Cosmologically Safe eV-Scale Sterile Neutrinos and Improved Dark Matter Structure

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We show that sterile neutrinos with masses $\gtrsim 1$ eV, as motivated by several short baseline oscillation anomalies, can be consistent with cosmological constraints if they are charged under a hidden sector force mediated by a light boson. In this case, sterile neutrinos experience a large thermal potential that suppresses mixing between active and sterile neutrinos in the early Universe, even if vacuum mixing angles are large. Thus, the abundance of sterile neutrinos in the Universe remains very small, and their impact on big bang nucleosynthesis, cosmic microwave background, and large-scale structure formation is negligible. It is conceivable that the new gauge force also couples to dark matter, possibly ameliorating some of the small-scale structure problems associated with cold dark matter.

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Introduction.—Several anomalies in short baseline neutrino oscillation experiments have spurred interest in models with more than three neutrino species. In particular, the excesses of electron neutrino events in LSND [1] and MiniBooNE [2], as well as the unexpected electron antineutrino disappearance at short baselines [3–6], could be explained in models with extra "sterile" neutrinos, i.e., light ($m \sim 1$ eV) new fermions that are uncharged under the standard model (SM) gauge group and mix with the three known neutrino species. On the other hand, a number of other neutrino oscillation experiments that did not observe any anomalous signals put such models under pressure [7–11], and a vigorous experimental program is currently under way to resolve the tension either by confirming the anomalies or by providing a definitive null result.

It is often argued that the tightest constraints on sterile neutrino models come from cosmology. Indeed, the simplest models—with just one or several sterile neutrinos, but no other new particles—are disfavored by the big bang nucleosynthesis (BBN) and Planck measurements of $N_{\rm eff}$, the number of relativistic particle species in the early Universe [12,13]. For sterile neutrino masses of ~ 1 eV or larger, even tighter constraints are obtained from large-scale structure formation [14], where the presence of extra neutrino species would lead to a washout of structure due to efficient energy transport by neutrinos.

In this Letter, we show that these constraints are evaded if sterile neutrinos have hidden interactions mediated, for instance, by a new gauge boson A', often called a dark photon, with a mass $M \lesssim$ MeV. As discussed below, gauge forces of this type are also interesting in dark matter (DM) physics and are probed in many cosmological and astrophysical searches [15]. We will show that at nonzero temperature the sterile neutrinos feel a Mikheyev-Smirnov-Wolfenstein (MSW) potential that suppresses mixing between active and sterile neutrinos in the early Universe, thus preventing sterile neutrino production in the early Universe. We will discuss constraints on this scenario from cosmology and particle physics. In the final part of the Letter, we will also discuss the possibility that A' also couples to the DM (χ) in the Universe, possibly easing the disagreement between small-scale structure observations and cold DM simulations.

Hidden sterile neutrinos.—We assume the SM is augmented by one extra species of light ($\sim$eV) neutrinos ν_s, which do not couple to the SM gauge bosons but are charged under a new $U(1)_\chi$ gauge symmetry. We assume that ν_s have relatively large ($\sim$10%) vacuum mixing with the active neutrinos and is thus capable of explaining the short baseline oscillation anomalies.

The sterile sector is expected to be coupled to the SM sector through high-scale interactions, and the two sectors decouple at temperatures $\gtrsim$TeV. Our results remain qualitatively correct even for decoupling temperatures as low as 1 GeV, i.e., just above the QCD phase transition. After decoupling, the temperature T_s of the sterile sector continues to drop as $T_s \sim 1/a$ (a being the scale factor of the Universe), while the temperature in the visible sector T_γ drops more slowly because of the entropy generated when heavier degrees of freedom (unstable hadrons, positrons, etc.) become inaccessible and annihilate or decay away. By the BBN epoch, the number of effective degrees of freedom of the visible sector g_* decreases from $\simeq$106.7 to $\simeq$10.75. Taking the sterile sector temperature as $T_s = (g_{*,\mathrm{T}_\gamma}/g_{*,\mathrm{TeV}})^{1/3} T_\gamma$, the additional effective number of fully thermalized neutrinos at BBN, for a single left-handed sterile neutrino (and its right-handed antineutrino) and a relativistic A', is

$$\Delta N_\nu \equiv \frac{\rho_{\nu_s} + \rho_{A'}}{\rho_\nu} = \frac{(g_{\nu_s} + g_{A'}) T_s^4}{g_\nu T_\nu^4} \tag{1}$$

$$= \frac{(\frac{7}{8} \times 2 + 3) \times (\frac{10.75}{106.7})^{4/3}}{(\frac{7}{8} \times 2) \times (\frac{4}{11})^{4/3}} \simeq 0.5, \tag{2}$$

which is easily consistent with the bound from BBN, viz., $\Delta N_\nu = 0.66^{+0.47}_{-0.45}$ [12]. Up to three generations of sterile neutrinos could be accommodated within $\simeq 1\sigma$. Note that we have conservatively taken T_ν at the end of BBN.

At lower temperatures, $T_s \lesssim 0.1$ MeV, A' becomes nonrelativistic, and decays to sterile neutrinos, heating them up by a factor of $\simeq 1.4$. However, these neutrinos with masses $m \gtrsim 1$ eV are nonrelativistic by the epoch of matter-radiation equality ($T_\gamma \simeq 0.7$ eV) and recombination ($T_\gamma \simeq 0.3$ eV). Thus, the impact of thermal abundances of A' and ν_s on the cosmic microwave background (CMB) and structure formation is negligible. See also Refs. [16–18] for alternate approaches. We will now show that oscillations of active neutrinos into sterile neutrinos, which are normally expected to bring the two sectors into equilibrium again, are also strongly suppressed due to "matter" effects.

The basic idea underlying our proposal is similar to the high-temperature counterpart of the MSW effect. Let us recall that at high temperatures, i.e., in the early Universe, an active neutrino with energy E experiences a potential $V_{\rm MSW} \propto G_F^2 E T_\gamma^4$ due to their own energy density [19]. This is not zero even in a CP symmetric universe. A similar, but much larger, potential can be generated at high temperature for sterile neutrinos if they couple to a light hidden gauge boson A'. There are two types of processes that can contribute to this potential—the sterile neutrino can forward scatter off an A' in the medium or off a fermion f that couples to A'.

These interactions of the sterile neutrino with the medium modify its dispersion relation through a potential $V_{\rm eff}$:

$$E = |\mathbf{k}| + \frac{m^2}{2E} + V_{\rm eff}, \tag{3}$$

where E and $|\mathbf{k}|$ are the energy and momentum of the sterile neutrino.

We calculated $V_{\rm eff}$ using the real-time formalism in thermal field theory (see Supplemental Material [20]). Physically, this potential is the correction to the sterile neutrino self-energy. In the low-temperature limit, i.e., T_s, $E \ll M$, we find $V_{\rm eff} \simeq -28\pi^3 \alpha_\chi E T_s^4/(45 M^4)$, similar to the potential for active neutrinos [19], with $\alpha_\chi \equiv e_\chi^2/(4\pi)$ being the $U(1)_\chi$ fine-structure constant. In the high-temperature limit T_s, $E \gg M$, we find $V_{\rm eff} \simeq +\pi\alpha_\chi T_s^2/(2E)$, similar to the result for hot QED [21]. We have assumed that there is no asymmetry in ν_s, which may be interesting to consider [16,22]. These analytical results are plotted in Fig. 1 (thick black lines). For comparison, we also calculated the potential numerically (thin black lines), and found excellent consistency with the analytical approximations in their region of validity. The potential is small only in a very small range of temperatures $T_s \approx M$, where the potential changes sign and goes through zero. Note that the potential is always smaller than $|\mathbf{k}|$ and vanishes at zero temperature.

In the presence of a potential, it is well known that neutrino mixing angles are modified. In the two-flavor approximation, the effective mixing angle θ_m in matter is given by [23]

$$\sin^2 2\theta_m = \frac{\sin^2 2\theta_0}{(\cos 2\theta_0 + \frac{2E}{\Delta m^2} V_{\rm eff})^2 + \sin^2 2\theta_0}, \tag{4}$$

where θ_0 is the vacuum mixing angle and $\Delta m^2 = m_s^2 - m_a^2$ is the difference between the squares of the mostly sterile mass eigenstate m_s and the active neutrino mass scale m_a. If the potential is much larger than the vacuum oscillation frequency, i.e.,

$$|V_{\rm eff}| \gg |\frac{\Delta m^2}{2E}|, \tag{5}$$

then θ_m will be tiny, and oscillations of active neutrinos into sterile ones are suppressed.

This is confirmed by Fig. 1, which summarizes our main results. For a typical neutrino energy $E \sim T_\gamma$ and $M \lesssim 10$ MeV, we see that condition (5) is well satisfied

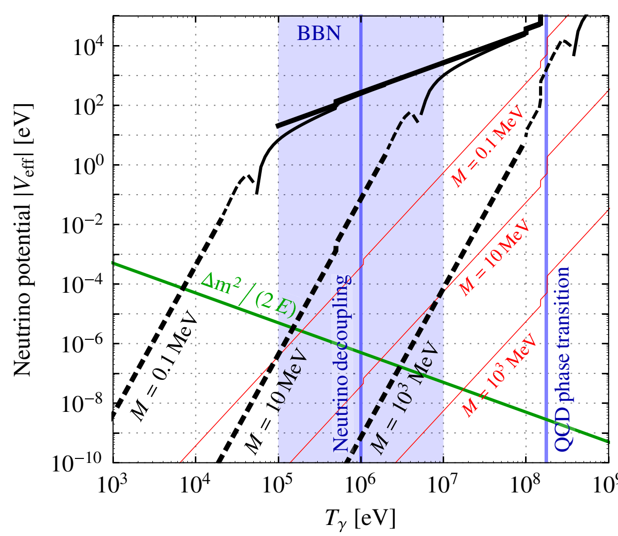

FIG. 1 (color online). Comparison of the effective matter potential $V_{\rm eff}$ for sterile neutrinos (black curves) to the active-sterile oscillation frequency $\Delta m^2/(2E)$ (green line) at $E \simeq T_\gamma$ and $\Delta m^2 = 1$ eV^2. As long as $|V_{\rm eff}| \gg \Delta m^2/(2E)$, oscillations are suppressed. Different black curves show $|V_{\rm eff}|$ for different values of the gauge boson mass M, with solid lines corresponding to $V_{\rm eff} > 0$ and dashed lines indicating $V_{\rm eff} < 0$. Thin (thick) lines show exact numerical (approximate analytical) results. The hidden sector fine-structure constant is taken as $\alpha_\chi \equiv e_\chi^2/(4\pi) = 10^{-2}/(4\pi)$. Red lines show the contribution to $V_{\rm eff}$ from an asymmetric DM particle with $m_\chi = 1$ GeV. The QCD phase transition and active neutrino decoupling epochs are annotated. The small kinks in the curves are due to changes in g_*, the effective number of degrees of freedom in the Universe.

down to temperatures $T_\gamma \lesssim 1$ MeV, i.e., until after the time of neutrino decoupling, when their thermal production becomes impossible. Thus, θ_m is suppressed and sterile neutrinos are not produced in significant numbers. There is also nonforward scattering of sterile neutrinos mediated by the hidden gauge boson, as well as the usual MSW potential for active neutrinos, which further suppress oscillations. A full numerical calculation using quantum kinetic equations [24] is consistent with our simple estimate using condition (5). Oscillations after decoupling reduce a small fraction, $\sin^2 2\theta_m \lesssim 0.1$, of the active neutrinos to steriles (which are nonrelativistic below 1 eV), consistent with $N_{\rm eff} = 3.30^{+0.54}_{-0.51}$ (95% limits) from cosmological data [13]. Note that in Fig. 1 we have conservatively taken sterile neutrino decoupling to occur at the same temperature, $T_\gamma \simeq 1$ MeV, as the decoupling of active neutrinos. In reality, sterile neutrino production ceases when $\Gamma_s \sim \sin^2\theta_s G_F^2 T_\gamma^5$ drops below the Hubble expansion rate $H \propto T_\gamma^2$, which happened at temperatures around 1 MeV$/(\sin^2\theta)^{1/3}$.

Even for M slightly larger than 1 MeV, sterile neutrino production remains suppressed until the BBN epoch, but it is interesting that in this case $V_{\rm eff}$ crosses zero while neutrinos are still in thermal equilibrium. This implies that there is a brief time period during which sterile neutrinos could be produced efficiently. However, as long as its duration is much shorter than the inverse of the sterile neutrino production rate $\Gamma_s^{-1} \sim [\sin^2\theta_s G_F^2 T_\gamma^5]^{-1}$, only partial thermalization of sterile neutrinos will occur. Interestingly, at the MSW resonance, i.e., $\Delta m^2 \simeq -2EV_{\rm eff}$, one may get some active-to-sterile neutrino (or antineutrino) conversion, depending on the adiabaticity of this resonance. This implies that, for $M \gtrsim 10$ MeV, we predict a fractional value of $\Delta N_{\rm eff}$ at BBN. A study of the detailed dynamics during this epoch is beyond the scope of our present work.

As a final remark, we would like to emphasize that, while Fig. 1 is for $E = T_\gamma$, it is important to keep in mind that active neutrinos follow a thermal distribution. We have checked that even for E different from T_γ, the value of $V_{\rm eff}$ does not change too much. Therefore, our conclusions regarding the suppression of sterile neutrino production remain valid even when the tails of the thermal distribution are taken into account.

Coupling to dark matter.—If a new gauge force of the proposed form exists, it is conceivable that not only sterile neutrinos but also DM particles χ couple to it. This, of course, leads to an additional contribution $2\pi\alpha_\chi(n_\chi - n_{\bar\chi})/M^2$ to $V_{\rm eff}$, through forward scattering off the net DM density (see Supplemental Material [20]). As long as DM is CP symmetric, we have $n_\chi - n_{\bar\chi} = 0$, and this extra contribution vanishes. Even for asymmetric DM [25], we see in Fig. 1 (red lines) that it is usually subleading for $m_\chi \gtrsim 1$ GeV.

The extra gauge interaction of DM does, however, lead to DM self-scattering, which has received considerable attention recently as a way of solving [26–28] the existing disagreement between the observed substructure of DM in the Milky Way and N-body simulations of galaxy formation. In particular, self-interacting DM can solve the "too big to fail" problem [29,30], i.e., the question of why very massive DM subhalos that are predicted to exist in a Milky Way–type galaxy have not been observed, even though one would expect star formation to be efficient in them and make them appear as luminous dwarf galaxies. Similarly, DM self-interactions could be the reason why the Milky Way appears to have fewer dwarf galaxies than expected from simulations (the "missing satellites" problem [31]). Finally, it may be possible to explain why the observed DM density distribution in Milky Way subhalos appears to exhibit a constant density core [32,33] rather than a steep cusp predicted in N-body simulations [34] ("cusp versus core problem"). While all of these problems could well have different explanations—for instance, the impact of baryonic feedback on N-body simulations is not yet well understood—it is intriguing that the self-scattering cross sections predicted in the scenario discussed here has exactly the right properties to mitigate these small-scale structure issues.

In our model, the "energy transfer cross section" in the center of the mass frame, $\sigma_T = \int d\Omega d\sigma/d\Omega(1-\cos\theta)$, is given in Born approximation by [35]

$$\sigma_T \simeq \frac{8\pi\alpha_\chi^2}{m_\chi^2 v_{\rm rel}^4}\left[\log(1+R^2) - \frac{R^2}{1+R^2}\right], \tag{6}$$

with $R \equiv m_\chi v_{\rm rel}/M$. Here, $v_{\rm rel}$ is the relative velocity of the two colliding DM particles. It is easy to see that σ_T is velocity independent for $v_{\rm rel} \ll M/m_\chi$ and drops roughly $\propto v_{\rm rel}^{-4}$ for larger $v_{\rm rel} \gg M/m_\chi$. This implies that the velocity-averaged cross section per unit DM mass, $\langle\sigma_T\rangle/m_\chi$, can be of order 0.1–1 cm^2/g in galaxies $[v_{\rm rel} \sim \mathcal{O}(100\ {\rm km/sec})]$, as required to mitigate the small-scale structure problems [27,28], while remaining well below this value in galaxy clusters $[v_{\rm rel} \sim \mathcal{O}(1000\ {\rm km/sec})]$, from which the most robust constraints are obtained [36]. The cross section given in Eq. (6) becomes inaccurate in the limit $\alpha_\chi m_\chi/M > 1$, and one needs to take nonperturbative or resonant effects into account. In computing $\langle\sigma_T\rangle$, we take the analytical expressions for σ_T for symmetric DM, as summarized in [37], and convolve with a DM velocity distribution, which we take to be of Maxwell-Boltzmann form, with velocity dispersion $v_{\rm rel}$.

As for the missing satellites problem, it was shown in [38–40] that DM-neutrino scattering can decrease the temperature of kinetic decoupling of DM, $T_{\rm kd}$, which can increase the cutoff in the structure power spectrum, $M_{\rm cut} \propto T_{\rm kd}^{-3}$, to the scales of the dwarf galaxies. $T_{\rm kd}$ is determined by equating the DM momentum relaxation rate $\sim(T_s/m_\chi)n_\chi\sigma_{\chi s}$ with the Hubble expansion rate. Here, $n_\chi \sim T_s^3$ is the DM number density, and $\sigma_{\chi s} \sim T_s^2/M^4$

is the DM-sterile neutrino scattering cross section. Quantitatively [39],

$$\frac{M_{\rm cut}}{M_{\rm Sun}} \simeq 3.2\times 10^{13}\alpha_\chi^{3/2}\left(\frac{T_s}{T_\gamma}\right)_{\rm kd}^{9/2}\left(\frac{\rm TeV}{m_\chi}\right)^{3/4}\left(\frac{\rm MeV}{M}\right)^3. \quad (7)$$

In previous literature, the exponent of T_s/T_γ in Eq. (7) is sometimes incorrectly given as $3/2$ [41]. We find the cutoff can be raised to $M_{\rm cut} = 10^9\text{–}10^{10} M_{\rm Sun}$, as required to solve the missing satellites problem. The number of sterile neutrino generations N_s, assumed to be 1 here, only weakly impacts the result as $M_{\rm cut} \propto N_s^{3/4}$. Note that in contrast to Ref. [39], we obtain a small T_s/T_γ, from decays of heavy standard model particles after the decoupling of the sterile sector.

In Fig. 2, we show the region of parameter space favored by these considerations. We see that it is possible to simultaneously mitigate the cusp versus core problem, too big to fail problem, as well as the missing satellites problem, while remaining consistent with the cluster constraint and simultaneously suppressing sterile neutrino production to evade BBN and CMB constraints. The potentially interesting solution to all of the enduring problems with small-scale structures was first shown in a scenario with active neutrinos [39], which has since been constrained using laboratory data, BBN, and large-scale structure [42–44]. A qualitative extension to sterile neutrinos was suggested therein, and we see here that such a scenario may be realized with no conflict with cosmology.

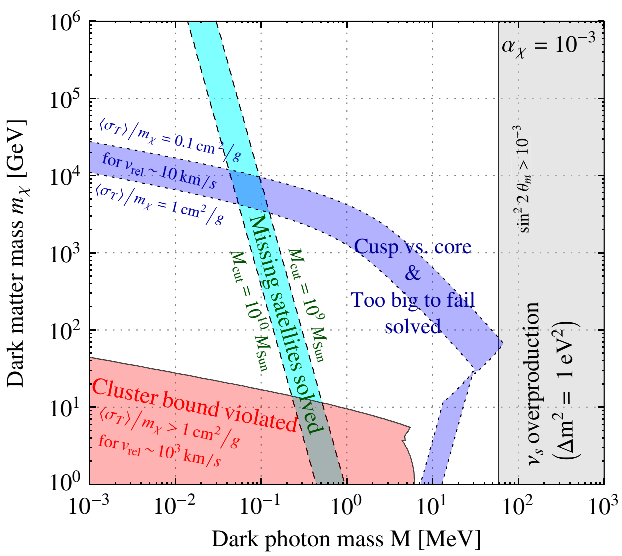

FIG. 2 (color online). Constraints on DM self-interactions from the requirements that the self-interaction in galaxy clusters is small, i.e., $\langle\sigma_T\rangle/m_\chi \lesssim 1$ cm^2/g, and that production of 1 eV sterile neutrinos is suppressed, i.e., $\sin^2 2\theta_m \lesssim 10^{-3}$ at $T_\gamma = 1$ MeV. We also show the favored parameter region for mitigating the cusp versus core and too big to fail problems, i.e., $\langle\sigma_T\rangle/m_\chi = 0.1\text{–}1$ cm^2/g in dwarf galaxies, and solve the missing satellites problem ($M_{\rm cut} = 10^{9\text{–}10} M_{\rm Sun}$). The kink in the σ_T contours is from an approximate treatment of the regime between the Born and classical limits.

The DM relic abundance may be produced by Sommerfeld-enhanced annihilations of DM into A' pairs that decay to sterile neutrinos, or alternatively through an asymmetry. However, unlike in Ref. [39], we do not use separate couplings of DM and ν to do this, so this should identify the preferred value for DM mass in the range $m_\chi \sim 1\text{–}100$ TeV. As long as DM chemical freeze-out happens well above $T_\gamma \sim$ GeV and the sterile neutrinos have time to rethermalize with ordinary neutrinos (and photons) via high-scale interactions, our scenario remains unaltered by DM annihilation.

Discussion and summary.—We now discuss the possible origin of a new gauge force in the sterile neutrino sector and on further phenomenological consequences. In Ref. [45], Pospelov has proposed a model with sterile neutrinos charged under gauged baryon number. He has argued that the model is consistent with low energy constraints, in particular, the one from $K \to \pi\pi\nu\nu$, even for $\kappa^2 \sin\theta/M^2 \sim 1000 G_F$. This is precisely the parameter region in which sterile neutrino production in the early Universe is suppressed, as we have demonstrated above. In Refs. [45–47], the phenomenological consequences of this model have been investigated, and it has been shown that strong anomalous scattering of solar neutrinos in DM detectors is expected. As an alternative to gauged baryon number, sterile neutrinos could also be charged under a gauge force that mixes kinetically with the photon [46]. In this case, $M \gtrsim 10$ MeV is preferred unless the coupling constants are extremely tiny. Once again, in this model interesting solar neutrino signals in DM detectors can occur. Finally, while we have focused here on new gauge interactions, it is also conceivable that the new interaction is instead mediated by a scalar [48,49]. However, in this case $\sigma_{\chi s} \propto m_{\nu_s}^2$, which is too small, and the missing satellite problem cannot be solved.

In summary, we have shown that eV-scale sterile neutrinos can be consistent with cosmological data from BBN, CMB, and large-scale structure if we allow them to be charged under a new gauge interaction mediated by a MeV-scale boson. In this case, sterile neutrino production in the early Universe is suppressed due to the thermal MSW potential generated by the mediator and by sterile neutrinos themselves. Our proposed scenario leads to a small fractional number of extra relativistic degrees of freedom in the early Universe, which may be experimentally testable in the future. If the considered boson also couples to DM, it could simultaneously explain observed departures of small-scale structures from the predictions of cold DM simulations.

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[1] A. Aguilar *et al.* (LSND Collaboration), Phys. Rev. D **64**, 112007 (2001).
[2] A. Aguilar-Arevalo *et al.* (MiniBooNE Collaboration), Phys. Rev. Lett. **110**, 161801 (2013).
[3] T. Mueller, D. Lhuillier, M. Fallot, A. Letourneau, S. Cormon *et al.*, Phys. Rev. C **83**, 054615 (2011).
[4] G. Mention, M. Fechner, Th. Lasserre, Th. A. Mueller, D. Lhuillier, M. Cribier, and A. Letourneau, Phys. Rev. D **83**, 073006 (2011).
[5] A. Hayes, J. Friar, G. Garvey, and G. Jonkmans, arXiv:1309.4146.
[6] M. A. Acero, C. Giunti, and M. Laveder, Phys. Rev. D **78**, 073009 (2008).
[7] J. Kopp, M. Maltoni, and T. Schwetz, Phys. Rev. Lett. **107**, 091801 (2011).
[8] J. Kopp, P. A. N. Machado, M. Maltoni, and T. Schwetz, J. High Energy Phys. 05 (**2013**) 050.
[9] J. Conrad, C. Ignarra, G. Karagiorgi, M. Shaevitz, and J. Spitz, Adv. High Energy Phys. **2013**, 163897 (2013).
[10] J. R. Kristiansen, O. y. Elgarøy, C. Giunti, and M. Laveder, arXiv:1303.4654.
[11] C. Giunti, M. Laveder, Y. F. Li, and H. W. Long, Phys. Rev. D **88**, 073008 (2013).
[12] G. Steigman, Adv. High Energy Phys. **2012**, 268321 (2012).
[13] P. Ade *et al.* (Planck Collaboration), arXiv:1303.5076.
[14] J. Hamann, S. Hannestad, G. G. Raffelt, and Y. Y. Wong, J. Cosmol. Astropart. Phys. 09 (**2011**) 034.
[15] J. Jaeckel and A. Ringwald, Annu. Rev. Nucl. Part. Sci. **60**, 405 (2010).
[16] R. Foot and R. R. Volkas, Phys. Rev. Lett. **75**, 4350 (1995).
[17] V. Barger, J. P. Kneller, P. Langacker, D. Marfatia, and G. Steigman, Phys. Lett. B **569**, 123 (2003).
[18] C. M. Ho and R. J. Scherrer, Phys. Rev. D **87**, 065016 (2013).
[19] D. Nötzold and G. Raffelt, Nucl. Phys. B **307**, 924 (1988).
[20] See Supplemental Material at http://link.aps.org/supplemental/10.1103/PhysRevLett.112.031803 for a derivation of the thermal potentials.
[21] H. A. Weldon, Phys. Rev. D **26**, 2789 (1982).
[22] N. Saviano, A. Mirizzi, O. Pisanti, P. D. Serpico, G. Mangano, and G. Miele, Phys. Rev. D **87**, 073006 (2013).
[23] E. K. Akhmedov, arXiv:hep-ph/0001264.
[24] S. Hannestad, R. S. Hansen, and T. Tram, preceding Letter, Phys. Rev. Lett. **112**, 031802 (2014).
[25] K. M. Zurek, arXiv:1308.0338.
[26] M. Vogelsberger, J. Zavala, and A. Loeb, Mon. Not. R. Astron. Soc. **423**, 3740 (2012).
[27] M. Rocha, A. H. Peter, J. S. Bullock, M. Kaplinghat, S. Garrison-Kimmel, J. Onorbe, and L. A. Moustakas, Mon. Not. R. Astron. Soc. **430**, 81 (2013).
[28] J. Zavala, M. Vogelsberger, and M. G. Walker, Mon. Not. R. Astron. Soc. **431**, L20 (2013).
[29] M. Boylan-Kolchin, J. S. Bullock, and M. Kaplinghat, Mon. Not. R. Astron. Soc. **415**, L40 (2011).
[30] M. Boylan-Kolchin, J. S. Bullock, and M. Kaplinghat, Mon. Not. R. Astron. Soc. **422**, 1203 (2012).
[31] A. A. Klypin, A. V. Kravtsov, O. Valenzuela, and F. Prada, Astrophys. J. **522**, 82 (1999).
[32] B. Moore, Nature (London) **370**, 629 (1994).
[33] R. A. Flores and J. R. Primack, Astrophys. J. **427**, L1 (1994).
[34] J. F. Navarro, C. S. Frenk, and S. D. M. White, Astrophys. J. **490**, 493 (1997).
[35] J. L. Feng, M. Kaplinghat, and H.-B. Yu, Phys. Rev. Lett. **104**, 151301 (2010).
[36] P. J. Fox and M. R. Buckley, Phys. Rev. D **81**, 083522 (2010).
[37] S. Tulin, H.-B. Yu, and K. M. Zurek, Phys. Rev. D **87**, 115007 (2013).
[38] T. Bringmann and S. Hofmann, J. Cosmol. Astropart. Phys. 04 (**2007**) 016.
[39] L. G. van den Aarssen, T. Bringmann, and C. Pfrommer, Phys. Rev. Lett. **109**, 231301 (2012).
[40] I. M. Shoemaker, Phys. Dark Univ. **2**, 157 (2013).
[41] T. Bringmann (private communication).
[42] R. Laha, B. Dasgupta, and J. F. Beacom, arXiv:1304.3460.
[43] B. Ahlgren, T. Ohlsson, and S. Zhou, Phys. Rev. Lett. **111**, 199001 (2013).
[44] F.-Y. Cyr-Racine, R. de Putter, A. Raccanelli, and K. Sigurdson, arXiv:1310.3278.
[45] M. Pospelov, Phys. Rev. D **84**, 085008 (2011).
[46] R. Harnik, J. Kopp, and P. A. Machado, J. Cosmol. Astropart. Phys. 07 (**2012**) 026.
[47] M. Pospelov and J. Pradler, Phys. Rev. D **85**, 113016 (2012).
[48] K. Babu and I. Rothstein, Phys. Lett. B **275**, 112 (1992).
[49] K. Enqvist, K. Kainulainen, and M. J. Thomson, Phys. Lett. B **280**, 245 (1992).