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New Target for Cosmic Axion Searches

Daniel Baumann,[1,2,*] Daniel Green,[3,4,†] and Benjamin Wallisch[1,‡]
[1]DAMTP, University of Cambridge, Cambridge, CB3 0WA, United Kingdom
[2]IoP, University of Amsterdam, Amsterdam, 1090 GL, Netherlands
[3]Department of Physics, University of California, Berkeley, CA 94720, USA
[4]Canadian Institute for Theoretical Astrophysics, Toronto, ON M5S 3H8, Canada
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Future cosmic microwave background experiments have the potential to probe the density of relativistic species at the subpercent level. This sensitivity allows light thermal relics to be detected up to arbitrarily high decoupling temperatures. Conversely, the absence of a detection would require extra light species never to have been in equilibrium with the Standard Model. In this Letter, we exploit this feature to demonstrate the sensitivity of future cosmological observations to the couplings of axions to photons, gluons, and charged fermions. In many cases, the constraints achievable from cosmology will surpass existing bounds from laboratory experiments and astrophysical observations by orders of magnitude.

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Introduction.—Most of what we know about the history of the Universe comes from the observations of light emitted at or after recombination. To learn about earlier times we rely either on theoretical extrapolations or the observations of relics that are left over from an earlier period. One of the most remarkable results of the Planck satellite is the detection of free-streaming cosmic neutrinos [1–3], with an energy density that is consistent with the predicted freeze-out abundance created one second after the Big Bang. Probing even earlier times requires detecting new particles that are more weakly coupled than neutrinos. Such particles arise naturally in many extensions of the Standard Model (SM) [4,5]. Particularly well-motivated are Goldstone bosons created by the spontaneous breaking of additional global symmetries.

Goldstone bosons are either massless (if the broken symmetry was exact) or naturally light (if it was approximate). Examples of light pseudo-Nambu-Goldstone bosons (pNGBs) are *axions* [6–8], *familons* [9–11], and *majorons* [12,13], associated with spontaneously broken Peccei-Quinn, family, and lepton-number symmetry, respectively. Below the scale of the spontaneous symmetry breaking, the couplings of the Goldstone bosons ϕ to the SM degrees of freedom can be characterized through a set of effective interactions

$$\frac{\mathcal{O}_\phi \mathcal{O}_{\rm SM}}{\Lambda^\Delta}, \tag{1}$$

where Λ is related to the symmetry breaking scale. Axion, familon, and majoron models are characterized by different couplings in Eq. (1). These couplings are constrained by laboratory experiments [5,14], by astrophysics [15,16], and by cosmology [17,18]. While laboratory constraints have the advantage of being direct measurements, their main drawback is that they are usually rather model specific and sensitive only to narrow windows of pNGB masses. Astrophysical and cosmological constraints are complimentary since they are relatively insensitive to the detailed form of the couplings to the SM and span a wide range of masses. The main astrophysical constraints on new light particles come from stellar cooling [15]. In order not to disrupt successful models of stellar evolution, any new light particles must be more weakly coupled than neutrinos. Moreover, since neutrinos couple to the rest of the SM through a dimension-six operator (suppressed by the electroweak scale), the constraints on extra particles are particularly severe for dimension-four and dimension-five couplings to the SM.

In this Letter, we will show that cosmology is remarkably sensitive to extra light particles. This is because interactions like Eq. (1) can bring these particles into equilibrium with the SM particles. Moreover, thermal equilibrium is democratic. Any new light field that was in thermal equilibrium in the past will have a number density that is comparable to that of photons. This is why neutrinos have been detected with high significance in the CMB [1–3] despite their weak coupling. Like astrophysical constraints, cosmology therefore requires any new light particles to be more weakly coupled than neutrinos. Given the Moore's law-like improvements in CMB detector sensitivity [19,20], cosmology will push the sensitivity to new light particles beyond the strength of weak scale interactions and has the potential to explore a fundamentally new territory of physics beyond the SM.

Preliminaries.—The total energy density in relativistic species is often defined as

$$\rho_r = \left[1 + \frac{7}{8}\left(\frac{4}{11}\right)^{4/3} N_{\rm eff}\right]\rho_\gamma, \tag{2}$$

where ρ_γ is the energy density of photons and the parameter $N_{\rm eff}$ is called the effective number of neutrinos, although there may be contributions that have nothing to do with neutrinos (see, e.g., Ref. [21]). The SM predicts

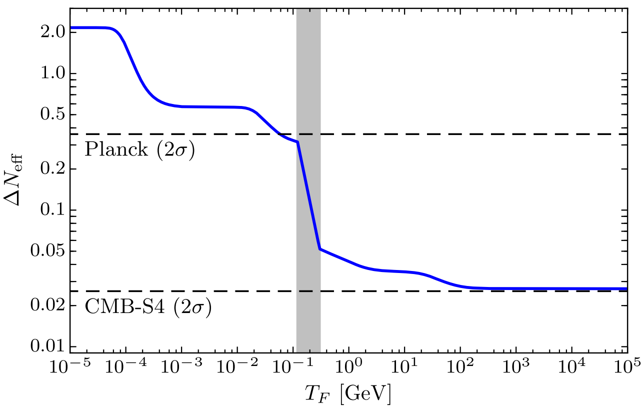

FIG. 1. Contribution of a single thermally decoupled Goldstone boson to the effective number of neutrinos, $\Delta N_{\rm eff}$, as a function of the freeze-out temperature T_F. Shown are also the current 2σ sensitivity of the Planck satellite [1] and an (optimistic) estimate of the sensitivity of a future CMB-S4 mission [3]. The gray band indicates the QCD phase transition.

$N_{\rm eff} = 3.046$ from neutrinos [22] and the current constraint from the Planck satellite is $N_{\rm eff} = 3.04 \pm 0.18$ [1]. Figure 1 shows the extra contribution to the radiation density of a thermally decoupled Goldstone boson as a function of its freeze-out temperature T_F. We see that particles that decoupled *after* the QCD phase transition are effectively ruled out by the observations of the Planck satellite [17]. On the other hand, the effect of particles that decoupled *before* the QCD phase transition is suppressed by an order of magnitude, $0.05 \geq \Delta N_{\rm eff} \geq 0.027$. Although Planck is blind to these particles, this regime is within reach of future experiments. In particular, the planned CMB Stage IV (CMB-S4) experiments have the potential to constrain (or detect) extra relativistic species at the level of $\sigma(N_{\rm eff}) \sim 0.01$ [3,19,20]. To simplify the narrative, we will assume that the sensitivity to $\Delta N_{\rm eff} = 0.027$ will be reached with CMB-S4, either on its own or in conjunction with other data [23,24]. Alternatively, our arguments could be viewed as strong motivation for reaching this critical level of sensitivity in future experiments.

Future CMB constraints.—Even the absence of a detection would be very informative, because it would strongly constrain the couplings between the extra light relics and the SM degrees of freedom. This is because a thermal abundance can be avoided if the reheating temperature of the Universe, T_R, is below the would-be freeze-out temperature, i.e., $T_R < T_F$. In that case, the extra particles have never been in thermal equilibrium and their densities therefore do not have to be detectable. In the absence of a detection, requiring $T_F(\Lambda) > T_R$ would place very strong bounds on the scale(s) in Eq. (1), i.e., $\Lambda > T_F^{-1}(T_R)$. We note that these constraints make no assumption about the nature of dark matter because the thermal population of axions arises independently of a possible cold population. On the other hand, we have to assume that the effective description of the pNGBs with interactions of the form of Eq. (1) holds up to $T_F \ll \Lambda$. This is equivalent to assuming that the UV completion of the effective theory is not too weakly coupled. Moreover, we also require the absence of any significant dilution of $\Delta N_{\rm eff}$ after freeze-out. In practice, this means that we are restricting to scenarios with $\Delta g_*(T_F) \lesssim g_*^{\rm SM}(T_F) \approx 10^2$. Finally, our results will be restricted to $m_\phi < 1$ MeV, so that the only possible decays of the pNGBs are to photons or neutrinos. In the remainder, we will derive future CMB constraints on the couplings of pNGBs to SM gauge fields (for axions) and charged fermions (for familons). Similar bounds for the couplings to neutrinos (for majorons) can be found in Refs. [25–27].

Constraints on axions.—Axions arise naturally in many areas of high-energy physics, the QCD axion being a particularly well-motivated example. They are a compelling example of a new particle that is experimentally elusive [5,14] because of its weak coupling rather than due to kinematic constraints. What typically distinguishes axions from other pNGBs are their unique couplings to the SM gauge fields. Below the scale of electroweak symmetry breaking (EWSB), we consider the following effective theory with shift-symmetric couplings of the axion:

$$\mathcal{L}_{\phi{\rm EW}} = -\frac{1}{4}\left(\frac{\phi}{\Lambda_\gamma} F_{\mu\nu}\tilde{F}^{\mu\nu} + \frac{\phi}{\Lambda_g} G^a_{\mu\nu}\tilde{G}^{\mu\nu,a}\right), \tag{3}$$

where $X_{\mu\nu} \equiv \{F_{\mu\nu}, G^a_{\mu\nu}\}$ are the field strengths for the photons and gluons, and $\tilde{X}^{\mu\nu} \equiv \frac{1}{2}\epsilon^{\mu\nu\rho\sigma}X_{\rho\sigma}$ are their duals. Axion models will typically include couplings to all SM gauge fields, but only the coupling to gluons is strictly necessary to solve the strong CP problem. At high energies, the rate of axion production through the gauge field interactions can be expressed as [28] (see also Refs. [29–32])

$$\Gamma(\Lambda_n, T) = \sum_n \gamma_n(T)\frac{T^3}{\Lambda_n^2}. \tag{4}$$

The prefactors $\gamma_n(T)$ have their origin in the running of the couplings and are only weakly dependent on temperature. For simplicity of presentation, we will treat these functions as constants, but take them into account in Ref. [33]. We see that the production rate, $\Gamma \propto T^3$, decreases faster than the expansion rate during the radiation era, $H \propto T^2$. The axions therefore *freeze out* when the production rate drops below the expansion rate, with the freeze-out temperature T_F determined by $\Gamma(T_F) = H(T_F)$. To avoid this thermal axion abundance requires $T_F > T_R$, or

$$\Gamma(\Lambda_n, T_R) < H(T_R) = \frac{\pi}{\sqrt{90}}\sqrt{g_{*,R}}\frac{T_R^2}{M_{\rm pl}}, \tag{5}$$

where $M_{\rm pl}$ is the reduced Planck mass and $g_{*,R} \equiv g_*(T_R)$ denotes the effective number of relativistic species at T_R. For a given reheating temperature, this is a constraint on the couplings Λ_n in Eq. (4). Treating the different axion couplings separately, we can write

$$\Lambda_n > \left(\frac{\pi^2}{90} g_{*,R}\right)^{-1/4} \sqrt{\gamma_{n,R} T_R M_{\rm pl}}, \tag{6}$$

where $\gamma_{n,R} \equiv \gamma_n(T_R)$.

The operator that has been most actively investigated experimentally is the coupling to photons. Photons are easily produced in large numbers in both the laboratory and in many astrophysical settings which makes this coupling a particularly fruitful target for axion searches. The electroweak couplings in the high-energy theory prior to EWSB are related to the photon coupling Λ_γ through the Weinberg mixing angle. In [33], we show in detail how the constraints (6) on the couplings to the electroweak gauge bosons map into a constraint on the coupling to photons. This constraint is a function of the relative size of the couplings to the $SU(2)_L$ and $U(1)_Y$ sectors. To be conservative, we will here present the weakest constraint which arises when the axion only couples to the $U(1)_Y$ gauge field. A specific axion model is likely to also couple to the $SU(2)_L$ sector and the constraint on Λ_γ would then be stronger (as can be seen explicitly in [33]). Using $\gamma_{\gamma,R} \approx \gamma_\gamma(10^{10}\ \text{GeV}) = 0.029$ and $g_{*,R} = 106.75 + 1$, we find

$$\Lambda_\gamma > 1.4 \times 10^{13}\ \text{GeV} \sqrt{T_{R,10}}, \tag{7}$$

where $T_{R,10} \equiv T_R/10^{10}$ GeV. For a reheating temperature of about 10^{10} GeV, the bound in (7) is three orders of magnitude stronger than the best current constraints (cf. Fig. 2). Even for a reheating temperature as low as 10^4 GeV the bound from the CMB would still marginally improve over existing constraints.

Massive axions are unstable to decay mediated by the operator $\phi F\tilde{F}$. However, for couplings compatible with the stellar cooling constraint, $\Lambda_\gamma > 1.3 \times 10^{10}$ GeV [35], and masses $m_\phi \lesssim 10$ keV, these decays occur after recombination and, hence, the axions are effectively stable [33]. The regime $10\ \text{keV} < m_\phi < 1\ \text{MeV}$ (where the axion decays between neutrino decoupling and recombination) is constrained by effects on the CMB and on big bang nucleosynthesis (BBN) [36–38].

The coupling to gluons is especially interesting for the QCD axion since it has to be present in order to solve the strong CP problem. The axion production rate associated with the gluon interaction in Eq. (3) is $\Gamma_g \simeq 0.41 T^3/\Lambda_g^2$ [28]. As before, we have dropped a weakly temperature-dependent prefactor, but account for it in Ref. [33]. The bound (6) then implies

$$\Lambda_g > 5.4 \times 10^{13}\ \text{GeV} \sqrt{T_{R,10}}. \tag{8}$$

Laboratory constraints on the axion-gluon coupling are usually phrased in terms of the induced electric dipole moment (EDM) of nucleons: $d_n = g_d \phi_0$, where ϕ_0 is the value of the local axion field. The coupling g_d is given for the QCD axion by [39,40]

$$g_d \approx \frac{2\pi}{\alpha_s} \times \frac{3.8 \times 10^{-3}\ \text{GeV}^{-1}}{\Lambda_g}. \tag{9}$$

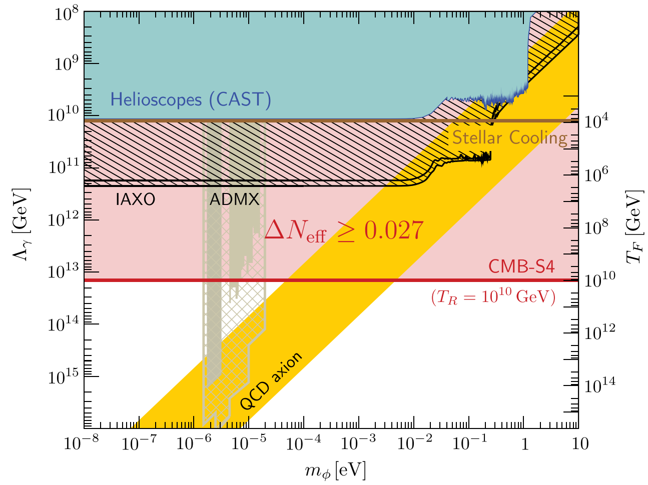

FIG. 2. Comparison between current constraints on the axion-photon coupling and the sensitivity of a future CMB-S4 mission (figure adapted from Ref. [34]). Future laboratory constraints (IAXO and ADMX) are shown as shaded regions. The yellow band indicates a range of representative models for the QCD axion. The future CMB bound is a function of the reheating temperature T_R. We note that ADMX assumes that the axion is all of the dark matter, while all other constraints do not have this restriction.

Constraints on g_d (and hence Λ_g) are shown in Fig. 3. We see that future CMB-S4 observations will improve over existing constraints on Λ_g by up to six orders of magnitude if $T_R = \mathcal{O}(10^{10}\ \text{GeV})$. Even if the reheating temperature is as low as 10^4 GeV, the future CMB constraints will be tighter by three orders of magnitude.

Constraints on familons.—Spontaneously broken global symmetries have also been envoked to explain the approximate $U(3)^5$ flavor symmetry of the Standard Model. The associated pNGBs—called *familons* [9–11]—couple to the SM through Yukawa couplings,

$$\mathcal{L}_{\phi\psi} = -\frac{\partial_\mu \phi}{\Lambda_\psi} \bar{\psi}_i \gamma^\mu (g_V^{ij} + g_A^{ij}\gamma^5)\psi_j$$
$$\to \frac{\phi}{\Lambda_\psi}\left(\mathrm{i}H\bar{\psi}_{L,i}\left[\sum_{I=V,A}(\lambda_i \mp \lambda_j)g_I^{ij}\right]\psi_{R,j} + \text{H.c.}\right), \tag{10}$$

where H is the Higgs doublet and $\psi_{L,R} \equiv \frac{1}{2}(1 \mp \gamma^5)\psi$. The $SU(2)_L$ and $SU(3)_c$ structures in Eq. (10) take the same form as for the SM Yukawa couplings [43], but this has been left implicit to avoid clutter. In the second line we have integrated by parts and used the equations of motion. The subscripts V and A denote the couplings to the vector and axial-vector currents, respectively, and $\lambda_i \equiv \sqrt{2}m_i/v$ are the Yukawa couplings, with $v = 246$ GeV being the Higgs vacuum expectation value. We note that the diagonal couplings, $i = j$, are only to the axial part, as expected from vector current conservation. In Table I, we have collected accelerator and astrophysics constraints on the effective couplings $\Lambda_{ij}^I \equiv \Lambda_\psi/g_I^{ij}$ and $\Lambda_{ij} \equiv \Lambda_\psi/[(g_V^{ij})^2 + (g_A^{ij})^2]^{1/2}$. We see that current data typically constrain the couplings to the first generation fermions much more than those to the

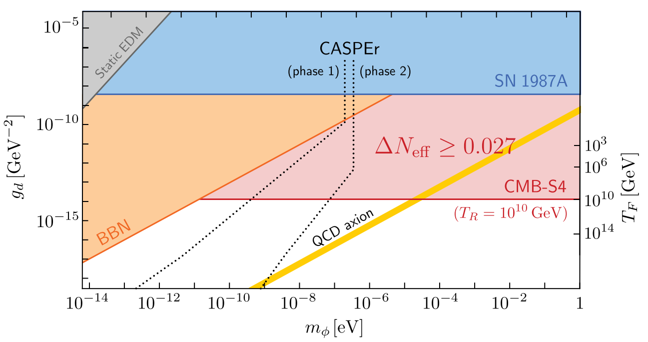

FIG. 3. Comparison between current constraints on the axion-gluon coupling and the sensitivity of a future CMB-S4 mission (figure adapted from Refs. [39,41]). The dotted lines are the projected sensitivities of the NMR experiment CASPEr [42]. We note that CASPEr, the static EDM [39], and BBN constraints [41] assume that the axion is all of the dark matter, while SN 1987A [15] and the future CMB constraint do not have this restriction.

second and third generations. We wish to compare these constraints to the reach of future CMB observations.

Above the EWSB scale, the production of the familon ϕ is determined by a four-point interaction. This allows the following processes: $\bar{\psi}_i + \psi_j \to H + \phi$ and $\psi_i + H \to \psi_j + \phi$. The total production rate is [33]

$$\Gamma_{ij}^I \simeq 0.37 N_\psi \frac{(\lambda_i \mp \lambda_j)^2}{8\pi} \frac{T^3}{(\Lambda_{ij}^I)^2}, \tag{11}$$

TABLE I. Current experimental constraints on Goldstone-fermion couplings [17,44,45] and future CMB constraints. The quoted freeze-out bounds are for $T_R = 10^{10}$ GeV and require that a future CMB experiment excludes $\Delta N_{\rm eff} = 0.027$. The freeze-in bounds, in contrast, do not depend on T_R and assume weaker exclusions $\Delta\tilde{N}_{\rm eff}$ [see the last column for estimates of the freeze-in contributions associated with the different couplings, $\Delta\tilde{N}_{\rm eff} \approx \Delta N_{\rm eff}(\frac{1}{4} m_i)$].

	Current constraints		Future CMB constraints		
Coupling	Bound [GeV]	Origin	Freeze-Out [GeV]	Freeze-In [GeV]	$\Delta\tilde{N}_{\rm eff}$
Λ_{ee}	1.2×10^{10}	White dwarfs	6.0×10^{7}	2.7×10^{6}	1.3
$\Lambda_{\mu\mu}$	2.0×10^{6}	Stellar cooling	1.2×10^{10}	3.4×10^{7}	0.5
$\Lambda_{\tau\tau}$	2.5×10^{4}	Stellar cooling	2.1×10^{11}	9.5×10^{7}	0.05
Λ_{bb}	6.1×10^{5}	Stellar cooling	9.5×10^{11}	$\cdots$	0.04
Λ_{tt}	1.2×10^{9}	Stellar cooling	3.5×10^{13}	$\cdots$	0.03
$\Lambda_{\mu e}^V$	5.5×10^{9}	$\mu^+ \to e^+\phi$	6.2×10^{9}	4.8×10^{7}	0.5
$\Lambda_{\mu e}$	3.1×10^{9}	$\mu^+ \to e^+\phi\gamma$	6.2×10^{9}	4.8×10^{7}	0.5
$\Lambda_{\tau e}$	4.4×10^{6}	$\tau^- \to e^-\phi$	1.0×10^{11}	1.3×10^{8}	0.05
$\Lambda_{\tau\mu}$	3.2×10^{6}	$\tau^- \to \mu^-\phi$	1.0×10^{11}	1.3×10^{8}	0.05
Λ_{cu}^A	6.9×10^{5}	$D^0 - \bar{D}^0$	1.3×10^{11}	2.0×10^{8}	0.05
Λ_{bd}^A	6.4×10^{5}	$B^0 - \bar{B}^0$	4.8×10^{11}	3.7×10^{8}	0.04
Λ_{bs}	6.1×10^{7}	$b \to s\phi$	4.8×10^{11}	3.7×10^{8}	0.04
Λ_{tu}	6.6×10^{9}	Mixing	1.8×10^{13}	2.1×10^{9}	0.03
Λ_{tc}	2.2×10^{9}	Mixing	1.8×10^{13}	2.1×10^{9}	0.03

where $N_\psi = 1$ for charged leptons and $N_\psi = 3$ for quarks. The "$-$" and "$+$" signs apply to $I = V$ and $I = A$, respectively. Deriving the freeze-out temperature and imposing $T_F > T_R$, we find

$$\Lambda_{ij}^I > \frac{m_i \mp m_j}{m_{\tau/t}} \sqrt{T_{R,10}} \begin{cases} 1.0\times10^{11}\ \text{GeV}, \\ 1.8\times10^{13}\ \text{GeV}, \end{cases} \tag{12}$$

where the first line applies to charged leptons with $m_\tau \approx 1.8$ GeV and the second to quarks with $m_t \approx 173$ GeV. In Table I, we show how these bounds compare to current laboratory and astrophysics constraints for a fiducial reheating temperature of 10^{10} GeV. Except for the coupling to electrons, the constraints from future CMB experiments are orders of magnitude stronger than existing constraints. For lower reheating temperatures the constraints would weaken proportional to $\sqrt{T_R}$. We note that, except for the top quark, laboratory and astrophysical constraints are considerably weaker for second and third generation particles because of kinematics, while the cosmological constraints are strengthened for the higher mass fermions due to the larger effective strength of the interactions.

Below the EWSB scale, the leading coupling of the familon to fermions becomes marginal after replacing the Higgs in Eq. (10) with its vacuum expectation value. The temperature dependence of the interaction rate is then weaker than that of the Hubble expansion rate, leading to a recoupling (i.e., freeze-in) of the pNGBs at low temperatures. To avoid a large density of pNGBs requires that the freeze-in temperature $T_{\tilde{F}}$ is smaller than the mass of the fermions participating in the interactions, $T_{\tilde{F}} < m_\psi$, so that the interaction rate becomes Boltzmann suppressed before freeze-in can occur. Again, this constraint can be expressed as a bound on the scales that couple the pNGBs to the SM fermions.

For the diagonal couplings in Eq. (10), the production rate is dominated by a Compton-like process, $\{\gamma, g\} + \psi_i \to \psi_i + \phi$, and by fermion-antifermion annihilation, $\bar{\psi}_i + \psi_i \to \{\gamma, g\} + \phi$, where $\{\gamma, g\}$ is either a photon or gluon depending on whether the fermion is a lepton or quark. Since freeze-in occurs at low temperatures, the quark production becomes sensitive to strong coupling effects. We therefore only present bounds for the lepton couplings. Above the lepton mass, the production rate is $\tilde{\Gamma}_{ii} \simeq 5.3\alpha|\tilde{\epsilon}_{ii}|^2/(8\pi)T$ [33], where $\tilde{\epsilon}_{ii} \equiv 2m_i/\Lambda_{ii}$. Deriving the freeze-in temperature and imposing $T_{\tilde{F}} < m_i$, we find

$$\Lambda_{ii} > 9.5\times10^7\ \text{GeV} \left(\frac{g_{*,\tau}}{g_{*,i}}\right)^{1/4} \left(\frac{\alpha_i}{\alpha_\tau}\frac{m_i}{m_\tau}\right)^{1/2}, \tag{13}$$

where $g_{*,i}$ and α_i are the effective number of relativistic species and the fine-structure constant at $T = m_i$. Except for the coupling to electrons, these new bounds are significantly stronger than the existing constraints. In particular, it is worth noting that the Planck constraint on the diagonal muon coupling, $\Lambda_{\mu\mu} > 3.4\times10^7$ GeV, improves on the current experimental bound by more than an order of magnitude.

For the off-diagonal couplings in Eq. (10), we have the possibility of a freeze-in population of the familon from the decay of the heavy fermion, $\psi_i \to \psi_j + \phi$. For $m_i \gg m_j$, the production rate associated with this process is $\tilde{\Gamma}_{ij} \simeq 0.31 N_\psi |\tilde{\epsilon}_{ij}|^2/(8\pi) m_i^2/T$ [33], where $\tilde{\epsilon}_{ij} \approx m_i/\Lambda_{ij}$. Requiring the corresponding freeze-in temperature to be below the mass of the heavier fermion, $T_{\tilde{F}} < m_i$, we get

$$\Lambda_{ij} > \left(\frac{g_{*,\tau/t}}{g_{*,i}}\right)^{1/4} \left(\frac{m_i}{m_{\tau/t}}\right)^{1/2} \begin{cases} 1.3 \times 10^8 \text{ GeV}, \\ 2.1 \times 10^9 \text{ GeV}, \end{cases} \tag{14}$$

where the first line applies to charged leptons and the second to quarks. We see that this improves over existing constraints for the third generation leptons and for the second and third generation quarks (except the top).

Conclusions.—In closing, we would like to re-emphasize that $\Delta N_{\mathrm{eff}} = 0.027$ is an important theoretical threshold (see Refs. [17,28,46,47] for related discussions). Remarkably, this target is within reach of future cosmological observations [23], including the planned CMB-S4 mission [19]. These observations therefore have the potential to probe for light thermal relics up to arbitrarily high decoupling temperatures. We consider this to be a unique opportunity to detect new particles, or place very strong constraints on their couplings to the Standard Model.

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[*]dbaumann@damtp.cam.ac.uk
[†]drgreen@cita.utoronto.ca
[‡]b.wallisch@damtp.cam.ac.uk

[1] P. A. R. Ade *et al.* (Planck Collaboration), Astron. Astrophys. **594**, A13 (2016).
[2] B. Follin, L. Knox, M. Millea, and Z. Pan, Phys. Rev. Lett. **115**, 091301 (2015).
[3] D. Baumann, D. Green, J. Meyers, and B. Wallisch, J. Cosmol. Astropart. Phys. 01 (2016) 007.
[4] J. Jaeckel and A. Ringwald, Annu. Rev. Nucl. Part. Sci. **60**, 405 (2010).
[5] R. Essig *et al.*, arXiv:1311.0029.
[6] R. D. Peccei and H. R. Quinn, Phys. Rev. Lett. **38**, 1440 (1977).
[7] S. Weinberg, Phys. Rev. Lett. **40**, 223 (1978).
[8] F. Wilczek, Phys. Rev. Lett. **40**, 279 (1978).
[9] F. Wilczek, Phys. Rev. Lett. **49**, 1549 (1982).
[10] D. Reiss, Phys. Lett. B **115**, 217 (1982).
[11] J. Kim, Phys. Rep. **150**, 1 (1987).
[12] Y. Chikashige, R. N. Mohapatra, and R. D. Peccei, Phys. Lett. B **98**, 265 (1981).
[13] Y. Chikashige, R. N. Mohapatra, and R. D. Peccei, Phys. Rev. Lett. **45**, 1926 (1980).
[14] P. W. Graham, I. Irastorza, S. Lamoreaux, A. Lindner, and K. van Bibber, Annu. Rev. Nucl. Part. Sci. **65**, 485 (2015).
[15] G. G. Raffelt, *Stars as Laboratories for Fundamental Physics* (University of Chicago Press, Chicago, IL, 1996).
[16] G. G. Raffelt, Proc. Int. Sch. Phys. Fermi **182**, 61 (2012).
[17] C. Brust, D. Kaplan, and M. Walters, J. High Energy Phys. 12 (2013) 058.
[18] D. Marsh, Phys. Rep. **643**, 1 (2016).
[19] K. Abazajian *et al.*, Astropart. Phys. **63**, 66 (2015).
[20] W. L. K. Wu, J. Errard, C. Dvorkin, C. L. Kuo, A. T. Lee, P. McDonald, A. Slosar, and O. Zahn, Astrophys. J. **788**, 138 (2014).
[21] S. Weinberg, Phys. Rev. Lett. **110**, 241301 (2013).
[22] G. Mangano, G. Miele, S. Pastor, T. Pinto, O. Pisanti, and P. Serpico, Nucl. Phys. **B729**, 221 (2005).
[23] A. Font-Ribera, P. McDonald, N. Mostek, B. Reid, H.-J. Seo, and A. Slosar, J. Cosmol. Astropart. Phys. 05 (2014) 023.
[24] A. Manzotti, S. Dodelson, and Y. Park, Phys. Rev. D **93**, 063009 (2016).
[25] Z. Chacko, L. J. Hall, T. Okui, and S. J. Oliver, Phys. Rev. D **70**, 085008 (2004).
[26] S. Hannestad and G. G. Raffelt, Phys. Rev. D **72**, 103514 (2005).
[27] A. Friedland, K. Zurek, and S. Bashinsky, arXiv:0704.3271.
[28] A. Salvio, A. Strumia, and W. Xue, J. Cosmol. Astropart. Phys. 01 (2014) 011.
[29] E. Braaten and T. C. Yuan, Phys. Rev. Lett. **66**, 2183 (1991).
[30] M. Bolz, A. Brandenburg, and W. Buchmüller, Nucl. Phys. **B606**, 518 (2001).
[31] E. Masso, F. Rota, and G. Zsembinszki, Phys. Rev. D **66**, 023004 (2002).
[32] P. Graf and F. D. Steffen, Phys. Rev. D **83**, 075011 (2011).
[33] See Supplemental Material at http://link.aps.org/supplemental/10.1103/PhysRevLett.117.171301 for derivations of the pNGB production rates and a discussion of the (negligible) effects of decays of pNGBs on the derived bounds.
[34] G. Carosi, A. Friedland, M. Giannotti, M. Pivovaroff, J. Ruz, and J. Vogel, arXiv:1309.7035.
[35] A. Friedland, M. Giannotti, and M. Wise, Phys. Rev. Lett. **110**, 061101 (2013).
[36] D. Cadamuro, S. Hannestad, G. Raffelt, and J. Redondo, J. Cosmol. Astropart. Phys. 02 (2011) 003.
[37] D. Cadamuro and J. Redondo, J. Cosmol. Astropart. Phys. 02 (2012) 032.
[38] M. Millea, L. Knox, and B. D. Fields, Phys. Rev. D **92**, 023010 (2015).
[39] P. W. Graham and S. Rajendran, Phys. Rev. D **88**, 035023 (2013).
[40] M. Pospelov and A. Ritz, Phys. Rev. Lett. **83**, 2526 (1999).
[41] K. Blum, R. D'Agnolo, M. Lisanti, and B. Safdi, Phys. Lett. B **737**, 30 (2014).
[42] D. Budker, P. W. Graham, M. Ledbetter, S. Rajendran, and A. O. Sushkov, Phys. Rev. X **4**, 021030 (2014).
[43] M. Peskin and D. Schroeder, *An Introduction to Quantum Field Theory* (Addison-Wesley, Reading, MA, 1995).
[44] J. L. Feng, T. Moroi, H. Murayama, and E. Schnapka, Phys. Rev. D **57**, 5875 (1998).
[45] B. Hansen, H. Richer, J. Kalirai, R. Goldsbury, S. Frewen, and J. Heyl, Astrophys. J. **809**, 141 (2015).
[46] M. Kawasaki, M. Yamada, and T. T. Yanagida, Phys. Rev. D **91**, 125018 (2015).
[47] Z. Chacko, Y. Cui, S. Hong, and T. Okui, Phys. Rev. D **92**, 055033 (2015).