ФИЗИКА ЭЛЕМЕНТАРНЫХ ЧАСТИЦ И АТОМНОГО ЯДРА
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LIGHT NEUTRINOS IN COSMOLOGY

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Ed. Instituts d'Investigació, València, Spain

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Neutrinos can play an important role in the evolution of the Universe, modifying some of the cosmological observables. We describe how the precision of present cosmological data can be used to learn about neutrino properties, in particular, their mass. We show how the analysis of current cosmological observations provides an upper bound on the sum of neutrino masses, with improved sensitivity from future cosmological measurements.

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INTRODUCTION

In this contribution I summarize the topics discussed in my two lectures at the School on Neutrino Cosmology, one of the best examples of the very close ties that have developed between nuclear physics, particle physics, astrophysics and cosmology. I tried to present the most interesting aspects, but many others that were left out can be found in the reviews [1–3] and, in particular, in [4]. A more general review on the connection between particle physics and cosmology can be found in [5].

We begin with a description of the properties and evolution of the background of relic neutrinos that fills the Universe. Then we review the influence of neutrinos on Primordial Nucleosynthesis. The largest part of this contribution is devoted to the impact of massive neutrinos on cosmological observables, that can be used to extract bounds on neutrino masses from present data. Finally we discuss the sensitivities on neutrino masses from future cosmological experiments.

Massive neutrinos could also play a role in the generation of the baryon asymmetry of the Universe from a previously created lepton asymmetry. In these leptogenesis scenarios, one can also obtain quite restrictive bounds on light neutrino masses, which are however model-dependent. We do not discuss this subject here, as it was covered by other lecturer at the School [6].

1. THE COSMIC NEUTRINO BACKGROUND

The existence of a relic sea of neutrinos is a generic feature of the standard hot Big Bang model, in number only slightly below that of relic photons that constitute the cosmic microwave background (CMB). This cosmic neutrino background

(CNB) has not been detected yet, but its presence is indirectly established by the accurate agreement between the calculated and observed primordial abundances of light elements, as well as from the analysis of the power spectrum of CMB anisotropies. In this section we will summarize the evolution and main properties of the CNB.

1.1. Relic Neutrino Production and Decoupling. Produced at large temperatures by frequent weak interactions, cosmic neutrinos of any flavour $(\nu_e, \nu_\mu, \nu_\tau)$ were kept in equilibrium until these processes became ineffective in the course of the expansion of the early Universe. While coupled to the rest of the primeval plasma (relativistic particles such as electrons, positrons, and photons), neutrinos had a momentum spectrum with an equilibrium Fermi–Dirac distribution with temperature T,

$$f_{\rm eq}(p,T) = \left[\exp\left(\frac{p-\mu_\nu}{T}\right)+1\right]^{-1}, \qquad (1)$$

which is just one example of the general case of particles in equilibrium (fermions or bosons, relativistic or nonrelativistic), as shown, e.g., in [7]. In the previous equation we have included a neutrino chemical potential μ_ν that would exist in the presence of a neutrino–antineutrino asymmetry.

As the Universe cools, the weak interaction rate Γ_ν falls below the expansion rate and neutrinos decouple from the rest of the plasma. An estimate of the decoupling temperature $T_{\rm dec}$ can be found by equating the thermally averaged value of the weak interaction rate, $\Gamma_\nu = \langle \sigma_\nu \, n_\nu \rangle$, where $\sigma_\nu \propto G_F^2$ is the cross section of the electron–neutrino processes, with G_F — the Fermi constant, and n_ν is the neutrino number density, with the expansion rate given by the Hubble parameter H, $H^2 = 8\pi\rho/3M_P^2$, where $\rho \propto T^4$ is the total energy density, dominated by relativistic particles, and $M_P = 1/G^{1/2}$ is the Planck mass. If we approximate the numerical factors to unity, with $\Gamma_\nu \approx G_F^2 T^5$ and $H \approx T^2/M_P$, we obtain the rough estimate $T_{\rm dec} \approx 1$ MeV. More accurate calculations give slightly higher values of $T_{\rm dec}$ which are flavour-dependent since electron neutrinos and antineutrinos are in closer contact with electrons and positrons, as shown, e.g., in [1].

Although neutrino decoupling is not described by a unique $T_{\rm dec}$, it can be approximated as an instantaneous process. The standard picture of *instantaneous neutrino decoupling* is very simple (see, e.g., [7] or [8]) and reasonably accurate. In this approximation, the spectrum in Eq. (1) is preserved after decoupling, since both neutrino momenta and temperature redshift identically with the expansion of the Universe. In other words, the number density of noninteracting neutrinos remains constant in a comoving volume since the decoupling epoch. We will see later that active neutrinos cannot possess masses much larger than 1 eV, so they were ultra-relativistic at decoupling. This is the reason why the momentum distribution in Eq. (1) does not depend on the neutrino masses, even after decoupling, i.e., there is no neutrino energy in the exponential of $f_{\rm eq}(p)$.

When calculating quantities related to relic neutrinos, one must consider the various possible degrees of freedom per flavour. If neutrinos are massless or Majorana particles, there are two degrees of freedom for each flavour, one for neutrinos (one negative helicity state) and one for antineutrinos (one positive helicity state). Instead, for Dirac neutrinos there are in principle twice more degrees of freedom, corresponding to the two helicity states. However, the extra degrees of freedom should be included in the computation only if they are populated and brought into equilibrium before the time of neutrino decoupling. In practice, the Dirac neutrinos with the «wrong-helicity» states do not interact with the plasma at temperatures of the MeV order and have a vanishingly small density with respect to the usual left-handed neutrinos (unless neutrinos have masses close to the keV range, as explained in Sec. 6.4 of [1], but such a large mass is excluded for active neutrinos). Thus, the relic density of active neutrinos does not depend on their nature, either Dirac or Majorana particles.

Shortly after neutrino decoupling the temperature drops below the electron mass, favouring $e^{\pm}$ annihilations that heat the photons. If one assumes that this entropy transfer did not affect the neutrinos because they were already completely decoupled, it is easy to calculate the change in the photon temperature before any $e^{\pm}$ annihilation and after the electron–positron pairs disappear by assuming entropy conservation of the electromagnetic plasma. The result is

$$\frac{T_\gamma^{\rm after}}{T_\gamma^{\rm before}} = \left(\frac{11}{4}\right)^{1/3} \simeq 1.40102, \tag{2}$$

which is also the ratio between the temperatures of relic photons and neutrinos $T_\gamma/T_\nu = (11/4)^{1/3}$. The evolution of this ratio during the process of $e^{\pm}$ annihilations is shown in Fig. 1, *a*, while one can see in the other plot how in this epoch the photon temperature decreases with the expansion less than the inverse of the scale factor a. Instead, the temperature of the decoupled neutrinos always falls as $1/a$.

It turns out that the standard picture of neutrino decoupling described above is slightly modified: the processes of neutrino decoupling and $e^{\pm}$ annihilations are sufficiently close in time so that some relic interactions between $e^{\pm}$ and neutrinos exist. These relic processes are more efficient for larger neutrino energies, leading to nonthermal distortions in the neutrino spectra at the per cent level and a slightly smaller increase of the comoving photon temperature, as noted in a series of works (see the full list given in the review by [1]). A proper calculation of the process of noninstantaneous neutrino decoupling demands solving the momentum-dependent Boltzmann equations for the neutrino spectra, a set of integro-differential kinetic equations that are difficult to solve numerically. The most recent analysis [9] of this problem included the effect of flavour neutrino oscillations on the neutrino decoupling process. One finds an increase in the neutrino energy densities with respect to the instantaneous decoupling approximation (0.73 and 0.52% for ν_es

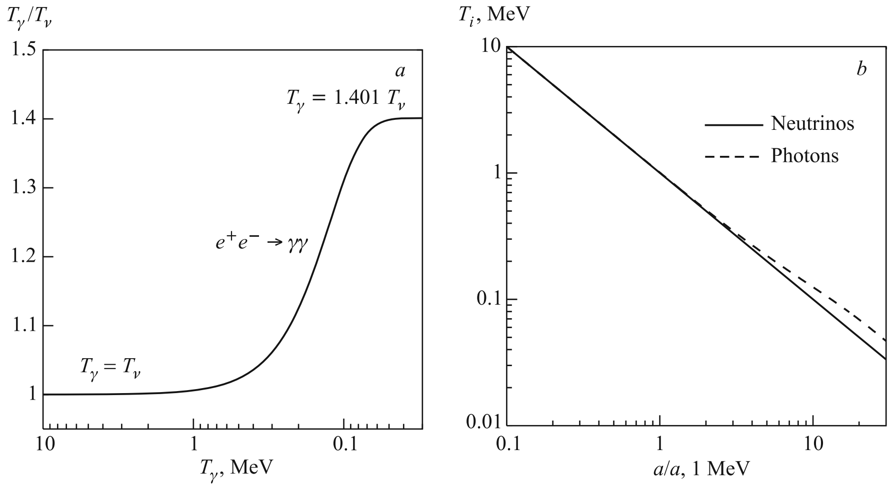

Fig. 1. Photon and neutrino temperatures during the process of $e^\pm$ annihilations: evolution of their ratio (*a*) and their decrease with the expansion of the Universe (*b*)

and $\nu_{\mu,\tau}$s, respectively) and a value of the comoving photon temperature after $e^\pm$ annihilations which is a factor of 1.3978 larger, instead of 1.40102. These changes modify the contribution of relativistic relic neutrinos to the total energy density which is taken into account using $N_{\rm eff} \simeq 3.046$, as defined later in Eq. (11). In practice, the distortions calculated in [9] only have small consequences on the evolution of cosmological perturbations, and for many purposes they can be safely neglected.

Any quantity related to relic neutrinos can be calculated after decoupling with the spectrum in Eq. (1) and T_ν. For instance, the number density per flavour is fixed by the temperature,

$$n_\nu = \frac{3}{11}\, n_\gamma = \frac{6\zeta(3)}{11\pi^2}\, T_\gamma^3, \tag{3}$$

which leads to a present value of 113 neutrinos and antineutrinos of each flavour per cm^3. Instead, the energy density for massive neutrinos should, in principle, be calculated numerically, with two well-defined analytical limits,

$$\rho_\nu(m_\nu \ll T_\nu) = \frac{7\pi^2}{120}\left(\frac{4}{11}\right)^{4/3} T_\gamma^4, \tag{4}$$

$$\rho_\nu(m_\nu \gg T_\nu) = m_\nu n_\nu. \tag{5}$$

1.2. Background Evolution. Let us discuss the evolution of the CNB after decoupling in the expanding Universe, which is described by the Friedmann–Robertson–Walker metric [8]

$$ds^2 = dt^2 - a(t)^2\,\delta_{ij}\,dx^i\,dx^j, \tag{6}$$

where we assumed negligible spatial curvature. Here $a(t)$ is the scale factor usually normalized to unity now ($a(t_0) = 1$) and related to the redshift z as $a = 1/(1+z)$. General relativity tells us the relation between the metric and the matter and energy in the Universe via the Einstein equations, whose time–time component is the Friedmann equation

$$\left(\frac{\dot{a}}{a}\right)^2 = H^2 = \frac{8\pi G}{3}\rho = H_0^2\frac{\rho}{\rho_c^0}, \tag{7}$$

that gives the Hubble rate in terms of the total energy density ρ. At any time, the critical density ρ_c is defined as $\rho_c = 3H^2/8\pi G$, and the current value H_0 of the Hubble parameter gives the critical density today

$$\rho_c^0 = 1.8788 \cdot 10^{-29}\,h^2\ \mathrm{g}\cdot\mathrm{cm}^{-3}, \tag{8}$$

where $h \equiv H_0/(100\ \mathrm{km}\cdot\mathrm{s}^{-1}\cdot\mathrm{Mpc}^{-1})$.

The different contributions to the total energy density are

$$\rho = \rho_\gamma + \rho_{\mathrm{cdm}} + \rho_b + \rho_\nu + \rho_\Lambda, \tag{9}$$

and the evolution of each component is given by the energy conservation law in an expanding Universe $\dot{\rho} = -3H(\rho + p)$, where p is the pressure. Thus the homogeneous density of photons ρ_γ scales like a^{-4}, that of nonrelativistic matter (ρ_{cdm} for cold dark matter and ρ_b for baryons) like a^{-3}, and the cosmological constant density ρ_Λ is of course time-independent. Instead, the energy density of neutrinos contributes to the radiation density at early times but behaves as matter after the nonrelativistic transition.

The evolution of all densities is shown in Fig. 2, *a*, starting at MeV temperatures until now. We also display the characteristic times for the end of Primordial Nucleosynthesis and for photon decoupling or recombination. The evolution of the density fractions $\Omega_i \equiv \rho_i/\rho_c$ is shown on the right panel, where it is easier to see which of the Universe components is dominant, fixing its expansion rate: first radiation in the form of photons and neutrinos (Radiation Domination or RD), then matter which can be CDM, baryons and massive neutrinos at late times (Matter Domination or MD) and finally the cosmological constant density takes over at low redshift (typically $z < 0.5$).

Massive neutrinos are the only particles that present a transition from radiation to matter, when their density is clearly enhanced (upper solid lines in Fig. 2).

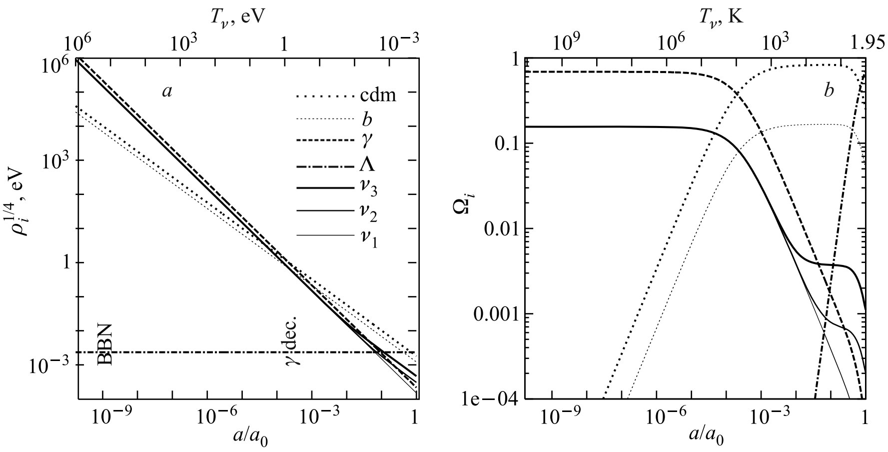

Fig. 2. Evolution of the background energy densities (a) and density fractions Ω_i (b) from the time when $T_\nu = 1$ MeV until now, for each component of a flat ΛMDM model with $h = 0.7$ and current density fractions $\Omega_\Lambda = 0.70$, $\Omega_b = 0.05$, $\Omega_\nu = 0.0013$, and $\Omega_{\rm cdm} = 1 - \Omega_\Lambda - \Omega_b - \Omega_\nu$. The three neutrino masses are $m_1 = 0$, $m_2 = 0.009$ eV, and $m_3 = 0.05$ eV

Obviously the contribution of massive neutrinos to the energy density in the nonrelativistic limit is a function of the mass (or the sum of all masses for which $m_i \gg T_\nu$), and the present value Ω_ν could be of order unity for eV masses (see Sec. 4).

2. NEUTRINOS AND PRIMORDIAL NUCLEOSYNTHESIS

In the course of its expansion, when the early Universe was only less than a second old, the conditions of temperature and density of its nucleon component were such that light nuclei could be created via nuclear reactions (for a recent review, see [10]). During this epoch, known as Primordial or Big Bang Nucleosynthesis (BBN), the primordial abundances of light elements were produced: mostly ^{4}He but also smaller quantities of less stable nuclei such as D, ^{3}He, and ^{7}Li. Heavier elements could not be produced because of the rapid evolution of the Universe and its small nucleon content, related to the small value of the baryon asymmetry which normalized to the photon density, $\eta_b \equiv (n_b - n_{\bar b})/n_\gamma$, is about a few times 10^{-10}. Measuring these primordial abundances today is a very difficult task, because stellar process may have altered the chemical compositions. Still, data on the primordial abundances of ^{4}He, D, and ^{7}Li exist and can be compared with the theoretical predictions to learn about the conditions of the Universe at such an early period. Thus BBN can be used as a cosmological test of any nonstandard physics or cosmology [11].

The physics of BBN is well understood, since in principle only involves the Standard Model of particle physics and the time evolution of the expansion rate as given by the Friedmann equation. In the first phase of BBN, the weak processes that keep the neutrons and protons in equilibrium,

$$n + \nu_e \leftrightarrow p + e^-, \quad n + e^+ \leftrightarrow p + \bar{\nu}_e \tag{10}$$

freeze and the neutron-to-proton ratio becomes a constant (later diminished due to neutron decays, $n \to p + e^- + \bar{\nu}_e$). This ratio largely fixes the produced ^{4}He abundance. Later, all the primordial abundances of light elements are produced and their value depends on the competition between the nuclear reaction rates and the expansion rate. These values can be quite precisely calculated with a BBN numerical code (see, e.g., [12]). At present there exists a nice agreement with the observed abundance of D for a value of the baryon asymmetry $\eta = (5.7 \pm 0.6) \cdot 10^{-10}$ [10], which also agrees with the region determined by CMB [13] and large-scale structure data (LSS). Instead, the predicted primordial abundance of ^{4}He tends to be a bit smaller than the observed value [14, 15]. However, it is difficult to consider this as a serious discrepancy, because the accuracy of the observations of ^{4}He is limited by systematic uncertainties.

There are two main effects of relic neutrinos at BBN. The first one is that they contribute to the relativistic energy density of the Universe (if $m_\nu \ll T_\nu$), thus fixing the expansion rate. This is why BBN gave the first allowed range of the number of neutrino species before accelerators (see the next section). On the other hand, BBN is the last period of the Universe sensitive to neutrino flavour, since electron neutrinos and antineutrinos play a direct role in the processes in Eq. (10).

3. EXTRA RADIATION AND THE EFFECTIVE NUMBER OF NEUTRINOS

Together with photons, in the standard case neutrinos fix the expansion rate during the cosmological era when the Universe is dominated by radiation. Their contribution to the total radiation content can be parametrized in terms of the effective number of neutrinos $N_{\rm eff}$, through the relation

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

where we have normalized to the photon energy density because its value today is known from the measurement of the CMB temperature. This equation is valid when neutrino decoupling is complete and holds as long as all neutrinos are relativistic.

We know that the number of light neutrinos sensitive to weak interactions (flavour or active neutrinos) equals three from the analysis of the invisible Z-boson width at LEP, $N_\nu = 2.984 \pm 0.008$ [16], and we saw in the previous section from the analysis of neutrino decoupling that they contribute as $N_{\rm eff} \simeq 3.046$. Any departure of $N_{\rm eff}$ from this last value would be due to nonstandard neutrino features or to the contribution of other relativistic relics. For instance, the energy density of a hypothetical scalar particle ϕ in equilibrium with the same temperature as neutrinos would be $\rho_\phi = (\pi/30)\, T_\nu^4$, leading to a departure of $N_{\rm eff}$ from the standard value of $4/7$. For detailed discussion of cosmological scenarios, where $N_{\rm eff}$ is not fixed to 3, see, e.g., [1, 11].

In the previous section, we saw that the expansion rate during BBN fixes the produced abundances of light elements, and in particular that of ^{4}He. Thus, the value of $N_{\rm eff}$ can be constrained at the BBN epoch from the comparison of theoretical predictions and experimental data on the primordial abundances of light elements. In addition, a value of $N_{\rm eff}$ different from the standard one would modify the transition epoch from a radiation-dominated to a matter-dominated Universe, which has some consequences on some cosmological observables such as the power spectrum of CMB anisotropies, leading to independent bounds on the radiation content. These are two complementary ways of constraining $N_{\rm eff}$ at very different epochs. Interestingly, recent data on the anisotropies of the CMB from WMAP [13] and the primordial ^{4}He abundance [14, 15] seem to favor a value of $N_{\rm eff} > 3$, although with large errorbars. The upcoming CMB measurements by the PLANCK satellite will soon pin down the radiation content of the Universe.

4. MASSIVE NEUTRINOS AS DARK MATTER

Nowadays the existence of Dark Matter (DM), the dominant nonbaryonic component of the matter density in the Universe, is well established. A priori, massive neutrinos are excellent DM candidates, in particular because we are certain that they exist, in contrast with other candidate particles. Together with CMB photons, relic neutrinos can be found anywhere in the Universe with a number density given by the present value of Eq. (3) of 339 neutrinos and antineutrinos per cm^3, and their energy density in units of the critical value of the energy density (see Eq. (8)) is

$$\Omega_\nu = \frac{\rho_\nu}{\rho_{\rm c}^0} = \frac{\sum\limits_i m_i}{93.14\, h^2 \ {\rm eV}}. \qquad (12)$$

Here $\sum_i m_i$ includes all masses of the neutrino states which are nonrelativistic today. It is also useful to define the neutrino density fraction f_ν with respect to the total matter density $f_\nu \equiv \rho_\nu/(\rho_{\rm cdm} + \rho_b + \rho_\nu) = \Omega_\nu/\Omega_m$.

In order to check whether relic neutrinos can have a contribution of order unity to the present values of Ω_ν or f_ν, we should consider which neutrino masses are allowed by noncosmological data. Oscillation experiments measure the differences of squared neutrino masses $\Delta m^2_{21} = m^2_2 - m^2_1$ and $\Delta m^2_{31} = m^2_3 - m^2_1$, the relevant ones for solar and atmospheric neutrinos, respectively, [17, 18]. As a reference, we take the following values of mixing parameters (3σ ranges or upper bounds) from an update of [19]:

$$\begin{gathered}\Delta m^2_{21} = (7.59^{+0.68}_{-0.56}) \cdot 10^{-5}\ \mathrm{eV}^2, \quad |\Delta m^2_{31}| = (2.40^{+0.35}_{-0.33}) \cdot 10^{-3}\ \mathrm{eV}^2, \\ s^2_{12} = 0.32^{+0.06}_{-0.05}, \quad s^2_{23} = 0.50^{+0.17}_{-0.14}, \quad s^2_{13} \leqslant 0.053.\end{gathered} \tag{13}$$

Here $s_{ij} = \sin\theta_{ij}$, where θ_{ij} ($ij = 12, 23$ or 13) are the three mixing angles. Unfortunately, oscillation experiments are insensitive to the absolute scale of neutrino masses, since the knowledge of $\Delta m^2_{21} > 0$ and $|\Delta m^2_{31}|$ leads to the two possible schemes shown in Fig. 1 of [4], but leaves one neutrino mass unconstrained. These two schemes are known as normal (NH) and inverted (IH) hierarchies, characterized by the sign of Δm^2_{31}, positive and negative, respectively. For small values of the lightest neutrino mass m_0, i.e., m_1 (m_3) for NH (IH), the mass states follow a hierarchical scenario, while for masses much larger than the differences all neutrinos share in practice the same mass and then we say that they are degenerate. In general, the relation between the individual masses and the total neutrino mass can be found numerically, as shown in Fig. 3.

There are two types of laboratory experiments searching for the absolute scale of neutrino masses, a crucial piece of information for constructing models of neutrino masses and mixings. The neutrinoless double-beta decay $(Z, A) \rightarrow (Z + 2, A) + 2e^-$ (in short $0\nu2\beta$) is a rare nuclear process where lepton number is violated and whose observation would mean that neutrinos are Majorana particles. If the $0\nu2\beta$ process is mediated by a light neutrino, the results from neutrinoless double-beta decay experiments are converted into an upper bound or a measurement of the effective mass $m_{\beta\beta}$

$$m_{\beta\beta} = |c^2_{12} c^2_{13}\, m_1 + s^2_{12} c^2_{13}\, m_2\, \mathrm{e}^{i\phi_2} + s^2_{13}\, m_3\, \mathrm{e}^{i\phi_3}| \tag{14}$$

where $\phi_{1,2}$ are the two Majorana phases that appear in lepton-number-violating processes. See [20, 21] for more details and the current experimental results.

Beta-decay experiments, which involve only the kinematics of electrons, are in principle the best strategy for measuring directly the neutrino mass [22]. The current limits from tritium beta decay apply only to the range of degenerate

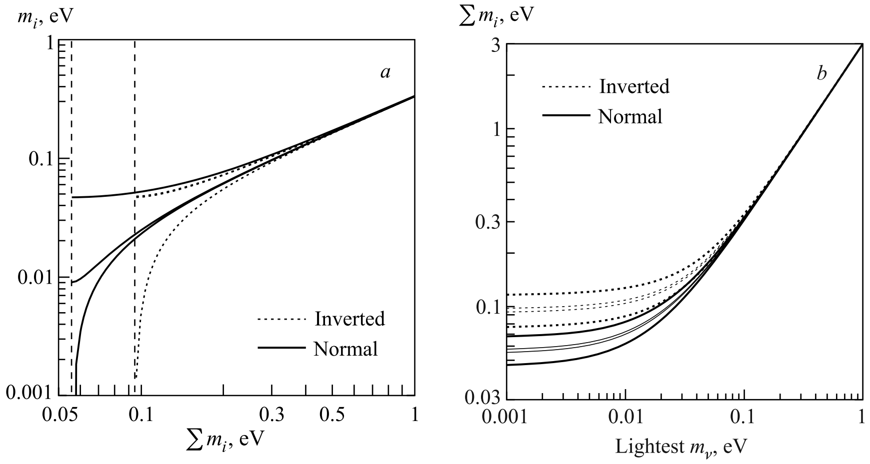

Fig. 3. Expected values of neutrino masses according to the values in Eq. (13). *a*) Individual neutrino masses as a function of the total mass for the best-fit values of the Δm^2. *b*) Ranges of total neutrino mass as a function of the lightest state within the 3σ regions (thick lines) and for a future determination at the 5% level (thin lines)

neutrino masses, so that $m_\beta \simeq m_0$, where

$$m_\beta = (c_{12}^2 c_{13}^2\, m_1^2 + s_{12}^2 c_{13}^2\, m_2^2 + s_{13}^2\, m_3^2)^{1/2} \tag{15}$$

is the relevant parameter for beta decay experiments. The bound at 95% CL is $m_0 < 2.05{-}2.3$ eV from the Troitsk and the Mainz experiments, respectively. This value is expected to be improved by the KATRI project to reach a discovery potential for 0.3–0.35 eV masses (or a sensitivity of 0.2 eV at 90% CL). Taking into account the present upper bound and the minimal values of the total neutrino mass in the normal (inverted) hierarchy, the sum of neutrino masses is restricted to the approximate range

$$0.06\,(0.1)\ \text{eV} \lesssim \sum_i m_i \lesssim 6\ \text{eV}. \tag{16}$$

As we discuss in the next sections, cosmology is at first order sensitive to the total neutrino mass $\sum\limits_i m_i$ if all states have the same number density, providing information on m_0 but blind to neutrino mixing angles or possible CP-violating phases. Thus cosmological results are complementary to terrestrial experiments. The interested reader can find the allowed regions in the parameter space defined by any pair of parameters $(\sum\limits_i m_i, m_{\beta\beta}, m_\beta)$ in [23].

Now we can find the possible present values of Ω_ν in agreement with the three neutrino masses shown in Fig. 3 and the approximate bounds of Eq. (16). Note that even if the three neutrinos are nondegenerate in mass, Eq. (12) can be safely applied, because we know from neutrino oscillation data that at least two of the neutrino states are nonrelativistic today, since both $(\Delta m^2_{31})^{1/2} \simeq 0.05$ eV and $(\Delta m^2_{21})^{1/2} \simeq 0.009$ eV are larger than the temperature $T_\nu \simeq 1.96$ K $\simeq 1.7 \cdot 10^{-4}$ eV. If the third neutrino state is very light and still relativistic, its relative contribution to Ω_ν is negligible and Eq. (12) remains an excellent approximation of the total density. One finds that Ω_ν is restricted to the approximate range

$$0.0013\,(0.0022) \lesssim \Omega_\nu \lesssim 0.13, \tag{17}$$

where we already included that $h \approx 0.7$. This applies only to the standard case of three light active neutrinos, while in general a cosmological upper bound on Ω_ν has been used since the 1970s to constrain the possible values of neutrino masses. For instance, if we demand that neutrinos should not be heavy enough to overclose the Universe ($\Omega_\nu < 1$), we obtain an upper bound $\sum\limits_i m_i \lesssim 45$ eV (again fixing $h = 0.7$). Moreover, since from the present analyses of cosmological data we know that the approximate contribution of matter is $\Omega_m \simeq 0.3$, the neutrino masses should obey the stronger bound $\sum\limits_i m_i \lesssim 15$ eV. We see that with this simple argument one obtains a bound which is roughly only a factor of 2 worse than the bound from tritium beta decay, but of course with the caveats that apply to any cosmological analysis. In the three-neutrino case, these bounds should be understood in terms of $m_0 = \sum\limits_i m_i/3$.

Dark matter particles with a large velocity dispersion such as that of neutrinos are called hot dark matter (HDM). The role of neutrinos as HDM particles has been widely discussed since the 1970s, and the reader can find a historical review in [24]. It was realized in the mid-1980s that HDM affects the evolution of cosmological perturbations in a particular way: it erases the density contrasts on wavelengths smaller than a mass-dependent free-streaming scale. In a Universe dominated by HDM, this suppression is in contradiction with various observations. For instance, large objects such as superclusters of galaxies form first, while smaller structures like clusters and galaxies form via a fragmentation process. This top-down scenario is at odds with the fact that galaxies seem older than clusters.

Given the failure of HDM-dominated scenarios, the attention then turned to cold dark matter (CDM) candidates, i.e., particles which were nonrelativistic at the epoch when the Universe became matter-dominated, which provided a better agreement with observations. Still in the mid-1990s it appeared that a small mixture of HDM in the Universe dominated by CDM fitted better the observational data on density fluctuations at small scales than a pure CDM model. However, within the presently favoured ΛCDM model dominated at late times by a cosmological constant (or some form of dark energy) there is no need for a significant contribution of HDM. Instead, one can use the available cosmological data to find how large the neutrino contribution can be, as we will see later.

Do neutrino oscillations have an effect on any cosmological epoch? In the standard picture, all flavour neutrinos were produced with the same energy spectrum, so no effect is expected from oscillations among these three states (up to small spectral distortions, see [9]). But there are two cases where neutrino oscillations could have cosmological consequences: flavour oscillations with nonzero relic neutrino asymmetries and active-sterile neutrino oscillations (for more details, see, e.g., Sec. 5 in [25]).

5. EFFECTS OF NEUTRINO MASSES ON COSMOLOGY

In this section we will briefly describe the main cosmological observables and the effects that neutrino masses cause on them. A more detailed discussion of the effects of massive neutrinos on the evolution of cosmological perturbations can be found in Subsecs. 4.5 and 4.6 of [4].

5.1. Brief Description of Cosmological Observables. Although there exist many different types of cosmological measurements, here we will restrict the discussion to those that are more important for obtaining an upper bound or eventually a measurement of neutrino masses.

First of all, we have the CMB temperature anisotropy power spectrum, defined as the angular two-point correlation function of CMB maps $\delta T/\bar{T}(\hat{n})$ ($\hat{n}$ being a direction in the sky). This function is usually expanded in Legendre multipoles

$$\left\langle \frac{\delta T}{\bar{T}}(\hat{n}) \frac{\delta T}{\bar{T}}(\hat{n}') \right\rangle = \sum_{l=0}^{\infty} \frac{(2l+1)}{4\pi} C_l \, P_l(\hat{n} \cdot \hat{n}'), \tag{18}$$

where $P_l(x)$ are the Legendre polynomials. For Gaussian fluctuations, all the information is encoded in the multipoles C_l which probe correlations on angular scales $\theta = \pi/l$. We have seen that each neutrino family can only have a mass of the order of 1 eV, so that the transition of relic neutrinos to the nonrelativistic regime is expected to take place after the time of recombination between electrons and nucleons, i.e., after photon decoupling. Since the shape of the CMB spectrum is related mainly to the physical evolution *before* recombination, it will be only marginally affected by the neutrino mass, except for an indirect effect through the modified background evolution. There exists interesting complementary information to the temperature power spectrum if the

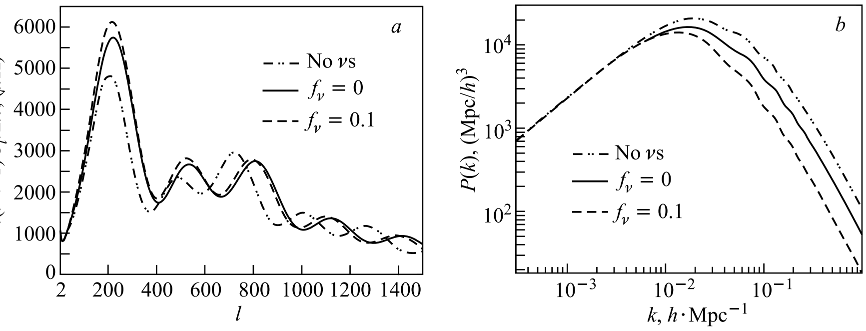

Fig. 4. CMB temperature anisotropy spectrum C_l^T and matter power spectrum $P(k)$ for three models: the neutrinoless ΛCDM model, a more realistic ΛCDM model with three massless neutrinos ($f_\nu \simeq 0$), and finally a ΛMDM model with three massive degenerate neutrinos and a total density fraction $f_\nu = 0.1$. In all models, the values of the cosmological parameters ($\omega_b = \Omega_b h^2$, $\omega_m = \Omega_m h^2$, Ω_Λ, A_s, n, τ) have been kept fixed

CMB polarization is measured, and currently we have some less precise data on the temperature $\times$ E-polarization (TE) correlation function and the E-polarization self-correlation spectrum (EE).

The current Large Scale Structure (LSS) of the Universe is probed by the matter power spectrum, observed with various techniques described in the next section (directly or indirectly, today or in the near past at redshift z). It is defined as the two-point correlation function of nonrelativistic matter fluctuations in Fourier space, $P(k,z) = \langle|\delta_m(k,z)|^2\rangle$, where $\delta_m = \delta\rho_m/\bar{\rho}_m$. Usually $P(k)$ refers to the matter power spectrum evaluated today (at $z = 0$). In the case of several fluids (e.g., CDM, baryons and nonrelativistic neutrinos), the total matter perturbation can be expanded as $\delta_m = \sum\limits_i \bar{\rho}_i\delta_i / \sum\limits_i \bar{\rho}_i$. Since the energy density is related to the mass density of nonrelativistic matter through $E = mc^2$, δ_m represents indifferently the energy or mass power spectrum. The shape of the matter power spectrum is affected by the free-streaming caused by small neutrino masses of $\mathcal{O}$(eV) and thus it is the key observable for constraining m_ν with cosmological methods.

We will show later in Fig. 4 the typical shape of both the CMB temperature anisotropy spectrum C_l and the matter power spectrum $P(k)$.

5.2. Neutrino Free-Streaming. After thermal decoupling, relic neutrinos constitute a collisionless fluid, where the individual particles free-stream with a characteristic velocity that, in average, is the thermal velocity $v_{\rm th}$. It is possible to define a horizon as the typical distance on which particles travel between time t_i and t. When the Universe was dominated by radiation or matter $t \gg t_i$, this horizon is, as usual, asymptotically equal to $v_{\rm th}/H$, up to a numerical factor of order one. Similar to the definition of the Jeans length (see Subsec. 4.4 in [4]), we can define the neutrino free-streaming wavenumber and length as

$$k_{\rm FS}(t) = \left(\frac{4\pi G \bar{\rho}(t) a^2(t)}{v_{\rm th}^2(t)} \right)^{1/2}, \quad \lambda_{\rm FS}(t) = 2\pi \frac{a(t)}{k_{\rm FS}(t)} = 2\pi \sqrt{\frac{2}{3}} \frac{v_{\rm th}(t)}{H(t)}. \quad (19)$$

As long as neutrinos are relativistic, they travel at the speed of light and their free-streaming length is simply equal to the Hubble radius. When they become nonrelativistic, their thermal velocity decays like

$$v_{\rm th} \equiv \frac{\langle p \rangle}{m} \simeq \frac{3T_\nu}{m} = \frac{3T_\nu^0}{m} \left(\frac{a_0}{a} \right) \simeq 150\,(1+z) \left(\frac{1 \text{ eV}}{m} \right) \text{km} \cdot \text{s}^{-1}, \quad (20)$$

where we used for the present neutrino temperature $T_\nu^0 \simeq (4/11)^{1/3} T_\gamma^0$ and $T_\gamma^0 \simeq 2.726$ K. This gives for the free-streaming wavelength and wavenumber during matter or Λ domination

$$\lambda_{\rm FS}(t) = 7.7 \frac{1+z}{\sqrt{\Omega_\Lambda + \Omega_m (1+z)^3}} \left(\frac{1 \text{ eV}}{m} \right) h^{-1}\,\text{Mpc}, \quad (21)$$

$$k_{\rm FS}(t) = 0.82 \frac{\sqrt{\Omega_\Lambda + \Omega_m (1+z)^3}}{(1+z)^2} \left(\frac{m}{1 \text{ eV}} \right) h\,\text{Mpc}^{-1}, \quad (22)$$

where Ω_Λ and Ω_m are the cosmological constant and matter density fractions, respectively, evaluated today. So, after the nonrelativistic transition and during matter domination, the free-streaming length continues to increase, but only like $(aH)^{-1} \propto t^{1/3}$, i.e., more slowly than the scale factor $a \propto t^{2/3}$. Therefore, the comoving free-streaming length $\lambda_{\rm FS}/a$ actually decreases like $(a^2 H)^{-1} \propto t^{-1/3}$. As a consequence, for neutrinos becoming nonrelativistic during matter domination, the comoving free-streaming wavenumber passes through a minimum $k_{\rm nr}$ at the time of the transition, i.e., when $m = \langle p \rangle = 3T_\nu$ and $a_0/a = (1+z) = 2.0 \cdot 10^3 (m/1 \text{ eV})$. This minimum value is found to be

$$k_{\rm nr} \simeq 0.018\ \Omega_m^{1/2} \left(\frac{m}{1 \text{ eV}} \right)^{1/2} h\,\text{Mpc}^{-1}. \quad (23)$$

The physical effect of free-streaming is to damp small-scale neutrino density fluctuations: neutrinos cannot be confined into (or kept outside of) regions smaller than the free-streaming length, for obvious kinematic reasons. There exists a gravitational back-reaction effect that also damps the metric perturbations on those scales. Instead, on scales much larger than the free-streaming scale the neutrino velocity can be effectively considered as vanishing and after the nonrelativistic transition the neutrino perturbations behave like CDM perturbations. In particular, modes with $k < k_{\rm nr}$ are never affected by free-streaming and evolve like in a pure ΛCDM model.

5.3. Impact of m_ν on the Matter Power Spectrum. The small initial cosmological perturbations in the early Universe evolve under the linear regime at any scale at early times and on the largest scales more recently, and produce the structures we see today. We cannot review here all the details (see [4] and references therein), but we will emphasize the main effects caused by massive neutrinos in the framework of the standard cosmological scenario: a Λ Mixed Dark Matter (ΛMDM) model, where Mixed refers to the inclusion of some HDM component.

First, let us describe the changes in the background evolution of the Universe. We have seen that massless neutrinos are always part of the radiation content, so in this case the present value of the matter contribution Ω_m^0 is equal to the contribution of CDM and baryons. Instead, massive neutrinos contribute to radiation at early times but to matter after becoming nonrelativistic. Thus, with respect to the massless neutrino case, massive neutrinos also contribute to Ω_m^0, reducing the values of $\Omega_{\rm CDM}^0$ and Ω_b^0. As a result, if these massive neutrinos have not yet become nonrelativistic at the time of radiation/matter equality (the epoch of the Universe when its starts to be dominated by matter and the contribution of radiation becomes subdominant), then this transition is delayed. The consequence of a late equality for the LSS matter power spectrum is the following: since on sub-Hubble scales the matter density contrast δ_m grows more efficiently during MD than during RD, the matter power spectrum is suppressed on small scales relatively to large scales.

At the perturbation level, we also saw that free-streaming damps small-scale neutrino density fluctuations. This produces a direct effect on the matter power spectrum (see Subsec. 4.5 of [4]), that depends on the value k with respect to $k_{\rm nr}$ in Eq. (23),

$$P(k) = \left\langle \left(\frac{\delta\rho_{\rm cdm} + \delta\,\rho_b + \delta\,\rho_\nu}{\rho_{\rm cdm} + \rho_b + \rho_\nu} \right)^2 \right\rangle = \\ = \left\langle \left(\frac{\Omega_{\rm cdm}\,\delta_{\rm cdm} + \Omega_b\delta_b + \Omega_\nu\,\delta_\nu}{\Omega_{\rm cdm} + \Omega_b + \Omega_\nu} \right)^2 \right\rangle = \\ = \begin{cases} \langle \delta_{\rm cdm}^2 \rangle & \text{for} \quad k < k_{\rm nr}, \\ [1 - \Omega_\nu/\Omega_m]^2 \, \langle \delta_{\rm cdm}^2 \rangle & \text{for} \quad k \gg k_{\rm nr}, \end{cases} \tag{24}$$

with $\Omega_m \equiv \Omega_{\rm cdm} + \Omega_b + \Omega_\nu$. Thus, the role of the neutrino masses would be simply to cut the power spectrum by a factor $[1-\Omega_\nu/\Omega_m]^2$ for $k \gg k_{\rm nr}$. However, it turns out that the presence of neutrinos actually modifies the evolution of the CDM and baryon density contrasts in such a way that the suppression factor is greatly enhanced, more or less by a factor four.

In conclusion, the combined effect of the shift in the time of equality and of the reduced CDM fluctuation growth during matter domination produces an attenuation of small-scale perturbations for $k > k_{\rm nr}$. It can be shown that for small values of f_ν this effect can be approximated in the large k limit by the well-known linear expression $P(k)^{f_\nu}/P(k)^{f_\nu=0} \simeq 1 - 8\,f_\nu$ [26]. For the comparison with the data, one could use instead some better analytical approximations to the full MDM or ΛMDM matter power spectrum, valid for arbitrary scales and redshifts, as listed in [4]. However, nowadays the analyses are performed using the matter power spectra calculated by Boltzmann codes such as CMBFAST [27] or CAMB [28], that solve numerically the evolution of the cosmological perturbations.

An example of $P(k)$ with and without massive neutrinos is shown in Fig. 4, where the effect of m_ν at large ks can be clearly visible. Such a suppression is probably better seen in Fig. 5, where we plot the ratio of the matter power spectrum for ΛMDM over that of ΛCDM, for different values of f_ν and three degenerate massive neutrinos, but for fixed parameters $(\omega_m,\, \Omega_\Lambda)$. For large ks, the linear approximation is a reasonable first-order approximation for $0 \;<\; f_\nu \;<\; 0.07$.

Is it possible to mimic the effect of massive neutrinos on the matter power spectrum with some combination of other cosmological parameters? If so, one would say that a parameter degeneracy exists, reducing the sensitivity to neutrino masses. This possibility depends on the interval $[k_{\rm min}, k_{\rm max}]$ in which the $P(k)$ can be accurately measured. Ideally, if we could have $k_{\rm min} \leqslant 10^{-2}h$ Mpc^{-1} and $k_{\rm max} \geqslant 1h$ Mpc^{-1}, the effect of the neutrino mass would be nondegenerate, because of its very characteristic step-like effect. In contrast, other cosmological parameters like the scalar tilt or the tilt running change the spectrum slope on all scales. The problem is that usually the matter power spectrum can only be

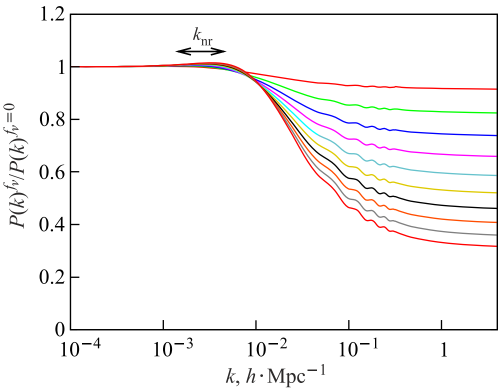

Fig. 5. Ratio of the matter power spectrum including three degenerate massive neutrinos with density fraction f_ν to that with three massless neutrinos. The parameters $(\omega_m, \Omega_\Lambda) = (0.147, 0.70)$ are kept fixed, and from top to bottom the curves correspond to $f_\nu = 0.01, 0.02, 0.03, \ldots, 0.10$. The individual masses m_ν range from 0.046 to 0.46 eV; and the scale $k_{\rm nr}$, from $2.1\cdot10^{-3}h$ Mpc^{-1} to $6.7\cdot10^{-3}h$ Mpc^{-1} as shown on the top of the figure

accurately measured in the intermediate region where the mass effect is neither null nor maximal: in other words, many experiments only have access to the transition region in the step-like transfer function. In this region, the neutrino mass affects the slope of the matter power spectrum in a way which can be easily confused with the effect of other cosmological parameters. Because of these parameter degeneracies, the LSS data alone cannot provide significant constraints on the neutrino mass, and it is necessary to combine them with other cosmological data, in particular, the CMB anisotropy spectrum, which could lift most of the degeneracies. Still, for exotic models with, e.g., extra relativistic degrees of freedom, a constant equation-of-state parameter of the dark energy different from -1 or a nonpower-law primordial spectrum, the neutrino mass bound can become significantly weaker (a factor 2 or more).

5.4. Impact of m_ν on the CMB Anisotropy Spectrum. For neutrino masses of the order of 1 eV (about $f_\nu \leqslant 0.1$) the three-neutrino species are still relativistic at the time of photon decoupling, and the direct effect of free-streaming neutrinos on the evolution of the baryon–photon acoustic oscillations is the same in the ΛCDM and ΛMDM cases. Therefore, the effect of the mass is indirect, appearing only at the level of the background evolution: the fact that the neutrinos account today for a fraction Ω_ν of the critical density implies some change either in the present value of the spatial curvature, or in the relative density of other species. If neutrinos were heavier than a few eV, they would already be nonrelativistic at decoupling. This case would have more complicated consequences for the CMB, as described in [29]. However, we will see later that this situation is disfavoured by current upper bounds on the neutrino mass.

Let us describe one example: we choose to maintain a flat Universe (the sum of all Ω_i is one) with fixed $(\omega_b = \Omega_b h^2, \omega_m = \Omega_m h^2, \Omega_\Lambda)$. Thus, while Ω_b and Ω_Λ are constant, $\Omega_{\rm cdm}$ is constrained to decrease as Ω_ν increases. The main effect on the CMB anisotropy spectrum results from a change in the time of equality. Since neutrinos are still relativistic at decoupling, they should be counted as radiation instead of matter around the time of equality, which is found by solving $\rho_b + \rho_{\rm cdm} = \rho_\gamma + \rho_\nu$. This gives $a_{\rm eq} = \Omega_r/(\Omega_b + \Omega_{\rm cdm})$, where Ω_r stands for the radiation density extrapolated until today *assuming that all neutrinos would remain massless*, given by Eq. (11) with $N_{\rm eff} \simeq 3.04$. So, when f_ν increases, $a_{\rm eq}$ increases proportionally to $[1 - f_\nu]^{-1}$: equality is postponed. This produces an enhancement of small-scale perturbations, especially near the first acoustic peak. Also, postponing the time of equality increases slightly the size of the sound horizon at recombination. These two features explain why in Fig. 4 the acoustic peaks are slightly enhanced and shifted to the left in the ΛMDM case.

Since the effect of the neutrino mass on CMB fluctuations is indirect and appears only at the background level, one could think that by changing the value of other cosmological parameters it would be possible to cancel exactly this effect (i.e., a parameter degeneracy). It can be actually shown that in the simplest ΛMDM model, with only seven cosmological parameters, one cannot vary the neutrino mass while keeping fixed $a_{\rm eq}$ and all other quantities governing the CMB spectrum. Therefore, it is possible to constrain the neutrino mass using CMB experiments alone [4, 30], although neutrinos are still relativistic at decoupling. This conclusion can be altered in more complicated models with extra cosmological parameters; for instance, allowing for an open Universe or varying the number of relativistic degrees of freedom. In such extended models the CMB alone is not sufficient for constraining the mass, but fortunately the LSS power spectrum can lift the degeneracy.

6. CURRENT BOUNDS ON NEUTRINO MASSES

Here we review how the available cosmological data is used to get information on the absolute scale of neutrino masses, complementary to laboratory experiments. Note that the bounds in the next subsections are all based on the Bayesian inference method, and the upper bounds on the sum of neutrino masses are given at 95% CL after marginalization over all free cosmological parameters. We refer the reader to [3,4] for a detailed discussion. Here it is assumed that the total neutrino mass is the only additional parameter with respect to a flat ΛCDM cosmological model characterized by 6 parameters, unless specified otherwise.

6.1. CMB Anisotropies. The experimental situation of the measurement of the CMB anisotropies is dominated by the seven-year release of WMAP data [13], which improved the already precise TT and TE angular power spectra of the previous releases, and included a detection of the E-polarization self-correlation spectrum (EE). On similar or smaller angular scales than WMAP, we have results from experiments that are either ground-based (ACBAR, VSA, CBI, DASI, ...) or balloon-borne (ARCHEOPS, BOOMERANG, MAXIMA, ...).

We saw in the previous section that the signature on the CMB spectrum of a neutrino mass smaller than 0.5 eV is small but does not vanish due to a background effect, proportional to Ω_ν, which changes some characteristic times and scales in the evolution of the Universe, and affects mainly the amplitude of the first acoustic peak as well as the location of all the peaks. Therefore, it is possible to constrain neutrino masses using CMB experiments only. In this framework, many analyses support the conclusion that a sensible bound on neutrino masses exists using CMB data only, of order of 2 eV for the total mass $M_\nu \equiv \sum\limits_i m_i$. This is an important result, since it does not depend on the uncertainties from LSS data discussed next.

6.2. Galaxy Redshift Surveys. We have seen that free-streaming of massive neutrinos produces a direct effect on the formation of cosmological structures. Actually, it is well known that the presence of neutrino masses leads to an attenuation of the linear matter power spectrum on small scales. In a seminal paper [26] it was shown that an efficient way to probe neutrino masses of order eV was to use data from large redshift surveys, which measure the distance to a large number of galaxies, giving us a three-dimensional picture of the Universe. At present, we have data from two large projects: the 2 degree Field (2dF) galaxy redshift survey and the Sloan Digital Sky Survey (SDSS).

One of the main goals of galaxy redshift surveys is to reconstruct the power spectrum of matter fluctuations on very large scales, whose cosmological evolution is described entirely by linear perturbation theory. However, the linear power spectrum must be reconstructed from individual galaxies which underwent a strongly nonlinear evolution. A simple analytic model of structure formation suggests that on large scales, the galaxy–galaxy correlation function should be, not equal, but proportional to the linear matter density power spectrum, up to a constant factor that is called the light-to-mass bias (b). This parameter can be obtained from independent methods, which tend to confirm that the linear biasing assumption is correct, at least in first approximation.

A conservative way to use the measurements of galaxy–galaxy correlations in an analysis of cosmological data is to take the bias as a free parameter, i.e., to consider only the shape of the matter power spectrum at the corresponding scales and not its amplitude (denoted as galaxy clustering data). An upper limit on M_ν between 0.8 and 1.7 eV is found from the analysis of galaxy clustering data (SDSS and/or 2dF, leaving the bias as a free parameter) added to CMB data. These values improve those found with CMB data only. The bounds on neutrino masses are more stringent when the amplitude of the matter power spectrum is fixed with a measurement of the bias, instead of leaving it as a free parameter. The upper limits on M_ν are reduced to values of order 0.5–0.9 eV although some analyses also add Lyman-α data (see the next subsection).

Finally, a galaxy redshift survey performed in a large volume can also be sensitive to the imprint created by the baryon acoustic oscillations (BAO) at large scales on the power spectrum of nonrelativistic matter. Since baryons are only a subdominant component of the nonrelativistic matter, the BAO feature is manifested as a small single peak in the galaxy correlation function in real space that was recently detected from the analysis of the SDSS luminous red galaxy (LRG) sample. The observed position of this baryon oscillation peak provides a way to measure the angular diameter distance out to the typical LRG redshift of $z = 0.35$, which in turn can be used to constrain the parameters of the underlying cosmological model.

6.3. Lyman-α Forest. The matter power spectrum on small scales can also be inferred from data on the so-called Lyman-α forest. This corresponds to the Lyman-α absorption of photons traveling from distant quasars ($z \sim 2{-}3$) by the neutral hydrogen in the intergalactic medium. As an effect of the Universe expansion, photons are continuously red-shifted along the line of sight, and can be absorbed when they reach a wavelength of 1216 Å in the rest-frame of the intervening medium. Therefore, the quasar spectrum contains a series of absorption lines, whose amplitude as a function of wavelength traces back the density and temperature fluctuations of neutral hydrogen along the line of sight. It is then possible to infer the matter density fluctuations in the linear or quasi-linear regime.

In order to use the Lyman-α forest data, one needs to recover the matter power spectrum from the spectrum of the transmitted flux, a difficult task that requires the use of hydrodynamical simulations for the corresponding cosmological model. Given the various systematics involved in the analysis, the robustness of Lyman-α forest data is still a subject of intense discussion between experts. In any case, the recovered matter power spectrum is again sensitive to the suppression of growth of mass fluctuations caused by massive neutrinos.

For a free bias, one finds that Lyman-α data help to reduce the upper bounds on the total neutrino mass to the level $M_\nu < 0.5{-}0.7$ eV. But those analyses that include Lyman-α data and a measurement of the bias do not always lead to a lower limit, ranging from 0.4 to 0.7 eV.

6.4. Summary and Discussion of Current Bounds. The upper bounds on M_ν from the previous subsections are representative of an important fact: a single cosmological bound on neutrino masses does not exist. Depending on the included set of data, the approximate ranges for the upper bounds are: 2–3 eV for CMB only, 0.9–1.7 eV for CMB and 2dF/SDSS-gal or 0.2–0.9 eV with the inclusion of a measurement of the bias and/or Lyman-α forest data and/or BAO data. For a discussion on the bounds on neutrino masses from different combinations of cosmological data, we refer the reader to [23], where they are compared with those coming from tritium beta decay and neutrinoless double-beta decay. In any case, current cosmological data probe the region of neutrino masses where the three neutrino states are degenerate, with a mass $M_\nu/3$.

7. FUTURE SENSITIVITIES ON m_ν FROM COSMOLOGY

In the near future we will have more precise data on cosmological observables from various experimental techniques and experiments. If the characteristics of these future experiments are known with some precision, it is possible to assume a «fiducial model», i.e., a cosmological model that would yield the best fit to future data, and to estimate the error bar on a particular parameter that will be obtained after marginalizing the hypothetical likelihood distribution over all the other free parameters. Technically, the simplest way to forecast this error is to compute a Fisher matrix, a technique has been widely used in the literature, for many different models and hypothetical datasets, now complemented by Monte Carlo methods. Here we will focus on the results for $\sigma(M_\nu)$, the forecast 68% CL error on the total neutrino mass, assuming various combinations of future observations:

CMB anisotropies measured with ground-based experiments or satellites such as PLANCK, galaxy redshift surveys, galaxy cluster surveys,... In particular, it has been emphasized the potentiality for measuring small neutrino masses of weak lensing experiments, which will look for the lensing effect caused by the large scale structure of the neighboring universe, either on the CMB signal [31] or on the apparent shape of galaxies (measured by cosmic shear surveys, see, e.g., [32]). We refer the reader to Sec. 6 of [4] for further details.

We give a graphical summary of the forecast sensitivities to neutrino masses of different cosmological data in Fig. 6, compared to the allowed values of neutrino masses in the two possible three-neutrino schemes. One can see from this figure that there are very good prospects for testing neutrino masses in the degenerate and quasi-degenerate mass regions above 0.2 eV or so. A detection at a significant level of the minimal value of the total neutrino mass in the inverted hierarchy scheme will demand the combination of future data from CMB lensing and cosmic shear surveys, whose more ambitious projects will provide a 2σ sensitivity to the minimal value in the case of normal hierarchy (of order 0.05 eV). The combination of CMB observations with future galaxy cluster surveys [33] or the measurement of the redshifted 21 cm signal from the epoch of reionization using low-frequency radio observations [34], should yield similar or even better sensitivities.

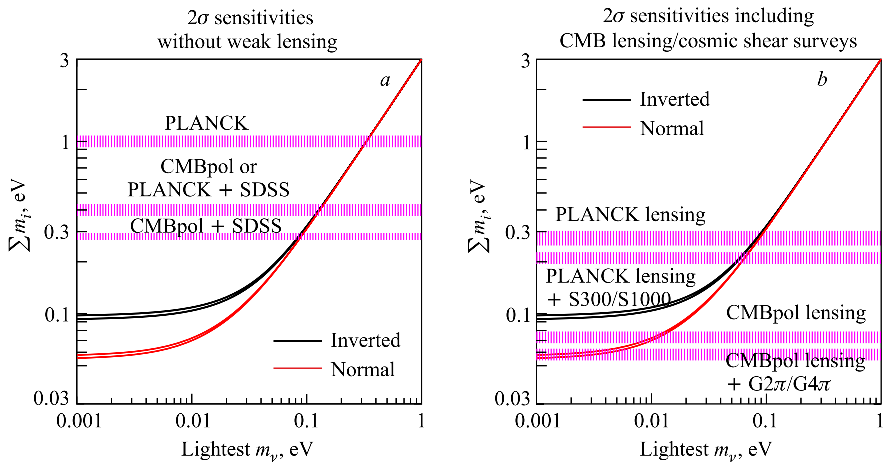

Fig. 6. Forecast 2σ sensitivities to the total neutrino mass from future cosmological experiments compared to the values in agreement with neutrino oscillation data (assuming a future determination at the 5% level). *a*) Sensitivities expected for future CMB experiments (without lensing extraction), alone and combined with the completed SDSS galaxy redshift survey. *b*) Sensitivities expected for future CMB experiments including lensing information, alone and combined with future cosmic shear surveys. Here CMBpol refers to a hypothetical CMB experiment roughly corresponding to the Inflation Probe mission

CONCLUSIONS

Neutrinos, despite the weakness of their interactions and their small masses, can play an important role in cosmology that we have reviewed in this contribution. In addition, cosmological data can be used to constrain neutrino properties, providing information on these elusive particles that complements the efforts of laboratory experiments. In particular, the data on cosmological observables have been used to bound the effective number of neutrinos (including a potential extra contribution from other relativistic particles).

But probably the most important contribution of cosmology to our knowledge of neutrino properties is the information it can provide on the absolute scale of neutrino masses. We have seen that the analysis of cosmological data can lead to either a bound or a measurement of the sum of neutrino masses, an important result complementary to terrestrial experiments such as tritium beta decay and neutrinoless double-beta decay experiments. In the next future, thanks to the data from new cosmological experiments, we could even hope to test the minimal values of neutrino masses guaranteed by the present evidences for flavour neutrino oscillations. For this and many other reasons, we expect that neutrino cosmology will remain an active research field in the next years.

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