

Boomerang mechanism explaining the excess radio background

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We propose a *boomerang mechanism* for the explanation of the excess radio background detected by ARCADE 2. In an early stage of the Universe, at a temperature T in the range ~ 0.1 keV–1 MeV, a fraction of relic neutrinos is resonantly converted into dark neutrinos by mixing induced by a preexisting lepton asymmetry. Dark neutrinos decay much later into a dark-standard photon state and a dark fermion, with a lifetime longer than the age of the Universe, as required by a solution to the excess radio background. This scenario circumvents the upper bound on the neutrino magnetic moment but still implies a testable lower bound.

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I. INTRODUCTION

The Absolute Radiometer for Cosmology, Astrophysics and Diffuse Emission (ARCADE 2) [1] has detected an excess radio background (ERB) in the 3–10 GHz frequency range with respect to the cosmic microwave background (CMB) thermal spectrum. The excess is statistically significant (more than 5σ) and cannot be explained by known population of sources since they give a contribution to the effective temperature that is 3–10 times smaller than the measured one [2]. Moreover, different observations place a strong upper limit on the anisotropy of the ERB that is, therefore, extremely smooth [3]. This represents a strong constraint for an astrophysical origin and, therefore, the ERB is currently regarded as a mystery [2]. The Tenerife Microwave Spectrometer (TMS) will soon take data in the 10–20 GHz frequency range [4,5] and might, therefore, help to shed light on this mystery. It was noticed that radiative relic neutrino decay can potentially explain the ARCADE 2 excess [6]. Recently, we have shown that such a solution indeed fits very well the six ARCADE 2 data points between 3 and 10 GHz giving rise to an excess [7].

It predicts an effective (radiometric) temperature for the nonthermal photons produced by the decays of one of the relic neutrino species, the lightest for definiteness, given by

$$T_{\gamma_{\text{nth}}}(E, 0) \simeq \frac{6\zeta(3)}{11\sqrt{\Omega_{\text{M}0}}} \frac{T_0^3}{E^{1/2} \Delta m_1^{3/2}} \frac{t_0}{\tau_1} \left(1 + \frac{a_{\text{D}}^3}{a_{\text{eq}}^3}\right)^{-\frac{1}{2}}, \quad (1)$$

where $t_0 = 4.35 \times 10^{17}$ s is the age of the Universe, $T_0 = (2.725 \pm 0.001)$ K is the photon temperature at the present time measured by the Far Infrared Absolute Spectrophotometer (FIRAS) instrument [8], $\Delta m_1 \equiv m_1 - m_0 \ll m_1$ is the mass difference between the lightest active neutrino and the new sterile state, τ_1 is the neutrino lifetime, and $E \leq \Delta m_1$ is the energy of the photon at the present time.¹ For the definition and values of the other cosmological parameters in Eq. (1), see Ref. [7]. This expression can be used to fit the six data points measured by ARCADE 2 in the 3–10 GHz frequency range. The best fit is obtained for $\Delta m_1 \simeq 4.0 \times 10^{-5}$ eV and $\tau_1 \simeq 1.46 \times 10^{21}$ s. As shown in Fig. 1, it provides an excellent fit to the ERB temperature spectrum, improving a simple power-law fit, with $\chi^2/\text{d.o.f.} \simeq 1$ with 4 degrees of freedom (d.o.f.) [7]. The result of the fit, expressed in terms of the quantity $\Delta m_1^{3/2} \tau_1$, gives at 99% confidence level (CL):

$$(\Delta m_1^{3/2} \tau_1)^{\text{ARCADE}} = 3.8_{-1.5}^{+7.2} \times 10^{14} \text{ eV}^{3/2} \text{ s}. \quad (2)$$

¹A discussion on the role of the absolute neutrino mass scale can be found in [7].

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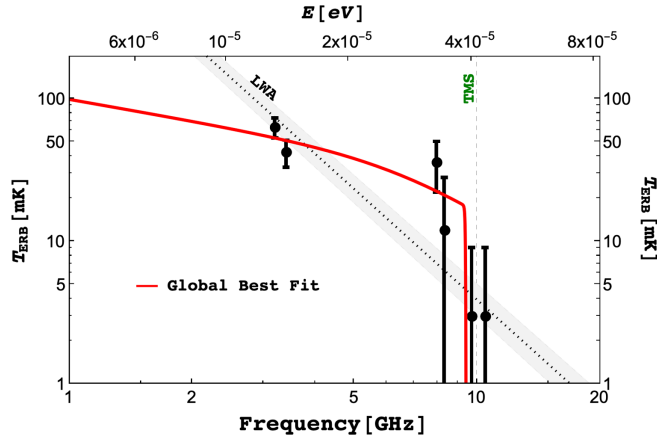


FIG. 1. Best-fit curve for T_{ERB} obtained with Eq. (1). The thick red curve corresponds to the best global fit obtained for $\Delta m_1 = 4.0 \times 10^{-5}$ eV and $\tau_1 = 1.46 \times 10^{21}$ s. The ARCADE 2 data points are taken from Ref. [1]. We also show the power-law fit $\beta = -2.58 \pm 0.05$ (dotted line with gray shade), obtained using the Long Wavelength Array (LWA) data at lower frequencies [9]. The vertical dashed line shows the TMS low-frequency threshold.

This region is shown in Fig. 2 (shaded regions with different shades for different frequency bands) in the plane Δm_1 versus τ_1 .

In a general way, we can write the radiative decay rate of the neutrino mass eigenstate ν_1 with mass m_1 into another neutrino mass eigenstate ν_0 with mass m_0 in terms of the *effective magnetic transition dipole moment* μ_{eff} , as [10]

$$\begin{aligned} \Gamma_{\nu_1 \rightarrow \nu_0 + \gamma} &= \frac{(m_1^2 - m_0^2)^3}{8\pi m_1^3} \mu_{\text{eff}}^2 \simeq \frac{\Delta m_1^3}{\pi} \mu_{\text{eff}}^2 \\ &\simeq 42.5 \text{ s}^{-1} \left(\frac{m_1 - m_0}{\text{eV}} \right)^3 \left(\frac{\mu_{\text{eff}}}{\mu_B} \right)^2. \end{aligned} \quad (3)$$

In the second expression we used $m_1 + m_0 \simeq 2m_1$, since we are interested in the case of quasidegenerate neutrinos. In the numerical expression we normalized μ_{eff} to the Bohr magneton $\mu_B = e\hbar/(2m_e) \simeq 296.3 \text{ GeV}^{-1}$. The effective magnetic transition dipole moment is defined as

$$\mu_{\text{eff}} \equiv \sqrt{|\mu_{1 \rightarrow 0}|^2 + |\epsilon_{1 \rightarrow 0}|^2}, \quad (4)$$

where $\mu_{1 \rightarrow 0}$ is the transition magnetic dipole and $\epsilon_{1 \rightarrow 0}$ is the transition electric dipole moment [11]. The most stringent upper bound on μ_{eff} comes from plasmon decays in globular cluster stars [12,13]:

$$\mu_{\text{eff}} \lesssim 1.2 \times 10^{-12} \mu_B. \quad (5)$$

In this case, from Eq. (3), one easily finds the following lower bound on the lifetime of ν_1 :

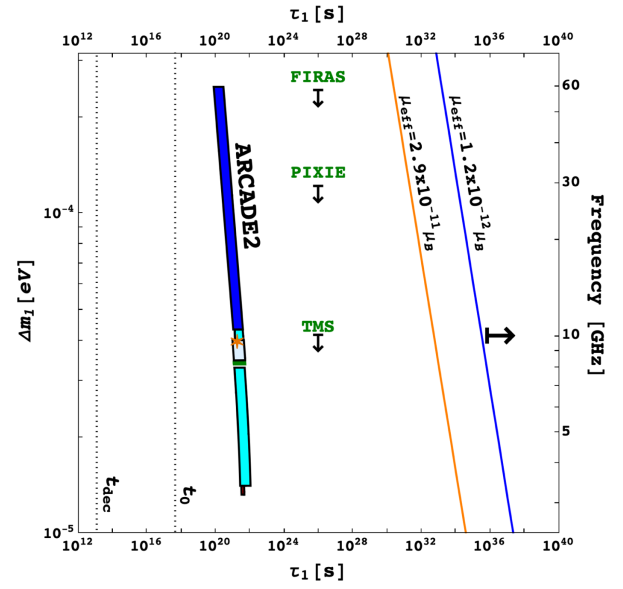


FIG. 2. Allowed region, shaded with different shades corresponding to different photon energy bands and with the best-fit point denoted by $\star$, explaining the ARCADE 2 excess radio background in the plane of Δm_1 versus τ_1 [7]. Lower bounds on the lifetime derived from the upper bound on the effective magnetic transition dipole moment [cf. Eq. (3)] are shown by the blue and orange lines. We also indicate the matter-radiation decoupling time t_{dec} and the current age of the Universe t_0 . The lowest frequency thresholds for FIRAS, PIXIE and TMS are also indicated.

$$\tau_{\nu_1 \rightarrow \nu_0 + \gamma} \gtrsim 2.5 \times 10^{21} \text{ s} \left(\frac{\text{eV}}{\Delta m_1} \right)^3. \quad (6)$$

One can see that for $\Delta m_1 \lesssim 10^{-4}$ eV, a necessary condition to address the ERB, one obtains $\tau_{\nu_1 \rightarrow \nu_0 + \gamma} \gtrsim 10^{33}$ s, an incredibly long lifetime yielding completely negligible contribution to the ERB. This conclusion remains valid even considering the less stringent laboratory upper bound from the GEMMA experiment measuring electron recoils induced by neutrino-electron elastic scattering [14,15],

$$\mu_{\text{eff}} \lesssim 2.9 \times 10^{-11} \mu_B. \quad (7)$$

The situation is depicted in Fig. 2. Therefore, irrespectively of whether the final neutrino is an active or a sterile neutrino species, neutrino radiative decay by itself cannot give any sizable contribution to the ERB.

The *boomerang mechanism* that we propose here provides a way to circumvent this bound, preserving at the same time the success of Eq. (1) in reproducing the ARCADE 2 data. While in radiative neutrino decays ordinary neutrinos directly decay into photons, in the boomerang mechanism neutrinos are first converted into dark (sterile) neutrinos and then these decay into dark photons and standard photons, with the latter constituting

the ERB.² We first discuss in Sec. II how active-dark neutrino mixing, in the presence of a sufficiently large lepton asymmetry, can convert a fraction of relic neutrinos into (quasidegenerate) dark neutrinos. In Sec. III we discuss how the dark neutrinos decay into a dark fermion species and into dark and standard photon state. In Sec. IV we combine together all constraints and determine an allowed region in the plane of active-dark neutrino mixing angle versus mass squared difference that maps into a corresponding region in the plane of the neutrino effective magnetic moment and resonance temperature. In Sec. V we compare our allowed region with current and future neutrino oscillation constraints. Finally, we draw some final remarks and conclude in Sec. VI.

II. ACTIVE-TO-DARK NEUTRINO CONVERSIONS IN THE EARLY UNIVERSE

We assume the existence of a light sterile dark neutrino field almost coinciding with mass eigenstate ν_0 with mass m_0 and quasidegenerate with the lightest neutrino with positive $\Delta m^2 \equiv m_1^2 - m_0^2$.³ We also assume, for definiteness, that it just mixes with the muon neutrino but all results are valid in general. The much higher values of $m_{2,3}^2 - m_0^2$ make in a way that the relevant mixing parameters are just Δm^2 and a very small muon–dark neutrino mixing angle θ_0 . With these assumptions and definitions the mixing in the early Universe is described by the effective Hamiltonian ΔH that in the interaction basis can be written as

$$\Delta H_{\mu\text{--dark}} = \frac{\Delta m^2}{4p} \begin{pmatrix} \cos 2\theta_0 - v(y, T_\nu, L) & -\sin 2\theta_0 \\ -\sin 2\theta_0 & -(\cos 2\theta_0 - v(y, T_\nu, L)) \end{pmatrix}, \quad (8)$$

where we introduced the dimensionless effective potential

$$v(y, T, L) = v_1(y, T_\nu) + v_2(y, T_\nu, L). \quad (9)$$

In this expression we denoted by T_ν the neutrino temperature, $y \equiv p/T_\nu$ and by L the effective (muonic) total asymmetry defined as

$$L \equiv 2L_{\nu_\mu} + L_{\nu_e} + L_{\nu_\tau} - \frac{1}{2}B_n, \quad (10)$$

²The name “boomerang” mechanism is justified by the fact that the visible sector first throws sterile neutrinos into the dark sector via mixing, and then these, decaying, throw back photons into the visible sector. The net result is that neutrinos are converted into photons. Trading off the effective magnetic moment of ordinary neutrinos with that of sterile neutrino to circumvent (5) is an ingredient, the only one, in common with the model proposed in Ref. [16] to explain the Experiment to Detect the Global EoR Signature (EDGES) anomaly.

³We adopt a different sign convention for Δm^2 than in [17,18].

where the neutrino asymmetries $L_{\nu_\alpha} \equiv (N_{\nu_\alpha} - N_{\bar{\nu}_\alpha})/N_\gamma^i$ and similarly for the neutron asymmetry B_n , with N_γ^i the photon abundance at some initial time t_i such that $m_\mu \gg T_i \gg m_e$, with $T_i \equiv T(t_i)$. The term $v_1(y, T_\nu)$ is the finite temperature contribution [19] and is given by

$$v_1(y, T_\nu) = \frac{eV^2}{\Delta m^2} \left(\frac{T_\nu}{T_\mu} \right)^6 y^2, \quad (11)$$

where $T_\mu \simeq 23.4$ MeV. For $L = 0$ and $\Delta m^2 > 0$, there would be a resonance, both for neutrinos and antineutrinos, at $T_\nu^{\text{res}}(y, 0) \simeq T_\mu (y^2 \Delta m^2 / eV^2)^{1/6}$. At this resonance relic neutrinos are not efficiently converted into dark neutrinos, though this could be used to trigger the generation of a large lepton asymmetry [18,20], as we discuss in Sec. II B.

A. Initial preexisting effective asymmetry

If we assume that there already exists an initial preexisting effective muon asymmetry L_i , for example generated by the decays of weakly coupled seesaw neutrinos as in leptogenesis [21], one has also to consider the term $v_2(y, T_\nu, L)$ given by [17]

$$v_2(y, T_\nu, L) \simeq \mp v_0 \frac{eV^2}{\Delta m^2} L \left(\frac{T_\nu}{\text{MeV}} \right)^4 y, \quad (12)$$

where the $- (+)$ sign holds for neutrinos (antineutrinos) and $v_0 = (4\sqrt{2}\zeta(3)/\pi^2)10^{12}G_F \text{ MeV}^2 \simeq 8$. If we assume that $|L_i| \gg L_\star(y) \simeq 0.4 \times 10^{-6} (y \Delta m^2 / eV^2)^{1/3}$, then the resonance condition is satisfied for $v_2(y, T_\nu, L_i) = 1$. This implies that for positive L_i and positive Δm^2 there is a resonance only for antineutrinos at a resonant temperature

$$T_\nu^{\text{res}}(y, L_i \gg L_\star) = \left(\frac{1}{v_0 L_i y} \frac{\Delta m^2}{eV^2} \right)^{\frac{1}{4}} \text{ MeV} < T_\nu^{\text{res}}(y, 0). \quad (13)$$

Notice that this resonance occurs at different times for different values of $y = y_{\text{res}}$. The asymmetry grows with time from the initial value and y_{res} spans all neutrino distribution starting from a small value $y_{\text{res}} \ll 1$ to large values $y_{\text{res}} \gg 1$, as discussed in detail in Ref. [18]. Since this process occurs, in general, during electron-positron annihilations, the neutrino temperature gets smaller than the photon temperature and, making use of entropy conservation, one has

$$T_\nu = T \left[\frac{g_S(m_e/T)}{g_S(0)} \right]^{\frac{1}{3}}, \quad (14)$$

where g_S is the number of entropy density ultrarelativistic degrees of freedom. For $T_\nu^{\text{res}}(y_{\text{res}} \ll 1, L_i) \lesssim 1$ MeV,

neutrino collisions can be neglected and at the resonance one has antineutrino conversions into dark neutrinos starting from small $y_{\text{res}} \ll 1$. If the resonance is crossed adiabatically, then all lightest antineutrinos are converted into dark neutrinos. Notice that if $L_i < 0$, then simply lightest neutrinos are converted into dark neutrinos instead of antineutrinos. For definiteness, we will usually refer to the case $L_i > 0$ in the following. More generally, to account also for nonadiabatic conversions, the fraction of converted neutrinos can be calculated using the Landau-Zener approximation:

$$f_{\nu_\mu \rightarrow \nu_{\text{dark}}} \simeq 1 - e^{-\frac{\pi}{2}\gamma_{\text{res}}}. \quad (15)$$

In this expression γ_{res} is the adiabaticity parameter at the resonance, given by

$$\gamma_{\text{res}} = \frac{|\Delta m^2| \sin^2 2\theta_0}{2y T_\nu^{\text{res}} H_{\text{res}}}, \quad (16)$$

where $H_{\text{res}} \simeq 0.2 \text{ s}^{-1} \sqrt{g_\rho(T_{\text{res}})} (T_{\text{res}}/\text{MeV})^2$ is the expansion rate at the resonance and g_ρ is the number of energy density ultrarelativistic degrees of freedom. The Landau-Zener approximation has been shown to reproduce quite well the numerical results obtained solving density matrix equation [18]. For simplicity, we can make use of a monochromatic approximation equivalent to say that all neutrinos are converted instantaneously at T_ν^{res} corresponding at $y_{\text{res}} = 3.15$. Using the prescription in Ref. [18], in this case one has to use $\langle L \rangle \simeq L_f/2 \simeq 0.2$ in the evaluation of the adiabaticity parameter, obtaining

$$\gamma_{\text{res}} \simeq 1.4 \times 10^9 \sqrt{\frac{10.75}{g_\rho(T_\nu^{\text{res}})}} \left(\frac{\Delta m^2}{\text{eV}^2} \right)^{\frac{1}{4}} \sin^2 2\theta_0. \quad (17)$$

A full momentum description would be able to describe the dependence of the adiabaticity parameter on momentum shown in Eq. (16). For values of y_{res} higher (lower) than 3.15, the adiabaticity parameter would be smaller than in the monochromatic approximation and so one can expect less (more) efficient conversion. However, bulk of neutrinos get converted for $y_{\text{res}} \simeq 3.15$ and one can expect that this value describes well the average effect. Therefore, an account of momentum dependence would produce just a slight correction, likely at the percent level, of the lower bound on the value of γ_{res} that in turn implies a lower bound on the active-dark neutrino mixing angle, as we derive in Sec. IV.⁴

⁴Within the current phenomenological status, such a small correction is very reliably negligible. If future upper bound on the neutrino effective magnetic moment will be improved by at least 2 orders of magnitude, then corrections of this kind might become worth to be included in the derivation of the allowed region. At this stage a full momentum description is not really needed.

B. Dynamical generation of the effective asymmetry

The effective lepton asymmetry that is necessary to create a resonance at low temperatures converting active neutrinos into dark neutrinos can also be generated dynamically by the same active-dark neutrino mixing [18,20]. This second case is quite attractive, since in this way one does not have to assume additional external mechanisms. If $L_i = 0$, the initial resonance at $T_\nu^{\text{res}}(0) \gtrsim 1 \text{ MeV}$ occurs in the collisional regime and triggers an initial exponential growth of the asymmetry up to a threshold value $L_t \simeq 0.4 \times 10^{-6} (\Delta m^2/\text{eV}^2)$ (not to be confused with $L_\star$). At this stage, collisions become ineffective, and the process enters a mixing-dominated regime, where antineutrinos get resonantly converted into dark neutrinos. If the process is fully adiabatic, a maximum final asymmetry $L_f = 0.375$ is ultimately generated, corresponding to the case when all antineutrinos ν_1 are converted into dark neutrinos. The exponential stage would then provide the initial value of the asymmetry, corresponding to L_i in Sec. II A, needed for the conversion of active to dark neutrinos. Notice that, since this asymmetry is generated after neutrino decoupling, such asymmetry is not constrained by CMB, since it does not result into any extra radiation compared to the standard case, unless the value of Δm^2 , and consequently of the resonant temperature, is large enough (above 0.1 eV^2) to produce some nonstandard big bang nucleosynthesis (BBN) and CMB effects [22]. This option might imply a reduced allowed region in the plane of Δm^2 versus $\sin^2 2\theta_0$ compared to the case of initial preexisting asymmetry discussed in Sec. II A. The reason is simply that the onset of the initial exponential growth stage prior to the resonant conversion is triggered by collisions and these become less efficient at values of $\sin^2 2\theta_0 \lesssim 10^{-10}$ [18]. This reduction would then affect just marginally the total (most conservative) allowed region that we show in Sec. IV.

III. DARK NEUTRINO DECAYS INTO DARK PHOTONS MIXED WITH PHOTONS

In order to explain the ERB, we assume that the dark neutrinos produced by active-to-dark neutrino conversions in the presence of a preexisting asymmetry decay into dark fermions ψ' with mass $m_{\psi'}$ and into a superposition of a dark and standard photon that we denote by γ' . The probability that γ' is detected as a photon, that can also be regarded as the decay branching ratio into photons, is denoted by ε . The dark photon is kinematically mixed with the standard photon but the kinetic mixing will not play any role, as we will comment. The dark fermion is assumed to be quasidegenerate with ν_0 , while we assume the dark photon mass $m_{\gamma'} \lesssim 10^{-15} \text{ eV}$ in order to evade the cosmological constraints from COBE/FIRAS for any value of the kinetic mixing parameter

χ_0 [23–26].⁵ We also assume that the decays can be described as fully nonrelativistically so that, at the decay, γ' has an energy $E_{\gamma'} \simeq \Delta m' \equiv m_0 - m_{\psi'}$. In this way the spectrum of nonthermal photons produced by the decays of the dark neutrinos and detected at the present time would be described by an effective temperature given exactly by the expression in Eq. (1) multiplied by ε and also divided by a factor $2/(1 - e^{-\frac{\varepsilon}{2}\gamma_{\text{res}}})$, taking into account that only relic antineutrinos are, adiabatically or nonadiabatically, converted into dark neutrinos.

The dark fermion radiative decay rate into γ' can be related to a dark neutrino effective magnetic moment μ'_{eff} by an expression analogous to Eq. (3):

$$\begin{aligned} \Gamma_{\nu_0 \rightarrow \psi' + \gamma'} &= \frac{(m_0^2 - m_{\psi'}^2)^3}{8\pi m_0^3} \mu_{\text{eff}}'^2 \simeq \frac{\Delta m'^3}{\pi} \mu_{\text{eff}}'^2 \\ &\simeq 42.5 \text{ s}^{-1} \left(\frac{\Delta m'}{\text{eV}} \right)^3 \left(\frac{\mu_{\text{eff}}'}{\mu_B} \right)^2. \end{aligned} \quad (18)$$

The important difference is that now μ'_{eff} does not have direct experimental constraints. However, because of the active-dark neutrino mixing, the active neutrino still has an effective neutrino magnetic moment $\mu_{\text{eff}} = \sqrt{\varepsilon} \sin^2 \theta_0 \mu'_{\text{eff}}$ and for this reason the experimental constraints on μ_{eff} still play a role.

IV. CONSTRAINTS AND ALLOWED REGION

Let us now combine all constraints and determine the allowed region in the parameter space of interest. First, we determine a convenient minimum set of parameters to display all the constraints. We start by imposing that the nonthermal photons produced by the dark neutrino decays, and mixed with the dark photons, can reproduce the ARCADE 2 data. We have then to impose that the dark neutrino lifetime is given by the lifetime determined in [7] in the case of direct relic neutrino decays (see best fit in Fig. 1), that we denote by $\tau_{\text{ARCADE}}(\nu_1 \rightarrow \nu_0 + \gamma)$, shortened by a factor $\varepsilon(1 - e^{-\gamma_{\text{res}}})/2$ to compensate the reduced photon production, explicitly:

$$\tau_{\nu_0 \rightarrow \psi' + \gamma'} = \frac{\varepsilon(1 - e^{-\gamma_{\text{res}}})}{2} \tau_{\text{ARCADE}}(\nu_1 \rightarrow \nu_0 + \gamma) \gtrsim t_0. \quad (19)$$

On the other hand, the lifetime cannot be shorter than t_0 , as also indicated in (19), since otherwise the exponential in the decay law kicks in and suppresses the photon effective temperature below the ARCADE 2 measured values.

⁵Here we do not assume the dark photon to be dark matter in which case the bounds on $m_{\gamma'}$ would be more stringent [27]. Moreover, we do not consider the Juno constraint on $m_{\gamma'}$ that extends down to $m_{\gamma'} \sim 10^{-18}$ eV [28], as we lack a detailed understanding of the Jovian magnetosphere and the associated uncertainties.

Since the lifetime is the inverse of the decay rate given in Eq. (3), the constraint (19) allows one to express μ'_{eff} in terms of γ_{res} . Therefore, for a fixed value of γ_{res} , one has a fixed value of μ'_{eff} . At the same time γ_{res} is expressed in terms of $\sin^2 2\theta_0$ and Δm^2 from Eq. (17).

We also have to impose the constraint $T_{\nu}^{\text{res}}(L_i) \lesssim 1$ MeV, corresponding to a neutrino collisionless regime.⁶ From Eq. (13) one can see that, for a fixed value of L_i , this constraint implies an upper bound on Δm^2 . Higher values of L_i correspond to higher allowed values of Δm^2 . However, one has to impose an upper bound $L_i \lesssim 10^{-2}$ from cosmological observations, since this would affect both BBN and CMB anisotropies [29]. It is then possible to express all other parameters in terms of $\sin^2 2\theta_0$, Δm^2 and L_i . Therefore, we can conveniently show all constraints in the plane Δm^2 versus $\sin^2 2\theta_0$, as shown in the left panel of Fig. 3 for the most conservative choice $\varepsilon = 1$, corresponding to $\gamma' = \gamma$, and highest value of $\tau_{\text{ARCADE}}(\nu_1 \rightarrow \nu_0 + \gamma)$ allowed at 99% CL [from Eq. (2)].

All constraints (shaded regions) are explicitly indicated. The horizontal dashed lines give the upper bound on Δm^2 for different values of L_i . One can also notice the experimental constraints on μ_{eff} that basically translate into an upper bound on $\sin^2 2\theta_0$. The orange line is the isocontour line for $\gamma_{\text{res}} = 1$, marking the border between the adiabatic and the nonadiabatic regime. In the right panel of Fig. 3 we also show constraints and allowed region in the plane μ_{eff} versus T_{ν}^{res} for the maximum allowed value of $L_i = 10^{-2}$. As one can see, we find a very interesting lower bound $\mu_{\text{eff}}/\mu_B \gtrsim 10^{-16}$ which might be testable [30]. Taking smaller values of L_i moves up this lower bound on μ_{eff} , making it even easier to be tested.

V. NEUTRINO OSCILLATION CONSTRAINTS

Having obtained an allowed region in the Δm^2 versus $\sin^2 2\theta_0$ plane, we compare it with the existing constraints from neutrino oscillation experiments. Active research is underway on sterile neutrinos with masses around the eV scale, motivated by the excess observed by LSND [31], and the intriguing results reported by MiniBooNE [32], BEST [33], and IceCube [34]. However, many other experiments have not found such evidence, leading to significant tension among datasets when combined [35,36].

Beyond the eV-scale sterile neutrino, searches have explored a wide range of mass for oscillations between active and sterile states [37–48], but no evidence has been found. For nearly degenerate active and sterile states, solar [42,49,50] and reactor data [51] constrain $\Delta m^2 \lesssim 10^{-12}$ eV² and $\sin^2 2\theta \lesssim 10^{-4}$, as shown in Fig. 4.

⁶Here we consider the less constrained option assuming the presence of an initial preexisting asymmetry L_i , as discussed in Sec. II A.

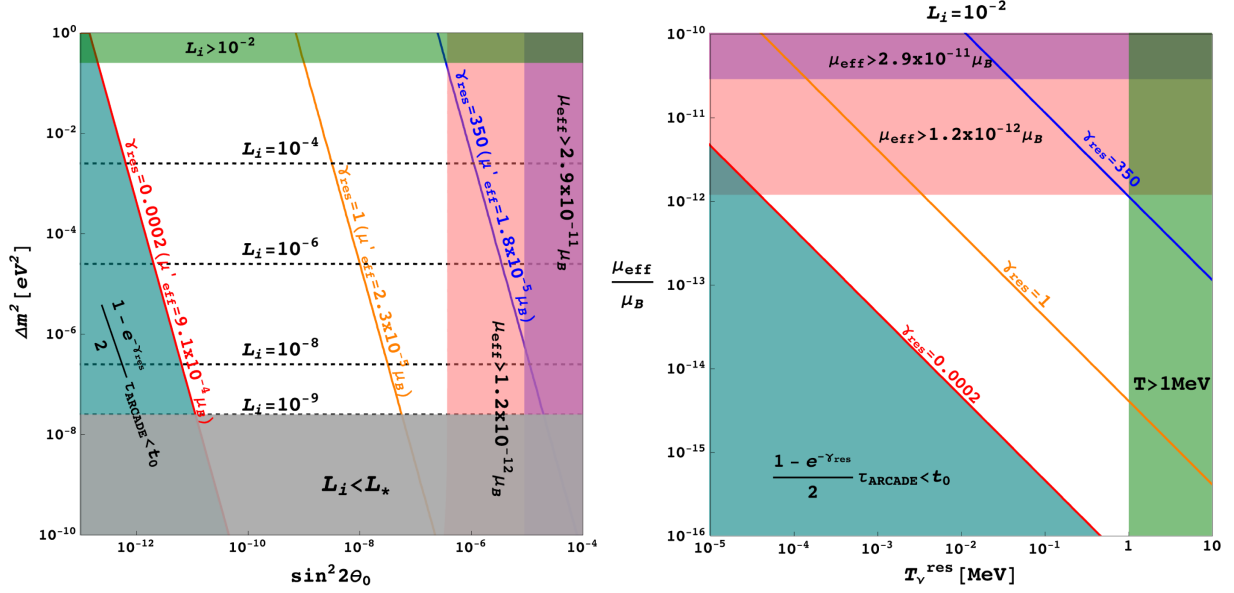


FIG. 3. Left panel: constraints (shaded) and allowed region (white) in the plane of Δm^2 versus $\sin^2 2\theta_0$. The horizontal dashed lines give an upper bound on Δm^2 from $T_\nu^{\text{res}} < 1$ MeV for the indicated values of L_i . Right panel: constraints (shaded) and allowed region (white) in the plane of μ_{eff} versus T_ν^{res} . We have used conservatively $\varepsilon = 1$.

The disappearance bounds of muons come from long-baseline experiments [52,53] and atmospheric experiments [54], $\Delta m^2 \lesssim 10^{-4}$ eV² and $\sin^2 2\theta_0 \lesssim 10^{-2}$. Astrophysical neutrinos, with large L/E , probe down to $\Delta m^2 \sim 10^{-20}$ eV² for maximal mixing [46,55–59].

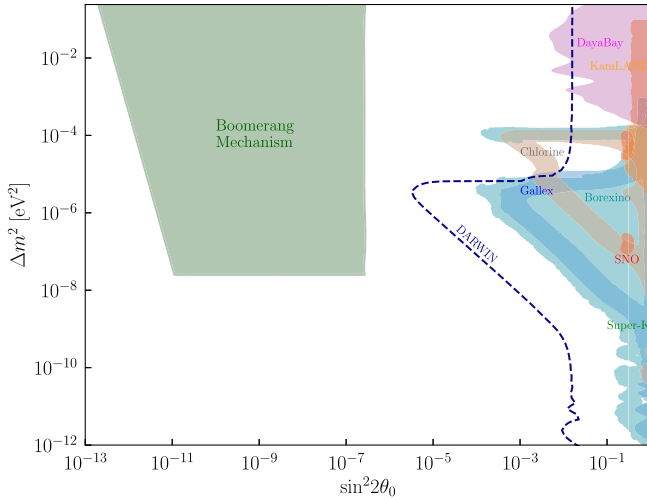


FIG. 4. Region allowed by the boomerang mechanism at 99% CL, along with the region excluded by oscillation experiments. We present some of the most stringent bounds on the mixing, including measurements from solar neutrino experiments such as Borexino [50], Gallex [60], Chlorine [61], SNO [62], and Super-Kamiokande [63] as well as reactor experiments like KamLAND [50] and Daya Bay [51]. We also show the future sensitivity from DARWIN [64].

Figure 4 compares the current bounds and future sensitivities with the region predicted by the boomerang mechanism (cf. the left panel of Fig. 3). Although this work focuses on muon-sterile oscillations, the mechanism applies to all flavors, so we adopt the strongest available limits on the mixing.

Upcoming dark matter experiments such as DARWIN [56,64] are expected to measure the solar neutrino flux from pp fusion via electron scattering. The presence of a light sterile neutrino would modify the electron-neutrino spectrum arriving at Earth. For very small mass splittings ($\Delta m^2 \sim 10^{-12}$ eV²), sterile neutrinos induce vacuum oscillations along the Earth-Sun baseline [56]. As Δm^2 increases, the oscillation length decreases and matter effects become increasingly relevant, enhancing the sensitivity to smaller mixing angles. For mass splittings around $\Delta m^2 \sim 10^{-6}$ eV², mixings as small as $\sin^2 2\theta_0 \sim 10^{-5}$ could be probed with a 300-ton-year exposure. Reactor experiments, such as JUNO [65], will also be able to explore the oscillation between active and sterile state for very small mixing angles.

VI. FINAL REMARKS

- (i) The boomerang mechanism effectively realizes a solution of the ERB mystery obtained in terms of radiative relic neutrino decays [7] but in two stages: in a first early stage relic antineutrinos (or neutrinos, depending on the sign of L_i) of one species are converted into dark neutrinos (the visible sector throws particles into the dark sector); in a second stage dark neutrinos decay into dark-standard photon

states (dark sector throws back particles into the visible one).⁷

- (ii) As Fig. 3 shows, it has a broad variety of phenomenological implications that make it testable in different ways. First of all the TMS experiment will soon verify the ARCADE 2 data, the existence of the ERB and the solution proposed in Ref. [7]. Notice also that this implies some nonstandard deviation in the 21 cm cosmological global signal (see [7] for details). Possible cosmological anomalies in BBN and/or CMB anisotropies might be addressed by the presence of a large lepton asymmetry [66]. Also notice that it implies nonstandard relic neutrino background properties that might be potentially measured [18]. Finally, the lower bound we found on μ_{eff} , 4 orders of magnitude below the current upper bound, will be tested by future experiments [30].
- (iii) The monochromatic approximation we used is well justified by numerical solution of density matrix equation with a full momentum description [18]. This also grasps the variation of the nonadiabaticity parameter with momentum. However, as discussed in Sec. IV, one would have small corrections to the allowed region we have derived.
- (iv) The allowed regions have been determined for the most conservative case $\varepsilon = 1$, corresponding to the extreme minimal case where existence of dark photons would be not necessary. They would clearly shrink for lower values until they would disappear for a lower bound $\varepsilon \sim 2 \times 10^{-4}$.
- (v) Having imposed $m_{\gamma'} < 10^{-15}$ eV suppresses the kinetic mixing since $m_{\gamma'}^2 \gg m_{\gamma}^2$, evading microwave background constraints [23]. Since kinetic mixing is suppressed, we cannot have $\gamma' = \gamma_{\text{dark}}$ (i.e., $\varepsilon = 0$) and a dynamical generation of a photon component due to a large kinetic mixing parameter χ_0 ; also, since kinetic mixing does not play a role, one could optionally have $m_{\gamma'} > 10^{-15}$ eV and negligible kinetic mixing without violating microwave background constraints.
- (vi) Notice that having ε and $\theta_0 \neq 0$ could induce a nonvanishing neutrino millicharge in addition to an effective magnetic moment and this might introduce further constraints to be taken into account [11]. However, these constraints are model dependent and we have assumed that the neutrino millicharge is negligible. In any case these could be circumvented replacing a constant θ_0 with a temperature-dependent effective mixing angle such that $\theta(T_{\text{res}}^\nu)$ coincides

with the required values for the mechanism to work shown in the left panel of Fig. 3, while for $T \ll T_{\text{res}}^\nu$ this is sufficiently small to generate a neutrino millicharge in agreement with all constraints. Such a temperature-dependent effective mixing angle can be generated introducing nonstandard neutrino interactions yielding off-diagonal (temperature-dependent) terms in the effective Hamiltonian [see Eq. (8)].

- (vii) We also have to take into account stellar cooling constraints from plasmon decays $\gamma^* \rightarrow \nu_0 + \psi'$. However, these can be circumvented coupling the dark fermions to a scalar in a way that either they get a mass higher than plasmon mass [67] or that the coupling to photons in stars is suppressed by having a symmetry restoration at an energy scale $\mathcal{O}(\text{keV})$ such that the scalar vacuum expectation value vanishes [68]. Interestingly, such a low-scale phase transition has been proposed with independent motivations, within a split-seesaw Majoron model [69].
- (viii) We have not discussed how the boomerang mechanism could be embedded within a full ultraviolet-complete model. For example, a possible direction is offered by models incorporating quasi-Dirac neutrinos [48,70]. Typically, we have maximal mixing for $\Delta m \ll m$, as $\tan 2\theta \sim 2m/\Delta m$, but arbitrary mixing is possible for certain textures of the Dirac and Majorana mass matrices, akin to the case of low-scale type-I seesaw [71]. Also notice that only within a definite model one could have more specific relation between ε and χ_0 and calculate neutrino millicharge and consequent constraints.

Before concluding, we would like to highlight that experimentally we will have first important results from the TMS experiment [4,5], that will test the ERB in the 10–20 GHz range within the next few years. If these should support a solution in terms of relic neutrino decays, then the identification of the mechanism responsible for the instability of the relic neutrino background would become a priority and this would justify dedicated numerical analyses, for example, including a full momentum dependence and dynamical asymmetry generation. An experimental improvement of the upper bound on the neutrino muon effective magnetic moment would then become very important and in case of a positive signal would strongly support the boomerang mechanism. Active-sterile neutrino oscillation experiments testing very small mixing angles would at this stage become strongly motivated, representing an ultimate test of the mechanism. In conclusion, the boomerang mechanism shows that a solution to the ERB in terms of relic radiative neutrino decays is possible. If the excess is confirmed, this might then provide a direct open window for the exploration of a new dark sector.

⁷It should be noticed how the second stage occurs much later (it is basically occurring now). That implies that dark neutrinos are fully nonrelativistic when they decay and their initial momentum distribution, after active-to-dark neutrino conversions, has no impact on the photon spectrum.

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DATA AVAILABILITY

The data are not publicly available. The data are available from the authors upon reasonable request.

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