Detecting a Secondary Cosmic Neutrino Background from Majoron Decays in Neutrino Capture Experiments

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We consider theories in which the generation of neutrino masses is associated with the breaking of an approximate global lepton number symmetry. In such a scenario the spectrum of light states includes the Majoron, the pseudo-Nambu Goldstone boson associated with the breaking of the global symmetry. For a broad range of parameters, the Majoron decays to neutrinos at late times, after the cosmic neutrinos have decoupled from the thermal bath, resulting in a secondary contribution to the cosmic neutrino background. We determine the current bounds on this scenario, and explore the possibility of directly detecting this secondary cosmic neutrino background in experiments based on neutrino capture on nuclei. For Majoron masses in the eV range or below, the neutrino flux from these decays can be comparable to that from the primary cosmic neutrino background, making it a promising target for direct detection experiments. The neutrinos from Majoron decay are redshifted by the cosmic expansion, and exhibit a characteristic energy spectrum that depends on both the Majoron mass and its lifetime. For Majoron lifetimes of order the age of the universe or larger, there is also a monochromatic contribution to the neutrino flux from Majoron decays in the Milky Way that can be comparable to the diffuse extragalactic flux. We find that for Majoron masses in the eV range, direct detection experiments based on neutrino capture on tritium, such as PTOLEMY, will be sensitive to this scenario with 100 gram-years of data. In the event of a signal, the galactic and extragalactic components can be distinguished on the basis of their distinct energy distributions, and also by using directional information obtained by polarizing the target nuclei.

I. INTRODUCTION

While oscillation experiments have conclusively established that neutrinos have small but non-vanishing masses, the dynamics that generates these masses remains a mystery. In particular, it is not known whether neutrinos are Dirac or Majorana particles, whether their spectrum is hierarchical, inverse hierarchical or quasi-degenerate, or why their masses are so small. Several ideas have been put forward to explain the origin of neutrino masses. These include the well-known seesaw mechanism [1–5], and the Majoron model [6–8].

Extensions of the SM based on the Majoron framework incorporate a lepton number as a global symmetry. Neutrinos acquire masses once the global lepton number symmetry is spontaneously broken. The Majoron is the pseudo-Nambu Goldstone boson associated with the breaking of the global symmetry. The form of its couplings to neutrinos is dictated by the non-linearly realized lepton number symmetry. In the case when neutrinos are Dirac, the couplings of the Majoron field $J$ take the schematic form,

$$i \frac{m_\nu}{f} J \nu \bar{\nu} + \text{h.c.}$$

(1)

Here $m_\nu$ represents the neutrino mass, while $f$ denotes the Majoron decay constant, which corresponds to the scale at which the global lepton number symmetry is broken. In the case when the neutrino masses are Majorona, the corresponding coupling takes the form

$$i \frac{m_\nu}{f} \frac{1}{2} J \nu \bar{\nu} + \text{h.c.}$$

(2)

Eqs. (1) and (2) represent the leading order interactions of the Majoron with neutrinos in an expansion in $1/f$. These couplings are diagonal in the basis in which the neutrino masses are diagonal. Hence the Majoron couples to each neutrino mass eigenstate with a strength proportional to the corresponding neutrino mass.

In general, the global symmetry need not be exact but only approximate, in which case the Majoron will acquire a mass. It then follows from the couplings of the Majoron to neutrinos, Eqs. (1) and (2), that any population of Majorons present in the early universe will eventually decay to neutrinos, unless such decays are forbidden by phase space considerations. If the Majoron lifetime is greater than the age of the universe, Majorons could constitute some or all of the observed dark matter of the universe [9–13]. The required population of Majorons could have survived till the present day as thermal relics, or have been produced nonthermally through, for example, the misalignment
mechanism. If, however, the Majoron lifetime is less than or comparable to the age of the universe, decays of the Majoron will result in a new contribution to the population of cosmic neutrinos. Provided the Majoron decays occur at late times, after the cosmic neutrinos have decoupled from the thermal bath, and at temperatures below the mass of the Majoron, this new population of neutrinos will be thermally decoupled and distinct from the cosmic neutrino background (CνB). Several authors have explored the possibility of detecting the neutrinos from Majoron decays in neutrino detectors such as Borexino, Super-Kamiokande and IceCube \cite{1416}, as well as in dark matter detectors (see, for example, \cite{17}). In this paper we determine the current bounds on this scenario, and explore the possibility of directly detecting this secondary CνB in neutrino capture experiments.

Cosmological observations can be used to place limits on the flux of neutrinos from Majoron decay. These limits depend on the Majoron mass and lifetime. Decays that occur prior to recombination result in a contribution to the energy density in radiation that is present during the epoch of acoustic oscillations. Precision measurements of the the cosmic microwave background (CMB) place severe limits on any such contribution, with the result that the number density of neutrinos from Majoron decays is always expected to be smaller than the CνB. In the case of Majorons that decay after recombination, the limits are much weaker. Current cosmological observations allow roughly 5% of the total energy density in dark matter to have decayed to dark radiation before today \cite{1820}. Then, if neutrinos are Majorana, the number density of neutrinos and antineutrinos from Majoron decay is given by

\begin{equation}
\frac{n_{\nu}}{\rho_c} = F \Omega_{\text{dm}} \frac{m_e}{m_J} \left(1 - e^{-t_0/\tau_J}\right). \tag{3}
\end{equation}

In this expression \(\tau_J\) represents the lifetime of the Majoron, and \(m_J\) its mass. \(F\) denotes the fractional contribution of Majorons to the total energy density in dark matter at early times, while \(\Omega_{\text{dm}}\) is the fractional contribution of dark matter to the total energy density \(\rho_c\). If the lifetime is much less than \(t_0\), the age of the universe, we have

\begin{equation}
\frac{n_{\nu_0}}{\rho_c} = \frac{n_{\bar{\nu}_0}}{\rho_c} = F \Omega_{\text{dm}} \frac{m_e}{m_J} = \frac{63 \text{ eV}}{5\% \text{ cm}^3 m_J}. \tag{4}
\end{equation}

In this limit, the number density of active neutrinos and antineutrinos is half of this. For comparison, note that the total number density of the three flavors of neutrinos in the CνB is \(168/\text{cm}^3\), with an equal number of antineutrinos. We see from this that the number density of neutrinos arising from Majoron decays can be very large. In particular, for Majoron masses below an eV, the flux from their decays can be comparable to or even exceed the flux from the cosmic background neutrinos. Therefore, the neutrinos from Majoron decays constitute a promising target for the direct detection experiments designed to detect the CνB.

Since there are only two particles in the final state, the neutrinos produced in Majoron decays are monochromatic. However, as a consequence of the cosmic expansion, these particles get redshifted, resulting in a characteristic energy distribution that depends not just on the Majoron mass, but also its width. If the Majoron lifetime is of order the age of the universe, we also expect a flux of neutrinos from the decay of Majorons in the Milky Way galaxy. Since the time taken to traverse the galaxy is small compared to the age of the universe, these neutrinos are monochromatic. Provided the population of Majorons is highly nonrelativistic during the matter dominated era, the distribution of Majorons in the galaxy is expected to follow the cold dark matter profile. Then the overdensity of Majorons in the Milky Way results in a flux of galactic neutrinos that is comparable to the diffuse extragalactic flux.

Experiments to detect the CνB based on neutrino capture on tritium were first proposed by Weinberg more than 50 years ago \cite{21}. Tritium undergoes beta decay with a half-life of 12.32 years,

\begin{equation}
^3\text{H} \rightarrow ^3\text{He} + e^- + \bar{\nu}_e. \tag{5}
\end{equation}

This process releases 18.6 keV of energy, which is distributed between the final state electron and neutrino. However, this decay can be also induced by neutrino capture,

\begin{equation}
^3\text{H} + \nu_e \rightarrow ^3\text{He} + e^- \tag{6}
\end{equation}

The final state electron now carries away the energy of the incident neutrino, in addition to the energy liberated in the decay. Therefore, given sufficient energy resolution, searches for electrons with energies above the beta decay endpoint can be used for direct detection of cosmic neutrinos.

Recently, this idea has taken concrete shape in the form of the PTOLEMY experiment \cite{22,23}, which aims to discover the CνB. PTOLEMY proposes to use 100g of tritium coated on graphene. A target of this size could detect

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\footnotesize{\textsuperscript{1} The bound on the fraction of decaying dark matter varies with the lifetime. The current bound indicates \(\tau/F \gtrsim 160\text{ Gyr}\) for \(\tau \gtrsim t_0\) \cite{18}, where \(\tau\) is the lifetime of the decaying dark matter while \(F\) denotes the fractional energy density in decaying dark matter to that of all dark matter. Since we focus on the case in which the lifetime is order of the age of universe, we take the bound to be \(\sim 5\%\).}
about 4 CνB events per year if neutrinos are Dirac, rising to 8 events per year if neutrinos are Majorana \[24\]. The energy resolution \(\Delta\) is at the level of 0.15 eV \[22\], which means that the experiment is not sensitive to neutrinos with energies below this threshold. The CνB neutrinos are non-relativistic. Therefore, in order to detect them, the neutrinos must be quasi-degenerate, with masses of order 0.1 eV.

In the case of Majorons with masses below an eV, the flux of neutrinos from Majoron decay can be comparable to or larger than the flux from the CνB. We find that for Majoron masses in the eV range, direct detection experiments based on neutrino capture on tritium will be sensitive to this scenario with 100 gram-years of data even if the SM neutrinos are lighter than 0.1 eV. Apart from an effect on the total event rate, we expect a distinctive energy spectrum of events inherited from the characteristic energy distributions of the neutrinos from extragalactic and galactic decays. These signals can be distinguished from the CνB, because the CνB neutrinos are all nonrelativistic today.

Interestingly, by employing polarized tritium nuclei as the target \[28\], we can also obtain directional information about the neutrino flux. Since galactic Majorons are expected to follow the dark matter distribution in the Milky Way, the monochromatic component of the neutrino flux encodes information about the profile of dark matter in the galaxy. Therefore, by observing the modulation of the event rate with the orientation of the tritium spin, we can shed light on the galactic dark matter distribution.

This paper is organized as follows. We introduce the Majoron model in Sec. II, and discuss the current bounds on this class of theories. In Sec. III we study the neutrino flux from Majoron decay, including both extragalactic and galactic sources. In Sec. IV, we estimate the signal rates at PTOLEMY based on two benchmark points. We also study the dependence of the signal on the orientation of the spin of the tritium nuclei in the case of a polarized target. We conclude in Sec. V.

II. THE MAJORON MODEL

As the Goldstone boson associated with the breaking of lepton number, the form of the Majoron couplings in the low energy theory is dictated by the non-linearly realized symmetry. If neutrinos are Dirac, below the symmetry breaking scale \(f\), the coupling of the Majoron to leptons takes the schematic form,

\[ i\frac{\lambda}{f} H L \nu^c + \text{h.c.} \tag{7} \]

In this equation the dimensionless parameter \(\lambda\) is related to the neutrino mass as \(\lambda = \sqrt{2} m_\nu/v_{\text{EW}}\), where \(v_{\text{EW}} = 246\) GeV is the electroweak VEV. In the case when neutrinos are Majorana, we have instead

\[ i\frac{J}{f} \frac{(HL)^2}{\Lambda} + \text{h.c.} \tag{8} \]

Here \(\Lambda\) is an ultraviolet scale that is related to the neutrino mass as \(\Lambda^{-1} = m_\nu/v_{\text{EW}}^2\). Eqs. (7) and (8) are the leading order interactions of the Majoron with neutrinos in an expansion in \(1/f\). At low energies these reduce to Eq. (1) and Eq. (2) respectively.

Assuming \(m_J \gg m_\nu\), we obtain the total decay width of the Majoron to neutrinos as

\[ \Gamma^D_J = \frac{m_\nu^2 m_J}{f^2 8\pi} \quad \text{or} \quad \Gamma^M_J = \frac{m_\nu^2 m_J}{f^2 16\pi}, \tag{9} \]

where \(\Gamma^D_J\) and \(\Gamma^M_J\) correspond to the Dirac and Majorana cases respectively. We focus on Majoron masses \(m_J\) of order an eV and Majoron lifetimes \(\tau_J\) of the order of the age of the Universe. This translates to a scale \(f\) of order \(10^6\) GeV.

The Majoron model is constrained by astrophysical, cosmological and collider data. These bounds can be translated into limits on the mass of the Majoron \(m_J\) and the decay constant \(f\). For light Majorons, with masses less than an MeV, the bounds from cosmology are the most severe. We therefore begin by considering the cosmological limits. Precision measurements of the CMB place limits on the total energy density in radiation during the epoch of acoustic oscillations. Additional energy density in radiation beyond the SM prediction is generally parametrized in terms of the effective number of neutrinos, \(\Delta N_{\text{eff}}\). The CMB constraint, \(\Delta N_{\text{eff}} < 0.3\ [29]\), translates into the requirement that the Majoron not have a thermal abundance at temperatures of order an MeV, when the weak interactions decouple. For low values of \(f\), Majoron-neutrino scattering will bring the Majoron into thermal equilibrium with the neutrinos.

\[2\] For models of new physics that alter the number density of nonrelativistic cosmic neutrinos, see for example [25–27].
To satisfy the bound, the decay constant $f$ must lie above 100 MeV for $m_J \lesssim 1$ MeV in the case of Majorana neutrinos \textsuperscript{30}. The limits in the Dirac case, though somewhat model dependent, are comparable. For Majorons with masses around eV, an even more severe constraint arises from the requirement that the Majoron not be in thermal equilibrium with the neutrinos during the CMB epoch. Inverse decays into Majorons prevent the neutrinos from free streaming, impacting the heights and locations of the CMB peaks. Current bounds indicate that neutrinos must be free streaming at temperatures $T$ of order an eV \textsuperscript{31, 32}. This translates into a constraint on the Majoron decay constant, $f \gtrsim 100$ GeV \textsuperscript{30}, in the case of Majorana neutrinos. Again, the limits in the Dirac case, though somewhat model dependent, are similar.

So far we have focused on the cosmological limits arising from the dominant decay channel of Majorons, the decay to two neutrinos. Other subdominant channels, such as decays charged leptons and photons, can also be used to set limits on $f$ because the bounds on fluxes of visible particles are usually stronger. Since our focus is on Majoron masses in the eV range, the only relevant channel is $J \rightarrow \gamma \gamma$. However, for light Majorons the branching ratio to photons is extremely small, and the resulting bounds are very weak \textsuperscript{15, 33, 34}.

The strongest astrophysical bounds arise from the effects of Majorons on supernovae. Neutrinos in a supernova acquire masses from matter effects, which are largest for electron neutrinos. Then, if the Majoron is light, electron neutrinos can decay into lighter flavors of neutrinos and Majorons. This can affect a supernova in two distinct ways. Prior to the bounce, these decays can deleptonize the core, preventing the explosion from taking place. After the bounce, the Majorons produced in these decays can free stream out, leading to overly rapid energy loss. While the resulting constraints depend on the details of supernova dynamics, they are at the level of $f \gtrsim 100$ keV \textsuperscript{35–39}, and therefore weaker than the cosmological bounds. Clearly, our benchmark values of $f \sim 10^6$ GeV and $m_J \sim 1$ eV are safe from all current bounds.

### III. THE NEUTRINO FLUX FROM MAJORON DECAYS

In this section we determine the flux and energy distribution of the neutrinos arising from Majoron decays, as a function of the initial abundance of the Majoron, its mass and lifetime. In the case of Majorons with lifetimes longer than or comparable to the current age of the universe we expect two distinct contributions to the flux, one from Majoron decays in the Milky Way galaxy, and the other from extragalactic decays. The galactic flux is expected to be monochromatic, with frequency equal to exactly half the Majoron mass. The diffuse extragalactic component of the flux, however, undergoes redshift as the universe expands, resulting in an energy distribution that, although not monochromatic, is highly distinctive.

#### A. Extragalactic Neutrinos

We first determine the number density of diffuse neutrinos arising from Majoron decaying outside the Milky Way. Assuming SM neutrinos are Majorana, the comoving number density of neutrinos and antineutrinos arising from Majoron decay in a time interval $dt$ is given by,

$$d(n_\nu a^3) = d(n_\bar{\nu} a^3) = -\frac{n_J a^3}{\tau_J} dt ,$$

where $a$ represents the scale factor and $n_J a^3$ the comoving number density of Majorons. Now $n_J$ evolves as

$$n_J = n_{J0} e^{-t/\tau_J} ,$$

where $n_{J0}$ is the initial number density of Majorons. If Majorons contribute a fraction $\mathcal{F}$ of the energy density in dark matter at early times, we have

$$n_{J0} a^3 = \mathcal{F} \Omega_{dm} \frac{\rho_c}{m_J}$$

Then, Eq. (10) becomes

$$d(n_\nu a^3) = -\frac{\mathcal{F} \Omega_{dm}}{\tau_J} \frac{\rho_c}{m_J} e^{-t/\tau_J} dt$$

At times when the universe dominated by dark matter and dark energy, the Hubble expansion at redshift $z$ can be approximated as

$$H(z) = H_0 \sqrt{\Omega_\Lambda + (1 + z)^3 \Omega_m} .$$
Here $\Omega_\Lambda$ and $\Omega_m$ denote the fractional contributions of dark energy and matter to the total energy density of the universe, and we are working in the limit $\Omega_\Lambda + \Omega_m = 1$. We can use Eq. (14) to obtain an expression for the age of the universe at redshift $z$,

$$t(z) = \frac{2}{3} \frac{1}{H_0 \sqrt{\Omega_\Lambda}} \ln \left( \sqrt{r(z)} + \sqrt{1 + r(z)} \right) \quad (15)$$

Here $r(z)$ represents the relative contributions of dark energy and matter to the total energy density of the universe at redshift $z$,

$$r(z) = \frac{\Omega_\Lambda}{\Omega_m} \frac{1}{(1 + z)^3}. \quad (16)$$

For neutrinos with energy $E$ today, we can infer that they were generated from Majorons that decayed at redshift $z_E$, where $z_E$ is given by

$$1 + z_E = \frac{p_{\text{max}}}{\sqrt{E^2 - m_\nu^2}} \quad \text{with} \quad p_{\text{max}} \equiv \sqrt{\frac{m_J^2}{4} - m_\nu^2}. \quad (17)$$

Using Eqs. (15) and (17) to eliminate $t$ in favor of $E$ in Eq. (13), we can obtain an expression for the energy distribution of the extragalactic neutrinos from Majoron decay (see for example [17]),

$$\frac{dn_{\nu}^{e,g}}{dE} = F_{\Omega_{\text{dm}}} \frac{\rho_c}{m_J} \frac{1}{\tau_J H_0 \sqrt{\Omega_\Lambda + \Omega_m (1 + z_E)^3}} e^{-\frac{E}{E^2 - m_\nu^2}} \frac{E}{E^2 - m_\nu^2} (E \in [m_\nu, m_J/2]). \quad (18)$$

### B. Galactic Neutrinos

If the Majoron lifetime is comparable to the age of universe, Majorons will clump inside the Milky Way galaxy. If the Majorons are sufficiently cold, their distribution in the galaxy will follow the cold dark matter profile. Since the time to traverse the galaxy is small compared to the age of the universe, any neutrinos we observe from Majoron decay in the galaxy must have arisen from decays in the recent past. The spectrum of the galactic neutrinos produced from Majoron decays is therefore monochromatic, with energy $m_J/2$. The flux of these neutrinos in the neighborhood of the earth is given by the line of sight integral,

$$n_{\nu}^{\text{gal}} = F_{r_\odot} r_\odot \frac{\rho_\odot}{m_J \tau_J} \frac{e^{-r_\odot/\tau_J}}{r_\odot \rho_\odot} \approx 2.19 \quad (20)$$

where $r_\odot = 8.33 \text{kpc}$ is the distance from the earth to the galactic center and $\rho_\odot = 0.3 \text{GeV/cm}^3$ is the dark matter density at the position of the Earth. Here the dimensionless factor $J$ corresponds to the $J$ factor for the NFW profile [40] averaged over the Milky Way galaxy,

$$J = \frac{1}{4\pi} \int_{-\pi/2}^{\pi/2} db \cos b \int_0^{2\pi} d\ell J^{\text{NFW}}(b, \ell) \approx 2.19 . \quad (21)$$

Here we have used galactic coordinates $(d, b, \ell)$ with distance from the earth $d$, galactic latitude $b$ and longitude $\ell$. In this coordinate, we can express any vector in the $(x, y, z)$ basis as:

$$x = d \cos b \cos \ell , \quad y = d \cos b \sin \ell , \quad z = d \sin b. \quad (21)$$

The earth lies at the center of this coordinate system while the galactic center is at $x = r_\odot$ and $y = z = 0$. This means the whole galactic plane corresponds to $z = 0$ plane. $J^{\text{NFW}}(b, \ell)$ is the $J$ factor for the NFW profile expressed as a function of $b$ and $\ell$ [41].

### IV. DIRECT DETECTION IN NEUTRINO CAPTURE EXPERIMENTS

The neutrinos from Majoron decay can be directly detected in experiments based on neutrino capture on nuclei. In this section we determine the size of the signal in experiments based on neutrino capture on tritium, with a focus on the PTOLEMY experiment.
The PTOLEMY experiment has been designed to detect the CνB. The target consists of 100 grams of tritium coated on a graphene substrate. Tritium undergoes beta decay to He$^3$ with a lifetime of 12.32 years, liberating 18.6 keV of energy. This decay can also be induced by incident electron neutrinos, in which case the liberated electrons will have an energy that exceeds the end point energy of the beta decay process, allowing them to be distinguished. In the case of a non-relativistic neutrino with mass $m_\nu$, the energy of the liberated electrons will exceed the end point energy by $2m_\nu$. Since PTOLEMY has an energy resolution of about 0.15 eV, and the CνB neutrinos are nonrelativistic, it is sensitive to CνB neutrinos with masses above 0.1 eV. In contrast, the neutrinos from the decays of Majorons with masses in the eV range and lifetimes of order the age of the universe are expected to be relativistic. We therefore expect that PTOLEMY will be able to detect these neutrinos, even if their masses lie well below 0.1 eV, provided their energies today are above 0.2 eV.

For low energy neutrinos the cross section $\sigma$ scales as the inverse of the relative velocity $v$, so that the product $\sigma v$ is a constant. The unpolarized capture rate on tritium for electron neutrinos is given by

$$ n_\nu \sigma v_\nu = \left[ (1 - v_\nu)n_{\nu_R} + (1 + v_\nu)n_{\nu_L} \right] \bar{\sigma}, $$

$$ \bar{\sigma} = 3.83 \times 10^{-45} \text{cm}^2, $$

(22)

where $v_\nu$, the magnitude of the neutrino velocity, is close to 1 for relativistic neutrinos. The number density of left(right)-helical neutrinos is denoted as $n_{\nu_{L(R)}}$. If neutrinos are Majorana (Dirac), the number density of neutrinos from Majoron decay is $n_{\nu_{L(R)}} = n_{\nu_{L(R)}} = n_\nu$ ($n_{\nu_{L(R)}} = \frac{1}{2}n_\nu$, $n_{\nu_{L(R)}} = \frac{1}{2}n_\nu$). Eq. (22) is valid for incident neutrinos with energies below a few keV. It is clear from this formula that, for a given Majoron mass and lifetime, the signal rate from Majoron neutrinos (denoted as $R_M$) is larger by a factor of two than the rate from Dirac neutrinos (denoted as $R_D$), $R_M = 2R_D$. Neutrinos from galactic and extragalactic Majoron decays exhibit characteristic energy distributions, and the signal event rate at PTOLEMY will maintain the relation $R_M = 2R_D$ throughout the entire spectrum.

Only neutrinos in the electron flavor can be captured on tritium. Since the neutrinos produced in Majoron decay are left-handed, the only neutrino that can be captured is the electron. The solar and atmospheric mass splittings are given by 0.0086 eV and 0.050 eV respectively [42]. Then, if the lightest neutrino is massless, the suppression factor is about 0.03 in the case of a normal hierarchy, and about 1/2 in the case of an inverted hierarchy. If neutrinos are quasi-degenerate, the suppression factor is about 1/3.

The number of extragalactic (galactic) neutrino events per year expected in PTOLEMY, in the case of Majorana neutrinos, is given by (see Fig. [1])

$$ R_{M}^{\text{gal}} = R_0 \left( 1 - e^{-t_0/\tau_J} \right), \quad R_{M}^{\text{gal}} = R_0 \frac{\epsilon_{\odot} \rho_{\odot}}{\epsilon_J} \Omega_{dm} \rho_c e^{-t_0/\tau_J}. $$

(24)

where $R_0$ is given by

$$ R_0 = N_H 2n_{\nu_0} \sum_i |U_{ei}^2| P_i \bar{\sigma} $$

$$ = \frac{3.1 \mathcal{F} \text{ eV}}{\text{year} \frac{50}{m_J} \text{cm}^2} \sum_i |U_{ei}^2| P_i \frac{M_T}{100 \text{g}} \frac{\bar{\sigma}}{100 g 3.83 \times 10^{-45} \text{cm}^2}. $$

(25)

Here $N_H$ is the number of tritium nuclei in the detector and $M_T$ is the mass of the tritium target. We see that in the case of quasi-degenerate and inverted spectra, we can expect to see as many as a few events per year, provided the neutrino energies are above the detector threshold. This falls to one event every few years in the case of a normal hierarchy. Note that this signal rate is comparable to the expected event rate for the standard CνB, which corresponds to 4(8) events at PTOLEMY if neutrinos are Dirac (Majorana) [4].

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3 We use the latest 2018 data for normal hierarchy from NuFIT 3.2 (2018), www.nu-fit.org.
4 Gravitational clustering of nonrelativistic neutrinos is expected to modify the CνB event rate, but only by order one factors [43-44].
A. Extragalactic neutrino spectrum

Since extragalactic neutrinos are approximately isotropic, we can use the unpolarized cross-section in Eq. (22). By combining Eq. (18) and Eq. (22), we obtain the event rate for Majorana neutrinos,

\[
\frac{1}{R_0} \frac{dR_{e.g.}^{e.g.}}{dE} = \frac{1}{\tau_J H_0 \sqrt{\Omega_\Lambda + \Omega_m (1 + z_E)^3}} e^{-\frac{(z_E)}{\tau_J}} \frac{E}{E^2 - m_{\nu}^2} \quad (E \in [m_{\nu}, m_J/2]),
\]

where \(z_E\) is defined in Eq. (17).

The spectrum of neutrinos from extragalactic Majoron decay is shown in Fig. 2. For a given Majoron lifetime, the spectrum has sharp edges at \(E = m_J/2\) and \(E = m_{\nu}\), with a continuous distribution connecting these two points. The peak at \(E = m_{\nu}\) corresponds to the fact that neutrinos from early decays of Majorons are all nonrelativistic with energies close to the neutrino mass.

By integrating Eq. (26), \(R_{e.g.}^{e.g.} = \int dE \frac{dR_{e.g.}^{e.g.}}{dE} = R_0 (1 - e^{-t_0/\tau_J})\), we reproduce the result for the total event rate. If the lifetime is smaller than the age of universe (\(\tau_J < t_0\)), most of the Majorons will have decayed. Then the event rate is expected to be close to \(R_0\), \(R_{e.g.}^{e.g.} \approx R_0\), provided the neutrinos retain enough energy to pass the detector threshold. If, however, the lifetime is longer than the age of the universe, the fraction of Majorons that have decayed before today is smaller. Then the ratio \(R_{e.g.}^{e.g.} / R_0\) gets smaller as \(\tau_J\) increases, as can be seen in Fig. 1.

B. Polarized tritium and the galactic neutrino spectrum

The capture cross section of neutrinos on tritium depends on the relative angle between the incoming neutrino velocity and the spin of the tritium nucleus. As discussed in [28], this means that if the tritium used in PTOLEMY is polarized, we can obtain information about the direction of the neutrino signal. This in turn can be used to obtain information about the distribution of the Majoron component of dark matter in the galaxy. Since the galactic neutrinos will have energies of order an eV, they can be treated as ultrarelativistic. Averaging over the angles of the outgoing electrons, one obtains the polarized capture rate for relativistic electron neutrinos as [28]

\[
n_{\nu} \sigma(\hat{s} \cdot \hat{v}) = 2n_{\nu L} \sigma (1 + B \hat{s} \cdot \hat{v}).
\]

Here the parameter \(B \approx 0.99\) incorporates the relevant couplings and form factors. The unit vector in the direction of the tritium spin is denoted as \(\hat{s}\), while the direction of the incoming neutrino is \(\hat{v}\). As mentioned earlier, \(n_{\nu L} = n_{\nu R}^J\) if neutrinos are Majorana, while \(n_{\nu L} = \frac{1}{2} n_{\nu R}^J\) if neutrinos are Dirac. Since the extragalactic source (discussed in section IIIA) is isotropic, the result for that case simply reduces to Eq. (22) after averaging the tritium spin in the relativistic limit. However, because the galactic distribution of Majorons follows an NFW profile, interesting directional information is encoded in term proportional to \(\hat{s} \cdot \hat{v}\) in Eq. (27).
FIG. 2: This plot shows \(dR_M/dE\) as a function of the neutrino energy \(E\) for two benchmark points, both with \(m_\nu = 0.05\,\text{eV}\) and 100 g tritium target. Solid lines represent the benchmark (\(m_J = 1\,\text{eV}, \tau_J = 10\,\text{Gyr}\) and \(F = 0.05\)) while the dashed lines show the benchmark (\(m_J = 4\,\text{eV}, \tau_J = 100\,\text{Gyr}\) and \(F = 1\)). We show the energy spectra from extragalactic neutrino decays (brown), galactic neutrino decays with tritium spin polarized towards (\(\theta = 0\), orange) and away (\(\theta = \pi\), blue) from the galactic center. Here we apply \(\sigma = 0.064\,\text{eV}\) (from the experimental resolution \(\Delta = \sqrt{8}\ln 2\pi = 0.15\,\text{eV}\)) Gaussian smearing of galactic neutrinos. These results are for Majorana neutrinos. The rates for Dirac neutrinos are half of those in the Majorana case.

We wish to determine the dependence of the signal on the orientation of the spin of the tritium nuclei relative to the galactic center. We begin by writing \(\hat{s} \cdot \hat{v}\) in terms of galactic coordinates,

\[
\hat{s} \cdot \hat{v} = \cos b_k \cos b_\ell \cos (\ell_\hat{s} - \ell_\hat{g}) + \sin b_k \sin b_\ell,
\]

where \(\hat{s}(\hat{v}) = \cos b_\xi(x) \cos \ell_\xi(x) \hat{e}_x + \cos b_\xi(x) \sin \ell_\xi(x) \hat{e}_y + \sin b_\xi(x) \hat{e}_z\) with \(\hat{e}_{x,y,z}\) being the unit vector along \(x, y, z\) direction. We also define \(\theta\) as the angle between the tritium spin and the unit vector pointing towards the galactic center \((\hat{e}_x)\),

\[
\cos \theta = \hat{s} \cdot \hat{e}_x = \cos b_k \cos \ell_\hat{g}.
\]

Since the form of the NFW distribution is invariant under rotations along the line connecting the earth and the galactic center, the final result should also be unchanged under such a rotation. This allows us to choose coordinates such that \(\ell_\hat{s} = 0\) or \(\pi\). Combining Eqs. (28) and (29), we obtain

\[
\hat{s} \cdot \hat{v} = \begin{cases} 
\cos b_k \cos b_\ell \cos \ell_\hat{s} + \sin b_k \sin b_\ell & (\ell_\hat{s} = 0) \\
- \cos b_k \cos b_\ell \cos \ell_\hat{s} + \sin b_k \sin b_\ell & (\ell_\hat{s} = \pi) 
\end{cases}
\]

\[
= \cos \theta \cos b_k \cos \ell_\hat{s} + \sin \theta \sin b_\ell.
\]

In this case the rate when the spin is oriented at an angle \(\theta\) with Majorana neutrinos is given by

\[
\frac{R_M^{\text{GM}}(\theta)}{R_0} = \frac{r_0 \rho_0}{\tau_J \Omega_{dm} \rho_c} e^{-t_0/\tau_J} \left[ J + \frac{B}{4\pi} \int_{-\pi/2}^{\pi/2} db \int_0^{2\pi} d\ell \cos b (-\cos \theta \cos b \cos \ell - \sin \theta \sin b) J_{\text{NFW}}(b, \ell) \right]
\]

\[
\approx \frac{r_0 \rho_0}{\tau_J \Omega_{dm} \rho_c} e^{-t_0/\tau_J} (2.19 - 0.60 \cos \theta).
\]

It is clear from Eq. (31) that the galactic signal varies by almost a factor of two depending on whether the spin of the polarized tritium points towards or away from the galactic center,

\[
\frac{R_M^{\text{GM}}(\theta = \pi)}{R_M^{\text{GM}}(\theta = 0)} \approx 1.8.
\]

This is a characteristic feature of the galactic neutrino signal. PTOLEMY with polarized tritium can test this prediction by measuring the modulation associated with the orientation of the spin.\footnote{Polarized tritium can also be used to measure the annual modulation of PTOLEMY signals, both in the context of the standard CoB \cite{49}, and of new physics \cite{50}.}
In principle, all galactic neutrinos have the same energy \( E = m_j/2 \). However, due to experimental resolution, the actual spectrum of galactic neutrinos will be smeared out. To obtain a realistic spectrum of galactic neutrinos, we apply a Gaussian smearing with the full width at half maximum (FWHM) \( \Delta \) defined as

\[
\Delta = \sqrt{8 \ln 2 \sigma},
\]

where \( \sigma \) is the standard deviation of the Gaussian. The expected energy resolution of PTOLEMY, \( \Delta = 0.15 \) eV, corresponds to a standard deviation \( \sigma = 0.064 \) eV. With this Gaussian smearing, the galactic neutrinos spectrum is given by

\[
\frac{1}{R_0} \frac{dR_{\text{gal}}^{\theta}}{dE} = \frac{1}{\sqrt{2\pi}\sigma} \frac{n_0^{\text{gal}}}{n_{\theta}^{\text{gal}}} (2.19 - 0.60 \cos \theta) \exp[-(E - \frac{m_j}{2})^2]/2\sigma^2],
\]

where

\[
\frac{n_0^{\text{gal}}}{n_{\theta}^{\text{gal}}} = \frac{R_{\odot} \rho_{\odot}}{\tau_j \Omega_{\text{dm}} c} e^{-t_0/\tau_j}.
\]

The galactic neutrino spectrum is shown in Fig. 2. From the above expression, we see that the total signal rate from galactic neutrinos is proportional to \( e^{-t_0/\tau_j} \). Therefore, \( R_{M, \text{gal}}^{\theta} \) is peaked at \( \tau_j = t_0 \), as can be seen in Fig. 4. If \( \tau_j \ll t_0 \) the event rate is small due to the galactic neutrino density being suppressed by \( e^{-t_0/\tau_j} \). On the other hand, if \( \tau_j \gg t_0 \), the signal is suppressed because of the small decay rate.

Comparing the signals from extragalactic and galactic Majoron decay, we see that both signals can be easily distinguished from each other, and from \( C_{\nu B} \) signals. \( C_{\nu B} \) signals accumulate close to \( E = m_e \), while the energy distribution of neutrinos from extragalactic decays exhibits a distinctive spectrum, as shown in Fig. 2. Galactic signals can be distinguished by their monochromatic nature, and, if the tritium target is polarized, by employing directional information.

The directional dependence of the signal also encodes information about the dark matter distribution in the Milky Way galaxy. Once \( R_{M, \text{gal}}^{\theta} \) is measured, the ratio of the \( \theta \) independent term to the coefficient of the \( \cos \theta \) term captures the ratio of the average of the dark matter profile to the weighted average along the line sight [see Eq. (31)]. Measuring this ratio can help discriminate between models of the dark matter profile in the Milky Way.

C. Benchmark points

For concreteness we consider two benchmark points. We first choose \( m_j = 1 \) eV and a lifetime that is lightly shorter than the age of the universe, \( \tau_j = 10 \) Gyr. In this case the current bounds from decaying dark matter indicate that Majorons can only contribute about 5% of the total dark matter abundance [18]. Therefore, the first benchmark point we choose is BP1: \( m_j = 1 \) eV, \( \tau_j = 10 \) Gyr and \( F = 0.05 \). From the analysis in [18], if the lifetime of the Majoron is sufficiently longer than the age of the universe, say \( \tau_j = 100 \) Gyr, it can constitute all of the dark matter abundance, so that \( F = 1 \). Therefore we consider another benchmark BP2: \( m_j = 4 \) eV, \( \tau_j = 100 \) Gyr and \( F = 1 \). For these two benchmark points with quasi-degenerate neutrino masses, we obtain the total signal rates from extragalactic decays [Eq. (24)] and galactic decays with a polarized tritium target [Eq. (31)] as

\[
\begin{align*}
\text{BP1}: & \quad R_{M}^{e, x} = 2.3/(100 \text{g} \cdot \text{yr}), \quad R_{M, \text{gal}}^{\theta}(\theta = \pi) = 1.4/(100 \text{g} \cdot \text{yr}) \\
\text{BP2}: & \quad R_{M}^{e, x} = 2.0/(100 \text{g} \cdot \text{yr}), \quad R_{M, \text{gal}}^{\theta}(\theta = \pi) = 2.4/(100 \text{g} \cdot \text{yr}).
\end{align*}
\]

The dominant background at PTOLEMY arises from regular beta decay events in which the electron carries away almost all the liberated energy. We expect that this background falls off very quickly for \( E > 0.2 \) eV in Fig. 2. The basis of this expectation is that PTOLEMY experiments have sensitivity to \( C_{\nu B} \) neutrinos with \( m_\nu \gg 0.1 \) eV, which means that the experiment must be able to discriminate events with electron energy 0.2 eV or more above the end point of the beta decay spectrum. Following this assumption, we obtain the sum of extragalactic and galactic event rates above \( E \geq 0.2 \) eV to be \( R_{M}^{\text{tot}} = 2.7 \) for BP1 and \( R_{M}^{\text{tot}} = 4.3 \) for BP2 with 100 gram tritium and one year exposure.

Therefore, we expect to have a few signal events with a specified energy spectrum in the small background range with just one year of data. Clearly, observing such a striking signal would provide strong evidence of Majoron decays.

V. CONCLUSION

In this paper we have analyzed the prospects for discovery of a secondary cosmic neutrino background originating from Majoron decays at future detectors based on neutrino capture on tritium, such as PTOLEMY. We have studied
the bounds on this scenario and determined the event rate and energy distribution of the signal as a function of the Majoron mass and lifetime. We find that for Majoron masses at the eV scale and lifetimes of order the age of the universe, these detectors will be sensitive to this class of models with 100 gram-years of data. We have considered two distinct sources of this signal - galactic and extragalactic, and find that each has a distinctive spectrum that can be distinguished using the data. Finally, we have explored the possibility of using polarized tritium to obtain directional information about the incoming neutrinos. This can be used to discriminate between neutrinos with galactic and extragalactic origins, and thereby shed light on the distribution of dark matter in our galaxy.

**Note added:** While this paper was being completed, we received Ref. [47], which discusses ideas similar to those considered here.

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