the constituent mass, i.e. the effective mass of \( X \) inside the composite [19, 20]. For the cosmological formation we consider hereafter, composites are saturated well before they finish forming, cf. Eq. (2). The constituent mass for a saturated composite is given by \( \bar{m}_X \approx m_X - E_X \), where \( E_X \) is the binding energy per \( X \). In terms of bare masses and couplings, in a saturated composite \( \bar{m}_X \approx [3 \pi m_X^2 m_\varphi^2/(2 \alpha_X)]^{1/4} \), where the \( \varphi - X \) coupling constant is \( \alpha_X \equiv g_\varphi^2/4\pi \). For parameters we consider, the binding energy is close to the unbound \( X \) mass, \( E_X \approx m_X \), and so \( \bar{m}_X \ll m_X \). This means the composite state of \( X \) particles has a total mass \( M_X \equiv N \bar{m}_X \), which is much less than the mass of unbound \( X \) particles, \( Nm_X \). As a consequence, after the composite is assembled, the mass density of dark matter in the universe decreases by a factor \( \bar{m}_X/m_X \), where the mass loss is accounted for by the emission of \( \varphi \) radiation.

Fermion composites will begin to assemble in the early universe by forming two-fermion bound states, where binding will occur so long as \( \alpha_X^2 m_X \gtrsim m_\varphi \) and \( \alpha_X \gtrsim 0.3 \left( m_X/10^7 \text{GeV} \right)^{2/5} \left( \zeta/10^{-6} \right)^{1/5} \) [15]. After two-fermion states form, composites will build up through processes like \( X_N + X_N \rightarrow X_{2N} + \varphi \), where \( X_N \) is a bound state formed from \( N \) fermions. At the temperature of composite assembly \( T_{ca} \), an estimate for the number of constituent particles in a typical composite can be obtained by comparing the \( X_N \) interaction rate to the Hubble rate [17, 20], \( n_{X_N} \sigma_{X_N} v_{X_N}/H \sim 1 \). Re-expressing this in terms of the \( X \) number density \( n_X = n_{X_N}/N \), the \( X \) composite cross-section \( \sigma_X = \sigma_{X_N}/N^{2/3} \) (where \( R_X \) scales as \( N^{1/3} \) in the saturation regime), and the \( X \) velocity \( v_{X_N} = v_X/N^{1/2} \), we arrive at an expression for the number of \( X \) particles in a typical composite,

\[
N_c = \left( \frac{2n_X \sigma_{X} v_X}{3H} \right)^{6/5} = \left( \frac{20 \sqrt{g_{ca} T_{ca}^3 c_a} M_{pl}^{3/2} \zeta}{m_X^{7/2}} \right)^{6/5} \left( \frac{5 \text{GeV}}{m_X} \right)^{21/5} \left( \frac{10^{-6}}{\zeta} \right)^{6/5},
\]

where in the first equality we have included a factor of \( 2/3 \) appropriate for composite assembly in a radiation-dominated universe [20], in the second equality we have used a composite cross-section \( \sigma_X = 4 \pi R_c^2 \) with \( R_c \equiv (3 \bar{m}_X/4 \pi \rho_c)^{1/3} = (9 \pi/4)^{1/3}/\bar{m}_X \), a velocity \( v_X = \sqrt{T/\bar{m}_X} \), the Friedmann relation is \( 3H^2 = g^* \pi^2 T^4/30 \), and \( T_r \approx 0.8 \text{ eV} \) is the temperature at matter-radiation equality. For \( m_X \gg \bar{m}_X \), the binding energy of these composites is \( E_X \approx m_X \), and composite assembly will finish around the temperature of \( X \) freeze-out, \( T_{ca} \sim m_X/10 \).

A few more facets of heavy composite cosmology are worth emphasizing. First, because the constituent mass \( \bar{m}_X \ll m_X \) determines the final density of DM, heavy asymmetric DM composites can account for the baryon-DM density coincidence: the present-day DM density approximately matches the baryon density, \( \Omega_{DM}/\Omega_B \approx 5 \). For asymmetric DM, this coincidence can be explained by having a single particle asymmetry that determines both the baryon and DM relic abundances. While a naive prediction for the constituent DM mass relative to the baryon mass is then \( \bar{m}_X \sim 5 m_b \sim 5 \text{ GeV} \), in the case that heavy asymmetric DM freezes out before the electroweak phase transition, electroweak sphalerons can dilute baryon number, leading to a looser prediction \( \bar{m}_X \sim 1 - 1000 \text{ GeV} \) [25]. Second, as previously discussed, in an asymmetric DM cosmology the symmetric DM component \( (XX) \) is depleted via annihilation. In the case of the heavy DM detailed above, \( XX \) annihilation to \( \varphi \) will deplete \( XX \) to a sub-DM
relic density, if the annihilation cross-section $\sigma_\text{v} \approx 3\alpha^2_{\text{e}} \langle \frac{m_X}{8m_X^2} \rangle \gtrsim 10^{-36}\text{cm}^2 / \xi [24]$, which corresponds to $\alpha_X \gtrsim 0.3 \left( \frac{m_X}{10^7 \text{GeV}} \right) \left( \frac{\xi}{10^{-6}} \right)^{1/2}$, although this restriction weakens if $X\bar{X}$ are depleted by additional annihilation channels or other mechanisms.

**NUCLEAR ACCELERATION, RADIATION, AND FUSION IN COMPOSITE DM**

Substantial energy can be released by nuclei accelerated inside large DM composites, both through fusion and bremsstrahlung processes. With the structure and cosmology of heavy composites previously laid out, we now turn to nuclear acceleration, radiation, and fusion inside large composites. We begin with the potential inside a saturated composite $\langle \phi \rangle \approx \frac{m_X}{g_X}$, obtained by requiring the composite’s internal potential (1) is minimized at equilibrium. Boundary conditions require that outside the composite the potential decays as

$$\phi(r) = \langle \phi \rangle e^{-m_\phi(r-R_X)} \left( \frac{R_X}{r} \right). \tag{3}$$

**Acceleration.** Nuclei with $A$ nucleons will have their momentum $p$ boosted to $p'$ as they enter the composite, according to $p'^2 + m_X^2 = p^2 + (m_X - V_n)^2$, where $V_n = A g_{\phi}/\langle \phi \rangle = A g_n m_X / g_X$. In the limit $V_n \ll m_N$, the second term can be expanded yielding $p'^2 - p^2 = 2m_N V_n$. Nuclei will accelerate over a time determined by the field gradient at the composite boundary and the velocity $v_X$ at which the composite moves, e.g., Eq. (3), $\tau_{\text{accel}} \approx (m_\phi v_X)^{-1} (1 + 2V_n/m_N v_X^2)^{-1/2}$.

**Ionization.** For parameters in Fig. 2 saturated composites crossing terrestrial material at speeds $v_X \approx 10^{-3}$ will accelerate nuclei on a timescale $\tau_{\text{ion}} \lesssim 10^{-18} \text{s}$ due to the sharp gradient of the potential. This timescale is shorter than both the electron orbital period ($10 \text{ eV}$) $\approx 10^{-17} \text{s}$ and $a_0/v_N \approx 10^{-17} \text{s}$ ($v_N/10^{-2} \text{eV}$) where $v_N$ is the nucleus final speed and $a_0$ is the Bohr radius. Such a perturbation is then non-adiabatic, i.e., electrons do not respond to the sudden nuclear motion in a similar timescale, resulting in excitation or ionization [40, 41], the so-called Migdal effect which has been recently considered to extend the sensitivity of direct detection experiments [42–48]. In particular, numerical results from [42] indicate that the probability of outer-shell electron ionization for C and O atoms, the most abundant elements in IceCube and SNO+, is of order $f_{\text{e}} \approx 10^{-2} - 10^{-1}$ for the nuclear kinetic energies considered here, with the probability peak located at ionized electron energies $\approx 1 - 10 \text{ eV}$. Hence after this impulsive motion, a sizeable fraction of atoms are partially ionized. However, further considerations indicate the atoms will be fully ionized. The atoms accelerated to relative energies $100 \text{ eV} - 1 \text{ MeV}$, will scatter with the free electrons, resulting in further ionization. The cross-section for ionizing atomic oxygen or carbon is $\sigma_i \approx 10^{-16} - 10^{-17} \text{cm}^2$ in the energy range of interest [49]. Ionization by electron-atom collisions will occur on a timescale given by $(f_{\text{e}} n_{\text{e}} v_N \sigma_i)^{-1} \lesssim 10^{-15} \text{s}$, where $n_{\text{e}}$ is the atom velocity and $v_N$ is the electron number density. This timescale is shorter than the composite crossing time ($2R_X / v_X$) $\approx 10^{-15} \text{s}$ ($R_X / \text{nm}$)($v_X / 10^{-3} \text{s}$), so long as composites are larger than a nm, and becomes even shorter as more electrons are ionized and $f_{\text{e}} \approx 1$. Hence atoms are fully ionized in the detection regions shown in Fig. 2. These estimates agree with [50], which finds order one ionization fractions for carbon and oxygen plasmas at $T \gtrsim 100 \text{ eV}$ and density $\approx 1 \text{ g cm}^{-3}$.

**Radiation.** The ion-electron plasma will have a photon opacity dominated by free-electron scattering, with a photon mean free path $(\alpha_T \sigma_T)^{-1} \approx 5 \text{ cm} \gg R_X$, where $\sigma_T \approx 10^{-24} \text{ cm}^2$ is the Thompson cross-section for electrons. Therefore, radiation never equilibrates with the plasma and we do not expect blackbody radiation. Instead, we expect thermal electron-ion bremsstrahlung, which has specific emissivity at frequency $\omega$ (see e.g. [51]) $j_\omega = \frac{16}{3} \left( \frac{m_e E_\text{brem}}{\pi} \right)^{1/2} \exp(-\omega/T)$, for electron mass $m_e$ and fine structure constant $\alpha = e^2/4\pi$. Since in Fig. 2 we require $T \gtrsim 100 \text{ eV}$, there is emission of ionizing radiation. The integrated emissivity over volume and frequency yields a radiated energy rate

$$\dot{E}_\text{brem} = \frac{64\pi^2 e^6}{9\sqrt{3m_e^3}} \left( \frac{2m_\text{e} T}{\pi} \right)^{1/2} n_e^2 R_X^3 \tag{4}$$

$$\approx 10^{10} \text{ GeV s}^{-1} \left( \frac{g_X}{9} \right)^{-2} \left( \frac{9n_\text{e}}{10^{-10}} \right)^{1/2} \left( \frac{m_X}{\text{TeV}} \right)^{1/2} \left( \frac{R_X}{\text{nm}} \right)^3.$$

At temperatures $T \approx 100 \text{ keV} - 1 \text{ MeV}$, we also expect a fraction of the ions to undergo thermonuclear fusion. In particular, we consider here the thermonuclear $^{16}\text{O}$ burning rateg tabulated in [52], since this is the most abundant isotope in the terrestrial crust and mantle [53–58]. We remark that this radiation rate dominates over ionization energy losses.

**Detection.** The large composites we have uncovered cannot be found by traditional dark matter experiments, which are flux limited to $M_X \lesssim 10^{19} \text{ GeV}$ [9, 59]. However, the copious energy released by fusion-capable composites make them observable at larger neutrino experiments like IceCube, Super-K, and large volume scintillators (LVS) like SNO+, Borexino, and JUNO; their enmity extends the $M_X$ mass reach to $3 \times 10^{25} \text{ GeV}$ in the case of IceCube (assuming 5 yrs and a km$^2$ detection area). To conservatively establish the sensitivity of IceCube and LVS to a flood of $\gtrsim \text{ eV}$ photons emitted from transiting composites, in Fig. 2 we require trigger threshold energy depositions of $\approx 10 \text{ TeV}$ and $1 \text{ MeV}$ per 100 ns respectively, which are an order of magnitude above the TeV [60] and 100 keV [61, 62] per 100 ns design thresholds of these experiments (this still underestimate
IceCube’s sensitivity, since our requirement implies \( \gtrsim 100 \) PeV radiated in a transit through IceCube). Comparing this to Eq. (4), we find that nucleon couplings as small as \( g_n \sim 10^{-14} \) at IceCube and \( g_n \sim 10^{-12} \) at LVS can be detected in the upper left portions of IceCube and LVS detection regions marked in Fig. 2. Smaller coupling values will result in too little radiation rate for composite detection. We also show where \(^{16}\text{O}\) fusion reactions occur at IceCube as a single composite crosses it, using the \(^{16}\text{O}\) burning rate in [52]. In this case, there will be additional gamma rays and byproducts with \(~\text{MeV} \) energies, e.g. \(^{32,33}\text{S}, ^{32}\text{P}, ^{28}\text{Si}, ^{24}\text{Mg}\) as well as p, n and \( \alpha \)'s [52].

**Capture on Earth and lack of X-nuclear scattering.**

Thus far we have not mentioned nuclear scattering against X fermions in composites. Compared to Eq. (4), composite energy loss from X-nuclear scattering will be negligible, in part because the fermi momentum of X is large, \( p_{fX} \sim \bar{m}_X \). Accounting for nuclear scattering with degenerate fermions [63–65], the scattering energy loss is \( \bar{E}_{X-N} \approx A^2 g_{Xf}^2 \bar{g}_X m_N \bar{m}_X^4 (m_N + 2\bar{m}_X)^4 \), which is tiny compared to bremsstrahlung in Fig. 2. On the other hand, energy loss in the form of radia-
tion, cf. Eq. (4), could result in stopping of composites before they reach detectors. This is relevant for lower mass composites with less initial kinetic energy. Using Eq. (4), a composite with an initial velocity $v_X$ will travel through the Earth’s mantle a distance $L_{cap} \simeq 2 \text{ km} \left( m_X/\text{TeV} \right)^{1/2} \left( 10^{-10}/g_n \right)^{1/2} \left( 1/\rho_1 \right)^{1/2} \times \left( v_X/200 \text{ km s}^{-1} \right)^3 \left( m_e/10 \text{ keV} \right)^2$ before being slowed below Earth’s escape velocity, where we have computed this distance considering the most abundant isotope $^{16}\text{O}$ and using element/density profiles from [66, 67] (see also [53–58, 68]). These scalings agree with the simple capture estimate $L_{cap} \sim \Delta E_{cap} v_X / \dot{E}_{\text{brem}}$, where $\Delta E_{cap}$ is the DM’s initial kinetic energy. Earth’s capture rate can be found using the method described in [67]. The captured composites may induce neutral reactions in the crust and mantle, resulting in a potential planetary heat signal relevant for future searches [65–67].

**WD explosions.** The transit of a large composite through a white dwarf (WD) can catalyze nuclear fusion reactions leading to a thermonuclear runaway and Type-Ia supernova explosion, similar to [69–74], although in this case fusion is initiated by nuclei accelerated inside the composite. As established in these references, WDs will ignite when certain ignition conditions are met as detailed in [75], where a set of critical temperatures and trigger masses are numerically computed for different white dwarf compositions and central densities. We conservatively require a critical temperature $T_{\text{crit}} \simeq 1 \text{ MeV}$ for a pure $^{12}\text{C}$ white dwarf. As they pass through a white dwarf, composites can lose kinetic energy to heat dissipation in the form of radiation, raising the possibility that they may be stopped before reaching the WD core. However, composites bounded by WDs in Fig. 2 are so massive that a negligible fraction of their kinetic energy is lost to this dissipative effect. Heat conduction out of the composite is dominated by relativistic white dwarf electrons, with a rate $Q_{\text{cond}} \simeq 4\pi^2 T_X^2 R_X/15 \kappa_e \rho_e \simeq 10^{27} \text{ GeV s}^{-1} \left( \rho_e/10^9 \text{ g cm}^{-3} \right)^{4/15} \left( R_X/\mu \text{m} \right)$, where $\kappa_e \simeq 10^{-3} \text{ cm}^2 \text{ g}^{-1} \left( T_e/10^7 \text{ K} \right)^{2.8} \left( 10^9 \text{ g cm}^{-3} / \rho_e \right)^{1.6}$ is the conductive opacity of the relativistic white dwarf electrons [76]. Composite radiation, on the other hand, is $\dot{Q}_{\text{rad}} = 4\pi R_X^2 \nabla \left( \langle \sigma T^3 \rangle / \kappa_e \rho_e \right) \simeq 16\pi R_X^2 \sigma T^4 m_e/\kappa_e \rho_e \simeq 10^{24} \text{ GeV s}^{-1} \left( R_X/\mu \text{m} \right) \left( m_e/\text{keV} \right)$, where $\kappa_e \simeq 10^7 \text{ cm}^2 \text{ g}^{-1} \left( T_e/10^7 \text{ K} \right)^{-2/7} \left( \rho_e/10^9 \text{ g cm}^{-3} \right)$ is the white dwarf radiative opacity dominated by free-electron transitions [74, 77]. We have assumed a black-body energy density since the stellar material is highly opaque to photons.

The rate of carbon fusion in dense WD matter is $R_{\text{fu}} \simeq 10^{12} \text{ cm}^{-3} \text{ s}^{-1} \left( \rho_e/10^9 \text{ g cm}^{-3} \right)^2$ at $T_{\text{crit}} \simeq 1 \text{ MeV}$, with an average energy release rate $Q \simeq +3 \text{ MeV}$ per reaction [78]. This yields a nuclear energy release rate $\dot{Q}_{\text{fu}} \simeq 4\pi Q R_{\text{fu}} R_X^3/3 \simeq 10^{28} \text{ GeV s}^{-1} \left( R_X/\mu \text{m} \right)^3$. Therefore, for composites with radii $R_X \gtrsim \mu \text{m}$, the heat release from nuclear fusion greatly exceeds conductive and radiative losses, setting the conditions for a sustained thermonuclear runaway. We remark that stellar masses contained within radii $\gtrsim \mu \text{m}$ are $\gtrsim 10^{-3} \text{ g}$, which are in agreement with the minimum trigger masses outlined in [75]. Fig. 2 shows the region where $V_n \sim \text{ MeV}$ composites ignite a WD by simply passing through. Since one encounter would occur for composite masses $M_X \lesssim 10^{42} \text{ GeV}$ in a $\sim \text{ Gyr}$ timescale, the survival of e.g. WD J160420.40 [72] implies constraints on nucleon couplings $g_n \lesssim 10^{-15} \left( m_X/10^8 \text{ GeV} \right)^{-1}$ in that region.

**Big Bang Nucleosynthesis.** It is natural to wonder whether BBN may constrain fusion-capable composites through over-production or disintegration of isotopes. An extensive analysis of fusion-capable DM composites on primordial abundances using relevant reaction rates (e.g. [79–81]) will be the subject of future work. Here we remark that in the IceCube and LVS detection regions shown in Fig. 2, early universe composites seem unlikely to alter standard BBN abundance predictions. Constraints on $g_n$ imply that even for the maximum coupling allowed, composites will not change the temperature of the primordial plasma until redshift $z_X \lesssim 10^8 (A/1)(g_n/10^{-10})(m_X/\text{TeV})$. However, by this redshift the baryon density will be significantly diluted according to $\Omega_b \rho_b (1 + z_X)^3$. The average number of baryons inside composites will then be $4\pi \Omega_b \rho_b (1 + z_X)^3 R_X^3 / 3 m_b \simeq 10^{-11} \left( m_X/\text{TeV} \right) \left( g_n/10^{-10} \right) \left( R_X / \text{nm} \right)^3$ where $m_b$ is the baryon mass. Comparing this to Fig. 2, parameter space where large neutrino experiments have sensitivity, corresponds to composite sizes too small to have more than one baryon per composite, by the time a baryon inside a composite would be substantially accelerated in the early universe. Similar estimates using Eq. (4) indicate that detectable fusion-capable composites do not observably alter the baryon-to-photon ratio after BBN, nor the ionization fraction after recombination.

**CONCLUSIONS**

We have studied the cosmology and detection of heavy composite DM that internally accelerates nuclei, resulting in copious collisional radiation and nuclear fusion. Prospects have been explored for detection of fusion-capable composites at IceCube and liquid scintillator experiments. There are many aspects of Standard Model particle acceleration in DM composites that remain. While here we considered composites that accelerate nuclei to MeV energies, if these were increased to relativistic energies, this would cause repulsive composite-SM scattering processes [82]. For smaller than $100 \text{ eV}$ acceleration energies, the Migdal effect and SM-SM collisional ionization should permit dark matter experiments to search for rather weakly-coupled composites. For liquid noble element experiments such as Xenon-1T, LUX, LZ or DEAP-3600, this will require a dedicated analysis...
of the scintillation signals produced and detection efficiencies [43, 83–85]. Given that asymmetric composites are often associated with SM asymmetries, similar acceleration effects should be explored for composites coupled to the SM through vector fields, and especially fields that couple to leptons, baryons, or a combination of these. Finally, it would be interesting to study whether fusion-capable composites could detectably alter isotopic abundances in the Earth over geological time periods. We leave these and other inquests into accelerative dark matter to future work.

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