Neutrinos in cosmology, with some significant digressions

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Abstract. Neutrinos play prominent roles in both particle physics and cosmology. In this talk, I will cover two broad topics. The first will be possible origins for neutrino masses and mixings and the implications of this physics for cosmology. Some non-cosmological digressions on the flavour problem in general will be made. The second topic will be Big Bang Nucleosynthesis (BBN) and bounds on active-sterile neutrino mixing.

1. SURVEY OF NEUTRINO COSMOLOGY

Two important issues come to mind when contemplating neutrinos in cosmology: the possible existence of additional neutrino or neutrino-like species, followed by the dynamical implications of neutrino mass and mixing [1].

1.1. Additional neutrino species

One may classify theoretically possible additional neutrino states into four categories:

• Light active. This option is clearly ruled out by the measured invisible width of the Z boson: we know that there are only three active neutrinos with masses less than about 45 GeV ($\nu_e, \nu_\mu, \nu_\tau$).

• Light sterile. The combined solar, atmospheric and LSND anomalies imply the existence of at least one light sterile flavour, if one demands that neutrino oscillations account for all of these data [2]. (The reason is simply that the three very different $\Delta m^2$ parameters required to furnish the appropriate oscillation lengths cannot be obtained from just the three known flavours of light active neutrino.)

• Heavy active. Additional active states with masses greater than 45 GeV are not precluded by present electroweak data. We will discuss how such states might arise in extensions of the standard model (SM) shortly.

• Heavy sterile. Various see-saw models of neutrino mass provide good particle physics motivations for this class of particle. In addition to the standard see-saw mechanism [3], I will also review the universal [4] and mirror see-saw [5] frameworks to show that several different patterns of light and heavy sterile neutrinos can arise. Perhaps the most important cosmological role proposed for heavy ster-
ile states is baryogenesis via sphaleron-reprocessed lepton asymmetries produced either through the decay or the interactions of the heavy neutral leptons.

1.1.1. The active/sterile distinction [6]

By an “active” neutrino we mean, of course, one that couples to the known left-handed weak interaction. The standard left-handed neutrinos, for example, sit in doublets of electroweak SU(2)$_L$ with weak-hypercharge $Y = -1$ [normalisation is such that electric charge $Q = (\tau_3 + Y)/2$]. Exotic active neutrinos are also theoretically possible, perhaps as components of higher-dimensional weak isospin representations.

The term “sterile neutrino” connotes a charge and colour neutral fermion that is also a singlet of SU(2)$_L$, so that it does not feel the weak interaction either. This terminology is appropriate in the context of how experiments are done: the weak force is the only known way of detecting neutrinos, so any neutrino-like species immune to this interaction will reveal itself only by its absence! However, in extensions of the SM it is quite common to find sterile neutrinos that do in fact feel some as yet hypothetical interaction, a right-handed weak force for example.

This motivates a division of sterile neutrinos into two further categories: “weakly” and “fully” sterile. Weakly sterile species are those that feel some hypothetical gauge interaction, while fully sterile states are those that couple to no gauge force, known or unknown. This distinction can be important, especially for theories of sterile neutrino mass and perhaps also for cosmology.\(^1\)

Examples of fully sterile neutrinos include the gauge singlets usually termed “right-handed neutrinos” that may be added to the minimal SM fermions, under the assumption of a gauge force desert up to the Planck scale. Another related example is the set of SU(5) singlets one may add to the 5$^* \oplus 10$ multiplets that compose the standard families in SU(5) grand unified theories (GUTs). These singlets are fully sterile if SU(5) is not embedded in a larger group such as SO(10) or E$_6$.

As is well known, fully sterile species allow gauge invariant bare Majorana mass terms. Since these parameters are not related to any symmetry breaking vacuum expectation values (VEVs), they have no natural scale. While many model builders set these masses to be large as a matter of course, it should be noted that taking them to zero increases the symmetry of the theory. This means that small bare Majorana masses for fully sterile neutrinos are technically natural (i.e. preserved under radiative corrections).

Weakly sterile neutrinos are probably more common in SM extensions. If the Glashow-Weinberg-Salam SU(2)$_L \otimes U(1)_Y$ electroweak model is extended to the left-right symmetric SU(2)$_L \otimes SU(2)_R \otimes U(1)_{B-L}$, then the right-handed neutrinos are in doublets of SU(2)$_R$ together with the right-handed charged leptons. This kind of right-handed neutrino couples to the right-sector $W$ and $Z$ bosons so it is weakly sterile. In many models, the Majorana mass for such a state is proportional to the right-handed

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\(^1\) Note that my definition indulges in “gauge interaction chauvinism” by ignoring spin-0 boson induced interactions. No slight is intended. Indeed, Higgs-induced active-sterile mixing is obviously a very important effect.
weak isospin symmetry breaking scale, which is experimentally constrained to be large. The left-right symmetric model may be embedded in SO(10) and $E_6$ GUTs. If so, then the (usually heavy) weakly-sterile neutrinos feel gauge interactions beyond the right-handed weak force, and they may pick up GUT symmetry-breaking scale masses. These examples illustrate how heavy sterile neutrinos typically arise in well-motivated SM extensions.

Mirror matter models [7] provide an interestingly different class of weakly sterile neutrino. Consider a model with gauge group $G_{SM} \otimes G_{SM}'$ with an exact discrete symmetry between the sectors. This is usually selected to be a non-standard parity symmetry so that the world is invariant under the full Poincaré group, including improper transformations (hence the designation “mirror matter”). Every ordinary particle has a mirror partner. The mirror neutrinos are immune to the ordinary weak interaction, but feel the mirror weak force, so they are weakly sterile. However, the exact discrete symmetry ensures that the mirror neutrinos are light for exactly the same reasons that ordinary neutrinos are light, whatever those reasons are! This illustrates how light (weakly) sterile neutrinos can arise in a simple and well-motivated extension of the SM. The mirror matter model provides a qualitatively different outcome for the sterile neutrino masses compared to most other models because the additional chiral interactions felt by these states are broken at a relatively low scale (the electroweak scale in fact).

1.1.2. Three see-saw scenarios

I now want to present three case studies of see-saw models of neutrino mass that can furnish qualitatively different varieties of additional neutrino-like states, both heavy and light. I will restrict myself to one-family examples for simplicity, with the realistic three-family extensions being (mostly) straightforward generalisations.

The standard see-saw mechanism arises when one adds a gauge-singlet neutrino $\nu_R$ to a minimal SM family and writes down all renormalisable gauge-invariant terms [3]. The neutrino mass matrix is contained within the mass term

$$
\left( \begin{array}{c} \nu_L \\ (\nu_R)^c \end{array} \right) \left( \begin{array}{cc} 0 & m \\ m & M \end{array} \right) \left( \begin{array}{c} (\nu_L)^c \\ \nu_R \end{array} \right) + H.c.
$$

where $m$ is the electroweak-scale Dirac mass and $M$ is a bare Majorana mass. In many obvious SM extensions, $M$ becomes proportional to a high gauge symmetry breaking scale. If indeed $M \gg m$, then the eigenvalues have approximate magnitudes $m^2/M$ and $M$, with the eigenstates being predominantly $\nu_L$ and $\nu_R$, respectively. Hence the standard active neutrino $\nu_L$ has a mass suppressed relative to the electroweak scale by the small hierarchy parameter $m/M$, while its right-handed fully sterile partner is very heavy. (In the extensions, the heavy partner becomes weakly sterile.)

My second case study is the universal see-saw scenario [4]. One adds heavy weak-isosinglet vector-like partners $U, D, N$ and $E$ to each of the standard fermions $u, d, \nu$ and $e$ and one works within the left-right symmetric augmentation of the electroweak sector. The idea is to explain why all fermions (top quark excluded), not just neutrinos, have small masses relative to the electroweak scale $\sim 200$ GeV.
Given a Higgs sector containing both SU(2)$_L$ and SU(2)$_R$ doublets but no bidoublets, the charge $-1/3$ quarks develop a mass matrix as given by

$$
\left( \begin{array}{cc}
\tilde d_L & D_L \\
0 & m_L \end{array} \right) \left( \begin{array}{c}
m_L \\
0 \end{array} \right) \left( \begin{array}{cc}
d_R & D_R \\
M & 0 \end{array} \right) + H.c.
$$

(2)

where $m_L (m_R)$ is proportional to the left-(right-)handed weak isospin breaking scale, and $M$ is the bare gauge-invariant Dirac mass for the vector-like $D$ (or it is induced by a Higgs singlet). Invoking the hierarchy $m_L \ll m_R \ll M$, the mass eigenvalues have approximate magnitudes $m_d \simeq m_L m_R / M$ and $m_D \simeq M$. The ratio $m_R / M$ is a hierarchy parameter, suppressing the down-quark charge relative to the electroweak scale. The light eigenstate has a predominant $d$ admixture, while the heavy eigenstate is mostly the exotic $D$. The other charged fermions develop similar mass matrices. (The electroweak scale mass for the top quark demands that it, at least, not be see-saw suppressed. I will come back to this point later, during a digression.)

The neutrino mass matrix has a very interesting form:

$$
\left( \begin{array}{ccc}
\nu_L & (\nu_R)^c & \bar N_L \\
(\nu_R)^c & (\nu_L)^c & \bar N_R \\
\bar N_L & (\bar N_R)^c & 0 \\
\bar N_R & (\bar N_L)^c & 0 \\
\end{array} \right) \left( \begin{array}{cccc}
0 & 0 & m_L & m'_L \\
0 & 0 & m_R & m'_R \\
m_L & m_R & M_{LL} & M_{LR} \\
m'_L & m'_R & M_{LR} & M_{RR} \\
\end{array} \right) \left( \begin{array}{ccc}
(\nu_L)^c \\
(\nu_R)^c \\
(\bar N_L)^c \\
(\bar N_R)^c \\
\end{array} \right) + H.c.
$$

(3)

The eigenvalues have approximate magnitudes

$$
\frac{m_L^2}{M} \ll m_{u,d,e}, \quad \frac{m_R^2}{M} \gg m_{u,d,e}, \quad M, \quad M,
$$

(4)

where $M$ denotes a generic large scale set by the parameters $M_{LL,LR,RR}$. The lightest eigenstate is the $\nu_L$ up to small admixtures. We see that the active neutrino mass is automatically more suppressed (by a factor $m_L / m_R$) than those of the charged fermions! This is the beauty of the universal see-saw mechanism.

But the point I will emphasise here is that the spectrum of heavy sterile neutrinos is quite interesting. There are two extremely heavy fully sterile states of mass set by the largest scale $M$. But there is also a heavy weakly-sterile state of mass $m_R^2 / M$ which feels the right-handed weak interactions. It might be interesting to analyse the leptogenesis implications of this kind of heavy sterile spectrum, and to compare the results to those derived from the standard see-saw model.

I come now to my third and final case study: the mirror see-saw mechanism [5]. We return to the $G_{SM} \otimes G_{SM}$ model introduced above. In addition to the standard left-handed neutrino $\nu_L$ and its mirror partner $\nu_R$, we introduce singlet states comprising the “standard” right-handed neutrino $\nu_R$ and its mirror partner $\nu_L'$. The discrete symmetry between the sectors includes the interchanges

$$
\nu_L \leftrightarrow \nu_R', \quad \nu_R \leftrightarrow \nu_L',
$$

(5)

where I have suppressed the Lorentz structure required for a parity transformation.
The neutrino mass matrix is given by
\[
( \bar{\nu}_L \ (\nu_R')^c \ (\nu_R)^c \ \bar{\nu}_L ) \begin{pmatrix}
0 & 0 & m_1 & m_2 \\
0 & 0 & m_2 & m_1 \\
m_1 & m_2 & M_1 & M_2 \\
m_2 & m_1 & M_2 & M_1
\end{pmatrix}
\begin{pmatrix}
(\nu_L)^c \\
\nu_R' \\
\nu_R \\
(\nu_L')^c
\end{pmatrix} + H.c. \tag{6}
\]

Superficially this looks similar to the universal see-saw matrix, but it is actually very different. The entries \(m_{1,2}\) are both of the electroweak scale, while \(M_{1,2}\) are large bare masses within the gauge singlet sector. Notice that \(m_1\) is proportional to the VEV of the standard Higgs doublet, whereas \(m_2\) is driven by the VEV of the mirror Higgs doublet. In the version of the model I am discussing, these two VEVs are equal: the discrete parity symmetry is not spontaneously broken, with improper Lorentz transformations consequently being exact symmetries of both the Lagrangian and the world.

The consequences of this matrix are best revealed by changing to the parity eigenstate basis,
\[
\begin{align*}
\nu_L^+ & \equiv \frac{\nu_L \pm (\nu_R')^c}{\sqrt{2}}, \\
\nu_R^+ & \equiv \frac{\nu_R \pm (\nu_L')^c}{\sqrt{2}}
\end{align*} \tag{7}
\]
which furnishes the mass matrix
\[
\begin{pmatrix}
0 & 0 & 0 & m_+ \\
0 & 0 & m_- & 0 \\
0 & m_- & M_- & 0 \\
m_+ & 0 & 0 & M_+
\end{pmatrix} \tag{8}
\]

The two \(2 \times 2\) blocks make it easy to see that the eigenvalues are of order
\[
\frac{m_+^2}{M_+}, \quad \frac{m_-^2}{M_-}, \quad M_+, \quad M_- \tag{9}
\]
with eigenstates that are predominantly \(\nu_L^+, \nu_L^-, \nu_R^+\) and \(\nu_R^-\), respectively.

The two light eigenstates therefore form a maximally-mixed active-mirror pair, with the mirror component being the weakly sterile \(\nu_R'\). This is reminiscent of pseudo-Dirac structure, but different from it. The important thing is that the mirror see-saw mechanism automatically provides light effectively-sterile neutrinos. The two heavy eigenstates are a pair of maximally mixed fully-sterile states. Some cosmological consequences of the heavy states have been explored in Ref.[8].

1.1.3. Heavy active neutrinos

We have seen that most candidates for heavy sterile neutrinos actually provide weakly sterile varieties which, of course, are not actually sterile at all. The active/sterile distinction is thus arguably less useful for heavy states. Nevertheless for completeness I will now review particle physics motivations for heavy neutral leptons that feel the standard left-handed weak interaction.
Three classes immediately spring to mind: a fourth standard family, but with a neutrino more massive than 45 GeV; an exotic family composed of higher weak-isospin representations that contain neutral leptons; and vector-like weak isospin doublet leptons.

A fourth family with a heavy neutrino is a possibility, but there does not appear to be much motivation for it, and its properties would in any case be quite constrained by precision electroweak phenomenology. I will consider this scenario no further. Exotic families are fun to think about [9], but, again, there does not at present seem to be sufficient motivation to spend much time on them. Apart from noting that compositeness and weak-isospin-cubed models provide some theoretical underpinning, we will move on.

Vector-like weak isospin doublet leptons have the SU(3)$_c$ $\otimes$ SU(2)$_L$ $\otimes$ U(1)$_Y$ quantum numbers

$$\psi_L \sim (1, 2)(-1), \quad \psi_R \sim (1, 2)(-1),$$

so the Dirac mass term $M\psi_L\psi_R + H.c.$ is invariant under $G_{SM}$. Such particles can arise in extended theories where the $\psi$’s transform chirally under a higher symmetry that spontaneously breaks (eventually) to $G_{SM}$. The heavy neutral leptons are the $\tau_3^L = +1$ components of $\psi$.

As an interesting example, consider the subgroup chain

$$E_6 \rightarrow SO(10) \rightarrow SU(5) \rightarrow SU(3)_c \otimes SU(2)_L \otimes U(1)_Y.$$  

The smallest nontrivial representation of $E_6$ decomposes as:

$$27 \rightarrow 16 \oplus 10 \oplus 1 \rightarrow (5^* \oplus 10 \oplus 1) \oplus (5^* \oplus 5) \oplus 1.$$  

The 16 of SO(10) contains a standard family, while the 10 of SO(10) contains a vector-like pair of lepton doublets (and charge $-1/3$ quarks). These SO(10) ten-plet fermions will pick up masses when $E_6$ breaks to SO(10), because they are chiral under the former but vector-like under the latter (the 27 of $E_6$ is a complex representation, whereas the 10 of SO(10) is real).

I do not know if heavy active neutrinos such as these have important cosmological applications. However, I would like to digress about an interesting role the exotic fermions provided by the 27 of $E_6$ could play in the flavour problem.

We have seen that the 27 of $E_6$ can supply heavy vector-like states to partner down quarks and charged leptons. It also contains four neutral leptons in addition to the standard neutrino. It is amusing that exotic partners for up quarks are absent by group theoretic necessity. The raw ingredients therefore exist to explain why

$$m_t \sim 200 \text{ GeV} \gg m_{b, \tau} \gg m_\nu.$$  

One invokes a not-quite-universal see-saw mechanism whereby the bottom and tau lepton masses are suppressed relative to the electroweak scale through mixing with their exotic partners, while the top quark, lacking a partner, cannot have a suppressed mass. One would expect to be able to doubly suppress the lightest active neutrino. This idea works for the third family; obviously some other physics must be invoked to explain why the charm and up masses are much less than the electroweak scale.
Davidson and I invented this framework in 1999 [10]; Rosner independently proposed similar ideas at about the same time [11]. While a complete model exploiting the above vision is still lacking, especially in the neutrino sector, it might be useful to sketch some work-in-progress.

The Higgs multiplets necessary for fermion mass generation can also be found in the 27 of $E_6$. Different components of the Higgs multiplet (or multiplets) must pick up hierarchical VEVs to implement the mechanism. The bottom and tau mass matrices are of the form

$$
\begin{pmatrix}
0 & \ell \\
x & M
\end{pmatrix}
$$

(14)

where $\ell$ is an electroweak scale mass and $x \ll M$ are given by high symmetry breaking scales, with $x/M$ being the intra-family hierarchy parameter.

We have explored three incarnations that provide different choices for which symmetries are broken at scales $x$ and $M$. The three versions utilise the three different electric charge embeddings existing within $E_6$: standard, flipped and double-flipped. To understand this feature, we need to return to the subgroups of $E_6$ and look at them in more detail. A chain of maximal subgroups is

$$
E_6 \rightarrow SO(10) \otimes U(1)^{''} \rightarrow SU(5) \otimes U(1)^{'} \otimes U(1)^{''}
$$

\rightarrow

$$
SU(3) \otimes SU(2) \otimes U(1)_{Y_{st}} \otimes U(1)^{'} \otimes U(1)^{''}.
$$

(15)

The standard hypercharge (and hence electric charge) embedding takes $Y = Y_{st}$; the flipped embedding sees $Y$ as a certain linear combination of $Y_{st}$ and the generator of U(1)$^{'}$; the double-flipped case requires also an admixture of the generator of U(1)$^{''}$. It is fairly easy to explicitly write down these three linear combinations, but I will refrain from doing so here.

The effect of flipping and double-flipping is to rearrange the exotic vector-like fermions within alternate SU(5) components of the 27, hence affecting the mass generation mechanism. For the flipped and double-flipped embeddings, the highest scale $M$ is associated with the breaking of U(1)$^{'} \otimes U(1)^{''}$ down to a U(1) subgroup, while $x$ is correlated with the breakdown of SU(5) times this leftover U(1) down to the standard model group $G_{SM}$. For the standard embedding, $M$ and $x$ are related to a two-step breakdown of U(1)$^{'} \otimes U(1)^{''}$ to a U(1) subgroup and then to nothing, with SU(5) remaining exact.

1.2. Oscillation effects

Neutrino disappearance and flavour transformations are now very well established experimentally [2]. Standard mass and mixing driven oscillations are most probably responsible for these effects. It is therefore important to explore the implications of oscillations for cosmology.²

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² The cosmological effects of absolute neutrino masses constitute a complementary area of study that I will not review here.
Oscillations between light-active and light-sterile neutrino flavours can give rise to dramatic effects for Big Bang Nucleosynthesis (BBN), the topic of the next section. Elsewhere in these proceedings Y. Wong describes how transformations amongst the active neutrinos themselves can lead to neutrino asymmetry equilibration, an important issue for BBN [12]. I will focus instead on how heavy-heavy neutrino oscillations can be responsible for the baryon asymmetry of the universe through a leptogenesis mechanism.\(^3\)

As we have seen, heavy sterile neutrinos are a generic consequence of see-saw mechanisms. In the leptogenesis scenario of Fukugita and Yanagida, out-of-equilibrium decays of heavy neutral leptons into Higgs bosons and doublet neutrinos are used to generate a lepton asymmetry that is then reprocessed by non-perturbative effects in the SM (sphaleron-induced transitions) into a baryon asymmetry [13]. P. di Bari discussed this mechanism at the Workshop. I will review an alternative idea, due to Akhmedov, Rubakov and Smirnov [14].

As in the Fukugita-Yanagida model, three heavy neutral lepton singlets \(N_{A,B,C}\) are produced in the cosmological plasma through Yukawa couplings to the standard doublet neutrinos and Higgs bosons. The dominant process involves top quarks because of their large Yukawa coupling constant. The production is CP symmetric, so equal numbers of \(N\)'s and \(\bar{N}\)'s are created.

The Yukawa coupling constants governing the \(N\)-Higgs-\(\nu\) interactions are chosen in a mild hierarchy so that one or two (but not all three) of the heavy singlets are brought into thermal equilibrium prior to the electroweak phase transition at a temperature \(T_{EW} \approx 100\) GeV. For definiteness, we will follow the inventors by supposing that \(N_A\) and \(N_B\) are equilibrated by \(T_{EW}\), but \(N_C\) is not.

CP violation exists in the mixing matrix relating the Higgs interaction \(N\)-eigenstates and the mass eigenstates. After their production, the \(N_{A,B}\) fermions begin oscillating in a CP asymmetric way while continuing to interact with the background plasma. The CP violating mixings serve to separate the overall zero lepton number into nonzero asymmetries \(L_{A,B,C}\), maintaining \(L_A + L_B + L_C = 0\). The fast Yukawa interactions of \(N_A\) and \(N_B\) transfer their asymmetries to standard doublet neutrinos during the epoch prior to the electroweak phase transition when sphaleron processes are rapid. A nonzero baryon asymmetry results. Crucially, the \(C\)-type asymmetry is never reprocessed into the baryon sector. After the electroweak phase transition, it gets transferred to the light neutrino sector, but by then sphaleron transitions have switched off.

The parameter ranges required for this mechanism to work are different from that of Fukugita-Yanagida. Relatively small singlet masses, in the GeV to 10’s of GeV range, are acceptable for instance. Small Yukawa coupling constants and mixing angles are also a feature according to Ref.[14].

\(^3\) Light-heavy oscillations are quickly decohered by collisions with the background plasma.
2. UPDATE ON BBN AND STERILE NEUTRINOS

Big Bang Nucleosynthesis begins at about $T \simeq 0.8$ MeV, when nuclear statistical equilibrium no longer obtains [15]. Neutrinos feature in two important ways: by contributing relativistic energy density to the plasma thus helping to drive the expansion of the universe, and through the beta equilibrium processes

$$n\nu_e \leftrightarrow pe^-, \quad p\bar{\nu}_e \leftrightarrow ne^+, \quad n \leftrightarrow pe^-\bar{\nu}_e,$$

that maintain the neutron to proton ratio at equilibrium values. Active to light-sterile oscillations can increase the expansion rate because of sterile neutrino production and alter beta-equilibrium by affecting the electron-neutrino and antineutrino distribution functions. The light elements produced through BBN have abundances that depend on the neutron to proton ratio and the relative sizes of nuclear reaction rates and the expansion rate. The primordial abundances thus provide a probe of active-sterile mixing (assuming that is the only modification to standard BBN).

The paradigm known as “standard BBN” is based on the following assumptions:

- There are three neutrino and antineutrino flavours and they are massless.
- All six have Fermi-Dirac distributions with zero chemical potentials.
- There is one free parameter, the baryon to photon ratio $\eta$. It also implicitly assumed that this parameter is homogeneous.

The outputs are the primordial abundances of $^4$He, D, $^3$He and $^7$Li. Before comparing standard BBN with observational data, I would like to comment on the assumptions listed above.

If terrestrial experiments establish the existence of significant active-sterile mixing, then the first assumption of standard BBN, as quoted above, will have to be modified. We will see very shortly what current measurements say about the relativistic energy density during BBN.

For temperatures above the neutrino decoupling temperature of about 1 MeV, Fermi-Dirac (FD) distributions are very well justified because standard weak interactions keep all active neutrino flavours in kinetic equilibrium. After decoupling, the FD forms are maintained unless oscillation effects are operative.

Chemical equilibrium is maintained between neutrinos and antineutrinos of each active flavour $\alpha = e, \mu, \tau$ by well-established electroweak processes such as $\nu_\alpha \bar{\nu}_\alpha \leftrightarrow e^+e^- \leftrightarrow \gamma\gamma$. This enforces the relation $\mu_{\nu_\alpha} = -\mu_{\bar{\nu}_\alpha}$ between the chemical potentials above the $\alpha$-flavour chemical decoupling temperature. Setting these chemical potentials to zero is an extra assumption adopted within standard BBN. The conditions $\mu_{\nu_e} = \mu_{\nu_\mu} = \mu_{\nu_\tau}$ and the analogous one for antineutrinos, of which the standard BBN zero chemical potential assumption is a special case, will be maintained under pure active-active oscillations, but are in general not dynamically consistent if active-sterile oscillations occur. The most dramatic manifestation of this is the exponentially fast creation of neutrino-antineutrino asymmetries through collision-affected active-sterile oscillations.
when certain conditions are satisfied [16].

The zero chemical potential assumption is adopted partly for simplicity and partly because the small observed baryon asymmetry, \( \eta \sim 10^{-10} \), encourages one to suppose the neutrino asymmetries have similar magnitudes. However, because relic neutrinos have not been directly detected, the observational constraints on their asymmetries are quite weak. At least two mechanisms for producing large neutrino asymmetries have been explored – through active-sterile oscillations (as mentioned above) and Affleck-Dine processes [17] – so there is no requirement for the neutrino asymmetries to be as small as their baryon counterpart.

Through the matter effect [18], nonzero asymmetries can suppress sterile neutrino production [19]. This must be taken into account when computing BBN relativistic energy density constraints on active-sterile mixing. An \( e \)-like asymmetry is a special case of distribution function modifications that affect beta-equilibrium and hence the \( ^4\text{He} \) abundance prediction.

Finally, standard BBN for simplicity assumes spatial homogeneity. As primordial abundance determinations improve, it will be interesting to see if reality reflects this simplicity. One interesting possibility in the inhomogeneous case is for a spatially varying \( e \)-like asymmetry to be generated by active-sterile oscillations, seeded by an inhomogeneous baryon distribution [20].

The prediction of standard BBN is summarised in Figure 1, taken from the recent review by G. Steigman [21]. One can see that \( ^4\text{He} \) is by far the most abundant species, and that it is quite insensitive to \( \eta \). For phenomenologically acceptable \( \eta \)’s, deuterium is the next most common isotope, falling fairly steeply with the baryon density (because it is easily destroyed). I will focus on these two species. Current measurements are summarised in Figs. 2-4, taken once again from the Steigman review [21]. One can see that the \( ^4\text{He} \) data clearly approach an asymptotic value with decreasing metallicity, as tracked by oxygen in this case. The deuterium data show more scatter, both as a function of redshift and metallicity.

The implications of these measurements have been recently summarised by di Bari [22], and even more recently by Steigman [21]. At the time of the di Bari paper, there were two basically incompatible \( ^4\text{He} \) abundance determinations, the so-called “high value”

\[
Y_p = 0.244 \pm 0.002
\]  

(17)

of Izotov and Thuan (IT) [23], and the “low value”

\[
Y_p = 0.234 \pm 0.003
\]  

(18)

of Olive and Steigman (OS) [24]. The treatment of systematic effects accounts for the quite serious disparity. The dispassionate may choose to explore the ramifications of these extractions separately, or adopt a “compromise” value, such as the

\[
Y_p = 0.238 \pm 0.005
\]  

(19)

\footnote{Nonzero chemical potentials imply unequal number densities for neutrinos and antineutrinos.}
figure suggested by OS which inflates the error to reflect the systematic uncertainties. I prefer to treat the two determinations separately.

The deuterium abundance range advocated by O’Meara et al [25] is

$$\frac{D}{H} = 3.0 \pm 0.4. \quad (20)$$

After considering all of the available data, Steigman argues for a larger systematic uncertainty and suggests the range

$$\frac{D}{H} = 3.0^{+1.0}_{-0.5} \quad (21)$$

as a safer alternative.

The internal test of standard BBN is whether a unique value for $\eta$ fits both the $^4\text{He}$ and deuterium data. The inferred $3\sigma$ baryon density ranges are:

$$\eta(\text{High Y}) = 2.0 - 7.0,$$
$$\eta(\text{Low Y}) = 0.6 - 3.5,$$
$$\eta(D/H) = 4.5 - 7.7, \quad (22)$$
where the deuterium range of Eq. (20) has been used. (The semi-analytical approach of di Bari [22] is a very convenient tool when faced with extracting $\eta$ from many different primordial abundance determinations.) One can see that the IT and D/H data yield a consistent $\eta$, whereas OS and D/H are inconsistent at a greater than 3$\sigma$ level.

The recent cosmic microwave background radiation (CMBR) anisotropy data now provide an important independent determination of the baryon density [26]. A comparison of the CMBR value with that (or those!) determined from primordial abundance considerations therefore provides a very interesting external test of standard BBN. The CMBR 3$\sigma$ range is

$$\eta(\text{CMBR}) = 3.6 - 9.3.$$  (23)

This is consistent with IT and D/H, but not with OS!

Clearly, the issue of systematic effects in the extraction of the primordial $^4$He abundance is critical. Interestingly, a reanalysis of the data used by IT to derive the High $Y$ range has recently been performed by Peimbert, Peimbert and Luridiana (PPL) [27] (see the discussion by Steigman [21]). PPL take into account systematic effects involving the temperature of different chemical components of the clouds which supply the helium data. They conclude that these effects lead to a lower $Y_p$ than derived by IT! By ignoring...
FIGURE 3. Deuterium abundance as a function of metallicity. Figure taken from Ref.[21] and reproduced by permission from Cambridge University Press.

FIGURE 4. Helium abundance (mass fraction) as a function of metallicity. Figure taken from Ref.[21] and reproduced by permission of Cambridge University Press.
another systematic effect involving the collisional excitation of hydrogen, they deduce

$$Y_p(PPL1) = 0.2356 \pm 0.0020,$$

(24)

whereas a preliminary result incorporating the collisional effect also is

$$Y_p(PPL2) = 0.2384 \pm 0.0025,$$

(25)

which again raises the figure slightly.

My summary of the situation is in Fig. 5. As you can see, the PPL reanalysis apparently strengthens the evidence that standard BBN is not completely consistent with the best elemental abundance and CMBR analyses we have at present. In a nutshell, the CMBR and D/H data are consistent with each other, but are inconsistent at a reasonably significant level with the helium data. An overlap is obtained only when extremes of the $3\sigma$ ranges are used. Given the reasonable doubts one may have about how well systematic effects are understood, I would not want to draw too strong a conclusion just yet. However, it is fair to say that there is a tension between the data sets quoted above. PPL are promising to update their preliminary $Y_p$ figure with the collisional effect incorporated. If it should move up, even a little, relative to what I have called PPL2, then the $\eta$ range will move towards consistency with deuterium and CMBR (recall that small changes in $Y_p$ lead to large changes in the inferred $\eta$). Perhaps the most dramatic new information we expect in the near future is a more precise CMBR figure for $\eta$. The MAP experiment promises to reduce the uncertainty to about 10%, while the future Planck mission may reduce it all the way to about 1%. It is easy to see from Fig. 5 that a high and precise CMBR $\eta$ could really create a crisis for standard BBN.

This rather interesting situation certainly provides some motivation for thinking about non-standard BBN. One way to dramatise the situation is to point out that a lower helium yield can be obtained by reducing the expansion rate of the universe. Parameterising the latter in terms of an effective number of neutrino flavours in equilibrium during BBN, one can fit the light element abundance data better by reducing this parameter from three to about two and a half. Of course, no known physics can actually prevent the full equilibration of all three active species. However, there is a well-known modification of standard BBN that can reduce the helium yield without affecting the expansion rate appreciably: the introduction of a nonzero chemical potential for electron neutrinos and antineutrinos [28]. Using the equilibrium value for the neutron to proton ratio

$$\frac{n_n}{n_p} = \exp\left[\frac{m_p - m_n}{T} - \frac{\mu_{\nu_e}}{T}\right],$$

(26)

we can estimate that $\xi_{\nu_e} \equiv \mu_{\nu_e}/T \sim 0.05$ will reduce $Y_p$ by about 0.01, and thus provide a good simultaneous fit to the helium, deuterium and CMBR data.

Interestingly, active-sterile neutrino oscillations can induce significant chemical potentials or asymmetries [16]. So it has been suggested that rather than being a cosmological liability, light sterile neutrinos might actually be required for a fully successful BBN, with various models explored in recent years [29]. (These scenarios may involve more than just nonzero chemical potentials – oscillation-induced spectral distortion away from
FIGURE 5. Summary of allowed $3\sigma$ ranges for the baryon-to-photon ratio $\eta_{10} \equiv \eta \times 10^{10}$ (horizontal axis) from different observations. Standard BBN would be successful were a single value of $\eta$ consistent with all the data. The different data sets are: CMBR (cosmic microwave background acoustic peaks), $D/H$ (primordial deuterium abundance), PPL+Hc (primordial Helium abundance from Peimbert et al with preliminary incorporation of collision effect), PPLnHc (same with collision effect correction removed), SO (Steigman and Olive “low-Y result” for Helium), Iz&T (Izotov and Thuan “high-Y result”).

Fermi-Dirac form for $T < 1$ MeV may also occur.) However, new experimental data, especially from SNO, and a much better understanding of active-active oscillations in the early universe [12] now require all of these scenarios to be re-analysed.

The active-active story is a rather interesting one. Neutrino oscillation dynamics in the early universe is intrinsically non-linear because neutrinos scatter off other neutrinos in the plasma. For the active-active case, it turns out that this tends to induce synchronisation [30]: neutrinos of all energies are driven to oscillate at the same rate, quite unlike the behaviour when neutrino-neutrino scattering is absent. Perhaps even more importantly, neutrino asymmetries in this case do not suppress oscillation amplitudes, unlike the better-studied active-sterile situation. Yvonne Wong will explain this more fully in her talk. If an active-active $\nu_\alpha \leftrightarrow \nu_\beta$ channel is governed by a large mixing angle, then the $\alpha-$ and $\beta-$like asymmetries will tend towards equilibration. If both the solar and atmospheric neutrino deficits are due to large-angle active-active mixing, and if the solar $\Delta m^2$ is at the upper end of the allowed range, about $10^{-5}$ eV$^2$, then all the asymmetries equilibrate. In these circumstances, the BBN constraints on $L_{\nu_e}$ apply to $L_{\nu_{\mu, \tau}}$ as well [12]. These developments are still new: no one has yet calculated neutrino distribution function and asymmetry evolution in a light sterile neutrino model with active-active transitions correctly accounted for.

Let us now turn to BBN bounds on the relativistic energy density and distribution function distortion: when would agreement between calculations and observation be unambiguously absent? Introduce the parameters $\Delta N^0_\nu$ and $\Delta N^I_\nu$, where the former parameterises the change in the relativistic energy density in terms of an effective change in the neutrino flavour count, while the latter is a coarse-grained measure of the effect of spectral alteration away from zero chemical potential FD form. Following di Bari [22],
we proceed by writing down approximate analytical fits (valid around $\eta \simeq \eta_{\text{CMB}}$) to the primordial abundance yields as a function of $\Delta N_{\nu}^p$:

$$
Y_p(\eta, \Delta N_{\nu}^p) \simeq Y_{p, SBBN}(\eta) + 0.0137 \Delta N_{\nu}^p,
$$

$$
\left( \frac{D}{H} \right)(\eta, \Delta N_{\nu}^p) \simeq \left( \frac{D}{H} \right)_{SBBN}(\eta)[1 + 0.135 \Delta N_{\nu}^p]^{0.8}
$$

(27)

Spectral distortion changes the helium yield but negligibly affects deuterium. The course-grained spectral distortion parameter may thus be defined by

$$
\Delta N_{\nu}^f \equiv \frac{Y_p(\eta, \Delta N_{\nu}^p, \delta f) - Y_p(\eta, \Delta N_{\nu}^p)}{0.0137}
$$

(28)

where $\delta f$ denotes the deviation of the actual $\nu_e$ and $\bar{\nu}_e$ distribution functions from zero chemical potential FD form and is in general a complicated function.

By comparing the measured yields with $Y_p(\eta_{\text{CMB}})$ one obtains $3\sigma$ allowed ranges for

$$
N_{\nu}^{\text{tot}} \equiv 3 + \Delta N_{\nu}^p + \Delta N_{\nu}^f.
$$

(29)

The preference for an effective neutrino number of less than three is clearly displayed in Fig. 6. Note that deuterium yields only weak constraints. Cosmic microwave background data also at present provide only weak constraints on $\Delta N_{\nu}^p$. The maximum increase in $N_{\nu}^{\text{tot}}$ one can tolerate before BBN fails unambiguously is about 0.3, taking the original IT high-$Y$ figures uncritically. Note in particular, that a fully equilibrated sterile flavour corresponding to $N_{\nu}^{\text{tot}} = 4$ would cause BBN failure by a comfortable margin (assuming the absence of a negative $\Delta N_{\nu}^f$ that compensates for $\Delta N_{\nu}^p$).

![Figure 6](image-url)

**FIGURE 6.** The $3\sigma$ ranges for the preferred effective total number of neutrino flavours implied by the different primordial helium abundance extractions. See the caption of Fig. 5 for an explanation of the labels on the vertical axis.

In the future, in addition to a baryon density at $10 - 1\%$ precision, CMBR acoustic peak data may be precise enough to allow a probe of relativistic energy density at the
\[ \pm 0.1 \text{ level as quantified by } \Delta N_\rho. \] If \( \Delta N_\rho + \Delta N_\nu < 0 \) continues to be supported by helium data, especially if better understanding of systematic effects increases confidence in the extracted primordial abundance figure, then a combination BBN and CMBR data could provide evidence for a nonzero \( e \)-type asymmetry!\(^5\)

As our final topic, we turn to the BBN cosmology of \( 2+2 \) and \( 3+1 \) neutrino models. This was analysed in great detail by di Bari [22] (see also [31]), who concluded that essentially all of the possible scenarios yield parameters in the range

\[ \Delta N_\rho = 0.9 - 1.0, \quad \Delta N_\nu \simeq 0. \] (30)

Basically, the oscillation parameter conditions required for the active-sterile asymmetry generation mechanism to work are not met, so negligible asymmetries result and the sterile flavour is fully or almost fully equilibrated. The only loophole is in the \( 2+2 \) model which has the \( \nu_{\mu,s} \) couplet lighter than the \( \nu_{e,\tau} \) couplet. If the mixing angles between \( \nu_{e,\tau} \) and the lighter \( \nu_s \) are very small, then some asymmetry generation may occur, but not enough to prevent BBN failure. In a couple of subcases, the addition of a second sterile neutrino flavour can allow one to engineer an acceptable outcome, but the resulting model looks rather forced.

These conclusions motivate a reconsideration of the “large pre-existing neutrino asymmetry” idea [19]. One supposes that some mechanism operating at higher temperatures produces asymmetries that are large enough to suppress active-sterile mixing via the matter effect and hence to also suppress sterile neutrino production. For instance, \( \nu_\mu - \nu_s \) mixing with atmospheric range oscillation parameters requires an asymmetry high than about \( 10^{-5} \). While this is five orders of magnitude higher than the baryon asymmetry, it is still very small by BBN standards. In particular, an \( L_{\nu_e} \sim 10^{-5} \) generated through the active-active asymmetry equilibration mechanism would have a negligible effect on the helium abundance. What happens when one moves away from pure \( \nu_\mu - \nu_s \) mixing is currently under investigation.

A particularly amusing possibility is that

1. MiniBooNE confirms the LSND effect, with its indirect evidence for a light sterile neutrino;
2. future long baseline experiments garner \textit{direct evidence} for a light sterile neutrino through neutral to charged current ratio measurements, thus confirming the oscillation explanation for the LSND anomaly;
3. BBN remains incompatible with four (or more) equilibrated species.

These circumstances would provide remarkable evidence for the existence of fairly large pre-existing lepton asymmetries. This would obviously be important information for high temperature models of leptogenesis and baryogenesis.

\(^5\) Beware that BBN and CMBR data probe the universe during different epochs, so one has to assume that the physics of the intervening period is understood when one draws conclusions using combined information.
3. CONCLUSIONS

The known neutrinos play an important role in cosmology, especially in Big Bang Nucleosynthesis through the relativistic energy density and hence expansion rate of the universe, and through the neutron-proton interconversion reactions mediated by $\nu_e$ and $\overline{\nu}_e$. Much more speculatively, heavy neutrino-like states might be responsible for producing the cosmological baryon asymmetry through the sphaleron reprocessing of a lepton asymmetry.

I have reviewed how additional neutrino-like states can arise in extensions of the Standard Model of particle physics. Three see-saw mechanisms – the standard, the universal, the mirror – were presented as case studies of how additional neutrino flavours may be required for understanding mass generation and mass hierarchies. All three furnish heavy neutrinos, and the mirror see-saw supplies possibly the best candidate theory for light sterile neutrinos.

The active/sterile distinction was critically reviewed, with sterility subdivided into full-blown and weak varieties, depending on whether or not those neutrino-like states couple to hypothetical gauge interactions beyond the known left-handed weak interaction. The additional interactions, if chiral, would provide a natural mass scale to the weakly sterile neutrinos through spontaneous symmetry breaking. Also, these interactions would affect the cosmology of the heavy neutral leptons, and hence also mechanisms of leptogenesis. I reviewed the Akhmedov-Rubakov-Smirnov leptogenesis model as an alternative to the better known Fukugita-Yanagida scenario.

I digressed on the subject of why the top quark has a weak scale mass while the other third family fermions have much smaller masses. I pointed out that the mathematics of the 27 of $E_6$ through a not-quite-universal see-saw mechanism might provide the answer.

The present status of Big Bang Nucleosynthesis was then discussed. We observed that there is a tension between the helium abundance data and both the deuterium and cosmic microwave background data. A recent reanalysis of helium data by Peimbert, Peimbert and Luridiana [27] strengthens the case for the inconsistency of standard BBN. The problem of systematic errors in the extraction of primordial abundances cautions against overly strong conclusions, however [21]. Future precise determinations of the baryon density and relativistic energy density from CMBR acoustic peak data promises to challenge standard BBN in important ways.

So, there is motivation to consider non-standard BBN. A nonzero $e$-like asymmetry can resolve the tension by lowering the helium abundance without affecting the deuterium abundance, the latter being in good agreement with the independent CMBR determination of the baryon density. Such an asymmetry can be generated by active-sterile oscillations. Previous specific models of this type need to be revised in light of the SNO solar neutrino results and our increased understanding of the early universe dynamics of active-active oscillation channels.

Finally, we observed that the much discussed $2 + 2$ and $3 + 1$ models lead almost inevitably to equilibration of the sterile flavour without the compensating generation of an $e$-like asymmetry. This argues for the existence of asymmetries generated at higher temperature scales of sufficient magnitude to suppress sterile neutrino production (and avoid being washed out) [19]. Should MiniBooNE confirm the LSND anomaly, and if primordial abundance data remain inconsistent with four equilibrated neutrino species,
then some remarkable evidence will exist for significant neutrino asymmetries generated at epochs significantly before that of BBN.

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REFERENCES

1. See A. D. Dolgov, Phys. Rep. 370, 333 (2002) for an extensive review and a thorough list of references.
2. Y. Fukuda et al., SuperKamiokande Collaboration, Phys. Rev. Lett. 81, 1158 (1998); 81, 4279(E) (1998); 82, 1810 (1999); hep-ex/0103032; J. N. Abdurashitov et al., SAGE Collaboration, Phys. Atom. Nucl. 63, 943 (2000); W. Hampel et al., GALLEX Collaboration, Phys. Lett. B447, 127 (1999); B. T. Cleveland et al., Homestake Collaboration, Ap. J. 496, 505 (1998); M. Altmann et al., GNO Collaboration, Phys. Lett. B490, 16 (2000); Y. Fukuda et al., Kamiokande Collaboration, Phys. Rev. Lett. 77, 1683 (1996); Q. R. Ahmed et al., SNO Collaboration, Phys. Rev. Lett. 87, 071301 (2001); Y. Fukuda et al., SuperKamiokande Collaboration, Phys. Rev. Lett. 82, 2644 (1999); Y. Fukuda et al., Kamiokande Collaboration, Phys. Lett. B335, 237 (1994); W. A. Mann (for the Soudan 2 Collaboration), Nucl. Phys. Proc. Suppl. 91, 134 (2000); R. Becker-Szendy et al., Nucl. Phys. Proc. Suppl. 38, 331 (1995); M. Ambrosio et al., MACRO Collaboration, Phys. Lett. B478, 5 (2000); C. Athanassopoulos et al., LSND Collaboration, Phys. Rev. Lett. 75, 2650 (1995); 77, 3082 (1996); 81, 1774 (1998); the LSND results will be soon checked by the MiniBooNE experiment: MiniBooNE Collaboration (R. Stefanski for the collaboration), Nucl. Phys. Proc. Suppl. 110, 420 (2002).
3. M. Gell-Mann, P. Ramond and R. Slansky, in Supergravity, eds. P. van Nieuwenhuizen and D. Z. Freedman (North Holland, 1979); T. Yanagida, in Proc. Workshop on Unified Theory and baryon Number of the Universe, eds. A. Sawada and H. Sugawara (KEK 1979); see also, R. N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. 44, 912 (1980).
4. A. Davidson and K. C. Wali, Phys. Rev. Lett. 59, 393 (1987); 60, 1813 (1988); S. Rajpoot, Phys. Lett. B191, 122 (1987).
5. R. Foot and R. R. Volkas, Phys. Rev. D52, 6595 (1995).
6. See R. R. Volkas, Prog. Part. Nucl. Phys. 48, 161 (2002), for additional discussion.
7. T. D. Lee and C. N. Yang, Phys. Rev. 104, 254 (1956); R. Foot, H. Lew and R. R. Volkas, Phys. Lett. B272, 67 (1991); Mod. Phys. Lett. A7, 2567 (1992); R. Foot, Mod. Phys. Lett. A9, 169 (1994); See also, I. Yu. Kobzarev, L. B. Okun and I. Ya. Pomeranchuk, Yad. Fiz. 3, 1154 (1966) [Sov. J. Nucl. Phys. 3, 837 (1966)]. For models with broken mirror symmetry see R. Foot and H. Lew, hep-ph/9411390; R. Foot, H. Lew and R. R. Volkas, JHEP 0007, 032 (2000); Z. G. Berezhiani and R. N. Mohapatra, Phys. Rev. D52, 6607 (1995); Phys. Lett. B375, 26 (1996); Z. Berezhiani, Acta Phys. Pol. B27, 1503 (1996).
8. N. F. Bell, Phys. Lett. B479, 257 (2000).
9. P. M. Fishbane, S. Meshkov and P. Ramond, Phys. Lett. B134, 81 (1984); P. M. Fishbane, S. Meshkov, R. E. Norton and P. Ramond, Phys. Rev. D31, 1119 (1985); R. Foot, H. Lew, R. R. Volkas and G. C. Joshi, Phys. Rev. D39, 3411 (1989).
10. A. Davidson and R. R. Volkas, unpublished notes.
11. J. Rosner, Phys. Rev. D61, 097303 (2000).
12. A. Dolgov, S. Hansen, S. Pastor, S. Petcov, G. Raffelt and D. Semikoz, Nucl. Phys. B632, 363 (2002); Y. Y. Y. Wong, Phys. Rev. D66, 025015 (2002); K. N. Abazajian, J. F. Beacom and N. F. Bell, Phys. Rev. D66, 013008 (2002).
13. M. Fukugita and T. Yanagida, Phys. Lett. B174, 45 (1986). For recent work see W. Buchmüller, P. Di Bari and M. Plümacher, Nucl. Phys. B643, 367 (2002); hep-ph/0209301.
14. E. Akhmedov, V. Rubakov and A. Smirnov, Phys. Rev. Lett. 81, 1359 (1998).
15. E. W. Kolb and M. S. Turner, *The Early Universe* (Addison-Wesley, 1990).
16. R. Foot, M. J. Thomson and R. R. Vokas, Phys. Rev. D53, 5349 (1996); X. Shi, Phys. Rev. D54, 2753 (1996); R. Foot and R. R. Vokas, D55, 5147 (1997); P. Di Bari, R. Foot, R. R. Vokas and Y. Y. Wong, Astropart. Phys. 15, 391 (2001). See M. Prakash et al., Annu. Rev. Nucl. Part. Sci. 51, 295 (2001) and Ref. [1] for reviews and additional references.
17. I. Affleck and M. Dine, Nucl. Phys. B249, 361 (1985).
18. L. Wolfenstein, Phys. Rev. D17, 2369 (1978); D20, 2634 (1979); S. P. Mikheyev and A. Yu. Smirnov, Nuovo Cimento C9, 17 (1986).
19. R. Foot and R. R. Vokas, Phys. Rev. Lett. 75 (1995) 4350.
20. P. Di Bari, Phys. Lett. B482, 150 (2000).
21. G. Steigman, “Primordial Alchemy: from the Big Bang to the present Universe” in *Cosmochemistry: the Melting Pot of the Elements*, eds. C. Esteban, R. J. Garcia Lopez, A. Herreo and F. Sanchez [Cambridge University Press, in press (April 2004)], astro-ph/0208186.
22. I. Izotov and T. X. Thuan, Ap. J. 500, 188 (1998).
23. K. Olive and G. Steigman, Ap. J. S. 97, 49 (1995).
24. M. O'Meara et al., Ap. J. 552, 718 (2001).
25. P. de Bernardis et al., Nature (London) 404, 955 (2000); S. Hanany et al., Ap. J. 545, L5 (2000); N. W. Halverson et al., Ap. J. 568, 38 (2002); A. E. Lange et al., Phys. Rev. D63, 042001 (2001); A. Balbi et al., Ap. J. 545, L1 (2000); C. B. Netterfield et al., Ap. J. 571, 604 (2002); C. Pryke et al., Ap. J. 568, 46 (2002); R. Stompor et al., Ap. J. 561, L7 (2001); T. J. Pearson et al., astro-ph/0205388; B. S. Mason et al., astro-ph/0205384.
26. A. Peimbert, M. Peimbert and V. Luridiana, Ap. J. 565, 668 (2002).
27. For a recent analysis see J. P. Kneller, R. J. Scherrer, G. Steigman and T. P. Walker, Phys. Rev. D64, 123506 (2001).
28. R. Foot and R. R. Vokas, Phys. Rev. D56, 6653 (1997); (E) D59, 029901 (1999); D61, 043507 (2000); N. F. Bell, R. Foot and R. R. Vokas, Phys. Rev. D58, 105010 (1998); R. Foot, Phys. Rev. D61, 023516 (2000).
29. S. Samuel, Phys. Rev. D48, 1462 (1993); V. A. Kostelecky and S. Samuel, Phys. Rev. D52, 621 (1995); S. Samuel, Phys. Rev. D53, 5382 (1996); J. Pantaleone, Phys. Rev. D58, 073002 (1998); S. Pastor, G. G. Raffelt and D. V. Semikoz, Phys. Rev. D65, 053011 (2002).
30. K. N. Abazajian, astro-ph/0205238.