

Implications of cosmological observables for particle physics: an overview

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Abstract. I review how precision data from observations of the cosmic microwave background anisotropies and the large-scale structure distribution can be used to probe particle physics. Some examples are the absolute neutrino mass scale, dark radiation, light sterile neutrinos, QCD axions, WIMP annihilation, and dark sector interactions.

1. Introduction: precision cosmological observations

The past two decades have witnessed tremendous advances in precision observations of the high-redshift universe, a progress that has finally enabled physical cosmology to be studied as an exact science. At the forefront of these observations are measurements of the cosmic microwave background (CMB) temperature and polarisation anisotropies, accomplished on large angular scales by space-based experiments such as COBE, WMAP, and most recently the Planck mission [1], and on small angular scales by an array of balloon and ground-based instruments.

The spatial distribution of matter in the universe on the largest scales has likewise been mapped out with great precision. The SDSS-III BOSS survey, for example, has measured the redshifts of 1.5 million luminous red galaxies out to a redshift of $z = 0.7$ over some 10,000 square degrees (about one third of the sky). These measurements have enabled the extraction of the baryon acoustic oscillations (BAO) scale, which is a powerful geometrical probe of the energy/matter content and geometry of the late-time universe [2].

Currently, the simplest cosmological model able to accommodate all available observational data is the concordance flat Λ CDM model. At leading order spacetime in this model is described by the Friedmann–Lemaître–Robertson–Walker (FLRW) metric with a flat spatial geometry, and the initial conditions are consistent with those of the simplest, single-field inflation models. The model’s energy content is usually reported in terms of each component’s present-day fractional contribution to the total *energy density*, i.e., 68% dark energy in the form of a cosmological constant, 27% cold dark matter (CDM), and 5% baryons, where the accompanying uncertainties are typically at the percent level [3]. In terms of the present-day *particle number densities* however, by far the dominant constituents are photons and three families of nearly massless neutrinos, both of which also dominated the universe’s energy density prior to CMB decoupling.

In view of the precision of these measurements, it is interesting to ponder what “non-standard” scenarios might be constrained by cosmological observations, and how well that goal might be achieved. Well-known examples of such investigations include massive neutrinos, and a non-canonical number of neutrino families; more recent explorations have also considered the



possibility of dark sector interactions such as elastic scattering, annihilation, and decay. While these scenarios usually go under the name of “non-standard” cosmologies, it is important to stress that many in fact arise quite naturally from the particle physics perspective, and cosmological constraints on such scenarios are often complementary to the reach of other experimental probes. It is in this sense that cosmological observations can be a powerful probe of particle physics.

In this article, I describe how cosmological observations can be used to probe particle physics.

2. Particle physics in cosmology

The concordance Λ CDM model has a well-defined particle content. The energy densities of its individual components have also been determined to great precision [3]. But what enables this determination in the first place? To answer this question, we note firstly that a modern-day cosmological observation typically maps out the distribution of an observable quantity (e.g., temperature and polarisation of the CMB photons, the number density of galaxies and clusters, etc.) as a function of redshift and/or angular position in the sky. To extract any meaningful information about the universe from this collection of data points requires that we have

- (i) A well-founded model of cosmology, i.e., a particle physics model embedded in an FLRW universe extended with inflation,
- (ii) A set of free model parameters associated with the cosmological model (e.g., particle number densities, particle masses, spectral index and amplitude of the initial fluctuations, etc.), and
- (iii) A theoretical framework—in this case, the Boltzmann equation in conjunction with general relativity—to compute and predict the observable quantities given (i) and (ii).

It is by comparing the predictions resulting from (iii) against observational data and applying the statistical techniques of parameter inference and model comparison that we learn about the properties of the universe.

The particle physics content of the cosmological model can influence these predictions in many interesting ways. At the level of homogeneous cosmology, the kinematic properties (i.e., ultrarelativistic, nonrelativistic, or somewhere in between) of a particle species strongly influence the time evolution of its energy density ρ . For example, a fully nonrelativistic species has $\rho(t) \propto a^{-3}(t)$, where $a(t)$ is the scale factor as a function of cosmic time t , while in the ultrarelativistic limit one expects $\rho(t) \propto a^{-4}(t)$. Because gravity couples universally to all forms of energy, the exact admixture of relativistic and nonrelativistic (and other) particle species has a strong impact on the expansion history of the universe, as described by the Friedmann equation, $H^2(t) = (8\pi G/3) \sum_i \rho_i(t)$, where the index i counts all forms of energy in the universe.

On the level of the inhomogeneities, both the kinematic properties and the non-gravitational interactions of a particle species can affect its response to perturbations in the spacetime metric (or, loosely speaking, the gravitational potentials). This can be understood in terms of the Boltzmann equation, which governs the evolution of a species’ phase space density $f(x, p, t)$:

$$\frac{\partial f}{\partial t} + \frac{d\mathbf{x}}{dt} \cdot \frac{\partial f}{\partial \mathbf{x}} + \frac{d\mathbf{p}}{dt} \cdot \frac{\partial f}{\partial \mathbf{p}} = C[f]. \quad (1)$$

On the l.h.s., the $d\mathbf{x}/dt$ term encodes the velocity of the phase space element at $\{\mathbf{x}, \mathbf{p}\} \rightarrow \{\mathbf{x} + d\mathbf{x}, \mathbf{p} + d\mathbf{p}\}$, while the $d\mathbf{p}/dt$ term is linked to the metric perturbations via the geodesic equation. On the r.h.s., the quantity $C[f]$ is called the collision integral, which incorporates non-gravitational particle interactions. Equation (1) therefore shows that the evolution of spatial inhomogeneities in a particle species depends strongly on what kind of particle it is.

Furthermore, these inhomogeneities feed back into the metric perturbations via the Einstein equation $G_{\mu\nu} = 8\pi G T_{\mu\nu}$, where $G_{\mu\nu}$ is the Einstein tensor and $T_{\mu\nu}$ is the stress–energy tensor encoding the matter/energy. This means that if we were to change the properties of one particle

species, e.g., make it interact more or less strongly, or give it a larger or smaller velocity dispersion, this change could in principle be felt in all other particle species present in the universe. How the photons and the baryons feel this change is particularly important, because they are ultimately what we observe as the CMB and the galaxies respectively. It is thus in this sense that cosmological observations can be sensitive to particle physics.

3. Λ CDM and beyond

Let us now consider in detail the particle physics of the Λ CDM model. Since only those particle species that have survived in some abundance down to roughly the keV temperature scale would impact on CMB and large-scale structure observables, particle physics of the standard model immediately stipulates the relevant particle content and interactions:

- (i) Photons: The photon number and energy densities have been established to 0.1% accuracy by the COBE FIRAS measurement of the CMB temperature and energy spectrum [4].
- (ii) Baryons: These are mainly in the form of hydrogen atoms (and 25% helium-4 by mass), and may be ionised or neutral depending on the abundance and energy of any ionising radiation that might be in the universe. In standard Λ CDM the CMB photons themselves play the role of the ionising radiation at temperatures above the eV scale. Furthermore, CMB photons can Compton-scatter off free electrons that result from the ionisation. The precise abundance of baryons relative to photons is not predicted by the standard model, and constitutes one of the free fitting parameters of Λ CDM.
- (iii) Neutrinos: Three flavours of light neutrinos are expected to be present, whereby the number density per flavour is fixed to 3/11 times the photon number density by standard model particle physics (with a minor, 1% correction in a detailed calculation). Although the discovery of neutrino flavour oscillations implies that some neutrinos have a minimum mass of 0.05 eV [5], in the cosmological context light neutrinos are conventionally assumed to be exactly massless. Standard weak interaction processes go out of equilibrium long before the CMB epoch; thus neutrinos are taken to be exactly non-interacting in the Λ CDM model.
- (iv) Cold dark matter (CDM): While clearly not part of the standard model of particle physics, we have nonetheless some hints from astrophysical observations that the dark matter should be nonrelativistic and should not interact with photons. The standard CDM paradigm further assumes a negligible velocity dispersion and the complete absence of non-gravitational interactions. The energy density residing in CDM is *a priori* unknown, and is thus one of the free fitting parameters of the Λ CDM model.

Particle physics extensions to Λ CDM can take the form of expanding the particle content of the model, and/or relaxing the assumptions about the kinematic and interaction properties of the existing particle content. I outline some of these possibilities below, as well as current observational constraints on these scenarios.

3.1. Neutrino mass

The discovery of neutrino oscillations has established with certainty that some neutrinos have a minimum mass of 0.05 eV. Comparing this number with the expected neutrino temperature today ($T_\nu = 1.95 \text{ K} \sim 10^{-4} \text{ eV}$), we see that cosmological neutrinos, although born ultrarelativistic, can become nonrelativistic as the universe cools down, and hence contribute to the matter content with an energy density $\Omega_\nu h^2 = \sum m_\nu / (94 \text{ eV})$ today.

If the neutrino mass is of order 0.1 eV, then the relativistic-to-nonrelativistic transition occurs close to the epoch of CMB decoupling. This leaves an interesting imprint on the CMB temperature anisotropies via potential decay induced by the transition. Furthermore, even after they have become nonrelativistic at low redshifts, the neutrinos still possess a large thermal

velocity dispersion. This dispersion causes neutrino clustering to be very inefficient on scales smaller than the instantaneous free-streaming scale λ_{FS} at redshift z ,

$$k_{\text{FS}} = \frac{2\pi}{\lambda_{\text{FS}}} \simeq 2.4 \sqrt{\frac{\Omega_m}{1+z}} \left(\frac{m_\nu}{\text{eV}} \right) h \text{ Mpc}^{-1}. \quad (2)$$

This inefficiency in neutrino clustering feeds back into the evolution of the metric perturbations and hence the growth of the dark matter and the baryon density perturbations. Thus, if we replace part of the CDM energy density with massive neutrinos, the end result is that the present-day matter power spectrum will show a suppression at wavenumbers larger than roughly

$$k_{\text{FS,min}} \equiv k_{\text{FS}}(z_{\text{nr}}) \simeq 0.03 \Omega_m^{1/2} \left(\frac{m_\nu}{\text{eV}} \right)^{1/2} h \text{ Mpc}^{-1}, \quad (3)$$

where z_{nr} is the redshift at which the neutrinos first become nonrelativistic. The fractional suppression is asymptotically $\Delta P/P \sim 8f_\nu$, where $f_\nu \equiv \Omega_\nu/\Omega_m$. Note that this relation implies that the suppression is sensitive primarily to the sum of the neutrino masses $\sum m_\nu$, less so to the individual masses themselves which affect only $k_{\text{FS,min}}$.

Using a combination of Planck CMB and BOSS BAO measurements, it is currently possible to constrain the sum of neutrino masses via these effects down to $\sum m_\nu < 0.23 \text{ eV}$ (95% C.I.) if we extend the Λ CDM model only with neutrino masses. The bound can deteriorate by up to 50% if assumptions about, e.g., spatial flatness, the nature of dark energy, etc., are relaxed. Getting away from the $0.2 \rightarrow 0.3 \text{ eV}$ ballpark is however difficult without some drastic revision of the basic cosmological framework. Previous attempts include breaking the scale-invariance of the initial conditions *and* the Copernican principle (using so-called “voids”) [6]. However, even such a nontrivial scenario has been found to be strongly constrained by observations of the kinetic Sunyaev–Zel’dovich effect [7, 8]. It is thus fair to conclude that neutrino mass limits from cosmology can be considered robust with respect to reasonable modifications of the Λ CDM model. See [9] for further discussions on model dependence.

It is interesting to observe that formally the Planck $\sum m_\nu$ bound is similar to those obtained in the WMAP era using CMB and galaxy power spectrum measurements. The post-Planck bounds are however arguably more robust for the following reasons: Firstly, with the detection of the lensing signal, Planck CMB measurements alone already yield $\sum m_\nu < 0.59 \text{ eV}$ (95% C.I.); this is good news, because CMB observables can be reliably calculated using linear perturbation theory. Secondly, the BAO features of the galaxy power spectrum suffer from milder nonlinear effects than its broad-band shape. Of course, if we are willing to go fully nonlinear, then as always the most aggressive constraint comes from using $\text{Ly}\alpha$ forest probes of the small-scale power spectrum amplitude: $\sum m_\nu < 0.15 \text{ eV}$ (95% C.I.) [10].

3.2. Dark radiation

The possible existence of extra relativistic particle species is often dealt with under the heading of “additional neutrinos”: the energy density of the new relativistic degrees of freedom ρ_x is added to that due to three families of neutrinos ρ_ν , and then parameterised with the parameter N_{eff} , defined via $\rho_x + \rho_\nu = N_{\text{eff}} \rho_\nu^0 = (3.046 + \Delta N_{\text{eff}}) \rho_\nu^0$, where ρ_ν^0 denotes the energy density residing in one species of thermalised effectively massless neutrinos with temperature $T_\nu = (4/11)^{1/3} T_\gamma$. Note that the Λ CDM value for N_{eff} is not exactly three: the 0.046 correction comes from taking into account QED finite temperature corrections, neutrino oscillations, and non-instantaneous neutrino decoupling. In terms of particle physics models, any thermalised particle species that decouples while ultrarelativistic and *remains* ultrarelativistic and non-interacting through the CMB epoch can be described by this simple extension of Λ CDM. Decay products of heavy particles can in principle also contribute to N_{eff} [11].

Deviation of N_{eff} from 3.046 has a non-trivial effect on the CMB temperature anisotropies and the matter power spectrum. First and foremost, increasing N_{eff} delays the matter–radiation equality epoch, which in turn contributes to the Integrated Sachs–Wolfe effect, measurable in the CMB temperature power spectrum through the odd-peak ratios. However, this effect can be offset by increasing simultaneously the matter density ω_m (i.e., N_{eff} and ω_m are degenerate parameters), and is thus not sufficient to pin down N_{eff} . Nonetheless, changing ω_m also alters the comoving distance D_* to the last scattering surface, which, along with the comoving sound horizon at CMB decoupling r_s , is the physical quantity that fixes the approximate locations of the CMB acoustic peaks. But even this is not sufficient information, because D_* can be tuned by adjusting the Hubble parameter h . Thus, what we have established from this discussion so far is that from the CMB acoustic peaks alone, there is a three-way parameter degeneracy between N_{eff} , ω_m , and h , and some orthogonal information is required to break it.

Obviously, a direct measurement of h in our local neighbourhood would do the trick. Alternatively, one could measure another distinct physical length scale, since all length scale measures in cosmology are sensitive to the Hubble parameter. The CMB temperature power spectrum itself offers such a possibility in its damping tail, where the damping scale k_d is determined by the diffusion of photons prior to CMB decoupling [12]; the damping tail was first measured by ACT [13] and SPT [14], and is now also probed by Planck. The equivalent scale can also be measured in the matter power spectrum, via the suppression of small-scale power due to diffusion damping of the baryon density perturbations [15].

Current constraints from the Planck mission on N_{eff} are remarkably consistent with its standard value; the combination of Planck CMB and BOSS BAO measurements yields $N_{\text{eff}} = 3.04 \pm 0.18$ (68% C.I.) [3]. It is however interesting to note that the Planck-inferred value of the Hubble parameter $H_0 = 100 h = 67.8 \pm 0.9 \text{ km s}^{-1} \text{ Mpc}^{-1}$ (68 % C.I.) is incompatible at $\sim 2\sigma$ with many direct measurements in the local neighbourhood, notably, the so-called HST value of $H_0 \equiv 100 h = 74.8 \pm 3.1 \text{ km s}^{-1} \text{ Mpc}^{-1}$ (68 % C.I.) [16]. In view of the parameter degeneracy discussed above, it is conceivable that a non-standard N_{eff} could ameliorate the incompatibility. Indeed, a combined fit of Planck CMB, BAO, and local measurements of H_0 finds $N_{\text{eff}} = 3.52_{-0.45}^{+0.48}$ (95% C.I.) [17]. However, one should be cautious when combining local and global measurements, since the former may suffer from fluctuations in our local neighbourhood. Thus it is perhaps prudent not to read too much into this 2σ discrepancy.

3.3. *eV-mass sterile neutrinos*

While the “dark radiation” prescription outlined in section 3.2 describes strictly the case of ultrarelativistic particle species, there are many well-motivated particle physics scenarios in which the light particle species does not remain ultrarelativistic throughout its entire evolution history. The phenomenology of these scenarios is essentially an amalgam of descriptions given above in section 3.1 and 3.2. I describe first in this section the case of an eV-mass sterile neutrino, and in section 3.4 the case of a QCD axion.

Several short-baseline (SBL) neutrino oscillations experiments have produced results that are inconsistent with the standard three-neutrino interpretation of the global neutrino oscillation data [18]. First amongst them is the LSND anomaly which saw the appearance of $\bar{\nu}_e$ in a $\bar{\nu}_\mu$ beam. This is followed by the more recent MiniBooNE ν_e excess, and the re-evaluation of the reactor neutrino fluxes which points to the disappearance of $\bar{\nu}_e$ on a distance of tens to hundreds of metres. If these anomalies are to be explained in terms of neutrino oscillations, then a fourth, sterile neutrino state must be introduced, which mix with the standard model neutrinos with mixing parameters in the ball park $\Delta m^2 \sim 1 \text{ eV}^2$ and $\sin^2 2\theta \sim 3 \times 10^{-3}$ [19].

Such a light sterile neutrino has important consequences for cosmology. Within the standard cosmological setting, this mixing and the ensuing flavour oscillations inevitably lead to full thermalisation of the sterile neutrinos to the same temperature as the standard model neutrinos,

and hence an excess of relativistic energy density $\Delta N_{\text{eff}} \simeq 1$ at early times [20]. Furthermore, because of their mass is “large”, $m_s > \sqrt{\Delta m^2} \sim 1$ eV, these sterile neutrinos become nonrelativistic at late times, and are then subject to the same sort of mass constraint discussed in section 3.1. In fact, extending the Λ CDM model with a free N_{eff} and m_s , the Planck collaboration finds the combined constraints $N_{\text{eff}} = 3.2 \pm 0.5$ (95% C.I.) and $m_s < 3.2$ eV (95% C.I.) [3], which are clearly incompatible with expectations for the SBL sterile neutrino.

Is there a way out? Assuming that the sterile neutrino interpretation of the SBL anomalies is correct, one possibility is to find a mechanism to suppress sterile neutrino thermalisation in the early universe. Some avenues appear promising, e.g., large lepton asymmetries [21], or hidden sterile neutrino self-interactions [22, 23]. Both mechanisms generate a new matter effect for active–sterile neutrino oscillations that suppresses the effective mixing angle and hence the sterile neutrino production rate; both, however, require new beyond-the-standard-model physics. Thus the question of the light sterile neutrino is still an open one, and the next generation of SBL oscillation experiments would be of crucial importance not only for neutrino physics.

3.4. *meV- to eV-mass axions*

The Peccei–Quinn solution of the strong CP problem predicts the existence of axions, low-mass pseudoscalars that are very similar to neutral pions, except that their mass and interaction strengths are suppressed relative to the pion case by a factor of order f_π/f_a , where $f_\pi \simeq 93$ MeV is the pion decay constant, and f_a is a large energy scale known as the axion decay constant or the Peccei–Quinn scale. This scale is related to axion mass via $m_a \simeq 6.0 (10^6 \text{ GeV}/f_a)$ eV.

Axions in the μeV mass range can serve as a CDM candidate. For heavier masses, however, a sizeable axion population can be produced by thermal processes. In particular, for $m_a > 60$ eV, axion scattering processes remain in equilibrium through the QCD phase transition at $T \sim 200$ MeV. When these axions finally decouple, they engender a relic axion background with number density $n_a = (1/2)(3.91/g_{*s}(T_{a,\text{dec}}))n_\gamma$ and temperature $T_a = (3.91/g_{*s}(T_{a,\text{dec}}))^{1/3}T_\gamma$, where $g_{*s}(T_{a,\text{dec}})$ denotes the effective entropy degrees of freedom at the time of axion decoupling, and n_γ and T_γ are the CMB photon number density and temperature respectively.

The thermalisation of hadronic axions via $a + \pi \leftrightarrow \pi + \pi$ and its decoupling have been calculated in [24]. The consequence of this thermal axion population on cosmological observations have also been explored in a series of works, most recently [25, 26]. The broad-brush phenomenology is similar to the case of eV-mass sterile neutrinos; how cosmological observations actually constrain axion physics and neutrino physics differ in the details [25]. Suffice it to say here that because the axion background is generally colder than the neutrino background, the potential decay induced by the axions’ relativistic-to-nonrelativistic transition occurs at too small a length scale to be observable in pre-Planck CMB data; Planck’s measurement of seven acoustic peaks constrains for the first time the axions mass to $m_a < 1.01$ eV (95% C.I.) using CMB data alone [25]. Adding the SDSS halo power spectrum improves the constraint to 0.86 eV.

This situation is of interest for solar axion searches by, e.g., the CAST experiment, which exploit the $a\gamma\gamma$ -interaction vertex both as a source of axions in the sun and for the back-conversion of axions into X-rays in a dipole magnet oriented towards the sun. Currently, CAST is able to exclude a small chunk of the QCD axion parameter space in m_a and $g_{a\gamma\gamma}$, where $g_{a\gamma\gamma}$ is the $a\gamma\gamma$ -coupling strength, around $m_a \sim 1$ eV [27]. In the future the proposed IAXO experiment will improve the sensitivity to $m_a < 0.2$ eV [28], while the ESA Euclid mission will easily pin down a relic axion population if the axion mass exceeds 0.15 eV [29]. Thus axion physics and cosmology will continue to go hand in hand for some time to come.

3.5. *WIMP annihilation*

The Weakly Interacting Massive Particle (WIMP) is a generic CDM candidate with a mass $m_\chi \simeq 10 \rightarrow 1000$ GeV, and is produced in the early universe via pair production from standard

model particles. An immediate consequence is that the WIMP must annihilate into standard model particles some time in the evolution history of the universe, which is the basis of indirect dark matter searches at high energy cosmic ray, gamma ray, and neutrino observatories [30, 31].

If annihilation happens frequently during the CMB epoch, then the energy injected into the universe can delay the hydrogen recombination process, and hence affect the free electron number density and subsequently CMB decoupling [32]. The energy injection rate per unit volume is

$$\frac{dE}{dt dV} = \rho_{\text{crit}}^2 \Omega_m^2 (1+z)^6 p_{\text{ann}}, \quad (4)$$

where $p_{\text{ann}} \equiv f(z)\langle\sigma v\rangle/m_\chi$ incorporates the annihilation cross-section $\langle\sigma v\rangle$, and $f(z)$ is an efficiency factor that determines what fraction of the energy eventually goes into heating and ionising the photon–baryon plasma at redshift z ; $f(z) = 0.01 \rightarrow 1$ for typical WIMP models.

Assuming a constant $f(z) = f_{\text{eff}}$, the Planck collaboration has placed a constraint of $p_{\text{ann}} < 4 \times 10^{-28} \text{ cm}^3 \text{ s}^{-1} \text{ GeV}^{-1}$ (95% C.I.) [3]. This can be translated into a constraint in the 2D $(f_{\text{eff}}\langle\sigma v\rangle, m_\chi)$ -parameter space. Interestingly, for certain WIMP models (as determined by the f_{eff} parameter), the constrained parameter space happens to coincide with the regions favoured by the purported annihilation signals of Fermi and AMS-02. See [33] for details.

3.6. Dark sector interactions

Extensions of Λ CDM to include dark sector interactions may appear at first glance exotic. However, such scenarios are in fact well-motivated from the particle physics perspective.

Dark matter elastic scattering: Interaction between the dark matter (DM) particle with standard model particles forms the basis of DM production and detection. While in the popular WIMP scenario freeze-out occurs at high temperatures (above the GeV scale), elastic scattering processes can continue to couple the WIMP to the standard model particle bath until a much later time. For the SUSY neutralino, for example, kinetic decoupling can occur at a temperature as low as $O(10)$ MeV, e.g., [34]. While this decoupling temperature is still too high to impact on CMB and large-scale structure observations, it is nonetheless instructive to ask, what is the lowest DM kinetic decoupling temperature that is compatible with observations.

The most relevant processes here are DM–neutrino and DM–photon scattering, since only neutrinos and photons survive in large abundances through the CMB epoch. In both cases the phenomenology is largely analogous to that of the photon–baryon system: elastic scattering of the DM with the ultrarelativistic particle species causes the two to form a tightly coupled fluid which exhibit acoustic oscillations on scales comparable to the Hubble length. Furthermore, because the coupling is never perfect, i.e., the ultrarelativistic species free-streams between scattering, diffusion wipes out the DM perturbations on scales below the diffusion length scale.

This scenario has been explored recently in, e.g., [35, 36], which find interesting constraints on the DM–neutrino and DM–photon cross-sections.

Neutrino self-interaction: Ultrarelativistic and non-interacting neutrinos free-stream. In the presence of an inhomogeneous spacetime, this free-streaming leads to an anisotropic stress in the neutrino fluid. If however we endow the neutrinos with a self-interaction that occurs sufficiently frequently, then the fluid can attain a state of local thermodynamic equilibrium, in which case only isotropic stress is present [37]. Thus, the absence of neutrino anisotropic stress can be taken as a sign of neutrino self-interaction. From the particle physics perspective, majoron-type neutrino mass models can give rise to such a self-interaction mediated by a light scalar.

The full details of how to incorporate neutrino self-interactions into the equations of motion for the neutrino inhomogeneities can be found in [38, 39], which improves upon earlier treatments in, e.g., [40]. Interesting constraints have been obtained in [38] on an effective 4-fermion self-coupling using the Planck data, and in [41] for mediation by a massless scalar. Curiously, the

former analysis also finds a second “peak” in the likelihood function corresponding to an effective coupling of $G_{\text{eff}} \simeq 8.6 \times 10^8 G_F$, where G_F is the Fermi constant. It would be interesting to see if this small preference for non-free-streaming neutrinos would survive further scrutiny.

4. Conclusions

There are many particle physics motivations for challenging the underlying assumptions of the flat concordance Λ CDM model, especially those assumptions we make in the dark matter and the neutrino sectors. In this article I have described several examples of how the Λ CDM model can be extended in the neutrino and dark matter sectors, and how these extensions can be probed using precision cosmological observations. Interesting constraints have already arisen from the current generation of CMB and large-scale structure observations. We expect further improvements with the advent of future probes such as the ESA Euclid mission [42].

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