On the robustness of cosmological axion mass limits

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We present cosmological bounds on the thermal axion mass in an extended cosmological scenario in which the primordial power spectrum of scalar perturbations differs from the usual power-law shape predicted by the simplest inflationary models. The power spectrum is instead modeled by means of a "piecewise cubic Hermite interpolating polynomial" (PCHIP). When using Cosmic Microwave Background measurements combined with other cosmological data sets, the thermal axion mass constraints are degraded only slightly. The addition of the measurements of \( \sigma_8 \) and \( \Omega_m \) from the 2013 Planck cluster catalogue on galaxy number counts relaxes the bounds on the thermal axion mass, mildly favouring a \( \sim 1 \) eV axion mass, regardless of the model adopted for the primordial power spectrum.

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I. INTRODUCTION

A possible candidate for an extra hot thermal relic component is the axion particle produced thermally in the early universe. Axions therefore can contribute to the hot dark matter component together with the standard relic neutrino background. Axions may be produced in the early universe via thermal or non thermal processes, and arise as the solution to solve the strong CP problem [1]. Axions are the Pseudo-Nambu-Goldstone bosons of a new global \( U(1)_{PQ} \) (Peccei-Quinn) symmetry that is spontaneously broken at an energy scale \( f_a \). The axion mass is given by

\[
m_a = \frac{f_\pi m_\pi}{f_a} \sqrt{\frac{R}{1 + R}} = 0.6 \text{ eV} \quad \frac{10^7 \text{ GeV}}{f_a},
\]

where \( f_a \) is the axion coupling constant, \( R = 0.553 \pm 0.043 \) is the up-to-down quark masses ratio and \( f_\pi = 93 \) MeV is the pion decay constant. Thermal axions will affect the cosmological observables in a very similar way to that induced by the presence of neutrino masses and/or extra sterile neutrino species. Massive thermal axions as hot relics affect large scale structure, since they will only cluster at scales larger than their free-streaming scale when they become non-relativistic, suppressing therefore structure formation at small scales (large wavenumbers \( k \)). Concerning Cosmic Microwave Background (CMB) physics, an axion mass will also lead to a signature in the CMB photon temperature anisotropies via the early integrated Sachs-Wolfe effect. In addition, extra light species as thermal axions will contribute to the dark radiation content of the universe, or, in other words, will lead to an increase of the effective number of relativistic degrees of freedom \( N_{\text{eff}} \), defined via

\[
\rho_{\text{rad}} = \left[ 1 + \frac{7}{8} \left( \frac{4}{11} \right)^{4/3} N_{\text{eff}} \right] \rho_\gamma,
\]

where \( \rho_\gamma \) refers to the present photon energy density. In the standard cosmological model in which a thermal axion content is absent, the three active neutrino contribution leads to the canonical value of \( N_{\text{eff}} = 3.046 \) [2]. The extra contribution to \( N_{\text{eff}} \) arising from thermal axions can modify both the CMB anisotropies (via Silk damping) and the light element primordial abundances predicted by Big Bang Nucleosynthesis. The former cosmological signatures of thermal axions have been extensively exploited in the literature to derive bounds on the thermal axion mass, see Refs. [3–8].

However, all the cosmological axion mass limits to date have assumed the usual simple power-law description for the primordial perturbations. The aim of this paper is to constrain the mass of the thermal axion using a non-parametric description of the Primordial Power Spectrum (PPS hereinafter) of the scalar perturbations, as introduced in Ref. [9]. While in the simplest models of inflation [10–12] the PPS has a scale-free power-law form, the PPS could be more complicated, presenting various features or a scale dependence. Several methods have been proposed in the literature to reconstruct the shape of the PPS (see the recent work of Ref. [13]). It has been shown [14, 15] that there are small hints for deviations from the power-law form, even when using different methods and different data sets.

The energy scales at which the PPS was produced during inflation can not be directly tested. We can only infer the PPS by measuring the current matter power spectrum in the galaxy distribution and the power spectrum of the CMB fluctuations. The latter one, measured with exquisite precision by the Planck experiment [16, 18], is the convolution of the PPS with the transfer function. Therefore, in order to reconstruct the PPS, the assumption of an underlying cosmological model is a mandatory first step in order to compute the transfer function.

Here we rather exploit a non-standard PPS approach,
which can allow for a good fit to experimental data even in models that deviates from the standard cosmological picture. In particular, we consider a thermal axion scenario, allowing the PPS to assume a more general shape than the usual power-law description. This will allow us to test the robustness of the cosmological thermal axion mass bounds (see Ref. [8] for a recent standard thermal axion analysis), as first performed in Ref. [19] for the neutrino mass case.

The structure of the paper is as follows. Section II describes the PPS modeling used in this study, as well as the description of the thermal axion model explored here and the cosmological data sets exploited to constrain such a model. In Sec. III we present and discuss the results arising from our Markov Chain Monte Carlo (MCMC) numerical analyses. We draw our conclusions in Sec. IV.

II. METHOD

In this section we focus on the tools used in the numerical analyses performed here. Subsection A describes the alternative model for the PPS of scalar perturbations used for the analyses here (see also Ref. [9]), while in Subsection B we introduce the cosmological model and the thermal axion treatment followed in this study. Finally, we shall present in Subsection C the cosmological data sets used in the MCMC analyses.

A. Primordial Power Spectrum Model

The primordial fluctuations in scalar and tensor modes are generated during the inflationary phase in the early universe. The simplest models of inflation predict a power-law form for the PPS of scalar and tensor perturbations (see e.g. [10][12] and references therein), but in principle inflation can be generated by more complicated mechanisms, thus giving a different shape for the PPS (see Refs. [20][21] and references therein). In order to study how the cosmological constraints on the parameters change in more general inflationary scenarios, we assume a non-parametric form for the PPS.

Among the large number of possibilities, we decided to describe the PPS of scalar perturbations using a function to interpolate the PPS values in a series of nodes at fixed position. The interpolating function we used is named “piecewise cubic Hermite interpolating polynomial” (PCHIP) [22] and it is a modified spline function, defined to preserve the original monotonicity of the point series that is interpolated. We use a modified version of the original PCHIP algorithm [23], detailed in Appendix A of Ref. [9].

To describe the scalar PPS with the PCHIP model, we only need to give the values of the PPS in a discrete number of nodes and to interpolate among them. We use 12 nodes which span a wide range of \(k\)-values:

\[
\begin{align*}
  k_1 &= 5 \times 10^{-6} \text{ Mpc}^{-1}, \\
  k_2 &= 10^{-3} \text{ Mpc}^{-1}, \\
  k_{j+1} &= k_2 \left(\frac{k_{j+1}}{k_2}\right)^{(j-2)/9} \quad \text{for} \quad j \in [3, 10], \\
  k_{11} &= 0.35 \text{ Mpc}^{-1}, \\
  k_{12} &= 10 \text{ Mpc}^{-1}. 
\end{align*}
\]

We choose equally spaced nodes in the logarithmic scale in the range \((k_2, k_{11})\), that is well constrained from the data [19], while the first and the last nodes are useful to allow for a non-constant behaviour of the PPS outside the well-constrained range.

The PCHIP PPS is described by

\[
P_s(k) = P_0 \times \text{PCHIP}(k; P_{s,1}, \ldots, P_{s,12}),
\]

where \(P_{s,j}\) is the value of the PPS at the node \(k_j\) divided by \(P_0 = 2.36 \times 10^{-9}\) [24].

B. Cosmological and Axion Model

The baseline scenario we consider here is the ΛCDM model, extended with hot thermal relics (the axions), together with the PPS approach outlined in the previous section. For the numerical analyses we use the following set of parameters, for which we assume flat priors in the intervals listed in Tab. I.

\[
\{\omega_b, \omega_c, \Theta_s, \tau, m_a, P_{s,1}, \ldots, P_{s,12}\},
\]

where \(\omega_b \equiv \Omega_b h^2\) and \(\omega_c \equiv \Omega_c h^2\) are, respectively, the physical baryon and cold dark matter energy densities, \(\Theta_s\) is the ratio between the sound horizon and the angular diameter distance at decoupling, \(\tau\) is the reionization optical depth, \(m_a\) is the axion mass in eV and \(P_{s,1}, \ldots, P_{s,12}\) are the parameters of the PCHIP PPS.

In order to compare the results obtained with the PCHIP PPS to the results obtained with the usual power-law PPS model, we describe the latter case with the following set of parameters:

\[
\{\omega_b, \omega_c, \Theta_s, \tau, m_a, n_s, \log[10^{10} A_s]\},
\]

where \(n_s\) is the scalar spectral index, \(A_s\) the amplitude of the primordial spectrum and the other parameters are the same ones described above. The flat priors we use are listed in Tab. I.

Concerning the contribution of the axion mass-energy density to the universe’s expansion rate, we briefly summarize our treatment in the following. Axions decoupled in the early universe at a temperature \(T_D\) given by the usual freeze out condition for a thermal relic:

\[
\Gamma(T_D) = H(T_D),
\]

where the thermally averaged interaction rate \(\Gamma\) refers to the \(\pi + \pi \rightarrow \pi^0 + a\) process. The freeze out equation
above can be numerically solved, obtaining the axion decoupling temperature \( T_D \) as a function of the axion mass \( m_a \). From the axion decoupling temperature it is possible to infer the present axion number density, related to the current photon density \( n_\gamma \) by

\[
n_a = \frac{g_{*S}(T_0)}{g_{*S}(T_D)} \times \frac{n_\gamma}{2},
\]

where \( g_{*S} \) represents the number of entropic degrees of freedom, with \( g_{*S}(T_0) = 3.91 \). As previously stated, the presence of a thermal axion will also imply an extra radiation component at the BBN period:

\[
\Delta N_{\text{eff}} = \frac{4}{7} \left( \frac{3 n_a}{2 n_\nu} \right)^{4/3},
\]

where \( n_a \) is given by Eq. (8) and \( n_\nu \) refers to the present neutrino plus antineutrino number density per flavour.

### C. Cosmological measurements

Our baseline data set consists of CMB measurements. These include the temperature data from the Planck satellite, see Refs. 18 [25], together with the WMAP 9-year polarization measurements, following [26]. We also consider high multipole data from the South Pole Telescope (SPT) [27] as well as from the Atacama Cosmology Telescope (ACT) [28]. The combination of all the above CMB data is referred to as the CMB data set.

Galaxy clusters represent an independent tool to probe the cosmological parameters. Cluster surveys usually report their measurements by means of the so-called cluster normalization condition, \( \sigma_8 \Omega_m^{1/2} \), where \( \gamma \approx 0.4 \) [29–31]. We shall use here the cluster normalization condition as measured by the Planck Sunyaev-Zeldovich (PSZ) 2013 catalogue [32], referring to it as the PSZ data set. The PSZ measurements of the cluster mass function provide the constraint \( \sigma_8(\Omega_m/0.27)^{0.3} = 0.764 \pm 0.025 \). As there exists a strong degeneracy between the value of the \( \sigma_8 \) parameter and the cluster mass bias, it is possible to fix the value of the bias parameter accordingly to the results arising from numerical simulations. In this last case, the error on the cluster normalization condition from the PSZ catalogue is considerably reduced: \( \sigma_8(\Omega_m/0.27)^{0.3} = 0.78 \pm 0.01 \). In our analyses, we shall consider the two PSZ measurements of the cluster normalization condition, to illustrate the impact of the cluster mass bias in the thermal axion mass bounds, as recently explored in Ref. [33] for the neutrino mass case.

Tomographic weak lensing surveys are sensitive to the overall amplitude of the matter power spectrum by measuring the correlations in the observed shape of distant galaxies induced by the intervening large scale structure. The matter power spectrum amplitude depends on both the \( \sigma_8 \) clustering parameter and the matter density \( \Omega_m \). Consequently, tomographic lensing surveys, via measurements of the galaxy power shear spectra, provide additional and independent constraints in the \( (\sigma_8, \Omega_m) \) plane. The Canada-France-Hawaii Telescope Lensing Survey, CFHTLenS, with six tomographic redshift bins (from \( z = 0.28 \) to \( z = 1.12 \)), provides a constraint on the relationship between \( \sigma_8 \) and \( \Omega_m \) of \( \sigma_8(\Omega_m/0.27)^{0.46} = 0.774 \pm 0.040 [34] \). We shall refer to this data set as CFHT.

We also address here the impact of a gaussian prior on the Hubble constant \( H_0 = 70.6 \pm 3.3 \text{ km/s/Mpc} \) from an independent reanalysis of Cepheid data [35], referring to this prior as the \( HST \) data set.

We have also included measurements of the large scale structure of the universe in their geometrical form, i.e., in the form of Baryon Acoustic Oscillations (BAO). These BAO wiggles, imprinted in the power spectrum of the galaxy distribution, result from the competition in the coupled photon-baryon fluid between radiation pressure and gravity. The BAO measurements that have been considered in our numerical analyses include the results from the WiggleZ [36], the 6dF [37] and the SDSS II surveys [38 [39], at redshifts of \( z = 0.44, 0.6, 0.73 \), \( z = 0.106 \) and \( z = 0.35 \), respectively. We also include in our analyses as well the Data Release 11 (DR11) of the BOSS experiment [40], which provides the most precise distant constraints [41] measuring both the Hubble parameter and the angular diameter distance at an effective redshift of 0.57.

### III. RESULTS

Table II depicts our results in the first scenario explored here, in which the axion mass is a free parameter and the PPS is described by the approach specified in Sec. II.A. Concerning CMB measurements only, the bounds on the thermal axion masses are largely relaxed in the case in which the PPS is not described by a simple power-law, as can be noticed after comparing the results depicted in Tab. II with those shown in Tab. III. This can be understood in terms of Fig. I which illustrates the degeneracies in the temperature anisotropies (left panel) and in the matter power spectrum (right panel) between the thermal axion mass and the PCHIP PPS. The left

| Parameter | Prior |
|-----------|-------|
| \( \Omega_b h^2 \) | [0.005, 0.1] |
| \( \Omega_{cdm} h^2 \) | [0.001, 0.09] |
| \( \Theta_s \) | [0.5, 10] |
| \( \tau \) | [0.01, 0.8] |
| \( m_a \) | [0.1, 3] |
| \( \sum m_\nu \) | [0.063] |
| \( P_{s,1}, \ldots, P_{s,12} \) | [0.01, 10] |
| \( n_{10} \) | [0.9, 1.1] |
| \( \log(10^{10} A_s) \) | [2.7, 4] |

TABLE I: Priors for the parameters used in the MCMC analyses.
\[
\begin{array}{ccccccccc}
\text{Parameter} & \text{CMB} & \text{CMB+HST} & \text{CMB+BAO} & \text{CMB+BAO} & \text{CMB+BAO} & \text{CMB+BAO} & \text{CMB+BAO} & \text{CMB+BAO} \\
& & & & +HST & HST-CFHT & +HST+PSZ (fixed bias) & +HST+PSZ \\
\Omega_m h^2 & 0.127^{+0.007}_{-0.007} & 0.122^{+0.006}_{-0.006} & 0.122^{+0.003}_{-0.003} & 0.121^{+0.003}_{-0.003} & 0.120^{+0.003}_{-0.003} & 0.118^{+0.002}_{-0.002} & 0.119^{+0.003}_{-0.004} \\
m_a [\text{eV}] & \text{Unconstrained} & < 1.31 & < 0.89 & < 0.91 & < 1.29 & 1.00^{+0.50}_{-0.48} & 0.93^{+0.70}_{-0.71} \\
\sigma_8 & 0.788^{+0.079}_{-0.086} & 0.821^{+0.052}_{-0.074} & 0.827^{+0.044}_{-0.057} & 0.825^{+0.045}_{-0.059} & 0.793^{+0.049}_{-0.058} & 0.760^{+0.023}_{-0.022} & 0.767^{+0.046}_{-0.044} \\
\Omega_m & 0.369^{+0.070}_{-0.065} & 0.314^{+0.045}_{-0.039} & 0.308^{+0.016}_{-0.015} & 0.304^{+0.016}_{-0.014} & 0.302^{+0.016}_{-0.015} & 0.304^{+0.016}_{-0.015} & 0.304^{+0.016}_{-0.016} \\
P_{s,1} & < 8.13 & < 8.17 & < 7.91 & < 8.06 & < 7.85 & < 8.09 & < 8.11 \\
P_{s,2} & 1.09^{+0.42}_{-0.35} & 1.04^{+0.41}_{-0.35} & 1.01^{+0.40}_{-0.32} & 0.99^{+0.42}_{-0.33} & 1.02^{+0.43}_{-0.34} & 1.01^{+0.42}_{-0.33} & 1.05^{+0.43}_{-0.38} \\
P_{s,3} & 0.68^{+0.39}_{-0.36} & 0.71^{+0.39}_{-0.39} & 0.71^{+0.39}_{-0.37} & 0.72^{+0.39}_{-0.38} & 0.69^{+0.39}_{-0.37} & 0.70^{+0.40}_{-0.38} & 0.69^{+0.40}_{-0.39} \\
P_{s,4} & 1.14^{+0.24}_{-0.22} & 1.15^{+0.24}_{-0.22} & 1.15^{+0.23}_{-0.21} & 1.15^{+0.23}_{-0.20} & 1.15^{+0.23}_{-0.21} & 1.15^{+0.23}_{-0.21} & 1.15^{+0.22}_{-0.21} \\
P_{s,5} & 1.02^{+0.11}_{-0.10} & 1.03^{+0.11}_{-0.11} & 1.00^{+0.11}_{-0.10} & 1.00^{+0.11}_{-0.10} & 0.99^{+0.11}_{-0.10} & 0.99^{+0.11}_{-0.10} & 0.99^{+0.11}_{-0.11} \\
P_{s,6} & 1.03^{+0.08}_{-0.07} & 1.06^{+0.08}_{-0.07} & 1.00^{+0.08}_{-0.07} & 1.00^{+0.08}_{-0.07} & 0.98^{+0.07}_{-0.06} & 0.98^{+0.07}_{-0.07} & 0.98^{+0.08}_{-0.07} \\
P_{s,7} & 0.99^{+0.07}_{-0.06} & 0.98^{+0.08}_{-0.07} & 0.98^{+0.07}_{-0.07} & 0.98^{+0.08}_{-0.07} & 0.96^{+0.07}_{-0.06} & 0.95^{+0.07}_{-0.06} & 0.96^{+0.07}_{-0.06} \\
P_{s,8} & 0.94^{+0.06}_{-0.06} & 0.95^{+0.08}_{-0.07} & 0.95^{+0.07}_{-0.06} & 0.95^{+0.08}_{-0.07} & 0.94^{+0.07}_{-0.06} & 0.94^{+0.07}_{-0.06} & 0.94^{+0.07}_{-0.06} \\
P_{s,9} & 0.92^{+0.06}_{-0.05} & 0.94^{+0.08}_{-0.06} & 0.94^{+0.07}_{-0.06} & 0.94^{+0.08}_{-0.06} & 0.93^{+0.07}_{-0.06} & 0.93^{+0.07}_{-0.06} & 0.94^{+0.07}_{-0.06} \\
P_{s,10} & 0.90^{+0.06}_{-0.06} & 0.91^{+0.08}_{-0.07} & 0.91^{+0.07}_{-0.06} & 0.91^{+0.08}_{-0.06} & 0.90^{+0.07}_{-0.06} & 0.90^{+0.07}_{-0.06} & 0.90^{+0.07}_{-0.07} \\
P_{s,11} & 1.25^{+0.30}_{-0.28} & 1.24^{+0.32}_{-0.31} & 1.23^{+0.31}_{-0.31} & 1.24^{+0.31}_{-0.31} & 1.22^{+0.30}_{-0.28} & 1.22^{+0.32}_{-0.28} & 1.23^{+0.31}_{-0.30} \\
P_{s,12} & \text{Unconstrained} & \text{Unconstrained} & \text{Unconstrained} & \text{Unconstrained} & \text{Unconstrained} & \text{Unconstrained} & \text{Unconstrained} \\
\end{array}
\]

TABLE II: 95% CL constraints on the physical cold dark matter density \( \Omega_m h^2 \), the axion mass \( m_a \) (in eV), the clustering parameter \( \sigma_8 \), the relative matter energy density \( \Omega_m \) and the \( P_{s,j} \) parameters for the PPS nodes from the different combinations of data sets explored here in the \( \Lambda \)CDM+\( m_a \) model, considering the PCHIP PPS modeling.

The addition to the CMB data of the HST prior on the Hubble constant provides a 95% CL upper limit on the thermal axion mass of 1.31 eV \(^1\), while the further addition of the BAO measurements brings this constraint down to 0.91 eV, as these last data sets are directly sensitive to the free-streaming nature of the thermal axion. Notice that these two 95% CL upper bounds are very similar to the ones obtained when considering the standard power-law power spectrum, which are 1.56 eV and 0.83 eV for the CMB+HST and CMB+HST+BAO data combinations, respectively.

Interestingly, when adding the CFHT bounds on the \( \sigma_8-\Omega_m \) relationship, the bounds on the thermal axion mass become weaker. The reason for that is due to the lower \( \sigma_8 \) values preferred by weak lensing measurements, values that can be achieved by allowing for higher ax-
ion masses. The larger the axion mass, the larger is the reduction of the matter power spectrum at small (i.e. cluster) scales, leading consequently to a smaller value of the clustering parameter $\sigma_8$.

If we instead consider now the PSZ data set with fixed cluster mass bias, together with the CMB, BAO and HST measurements, a non-zero value of the thermal axion mass of $\sim 1$ eV ($\sim 0.80$ eV) is favoured at $\sim 4\sigma$ ($\sim 3\sigma$) level, when considering the PCHIP (standard power-law) PPS approach. However, these results must be regarded as an illustration of what could be achieved with future cluster mass calibrations, as the Planck collaboration has recently shown in their analyses of the 2015 Planck cluster catalogue [33]. When more realistic approaches for the cluster mass bias are used, the errors on the so-called cluster normalization condition are larger, and, consequently, the preference for a non-zero axion mass of 1 eV is only mild in the PCHIP PPS case, while in the case of a standard power-law PPS such an evidence

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2 A similar effect when considering PSZ data for constraining either thermal axion or neutrino masses has also been found in Refs. [34,35,36].
completely disappears.

Figure 2 (left panel) shows the 68% and 95% CL allowed regions in the \((m_a, \Omega_m h^2)\) plane for some of the possible data combinations explored in this study, and assuming the PCHIP PPS modeling. Notice that, when adding BAO measurements, lower values of the physical cold dark matter density are preferred. This is due to the fact that large scale structure allows for lower axion masses than CMB data alone. The lower is the thermal axion mass, the lower is the amount of hot dark matter and consequently the lower should be the cold dark matter component. This effect is clear from the results shown in Tab. II and Tab. III, where the values of the physical cold dark matter density \(\Omega_c h^2\) and of the relative current matter density \(\Omega_m\) arising from our numerical fits are shown, for the different data combinations considered here.

The right panel of Fig. 2 shows the 68% and 95% CL allowed regions in the \((m_a, \sigma_8)\) plane in the PCHIP PPS scenario. The lower values of the \(\sigma_8\) clustering parameter preferred by PSZ data (see the results shown in Tab. I and Tab. II) are translated into a preference for non-zero thermal axion masses. Larger values of \(m_a\) will enhance the matter power spectrum suppression at scales below the axion free-streaming scale, leading to smaller values of the \(\sigma_8\) clustering parameter, as preferred by PSZ measurements. The evidence for non-zero axion masses is more significant when fixing the cluster mass bias in the PSZ data analyses.

Figure 3 shows the equivalent to Fig. 2 but for a standard power-law PPS. Notice that, except for the case in which CMB measurements are considered alone, the thermal axion mass constraints do not change significantly, if they are compared to the PCHIP PPS modeling. This fact clearly states the robustness of the cosmological bounds on thermal axion masses and it is applicable to the remaining cosmological parameters, see Tabs. I and III. Note that, for the standard case of a power-law PPS, the preference for non-zero axion masses appears only when considering the (unrealistic) PSZ analysis with a fixed cluster mass bias. When more realistic PSZ measurements of the cluster normalization condition are exploited, there is no preference for a non-zero thermal axion mass.

Besides the results concerning the thermal axion mass and the standard ΛCDM parameters, we also obtain constraints on the form of the PPS when modeled accordingly to the PCHIP scenario. The 95% CL limits for the \(P_{s,j}\) parameters are shown in Tab. I while an example of the reconstructed PPS is given in Fig. 3, where we show the 68%, 95% and 99% CL allowed regions arising from a fit to CMB data of the PCHIP PPS scale dependence, in the context of a ΛCDM+\(m_a\) model. We do not show the corresponding figures obtained from all the other data combinations since they are equivalent to Fig. 3 as one can infer from the very small differences in the 95% CL allowed ranges for the \(P_{s,j}\) parameters arising from different data sets, see Tab. I. Note that both \(P_{s,1}\) and \(P_{s,12}\) are poorly constrained at this confidence level: the reason for that is the absence of measurements at their corresponding wavenumbers. All the remaining \(P_{s,j}\), with \(j = 2, \ldots, 11\) are well-constrained. In particular, in the range between \(k_5\) and \(k_{10}\) (see Eq. (3)), the \(P_{s,j}\) are determined with few percent accuracy. Indeed, in the range covered between these nodes, the PPS does not present features and can be perfectly described by a power-law parametrization. Among the interesting features outside the former range, we can notice in Fig. 3 a significant dip at wavenumbers around \(k = 0.002\) Mpc\(^{-1}\), that comes from the dip at \(\ell = 20 – 30\) in the CMB temperature power spectrum and a small bump around \(k = 0.0035\) Mpc\(^{-1}\), corresponding to the increase at \(\ell \approx 40\). These features have been obtained in previous works [9, 14, 15] using different methods and data sets. In addition, we obtain an increase of power at \(k \approx 0.2\) Mpc\(^{-1}\), necessary to compensate the effects of the thermal axion mass in both the temperature anisotropies and the large scale structure of the universe.

IV. CONCLUSIONS

Axions provide the most elegant scenario to solve the strong CP problem, and may be produced in the early universe via both thermal and non-thermal processes. While non thermal axions are highly promising cold dark matter candidates, their thermal companions will contribute to the hot dark matter component of the universe, together with the (light) three active neutrinos of the standard model of elementary particles. Therefore, the cosmological consequences of light massive thermal axions are very much alike those associated with neutrinos, as axions also have a free-streaming nature, suppressing structure formation at small scales. Furthermore, these light thermal axions will also contribute to the dark radiation background, leading to deviations of the relativistic degrees of freedom \(N_{\text{eff}}\) from its canonically expected value of \(N_{\text{eff}} = 3.046\). Based on these signatures, several studies have been carried out in the literature deriving bounds on the thermal axion mass [3, 8].

Nevertheless, these previous constraints assumed that the underlying primordial perturbation power spectrum follows the usual power-law description governed, in its most economical form, by an amplitude and a scalar spectral index. Here we have relaxed such an assumption, in order to test the robustness of the cosmological axion mass bounds. Using an alternative, non-parametric description of the primordial power spectrum of the scalar perturbations, named PCHIP and introduced in Ref. [9], we have shown that, in practice, when combining CMB measurements with low redshift cosmological probes, the axion mass constraints are only mildly sensitive to the primordial power spectrum choice and therefore are not strongly dependent on the particular details of the underlying inflationary model. These results agree with the findings of Ref. [19] for the neutrino mass case. The tight-
est bound we find in the PCHIP primordial power spectrum approach is obtained when considering BAO measurements together with CMB data, with $m_a < 0.89$ eV at 95% CL. In the standard power-law primordial power spectrum modeling, the tightest bound is $m_a < 0.83$ eV at 95% CL, obtained when combining BAO, CMB and HST measurements. Notice that these bounds are very similar, confirming the robustness of the cosmological axion mass measurements versus the primordial power spectrum modeling.

Interestingly, both weak lensing measurements and cluster number counts weaken the thermal axion mass bounds. The reason for that is due to the lower $\sigma_8$ values preferred by these measurements, which could be generated by a larger axion mass. More concretely, Planck cluster measurements provide a measurement of the so-
called cluster normalization condition, which establishes a relationship between the clustering parameter $\sigma_8$ and the current matter mass-energy density $\Omega_m$. However, the errors on this relationship depend crucially on the knowledge of the cluster mass bias. A conservative approach for the cluster mass calibration results in a mild (zero) evidence for a non-zero axion mass of $1\,\text{eV}$ in the \textsc{pchip} (power-law) PPS case. We also illustrate a case in which the cluster mass bias is fixed, to forecast the expected results from future cosmological measurements. In this case, a non-zero value of the thermal axion mass of $\sim 1\,\text{eV}$ ($\sim 0.80\,\text{eV}$) is favoured at $\sim 4\sigma$ ($\sim 3\sigma$) level, when considering the \textsc{pchip} (power-law) PPS approach. Precise cluster mass calibration measurements are therefore mandatory, as the cluster mass bias is highly correlated with the clustering parameter $\sigma_8$, which, in turn, is highly affected by the free-streaming nature of a hot dark matter component, as thermal axions.

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[1] R. D. Peccei and H. R. Quinn, Phys. Rev. Lett. \textbf{38}, 1440 (1977); R. D. Peccei and H. R. Quinn, Phys. Rev. D \textbf{16}, 1791 (1977).
[2] G. Mangano, G. Miele, S. Pastor, T. Pinto, O. Pisanti and P. D. Serpico, Nucl. Phys. B \textbf{729}, 221 (2005) [hep-ph/0506164].
[3] A. Melchiorri, O. Mena and A. Slosar, Phys. Rev. D \textbf{76}, 041303 (2007) [arXiv:0705.2695 [astro-ph]].
[4] S. Hannestad, A. Mirizzi, G. G. Raffelt and Y. Y. Y. Wong, JCAP \textbf{0708}, 015 (2007) [arXiv:0706.4198 [astro-ph]].
[5] S. Hannestad, A. Mirizzi, G. G. Raffelt and Y. Y. Y. Wong, JCAP \textbf{0804}, 019 (2008) [arXiv:0803.1585 [astro-ph]].
[6] S. Hannestad, A. Mirizzi, G. G. Raffelt and Y. Y. Y. Wong, JCAP \textbf{1008}, 001 (2010) arXiv:1004.0695 [astro-ph.CO].
[7] M. Archidiacono, S. Hannestad, A. Mirizzi, G. Raffelt and Y. Y. Y. Wong, JCAP \textbf{1310}, 020 (2013) [arXiv:1307.0615 [astro-ph.CO]].
[8] E. Giusarma, E. Di Valentino, M. Lattanzi, A. Melchiorri and O. Mena, Phys. Rev. D \textbf{90}, 043507 (2014) [arXiv:1403.4852 [astro-ph.CO]].
[9] S. Gariazzo, C. Giunti and M. Laveder, arXiv:1412.7405 [astro-ph.CO].
[10] D. H. Lyth and A. Riotto, Phys. Rept. \textbf{314}, 1 (1999)
