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Cosmological constraints on neutrino plus axion hot dark matter

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Abstract. We use observations of the cosmological large-scale structure to derive limits on two-component hot dark matter consisting of mass-degenerate neutrinos and hadronic axions, both components having velocity dispersions corresponding to their respective decoupling temperatures. We restrict the data samples to the safely linear regime, in particular excluding the Lyman-α forest. Using standard Bayesian inference techniques we derive credible regions in the two-parameter space of m_a and $\sum m_\nu$. Marginalising over $\sum m_\nu$ provides $m_a < 1.2$ eV (95% C.L.). In the absence of axions the same data and methods give $\sum m_\nu < 0.65$ eV (95% C.L.). We also derive limits on m_a for a range of axion–pion couplings up to one order of magnitude larger or smaller than the hadronic value.

1. Introduction

The masses of the lightest known particles (neutrinos) are best constrained by the largest known scales (the entire universe). The well-established method of using cosmological precision data to constrain the cosmic hot dark matter fraction [1, 2] has been extended to hypothetical low-mass particles, notably to axions, in several papers [3–5].

We return to this topic to extend previous studies by some of us [3, 4] in several ways. First, we update the cosmological data sets to include the Wilkinson Microwave Anisotropy Probe 3-year data as well as the baryon acoustic oscillations measurements from the Sloan Digital Sky Survey that have since become available. Second, we use standard Bayesian inference techniques to construct credible regions in parameter space, in contrast to the likelihood maximisation method used before [3, 4]. Most importantly, we consider a two-component hot dark matter fraction consisting of axions and neutrinos. Since neutrinos are known to have nonvanishing masses, their hot-dark matter contribution is an unavoidable cosmological fit parameter. Axions and neutrinos decouple at different epochs and thus have different velocity dispersions that we implement self-consistently. In this regard our work parallels a recent study by another group [5].

We begin in section 2 with a brief summary of the relevant axion parameters and their decoupling conditions. In section 3 we describe the cosmological model and the parameter space we use. In section 4 we summarise the included data sets and briefly discuss our reasons for limiting the analysis to data in the safely linear regime of structure formation. We derive our new constraints in section 5 before concluding in section 6.

2. Hot dark matter axions

The Peccei–Quinn solution of the CP problem of strong interactions predicts the existence of axions, low-mass pseudoscalars that are very similar to neutral pions, except that their mass and interaction strengths are suppressed by a factor of order f_π/f_a, where $f_\pi \approx 93$ MeV is the pion decay constant, and f_a a large energy scale, the axion decay constant or Peccei–Quinn scale [6]. In more detail, the axion mass is

$$m_a = C_a \frac{z^{1/2}}{1+z} \frac{f_\pi m_\pi}{f_a} = C_a \frac{6.0 \text{ eV}}{f_a/10^6 \text{ GeV}}, \qquad (2.1)$$

where $z = m_u/m_d$ is the mass ratio of the up and down quarks. We will follow the previous axion literature and assume a value $z = 0.56$ [7, 8], but we note that it could vary in the range 0.3–0.6 [9]. Because of this uncertainty and to cover more general cases we will sometimes include a fudge factor C_a with the standard value 1. We will consider cases with $-1 < \log_{10}(C_a) < +1$.

A large range of f_a values (or, equivalently, m_a values) can be excluded by experiments and by astrophysical and cosmological arguments [10]. Axions with a mass of order 10 μeV could well be the cold dark matter of the universe [11] and if so will be found eventually by the ongoing ADMX experiment provided that

$1\ \mu\text{eV} < m_a < 100\ \mu\text{eV}$ [12]. In addition, a hot axion population is produced by thermal processes [13–15]. Axions attain thermal equilibrium at the QCD phase transition or later if $f_a \lesssim 10^8$ GeV, erasing the cold axion population produced earlier and providing a hot dark matter component instead.

In principle, $f_a \lesssim 10^9$ GeV is excluded by the supernova SN 1987A neutrino burst duration [10]. However, the sparse data sample, our poor understanding of the nuclear medium in the supernova interior, and simple prudence suggest that one should not base far-reaching conclusions about the existence of axions in this parameter range on a single argument or experiment alone. Therefore, it remains important to tap other sources of information, especially if they are easily available.

For those axion models with nonvanishing couplings to charged fermions, there exist stellar energy loss limits based on the axion–electron coupling that are competitive with the SN 1987A constraints so that here one does not rely on a single argument to exclude axions in the $f_a \lesssim 10^9$ GeV range. Therefore, we focus on hadronic models where axions do not directly couple to ordinary quarks and leptons. In this class of models all axion properties depend on f_a alone and not on model-dependent Peccei–Quinn charges of the ordinary quarks and leptons.

If axions do not couple to charged leptons, the main thermalisation process in the post-QCD epoch is [13]

$$a + \pi \leftrightarrow \pi + \pi \,. \tag{2.2}$$

The axion–pion interaction is given by a Lagrangian of the form [13]

$$\mathcal{L}_{a\pi} = \frac{C_{a\pi}}{f_\pi f_a} \left(\pi^0\pi^+\partial_\mu\pi^- + \pi^0\pi^-\partial_\mu\pi^+ - 2\pi^+\pi^-\partial_\mu\pi^0\right)\partial_\mu a \,. \tag{2.3}$$

In hadronic axion models, the coupling constant is [13]

$$C_{a\pi} = \frac{1-z}{3\,(1+z)} \,. \tag{2.4}$$

We note that in general the chiral symmetry-breaking Lagrangian gives rise to an additional piece for $\mathcal{L}_{a\pi}$ proportional to $(m_\pi^2/f_\pi f_a)\,(\pi^0\pi^0 + 2\pi^-\pi^+)\ \pi^0 a$. However, for hadronic axion models this term vanishes identically, in contrast, for example, to the DFSZ model (Roberto Peccei, private communication).

Based on the axion–pion interaction, the axion decoupling temperature in the early universe was calculated by some of us in Ref. [4], where all relevant details are reported. In our standard case we use the axion mass m_a as our primary parameter from which we derive the corresponding axion–pion interaction strength by virtue of equations (2.1) and (2.2). Noting that even in hadronic axion models there is some uncertainty in this relationship due to the uncertain quark-mass ratio z, we consider also more general cases in which we include a fudge factor C_a as defined in equation (2.1), thus allowing for a more general relationship between m_a and $C_{a\pi}$.

Table 1. Priors and standard values for the cosmological fit parameters considered in this work. All priors are uniform in the given intervals (i.e., top hat).

Parameter	Standard	Prior
$\omega_{\rm dm}$	—	0.01–0.99
$\omega_{\rm b}$	—	0.005–0.1
h	—	0.4–1.0
τ	—	0.01–0.8
$\ln(10^{10}A_s)$	—	2.7–4.0
n_s	—	0.5–1.5
$\sum m_\nu$ [eV]	0	0–20
m_a [eV]	0	0–20
$\log_{10}(C_a)$	0	−1–1

3. Cosmological model

We consider a cosmological model with vanishing spatial curvature and adiabatic initial conditions, described by nine free parameters,

$$\boldsymbol{\theta} = \{\omega_{\rm dm}, \omega_b, H_0, \tau, \ln(10^{10}A_s), n_s, \sum m_\nu, m_a, \log_{10}(C_a)\}. \tag{3.1}$$

Here, $\omega_{\rm dm} = \Omega_{\rm dm}h^2$ is the physical dark matter density, $\omega_b = \Omega_b h^2$ the baryon density, $H_0 = h$ 100 km s^{-1} Mpc^{-1} the Hubble parameter, τ the optical depth to reionisation, A_s the amplitude of the primordial scalar power spectrum, and n_s its spectral index. These six parameters represent the simplest parameter set necessary for a consistent interpretation of the currently available data.

In addition, we allow for a nonzero sum of neutrino masses $\sum m_\nu$, a nonvanishing axion mass m_a, and a fudge factor C_a relating m_a to f_a as defined in equation (2.1). These extra parameters will be varied one at a time, as well as in combination. Their "standard" values are given in table 1, along with the priors for all cosmological fit parameters considered here.

4. Data

4.1. Cosmic microwave background (CMB)

We use CMB data from the Wilkinson Microwave Anisotropy Probe (WMAP) experiment after three years of observation [16–18]. The data analysis is performed using version 2 of the likelihood calculation package provided by the WMAP team on the LAMBDA homepage [19].

4.2. Large scale structure (LSS)

We use the large-scale galaxy power spectra $P_{\rm g}(k)$ inferred from the luminous red galaxy (LRG) sample of the Sloan Digital Sky Survey (SDSS) [20, 21] and from the Two-degree Field Galaxy Redshift Survey (2dF) [22]. These power spectra are related to the underlying matter power spectrum $P_{\rm m}(k)$ via $P_{\rm g}(k) = b^2(k)P_{\rm m}(k)$, where the galaxy bias $b(k)$ is conventionally assumed to be constant with respect to k over the scales probed by galaxy clustering surveys.

However, recent studies suggest that this assumption may break down beyond $k \sim 0.1\ h\ \mathrm{Mpc}^{-1}$, and may source the apparent tension between the SDSS and the 2dF-inferred galaxy power spectra [20, 23]. To model the effects of scale-dependent biasing when extracting cosmological parameters from galaxy clustering data, both the SDSS-LRG and the 2dF teams advocate the use of the $Q_{\rm nl}$ fitting formula developed in Ref. [22] for ΛCDM cosmologies. See Ref. [24] for a detailed discussion.

In the present work, however, we take the view that fitting formulae developed for standard cosmologies may not be applicable in nonstandard scenarios, particularly those involving new length scales arising from, e.g., axion and neutrino free-streaming. Developing an alternative formula to properly handle these nonstandard effects on the galaxy bias is also beyond our present scope. We therefore adopt a conservative approach, and use only power spectrum data well below $k \sim 0.1\ h\ \mathrm{Mpc}^{-1}$, where a scale-*independent* bias is likely to hold true:

- 2dF, $k_{\rm max} \sim 0.09\ h\ \mathrm{Mpc}^{-1}$ (17 bands),
- SDSS-LRG, $k_{\rm max} \sim 0.07\ h\ \mathrm{Mpc}^{-1}$ (11 bands).

The combined set of these data is denoted LSS. We assume a scale-independent bias for each data set, and marginalise analytically over each bias parameter b^2 with a flat prior.

4.3. Baryon acoustic oscillations (BAO)

The baryon acoustic oscillations peak has been measured in the SDSS luminous red galaxy sample [25]. We use all 20 points in the two-point correlation data set supplied in Ref. [25] and the analysis procedure described therein, including power spectrum dewiggling, nonlinear corrections with the HALOFIT package [26], corrections for redshift-space distortion, and analytic marginalisation over the normalisation of the correlation function.

4.4. Type Ia supernovae (SNIa)

We use the luminosity distance measurements of distant type Ia supernovae provided by Davis et al. [27]. This sample is a compilation of supernovae measured by the Supernova Legacy Survey (SNLS) [28], the ESSENCE project [29], and the Hubble Space Telescope [30], as well as a set of 45 nearby supernovae. In total the sample contains 192 supernovae.

4.5. Lyman-α forest (Lyα)

Measurements of the flux power spectrum of the Lyman-α forest has been used to reconstruct the matter power spectrum on small scales at large redshifts. By far the

Table 2. 1D marginal 95% upper bounds on $\sum m_\nu$ and m_a for several different choices of data sets and models.

Data set	C_a prior	$\sum m_\nu$ [eV]	m_a [eV]
WMAP+LSS+SNIa	$\log_{10}(C_a) = 0$	0.63	2.0
WMAP+LSS+SNIa+BAO		0.59	1.2
Fixed $\sum m_\nu = 0$			
WMAP+LSS+SNIa+BAO		—	1.4
Fixed $m_a = 0$			
WMAP+LSS+SNIa+BAO		0.65	—
WMAP+LSS+SNIa	$-1 < \log_{10}(C_a) < 1$	0.61	2.2
WMAP+LSS+SNIa+BAO		0.60	1.1

largest sample of spectra comes from the SDSS survey. This data set was carefully analysed in McDonald et al. [31] and used to constrain the linear matter power spectrum. The derived linear fluctuation amplitude at $k = 0.009$ km s^{-1} and $z = 3$ is $\Delta^2 = 0.452^{+0.07}_{-0.06}$, and the effective spectral index $n_{\rm eff} = -2.321^{+0.06}_{-0.05}$. These results were derived using a very elaborate model of the local intergalactic medium in conjunction with hydrodynamic simulations.

While the Lyα data provide in principle a very powerful probe of the fluctuation amplitude on small scales, the question remains as to the level of systematic uncertainty in the result. The same data have been reanalysed by Seljak et al. [32] and Viel et al. [33–35], with somewhat different results. Specifically, the normalisation found in Refs. [33–35] is lower than that reported in Ref. [31].

This question of normalisation is particularly important for bounds on the hot dark matter content of the universe. Since the free-streaming scale of light neutrinos or axions is larger than the length scale probed by Lyα, their effect on the Lyα data amounts to an overall change in the normalisation that is completely degenerate with any possible shift due to systematics. The Lyα analysis in Ref. [31] already points to a higher fluctuation amplitude Δ^2 than that derived from the WMAP 3-year data; the addition of a hot dark matter component will render the two data sets even less compatible. This incompatibility in turn leads to a much stronger formal bound on the mass of the hot dark matter particle than would be expected considering the sensitivity of the present data (this is true for both neutrinos and other types of hot dark matter, such as axions).

These considerations suggest that the Lyα data are at present dominated by systematic effects. We therefore refrain from using them in the present analysis.

5. Results

We use standard Bayesian inference techniques, and explore the model parameter space with Monte Carlo Markov Chains (MCMC) generated using the publicly available

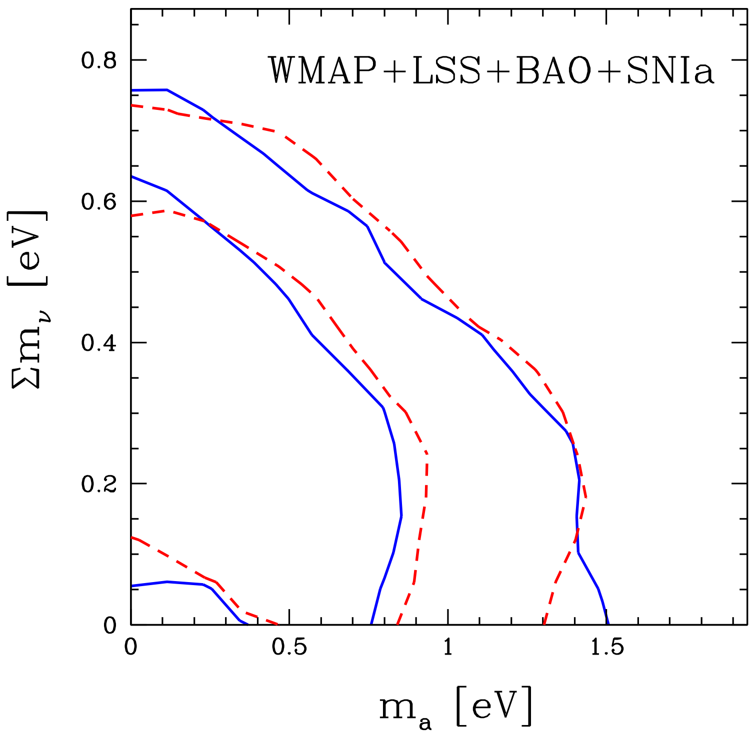

Figure 1. 2D marginal 68% and 95% contours in the $\sum m_\nu$-m_a plane derived from the full data set WMAP+LSS+BAO+SNIa. The blue/solid lines correspond to the fudge factor being fixed at $\log_{10}(C_a) = 0$, while the red/dashed lines indicate a top-hat prior on C_a in the interval $-1 < \log_{10}(C_a) < 1$.

CosmoMC package [36, 37]. Our results are summarised in tables 2 and 3, and figures 1 and 2.

For the case of a standard hadronic axion our bounds on $\sum m_\nu$ and m_a are tabulated in table 2. When BAO data are included our bounds are almost identical to those recently derived in Ref. [5] based on their conservative data set. The main difference is that we do not use the HST prior on h, but instead include the SDSS-BAO data. Since the BAO data break the Ω_m-h degeneracy, their inclusion has much the same effect as adding the HST prior. The importance of BAO data for the bound can be seen by the fact that the 95% upper bound is reduced from 2.0 eV to 1.2 eV. Our complete SNIa data set is also somewhat larger than the SNLS data set used in Ref. [5], containing in addition data from the GOODS and ESSENCE surveys. However, this has only a very modest impact on our results.

Our neutrino mass bound $\sum m_\nu < 0.65$ eV (95% C.L.) in the absence of axions is identical to that derived by some of us in Ref. [4], whereas the axion mass limit $m_a < 1.4$ eV in the absence of neutrino masses found here is significantly weaker than the 1.05 eV limit found earlier [4]. The agreement of the neutrino mass limits is coincidental because here we use different data, notably excluding the Lyman-α forest, and a different statistical methodology (marginalisation instead of maximisation). The relative difference between the limits can be interpreted such that the axion bound

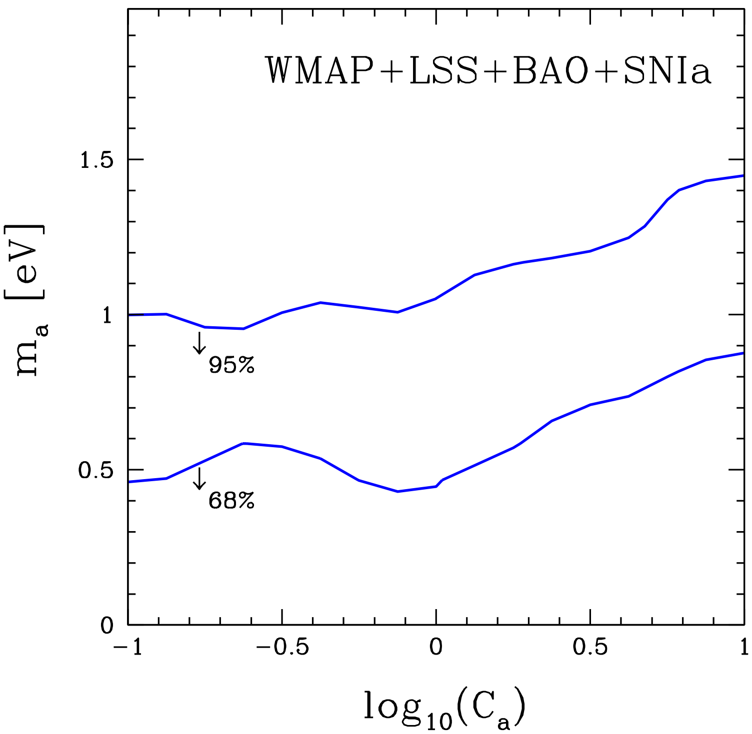

Figure 2. 2D marginal 68% and 95% contours in the m_a-$\log_{10}(C_a)$ plane, assuming a top-hat prior on the fudge factor in the interval $-1 < \log_{10}(C_a) < 1$.

benefits more from the inclusion of small-scale data than the neutrino mass bound, presumably because axions freeze out earlier and thus have a smaller velocity dispersion. Including the Lyα data here would strongly improve both limits as can be gleaned from the results of Ref. [5]. Adding Lyα to their conservative data set, the analysis of Ref. [5] finds that the marginalised axion mass limit improves by a factor 0.30, whereas the marginalised neutrino mass limit improves only by a factor 0.36, i.e., the relative gain for axions is 20% stronger. The changes in our new limits relative to those of Ref. [4] are in agreement with this picture.

Returning to our new limits, an important observation is that the upper bound on the sum of neutrino masses is largely independent of whether or not massive axions are present. The 95% upper limit on $\sum m_\nu$ is in either case approximately 0.6 eV, a bound very close to that found in previous studies using roughly the same data combination [38–41].

In figure 2 we show how the bound on the axion mass changes as C_a is allowed to vary up or down by up to a factor of 10, assuming a uniform prior on $\log_{10}(C_a)$ between -1 and $+1$. Note that the figure shows the 2D marginal contours, i.e., C_a and m_a are fitted simultaneously. For $\log_{10}(C_a) \leq 0$, the bound on m_a does not depend on C_a because the number of degrees of freedom at decoupling, $g_*(T_D)$, is approximately constant for a large range of f_a values (see table 2 of Ref. [4]). For $\log_{10}(C_a) > 0$, the value of $g_*(T_D)$ increases significantly with increasing C_a for a given m_a. This increase

Table 3. 1D marginal 68%/95% upper bounds on $\sum m_\nu$ and m_a for fixed fudge factors C_a. The data set used is WMAP+LSS+SNIa+BAO.

$\log_{10}(C_a)$	$\sum m_\nu$ [eV]	m_a [eV]
-1.0	0.39/0.64	0.51/0.98
0.0	0.37/0.59	0.60/1.2
1.0	0.40/0.63	0.69/1.4

in $g_*(T_D)$ leads to a corresponding drop in the present axion number density,

$$n_a = \frac{g_*(\text{today})}{g_*(T_D)} \times \frac{n_\gamma}{2}. \tag{5.1}$$

For a fixed m_a this amounts to a decrease in the ratio Ω_a/Ω_m with increasing C_a. The bound on m_a therefore becomes correspondingly weaker.

As can be seen in figure 2, the 2D marginal 95% upper limit on m_a stays roughly constant at $m_a \lesssim 1.1$ eV when $\log_{10}(C_a) \leq 0$, and increases roughly linearly with $\log_{10}(C_a)$ to about 1.4 eV at $\log_{10}(C_a) = 1$. We stress again that the m_a bounds in this figure are 2D bounds, and are formally—and often also in practice—not equivalent to 1D bounds on m_a derived under the assumption of a fixed C_a. For instance, the 2D bound on m_a at $C_a = 1$ in figure 2 is not exactly identical to the 1D bound quoted in table 2 for a fixed $C_a = 1$. However, despite this formality, the m_a-$\log_{10}(C_a)$ trend observed in figure 2 is also evident in table 3, which shows the 1D marginal 68% and 95% bounds on m_a for *fixed* values of C_a.

Finally, we note that if C_a is increased much beyond 10, the bound will deteriorate rapidly because axion decoupling will have occurred beyond the QCD phase transition.

6. Conclusions

We have updated previous limits from cosmological structure formation on the mass of hot dark matter axions. This limit applies to axions which were thermalised, mainly by axion–pion interactions, in the early universe, and which subsequently decoupled from the thermal plasma while still relativistic.

In the present study we investigate both the case where the neutrinos can be regarded as massless, as well as the case in which massive neutrinos are also allowed to contribute significantly to the hot dark matter fraction. In both cases we find an upper 95% limit on the mass of hadronic axions of 1.1–1.2 eV when all available cosmological data, except the Lyman-α forest, are used. Reassuringly, we find that the bound on the sum of neutrino masses is almost completely unaffected by the presence of hot dark matter axions.

Because of the uncertainty in the relation between the axion mass m_a and the energy scale f_a, we have also studied the case in which the relation $m_a = 6.0\ \text{eV}/(f_a/10^6\,\text{GeV})$ is modified by a fudge factor C_a. We have studied C_a in the range 0.1–10, which is fairly representative of the model uncertainties. We find that the axion mass bound is largely stable with respect to varying C_a in this range. For fixed values of C_a, the 1D marginal bound on m_a goes from 0.98 eV at $C_a = 0.1$ to 1.4 eV at $C_a = 10$. Essentially, this means that for hadronic axion models the uncertainty of the light quark mass ratios have a negligible impact on the axion mass limit.

Experimental and astrophysical limits on m_a or f_a are always derived from limits on the axion coupling to different particles. The cosmological hot dark matter limit, in contrast, primarily constrains the axion mass, with a very weak dependence on the axion–pion coupling. The hot dark matter limit of $m_a \lesssim 1$ eV is very similar to the limit derived from globular cluster stars based on the axion–photon coupling. However, this coupling is quite uncertain even in hadronic models because even there it depends on the unknown electric charge of the heavy quark in KSVZ-type models. The hot dark matter limit implies that it is very difficult to escape the limit $m_a \lesssim 1$ eV. One consequence is that in typical models, axions in the remaining allowed mass range necessarily escape freely from a supernova core. By courtesy of the SN 1987A neutrino burst duration, it follows that one can advance by another rung in the ladder of different limits and conclude that $m_a \lesssim 10^{-2}$ eV [10]. While this SN 1987A energy-loss limit does not have an obvious loophole, we repeat that it is based on a very small sample of detected neutrinos and is subject to various nuclear-physics and axion-model uncertainties.

Our results largely agree with those of Ref. [5] for their conservative data set. In contrast to Ref. [5] and to a previous study by some of us [4], we have not included the Lyman-α forest data which could formally improve both the neutrino and axion mass limits roughly by a factor of 3. We have explained in section 4.5 that using the Lyman-α forest exposes one to the risk that large systematic uncertainties in the normalisation of the power spectrum at small scales may dominate the final result.

The CAST experiment at CERN searches for axion-like particles emitted by the Sun by virtue of their coupling to photons [42, 43]. By including a helium filling of the magnet bores with variable pressure one can "adjust the photon mass," thereby allowing one to probe realistic combinations of m_a and axion–photon coupling. The completed runs with ^{4}He filling have already extended the experimental sensitivity to $m_a \sim 0.4$ eV. Further extensions to up to $m_a \sim 1.16$ eV with the forthcoming ^{3}He runs over three years are on the agenda [44]. This search range is not excluded by our limits, particularly as we believe that more restrictive limits derived from the Lyman-α forest may be dominated by systematic effects that are not reliably controlled.

One further caveat is that limits inferred from cosmological observations are by and large model-dependent. Additional free parameters not considered in this work, such as a nonstandard dark energy equation of state parameter, running in the primordial scalar spectral index, or a nonzero component of isocurvature modes in the initial conditions, could conceivably loosen the axion mass bound, as they have done many times before for the neutrino mass limit [38, 41, 45]. A significant and reliable improvement of cosmological hot dark matter limits is not immediately forthcoming. However, once data from the Planck CMB experiment [46] combined with other probes such as weak lensing surveys of galaxies [47–49] or of 21-cm emissions [50], or high-redshift galaxy surveys [51, 52] become available, the sensitivity will be pushed down by as mush as an order of magnitude even in the face of more complicated cosmological model frameworks. In that event, a detection of axions by CAST in the vicinity of $m_a \sim 1$ eV will have important ramifications for observational cosmology.

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