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Ultra-light scalar saving the $3+1$ neutrino scheme from the cosmological bounds

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ABSTRACT

The LSND and MiniBooNE results as well as the reactor and Gallium anomalies seem to indicate the presence of a sterile neutrino with a mass of ~ 1 eV mixed with active neutrinos. Such sterile neutrino can be produced in the early universe before the neutrino decoupling, leading to a contribution to the effective number of neutrinos (N_{eff}) as well as to a contribution to the sum of neutrino masses which are in tension with cosmological observations. We propose a scenario to relax this tension by a Yukawa coupling of the sterile neutrinos to ultra-light scalar particles which contribute to the dark matter in the background. The coupling induces an effective mass for ν_s which prevents its production in the early universe. We discuss the implications for the upcoming KATRIN experiment and future relic neutrino search experiments such as PTOLEMY. We also briefly comment on certain non-renormalizable forms of interaction between ν_s and the scalar and their consequences for the ν_s production in the early universe.

1. Introduction

The 3 neutrino mass and mixing scheme has been established as the standard paradigm to explain the results from various solar, atmospheric, long baseline and reactor neutrino experiments. However, there are a few hints that may point out the existence of a fourth sterile neutrino (ν_s) with a mass of ~ 1 eV mixed with active neutrinos. This is the essence of the so-called $3+1$ neutrino scheme which has been invoked to explain the LSND [1] and MiniBooNE [2] as well as the Gallium and reactor neutrino anomalies [3–5,8]. To explain the LSND and MiniBooNE anomalies, ν_μ should partially convert en-route into ν_e which implies that the sterile neutrino has to be mixed with ν_e and ν_μ, simultaneously. From ICECUBE and MINOS+, strong bounds are derived on the ν_s mixing with ν_μ shedding doubt on the $3+1$ oscillation solution to the LSND and MiniBooNE anomalies [6,7,9]. However, the reactor [4,5,8] and Gallium [3] anomalies (the observation that at short baselines $P(\nu_e \to \nu_e)$, $P(\bar{\nu}_e \to \bar{\nu}_e) < 1$) can be explained even if ν_s mixes only with ν_e so this solution is not ruled out by the ICECUBE or MINOS+ results which are based on the $\nu_\mu \to \nu_\mu$ observation. A recent analysis shows that the solutions to the Gallium and reactor anomalies are compatible with each other within the $3+1$ neutrino mixing scheme with $|U_{e4}|^2 \sim 0.01-0.02$ [10]. What makes this possibility even more exciting is that a sterile neutrino mixed with ν_e will lead to observable kinks in the spectrum of beta decay [11]. The upcoming KATRIN results can test the $3+1$ solution to reactor and Gallium anomalies [12].

On the other hand, the mixing of ν_s with ν_a implies that in the early universe before neutrinos decouple from the plasma, the neutrino oscillation brings the sterile neutrinos to thermal equilibrium with the active ones [13]. This means the effective relativistic degrees of freedom will increase by 1 unit ($N_{eff}=4$) which is disfavored by the CMB data [14] as well as by the Big Bang Nucleosynthesis (BBN). Moreover, the production of ν_s with a mass of 1 eV in the early universe will violate the upper bound on the sum of neutrino masses which are derived by combining the CMB and BAO results [14].

To avoid the bounds from cosmology, various models for self-interaction of ν_s has been proposed [15–19]. (See, however, [20–22].) The essence of all these scenarios is that the self-interaction of ν_s will induce an effective mass for ν_s at $T >$ MeV which will suppress the effective mixing and will therefore decrease the ν_a to ν_s oscillation probability, $P(\nu_a \to \nu_s)$. Notice that within these scenarios, the generated effective mass itself is given by the ν_s density. That is in order for the mechanism to be efficient, a nonzero ν_s density is required in the first place. This way the bound on N_{eff} can be satisfied but the bound on the sum of

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masses cannot be avoided. A recent thorough study shows that by adding the BAO data, the self-interaction scenario of ν_s characterized by an effective four-Fermion interaction will be still ruled out by the BAO data [23]. However, if the scenario involves light states coupled to ν_4 which open up the possibility of the removal of ν_4 by annihilation [24] or decay before the onset of structure formation (before the matter radiation equality) the BAO+CMB bound can be avoided, too. An alternative remedy is the late phase transition scenario proposed in [25]. A non-renormalizable coupling of ν_s to background scalar is suggested in [26] and its consequences for the ν_s abundance is discussed.

Ref. [27] proposes a $U(1)$ gauge model with a gauge boson of mass 10 eV coupled to ν_s as well as to asymmetric dark matter. The coupling creates an effective mass for ν_s proportional to dark matter density which is sizable even for vanishing ν_s density. The effective mixing at $T >$ MeV will be then suppressed, preventing the ν_s production. Moreover, the new gauge interaction opens up the possibility of relatively fast decay of the ν_4 components of the active neutrinos before the onset of structure formation. As a result, both the bound on N_{eff} from CMB and BBN and the bound on the sum of neutrino masses from BAO and CMB can be satisfied.

In this letter, we propose a scenario for making the $3+1$ solution to the short baseline anomalies compatible with cosmology based on a Yukawa coupling of ν_s to ultra-light real scalar which may be considered as dark matter. In section 2, we describe the scenario and demonstrate how it solves the tensions with cosmology. In section 3, we discuss the implications for KATRIN and PTOLEMY and formulate strategies to combine various observations to eventually elucidate the mechanism behind the absence of ν_s in the early universe. Section 4 summarizes our results.

2. The model

It is well-known that if dark matter (or a component of it) is of bosonic type, it can be as light as $\sim 10^{-21}$ eV. Despite its small mass, the ultra-light dark matter is considered to be cold because its production is non-thermal. These particles can be non-relativistic even at high temperatures. Recently such dark matter has gained popularity in the literature as it has been advocated as a solution to the small scale structure tensions within the WIMP scenario [28]. As long as their de Broglie wavelength is larger than their average distance with each other, they can be described by a classical field. In particular, a real ultra-light scalar dark matter can be described as

$$\phi = \frac{\sqrt{2\rho_\phi}}{m_\phi}\cos(m_\phi t - \vec{p}_\phi \cdot \vec{x}) \qquad (1)$$

where $|\vec{p}_\phi| \ll m_\phi$. For $t \gg 1/m_\phi$, ρ_ϕ (like in the case of other non-relativistic relics) scales as T^3. For $t \ll 1/m_\phi$, it can be shown that ρ_ϕ (the 00 element of the energy-momentum tensor, $T_{\mu\nu}$) is equal to minus the pressure, $-p_\phi$ (T_{ii}). Thus, the relation $T^{\mu\nu}_{;\nu} = 0$ or equivalently $\dot{\rho}_\phi + 3H(\rho_\phi + p_\phi) = 0$ implies that for $t \ll 1/m_\phi$, ρ_ϕ and therefore the amplitude of ϕ remains constant.

There is a vast literature discussing the production of such light particles in the early universe with non-relativistic velocities; for a review see [29]. In these mechanisms, ϕ is taken to be the phase of a complex scalar, $\Phi = |\Phi| e^{i\phi/f_0}$ with a $U(1)$ symmetry similar to the axion of the Peccei-Quinn symmetry [30]. Once the symmetry becomes spontaneously broken, ϕ obtains a random value in the range $(-f_0\pi, f_0\pi)$. If the symmetry breaking takes place during the inflation, the whole patch within our horizon will have the same value of initial ϕ (plus small fluctuations). The topological defects between the patches of different constant ϕ will safely be diluted away during inflation. It has been demonstrated in the literature that with this mechanism ϕ (consequently, ρ_ϕ) can be large enough to account for all DM. As we shall see, our scenario works even for smaller values of ϕ. The quantum fluctuations during inflation can provide the small variation in ρ_ϕ which provides the seeds for structure formation during the course of the history of the universe. Such small variation of ρ_ϕ is irrelevant to our discussion.

Refs. [31–33] assume a Yukawa coupling between ϕ and active neutrinos and discuss the implication of the oscillatory behavior of ϕ with time on the temporal modulation of various neutrino beams. Ref. [34] assumes a gauge interaction between complex ultra-light scalar DM and leptons and studies its impact on the flavor ratios of cosmic neutrinos detected by ICECUBE. Here, we assume a Yukawa coupling of the following form between real ϕ and ν_s

$$\lambda \phi \nu_s^T c \nu_s + \text{H.}c. \qquad (2)$$

where c is an asymmetric 2×2 matrix with components equal to ± 1. Notice that this coupling is renormalizable and invariant under the SM gauge group. As long as ϕ is lighter than the lightest neutrino mass eigenstate, ϕ remains stable and therefore a suitable dark matter candidate. Notice that we could write the interaction of type $\phi \bar{\nu}_s \nu_s$ with similar results but to avoid adding new degrees of freedom, we stick to this Majorana form which does not require right-handed component for ν_s.

The coupling in Eq. (2) induces an effective mass for ν_s given by

$$m_{eff} = \lambda \frac{\sqrt{2\rho_\phi}}{m_\phi}\cos(m_\phi t). \qquad (3)$$

Taking $m_\phi < 5 \times 10^{-17}~\text{eV} = \frac{1}{13~\text{sec}}$, for up to after neutrino decoupling (to be precise until $T \sim 0.22~\text{MeV}(m_\phi/(5\times 10^{-17}~\text{eV}))^{1/2}$), m_{eff} remains almost constant and equal to $m_{eff} = \lambda\sqrt{2\rho_\phi^{int}}/m_\phi$ where ρ_ϕ^{int} is the value of ρ_ϕ at $t \ll 1/m_\phi$. Taking for example $\rho_\phi^{int} = \rho_{DM}^0(0.22~~\text{MeV}\sqrt{m_\phi/(5\times 10^{-17}~\text{eV})}/T^0)^3$ (where the 0 superscript denotes the values today), we find $m_{eff} = 2.3 \times 10^{24}~\text{eV}\lambda(5\times 10^{-17}~\text{eV}/m_\phi)^{1/4}$. Notice that the format of the effective mass that ν_s receives is of the Lorentz invariant Majorana type which should be summed with the ν_s mass in vacuum (m_{ν_s}) to obtain dispersion relation *i.e.*, $E_{\nu_s}^2 - |\vec{p}_{\nu_s}|^2 = (m_{eff} + m_{\nu_s})^2$. Using the superradiance argument, a vector dark matter with a mass of $6\times 10^{-20} - 2\times 10^{-17}$ eV is constrained [35] but these bounds do not apply for the scalar dark matter. The superradiance bound from M87* rules out only the scalar dark matter of mass of 10^{-21} eV and lower [36].

Remember that in the case of the propagation of the active neutrinos in matter, the Lorentz violating effective mass of active neutrinos in medium (*e.g.*, $2\sqrt{2}G_F n_e \nu_e^\dagger \nu_e$) is added to $m_\nu^2/(2E_\nu)$ to obtain the Hamiltonian governing the neutrino flavor evolution. Here, the Lorentz conserving m_{eff} should be added to m_{ν_s} rather than to m_ν^2/E_ν. In the presence of $m_{eff} \gg m_{\nu_s}$, we can write the effective active sterile mixing angle as

$$\sin 2\theta_m|_T = \sin 2\theta \frac{m_{\nu_s}}{m_{eff}} = \begin{cases} \sin 2\theta_m^{int} & \text{at} \quad t \ll m_\phi^{-1}, \\ \sin 2\theta_m^{int}\left(\frac{0.22~\text{MeV}\sqrt{m_\phi/5\times 10^{-17}~\text{eV}}}{T}\right)^{3/2} & \text{at} \quad t \gg m_\phi^{-1}, \end{cases} \qquad (4)$$

where θ is the mixing angle in vacuum. At early universe when $T > 1$ MeV, the active neutrinos undergo scattering off the neighboring neutrinos and electrons. Each electroweak scattering will convert them to coherent active states without any ν_s component. To compute the oscillation probabilities, the evolution of full density matrix has to be computed [13] which is beyond the scope of the present paper. However, for $\sin^2 2\theta_m^{int} \ll 1$, a simplified estimate can be made as follows [37]: The rate of ν_a to ν_s conversion, $\Gamma_{\nu_a\to\nu_s}$, can be estimated as

$$\Gamma_{\nu_a\to\nu_s} = \frac{\sin^2 2\theta_m^{int}}{4\tau_\nu}$$

where τ_ν^{-1} is the interaction rate of neutrinos $\tau_\nu^{-1} \sim G_F^2 T^5$. Thus, the contribution to N_{eff} can be evaluated as

$$\delta N_{eff} = \int_{T_{min}}^{T_{max}} \Gamma_{\nu_a\to\nu_s} dt = \frac{\sin^2 2\theta_m^{int}}{4} \int_{T_{min}}^{T_{max}} \frac{1}{\tau_\mu} dt \, ,$$

where T_{min} is the neutrino decoupling temperature and T_{max} is the maximum temperature for which $(\Delta m^2/T)t \stackrel{>}{\sim} 1$. Notice that we use the fact that for $m_\phi \stackrel{<}{\sim} 5\times 10^{-17}$ eV up until T_{min}, m_{eff} and therefore $\sin^2 2\theta_m^{int}$ remain constant. Within the canonical $3+1$ scheme (in the limit $\sin 2\theta_m = \sin 2\theta$), Ref. [13] shows that for $\sin^2 2\theta \sim 4\times 10^{-4}$ and $\Delta m^2 \sim 1$ eV^2, the contribution to N_{eff} is reduced to 0.1. Scaling these results, we conclude that taking $\sin^2 2\theta_m^{int} = 4\times 10^{-5}$, the contribution will be less than $O(0.01)$ and therefore negligible. For $\Delta m^2 \sim 3$ eV^2 and $|U_{e4}|^2 \sim 2\times 10^{-2}$ (a typical solution to the Gallium and reactor neutrino anomalies [10] which is consistent with the most recent DANSS and STEREO bounds [38]), $\sin^2 2\theta_m^{int} = 4\times 10^{-5}$ can be achieved with $m_{eff}^{int} > 40$ eV which for ρ^{int} corresponding to ρ_{DM} implies $\lambda > 2\times 10^{-23}$. That is taking $\lambda \stackrel{>}{\sim} 2\times 10^{-23}(m_{\nu_s}/1 \text{ eV})$, the bound on N_{eff} can be safely relaxed but below $T \sim 0.01$ MeV (well above the matter radiation equality era) as well as in the Milky Way, m_{eff} can be neglected because ρ_ϕ and therefore the amplitude of ϕ will be suppressed.

Let us now discuss how the bounds from BAO and CMB on the sum of neutrino masses can be avoided. As we discussed, by choosing $\lambda > 10^{-23}$, the density of the ν_s particles produced at $T \stackrel{>}{\sim}$ MeV can be reduced to an arbitrarily small value. The contribution of them to the sum of the neutrino masses can be estimated as $\delta N_{eff} m_{\nu_s}$. Thus, as long as $\delta N_{eff} \stackrel{<}{\sim} 0.01$, the contribution is well below the bound on the sum of neutrino masses, $\sum_\nu m_\nu$ [14].

For $t \stackrel{>}{\sim} 1/m_\phi$, the ϕ field will start oscillating so m_{eff} can be even negative. This will have two dramatic consequences: (1) At certain epochs, ν_s can become lighter than even m_ϕ, opening the possibility of decay of ϕ to ν_s; (2) ν_s can become degenerate with active neutrinos[1] paving the way for non-adiabatic conversion of active neutrinos to ν_s despite the fact that $\dot{m}_{eff} = m_\phi m_{eff} \tan(m_\phi t) \ll m_{\nu_s}$. Let us discuss the consequences of each case.

Even when ν_s becomes lighter than ϕ, the perturbative lifetime of ϕ (*i.e.*, $4\pi/\lambda^2 m_\phi$) will be greater than the age of universe, however; as shown in [40], the ϕ field can convert into ν_s and $\bar{\nu}_s$ pairs through a mechanism known as parametric resonance production. During the epoch of our interest, the radiation dominates so $\rho_\phi \ll \rho_{\nu_a}$. Thus, even if ϕ completely decays into ν_s, the effects of the produced ν_s on cosmological observation will be negligible. If ϕ completely decays, another particle should play the role of dark matter.

The non-adiabatic conversion of ν_1, ν_2 and ν_3 to ν_4 can lead to a tension with the bounds from BAO and CMB on the sum of neutrino masses. However, such tension can be solved by opening up decay modes for ν_4. In fact, the λ coupling itself opens up a decay mode but the lifetime (being given by $(\sum_{i=1}^3 |U_{s4}|^2|U_{si}|^2\lambda^2\Delta m_{4i}^2/(4\pi E_4))^{-1}$ will be longer than the age of the universe. We can however introduce a new singlet ϕ' with an $SU(3)\times SU(2)\times U(1)$ invariant coupling of $\lambda'\phi'\nu_s^T c\nu_s$ which can result in decay $\nu_i \to \bar{\nu}_j\phi'$ with a rate of $\lambda'^2\Delta m_{ij}^2|U_{sj}|^2|U_{si}|^2/(4\pi E_i)$ where we have neglected the ϕ' mass. Notice that the lifetime of ν_i relative to that of ν_4 will be longer by a factor of $(\Delta m_{4j}^2|U_{s4}|^2)/(\Delta m_{ij}^2|U_{si}|^2)$. This means if we choose λ' in a range that ν_4 decays during $T =$ few eV-10 eV, the rest of ν_i will be free streaming at the matter radiation equality era [41][2] but they can decay after recombination era. With a single nonzero $U_{\alpha 4}$ ($\alpha \in \{e,\mu,\tau\}$), the unitarity of the neutrino 4×4 mixing matrix implies that all elements U_{s1}, U_{s2} and U_{s3} should be nonzero so the coupling $\lambda'\phi'\nu_s^T c\nu_s$ leads to eventual decay of all neutrinos to the lightest mass eigenstate; *i.e.*, $j=1$ for normal ordering and $j=3$ for inverted ordering. This also means that the decay of lighter ν_i does not take place for cosmic neutrinos [41]. The required lifetime of ν_4 can be achieved with $\lambda' = 3.5\times 10^{-12}$. With such small λ', the lighter neutrino mass eigenstates, ν_1, ν_2 and ν_3 as well as the produced ϕ' will be free streaming during $0.1 \text{ eV} < T \ll$ MeV. Taking ϕ' massless or with a mass much smaller than that of ν_1, its contribution to the $\sum m_\nu$ measurement from CMB and BAO will be negligible. The total energy in the form of ν_1, ν_2, ν_3 and ϕ will be equal to that of three neutrinos within the standard scheme leading to the same signatures as the standard 3ν scheme. Thus, the bounds from CMB and BAO can be safely satisfied. With $\lambda' = 3.5\times 10^{-12}$, the lifetime of ν_4 will be too long to be relevant for terrestrial, solar and even galactic supernova neutrinos.

If instead of the renormalizable Yukawa coupling in Eq. (2), we had taken a non-renormalizable coupling of form $\phi^2\nu_s^T c\nu_s$ as [26] or of form $i(\phi^*\partial_\mu\phi - \phi\partial_\mu\phi^*)\bar{\nu}\gamma^\mu\nu$ as [34], m_{eff} could not have become negative so the two consequences of $m_{\nu_s} + m_{eff} \to 0$ enumerated above would not have applied. The active background neutrinos at the start of the ϕ oscillation are mainly composed of $\tilde{\nu}_1$, $\tilde{\nu}_2$ and $\tilde{\nu}_3$, with a small contribution given by $\sin\theta_m$ of $\tilde{\nu}_4$ where "$\sim$" emphasizes that these are the energy eigenstates inside the (dark) matter medium. The neutrino propagation for the aforementioned non-renormalizable coupling will remain adiabatic so after the amplitude of ϕ diminishes due to the expansion, the background will mainly consist of the vacuum mass eigenstates ν_1, ν_2 and ν_3 with a small contribution from ν_4 given by $\sin^2\theta_m^{int} \sim 10^{-5}$. As a result, the contribution from ν_4 to $\sum m_\nu$ will be negligible so satisfying the bounds from CMB and BAO on $\sum m_\nu$ will not require a ν_4 decay mechanism.

Let us now discuss the stability of the ϕ mass in the presence of the λ coupling. This coupling is similar to the top Yukawa coupling in the SM and will similarly induce a quadratically divergent mass for ϕ. Like the standard model, we assume that there is a yet unknown mechanism (*e.g.*, SUSY) which cancels this divergent contribution. Still to have a "natural model", we should check whether the finite part of the contribution, $\lambda m_{\nu_s}/(4\pi)$ is smaller than m_ϕ. Taking $\lambda \sim 10^{-23}$, we see that this condition is readily satisfied. At

[1] To be precise, the degeneracy between mass eigenstates will be broken by $\Delta E = V_a$ where V_a is the effective potential for active neutrinos which at $T < m_e$ is composed of a contribution from symmetric neutrino background $25 G_F^2 T^5$ plus a contribution from the asymmetric electron background $\sqrt{2} G_F(n_e - n_{\bar{e}})$ [39]. However, at these temperatures both of these quantities are very small, satisfying the non-adiabaticity condition: $4\dot{\theta}_m/\Delta E|_{resonance} = \lambda\sqrt{2\rho_\phi^{int}}/(\sin\theta\cos\theta V_a) \gg 1$.

[2] However, see [42] which argues that the condition of free streaming might be relaxed.

$T \sim 10-20$ MeV when $\nu_a \to \nu_s$ may start, a "thermal" mass of $\sim \lambda n_{\nu_s}^{1/3}/\sqrt{12}$ is induced in which n_{ν_s} is the number density of the produced ν_s. Remembering that $\sqrt[3]{n_{\nu_s}} \ll T \sim 10-20$ MeV, we find the contribution is much smaller than $(\lambda/10^{-23})10^{-17}$ eV which is smaller than our benchmark value for m_ϕ. Similar consideration holds valid for $T < 0.1$ MeV where ν_s can be resonantly populated. As a result the thermal stability is guaranteed. For the larger values of λ, the stability can be jeopardized and a more careful study is required.

3. Prospects for KATRIN and PTOLEMY

The KArlsruhe TRItium Neutrino (KATRIN) experiment is designed to measure the neutrino mass by studying the endpoint of the spectrum of the emitted electron in the beta decay of Tritium. The experiment, which will soon release its first data, can be sensitive to the neutrino mass (or to be more precise to $m_{\nu_e} \equiv m_1|U_{e1}|^2 + m_2|U_{e2}|^2 + m_3|U_{e3}|^2$ [43]) down to 0.2 eV [44]. On the other hand, in the framework of the ΛCDM and the standard model of particles (including neutrino mass) combining the CMB and BAO [14] implies that neutrino mass should be smaller than this threshold and KATRIN cannot therefore discern the shift of the endpoint of the spectrum. However as shown in [45,46], there are ways to relax the bounds from cosmology on the sum of the neutrino masses opening up the hope for KATRIN to resolve a sizable shift of the endpoint and to measure m_{ν_e}.

If ν_e has a ν_4 component with a mass of ~ 1 eV, it will show up as a kink [11,12,47,48] in the spectrum of the emitted electron at $E_e = Q - m_{\nu_4}$ where Q is the mass difference between the mother and daughter nuclei. The height of the kink will be characterized by $|U_{e4}|^2$. Within the $3+1$ solution to the LSND and MiniBooNE anomalies or the $3+1$ solution to the reactor and Gallium anomalies, the size of the kink can be large enough to be resolved [12]. Let us discuss the implications of KATRIN observations combined with other observations within our scenario.

If future studies establish a deficit in the reactor $\bar{\nu}_e$ flux compatible with the $3+1$ scheme with $\Delta m^2 \sim 1$ eV^2 and $|U_{e4}|^2 \sim 0.01$ and on the other hand if KATRIN observes a kink with the corresponding position and amplitude, this will be a strong hint in favor of the $3+1$ scheme. There is a similar concept to detect relic neutrinos by the ν_e capture on Tritium. The PTOLEMY experiment is proposed to search for relic neutrinos invoking this concept [49]. Similarly to KATRIN, it can also study the beta decay spectrum. Within the 3ν mass scheme, we expect a peak further away from the endpoint at $E_e = Q + m_{\nu_e}$ due to the ν_e capture on Tritium. (To be more precise, we expect three peaks at $Q + m_{\nu_1}$, $Q + m_{\nu_2}$ and $Q + m_{\nu_3}$ which overlap with each other, looking like a single peak. Since within the 3ν scheme, we expect $F_{\nu_1} : F_{\nu_2} : F_{\nu_3} = 1:1:1$, the heights of these three overlapping peaks are given by $|U_{e1}|^2$, $|U_{e2}|^2$ and $|U_{e3}|^2$.) Within the $3+1$ scheme in addition to this peak, there will be another peak at $E_e = Q + m_{\nu_4}$ but with a height suppressed by $|U_{e4}|^2$. As we discussed in the previous section, within our scenario neutrinos will eventually decay into the lightest mass eigenstate; *i.e.*, ν_1 for the normal mass ordering and ν_3 for the inverted mass ordering. Thus, we expect a single peak at PTOLEMY-like setups at $Q + m_{\nu_1}$ ($Q + m_{\nu_3}$) with a height enhanced (suppressed) by a factor of $3|U_{e1}|^2$ (a factor of $3|U_{e3}|^2$) relative to the peak for the 3ν scheme for normal (inverted) mass ordering scheme. As a result, the presence of the kink at KATRIN results (or at PTOLEMY itself) but the absence of a second peak in the ν_e capture experiments might be taken as an indication for the ν_4 decay with a lifetime shorter than the age of the universe. We should however notice that in a scenario with a non-renormalizable coupling between ν and ϕ, as discussed in the previous section, the contribution of ν_4 to the background will be also negligible so similarly to our scenario, there will be no second peak. In this case however the height of the peak will be the same as that in the standard 3ν scheme rather than being enhanced (suppressed) by $3|U_{e1}|^2$ (or by $3|U_{e3}|^2$).

4. Summary

We have proposed a scenario to make the $3+1$ solution to the short baseline neutrino anomalies compatible with the cosmological observations. The scenario is based on a small Yukawa coupling between the sterile neutrino and ultra-light background scalar with a mass of $m_\phi < 5 \times 10^{-17}$ eV. This coupling will induce an effective mass for ν_s in the early universe when active neutrinos are still in thermal equilibrium with plasma, suppressing the effective active sterile mixing and therefore the ν_s production. This way the bound on N_{eff} is satisfied. Below $T \sim 0.01$ MeV (and for sure at present) the effective mass induced by the coupling to the dark matter is negligible.

After the neutrino decoupling when the ϕ field starts oscillating, the effective mass induced for ν_s (m_{eff}) can become negative, canceling the vacuum mass, m_{ν_s}. During the instants of total cancellation, active neutrinos can be resonantly converted to ν_s, causing a tension with the total neutrino mass bounds from BAO and CMB. A remedy is to open up the possibility of ν_4 decay before matter radiation equality. This rather fast decay can be achieved by coupling ν_s to another singlet scalar which is lighter than ν_4. We have discussed the interpretation of possible results from future observations of KATRIN and PTOLEMY within the framework of the present scenario and compare it with the predictions of certain alternative frameworks.

Throughout this letter, our main focus was on the $3+1$ solution to the short baseline anomalies but these results can be applied to even the $3+1$ solution to the ANITA events [50] which relies on nonzero $|U_{\tau 4}|$ instead of nonzero $|U_{e4}|$ or $|U_{\mu 4}|$. After this work was submitted to the archive, Ref. [51] appeared which confirms the suppression of δN_{eff} by the present scenario.

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