Constraining invisible neutrino decays with the cosmic microwave background

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Precision measurements of the acoustic peaks of the cosmic microwave background indicate that neutrinos must be freely streaming at the photon decoupling epoch when $T \approx 0.3$ eV. This requirement implies restrictive limits on "secret neutrino interactions," notably on neutrino Yukawa couplings with hypothetical low-mass (pseudo)scalars ϕ. For diagonal couplings in the neutrino mass basis we find $g \lesssim 1 \times 10^{-7}$, comparable to limits from supernova 1987A. For the off-diagonal couplings and assuming hierarchical neutrino masses we find $g \lesssim 1 \times 10^{-11}(0.05\ \mathrm{eV}/m)^2$ where m is the heavier mass of a given neutrino pair connected by g. This stringent limit excludes that the flavor content of high-energy neutrinos from cosmic-ray sources is modified by $\nu \to \nu' + \phi$ decays on their way to Earth.

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I. INTRODUCTION

The observed acoustic peaks in the temperature distribution of the cosmic microwave background (CMB) by WMAP [1,2] and other experiments such as CBI [3], DASI [4], ACBAR [5] and BOOMERANG [6–8] have provided a plethora of detailed information about our universe. In particular, several authors have independently realized that neutrino free-streaming affects the CMB acoustic peaks in a very characteristic way [9–11] and that the CMB observations therefore imply that neutrinos must be freely streaming around the photon decoupling epoch at $T \approx 0.3$ eV [12,13]. While ordinary weak interactions freeze out at $T \approx 1$ MeV, neutrinos could have "secret interactions" [14] that are still in equilibrium at late times or that actually recouple at late times. In particular, this applies to neutrino interactions with new massless or low-mass scalars or pseudoscalars. Typically, these particles would be the Nambu-Goldstone bosons of a new symmetry that is broken at some low energy scale. The Majoron model and its variants is one possible realization of this idea [15–18]. Even though the triplet Majoron model [16] is experimentally excluded, other Majoron models are still viable.

The late free-streaming requirement allows one to test neutrino interactions at eV energies, far below what is possible in the laboratory. It was previously recognized that signatures for neutrino mass generation in Majoron type models may show up in future CMB observations [10] and conversely, that existing observations exclude the "neutrinoless universe" [12] that had been invoked to escape the cosmological neutrino mass limit [19].

The purpose of our paper is to show that the free-streaming requirement translates into very stringent limits on the neutrino-(pseudo)scalar Yukawa couplings. Our limits suggest that interactions of this sort play no significant role for supernova physics [20–22]. Perhaps more interestingly, our limits exclude scenarios where the flavor content of high-energy cosmic-ray neutrinos is modified by decays $\nu \to \nu' + \phi$ in addition to the standard modification caused by flavor oscillations [23,24].

In Sec. II we first consider binary interactions among neutrinos and bosons that lead to new limits on the diagonal and off-diagonal Yukawa couplings. In Sec. III we consider decay and coalescence processes of the type $1 \leftrightarrow 2 + 3$ that provide new and very restrictive limits on the off-diagonal interactions. We summarize our findings in Sec. IV.

II. BINARY INTERACTIONS

We study ordinary neutrinos interacting with low-mass pseudoscalars ϕ via the coupling

$$L = -i\phi \sum_{jk} g_{jk} \bar{\nu}_j \gamma_5 \nu_k. \tag{1}$$

If the pseudoscalars are Nambu-Goldstone bosons of a new symmetry, as one should expect, then a derivative coupling is more appropriate, but for our most interesting process, neutrino decay $\nu \to \nu' + \phi$, the pseudoscalar and derivative couplings are equivalent. For binary processes it turns out to be conservative to use the pseudoscalar coupling, a point discussed in more detail at the end of this section. We do not explicitly study scalar interactions, but the results would be quantitatively similar.

The interaction Eq. (1) allows for binary processes of the type $\nu + \bar{\nu} \leftrightarrow \phi + \phi$, $\nu + \phi \leftrightarrow \nu + \phi$, and finally $\nu + \nu \to \nu + \nu$ with a ϕ-exchange. For this latter process the pseudoscalar and derivative couplings are equivalent because each fermion line has only one Nambu-Goldstone boson attached to it. Apart from numerical factors the scattering rate in a thermal environment of relativistic

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neutrinos is

$$\Gamma_{1+2\leftrightarrow 3+4} \approx g^4 T, \tag{2}$$

where g is the largest entry of the Yukawa coupling matrix. To avoid acoustic oscillations of the neutrino-Majoron fluid we require that the cosmic expansion rate H exceeds $\Gamma_{1+2\leftrightarrow 3+4}$ at the time of photon decoupling. Of course, this criterion is somewhat schematic, but to be numerically specific we will use the following conditions. The photon temperature at decoupling is $T_{\gamma,\mathrm{dec}} = 0.256$ eV. The corresponding neutrino temperature relevant for our estimate is $T_{\nu,\mathrm{dec}} = (4/11)^{1/3} T_{\gamma,\mathrm{dec}} = 0.18$ eV. At this epoch the universe is matter dominated so that the expansion rate is $H_{\mathrm{dec}} = 100\ \mathrm{kms}^{-1}\ \mathrm{Mpc}^{-1} (\Omega_{\mathrm{M}} h^2)^{1/2} (z_{\mathrm{dec}} + 1)^{3/2}$. With $\Omega_{\mathrm{M}} h^2 = 0.134$ and $z_{\mathrm{dec}} = 1088$ for the cosmic matter density and the redshift at decoupling, respectively [2], we find

$$H_{\mathrm{dec}} = 4.27 \times 10^{-14} \mathrm{s}^{-1} = 2.81 \times 10^{-29}\ \mathrm{eV}. \tag{3}$$

The free-streaming requirement $\Gamma_{1+2\leftrightarrow 3+4} < H_{\mathrm{dec}}$ thus translates into

$$g \lesssim 1.1 \times 10^{-7}. \tag{4}$$

Because of the g^4 dependence of the interaction rate, corrections to this limit from exact numerical factors are minor, and the limit is correspondingly robust. For example, the two independent studies [12,13] which find that neutrinos must free-stream around recombination both use approximations to solve the exact Boltzmann equation for neutrino perturbations. However, even though this does introduce a factor of order unity uncertainty in the $\Gamma_{1+2\leftrightarrow 3+4} < H_{\mathrm{dec}}$ requirement, the uncertainty in the bound on g is much smaller.

Our limit Eq. (4) on the largest of the Yukawa couplings is nominally more restrictive than the limit obtained from the energy-loss argument of supernova (SN) 1987A [20–22]. However, since our limit is based on a dimensional analysis with uncertain numerical factors, one can only claim that the two limits are comparable.

If the bosons are Nambu-Goldstone bosons of a new broken symmetry, there is a relation $g = m/f$ between the diagonal Yukawa couplings, the neutrino masses, and the symmetry breaking scale f. For processes with two Nambu-Goldstone boson lines attached to one fermion line, the coupling is derivative rather than pseudoscalar. The relevant cross section will be proportional to $m^2 E^2 / f^4$ with E a typical energy of the process rather than proportional to $g^4 = m^4/f^4$ that we used in our estimate. If neutrinos are relativistic at $T_{\nu,\mathrm{dec}}$, evidently the correct derivative structure for the interaction would lead to more restrictive limits. We note that current cosmological limits on the neutrino masses are $\sum m \lesssim 1$–1.5 eV, and even $\sum m \lesssim 0.4$ eV has been claimed [25]. Typical neutrino energies at photon decoupling are $E \approx 3T_{\nu,\mathrm{rec}} \approx 1$ eV so that even in the case of degenerate neutrino masses it is justified to treat neutrinos as relativistic. Note that the mass bounds have been obtained assuming noninteracting neutrinos, but since strongly interacting neutrinos are excluded by CMB the bound is self-consistent.

Of course, if the boson masses were much larger than the eV-scale, our considerations would change and these restrictive limits could be avoided.

III. DECAY AND COALESCENCE

Of greater interest is the decay process $\nu \rightarrow \nu' + \phi$ and the coalescence process $\nu' + \phi \rightarrow \nu$ as these are only of first order and thus the rate is proportional to g^2 rather than g^4. Here it is exact to use the nonderivative form of the interaction. The sum of the decay rates for $\nu \rightarrow \nu' + \phi$ and $\nu \rightarrow \bar{\nu}' + \phi$ in the rest frame of the parent neutrino with mass $m \gg m'$ is [26,27]

$$\Gamma_{\mathrm{decay}} = \frac{g^2}{16\pi} m. \tag{5}$$

In the frame of the thermal medium, a typical neutrino energy is $E = 3T$ so that the rate is reduced by a typical Lorentz factor $m/3T$.

The decay and coalescence processes are kinematically constrained, for relativistic particles, to couple nearly collinear modes of the interacting particles. Therefore, even if the decay is isotropic in the rest frame of the parent particle, the decay products will have directions within an approximate angle corresponding to the parent's Lorentz factor m/E. Therefore, to randomize the direction of the original neutrino requires a random walk of small angular steps. Therefore, we must include a factor $(m/E)^2$ in the medium-frame interaction rate to obtain the relevant "transport rate" rather than a naive interaction rate. The same argument was made in Ref. [10]. We also note that the usual transport cross section is the ordinary cross section times $(1 - \cos\theta)$ with θ the scattering angle. Therefore, for small angles the transport cross section is the ordinary cross section times $\frac{1}{2}\theta^2$, resulting in the same approximate $(m/E)^2$ correction factor.

We conclude that we should compare the "transport rate"

$$\Gamma_{\mathrm{T}} \approx \frac{g^2}{16\pi} m \left(\frac{m}{E}\right)^3 \tag{6}$$

with the expansion rate at photon decoupling Eq. (3). Using $E = 3T_{\nu,\mathrm{dec}}$ we thus find

$$\begin{aligned} g &\lesssim 12(3\pi)^{1/2} (H_{\mathrm{dec}} T_{\nu,\mathrm{dec}}^3)^{1/2} m^{-2} \\ &= 0.61 \times 10^{-11} \left(\frac{50\ \mathrm{meV}}{m}\right)^2, \end{aligned} \tag{7}$$

where as a mass scale we have used the largest neutrino mass of about 50 meV implied by oscillation data in a hierarchical mass scenario. This is by far the most restrictive limit on the off-diagonal neutrino-Majoron couplings.

Translating this constraint into a limit on the rest-frame lifetime yields

$$\tau \gtrsim 2 \times 10^{10}\mathrm{s}\left(\frac{m}{50\ \mathrm{MeV}}\right)^3. \tag{8}$$

Therefore, neutrinos could still be rather short-lived on cosmological time scales.

If the coupling g is between two nearly degenerate mass eigenstates, the decay rate acquires an additional factor of approximately $(\delta m^2)^3/m^6$ [27] so that the limit on g is relaxed by a factor $m^3/(\delta m^2)^{3/2}$.

IV. CONCLUSIONS

We have shown that the neutrino free-streaming requirement at the photon decoupling epoch implies new limits on the neutrino Yukawa couplings with hypothetical low-mass bosons. For the diagonal couplings we find $g \lesssim 10^{-7}$, comparable to the energy-loss argument of SN 1987A [20–22]. The SN limits suffer from the complication that one needs to distinguish between the free-streaming regime where Majoron emission simply provides a new channel of energy loss and the trapping regime where neutrino-Majoron interactions are so strong that Majorons are trapped. Unless the Majorons or other new bosons have masses which are large compared to the eV scale, our limits imply that Majorons and similar particles interact too weakly to be trapped in a SN core. Given the uncertainty of our limit, Majoron emission could perhaps still provide a non-negligible channel of energy loss, but a dominant dynamical role appears to be excluded.

We obtain a much more restrictive limit on the off-diagonal couplings of $g \lesssim 10^{-11}(50\ \mathrm{MeV}/m)^2$ with m the heavier mass of a given pair of neutrinos with nondegenerate masses. This is by far the most restrictive limit on such interactions.

It is of interest in the context of scenarios where high-energy neutrinos from cosmic-ray sources may decay on their way to Earth. These neutrinos are produced by the decay of charged pions that in turn are produced as secondary products by the high-energy cosmic-ray protons that must be produced somewhere in the universe. Standard flavor oscillations imply an expected flavor ratio at Earth of $\nu_e{:}\nu_\mu{:}\nu_\tau \approx 1{:}1{:}1$. Therefore, if future large-scale neutrino telescopes were to observe significant deviations from this expected flavor composition one would be tempted to conclude that these modifications are caused by $\nu \to \nu' + \phi$ decays [23,24]. Assuming a source distance of $D = 100$ Mpc and a neutrino energy $E = 10$ TeV, a strong decay effect obtains if $\Gamma_{\mathrm{decay}}(m/E) \gtrsim D^{-1}$. With Eq. (5) this implies the requirement

$$g \gtrsim 1.1 \times 10^{-7}\left(\frac{50\ \mathrm{meV}}{m}\right)\left(\frac{E}{10\ \mathrm{TeV}}\right)^{1/2}\left(\frac{100\ \mathrm{Mpc}}{D}\right)^{1/2}. \tag{9}$$

Therefore, g would need to be about 4 orders of magnitude larger than our new limit. Even if we consider decays between the second-lightest and lightest neutrinos with $m \approx 10$ meV and/or larger distances or smaller energies, this conclusion is not changed. For degenerate neutrino masses, the decay rate of cosmic-ray neutrinos and the early-universe reaction rate get penalized with the same factor so that, again, our conclusion remains unchanged.

Therefore, it appears that possible future deviations from the expected 1:1:1 flavor content of high-energy cosmic-ray neutrinos would have to be ascribed to effects other than Majoron decays. By the same token, invisible neutrino decays of the Majoron type can not affect solar or atmospheric neutrino observations.

In essence the reason why our bound is so much stronger is that neutrinos are almost nonrelativistic at decoupling, having energies in the sub-eV range. The Lorentz factor is therefore enormously smaller than the one considered for TeV neutrinos. On the other hand the effective ''baseline'' for decays is roughly $H_{\mathrm{dec}}^{-1} \sim 0.2$ Mpc, a number which is only 3 orders of magnitude smaller than the 100 Mpc considered typical for high-energy neutrinos observable in neutrino telescopes.

The diffuse cosmic neutrino background from all core-collapse supernovae in the universe is another potentially detectable neutrino flux from a cosmological distance. This flux can also be affected by invisible neutrino decays [28,29]. Typical energies are in the range of tens of MeV, i.e. 6 orders of magnitude smaller relative to our discussion of high-energy neutrinos. The typical source distance is at least a factor of 10 larger. Overall the required coupling strength for significant decay effects to occur is 3–4 orders of magnitude smaller than Eq. (9). Therefore, given the crude nature of our limit we can not claim with complete confidence that the diffuse supernova neutrinos are not affected by Majoron-type decays.

Either way, the universe as a neutrino laboratory once again provides information on these no-longer-so-elusive particles which is of direct relevance to other experimental directions in neutrino research.

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