The Sun: Light Dark Matter and Sterile Neutrinos

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Abstract

Next-generation experiments allow for the possibility of testing the neutrino flavor oscillation model to very high levels of accuracy. Here, we explore the possibility that the dark matter in the current universe is made of two particles, a sterile neutrino and a very light dark matter particle. By using a 3+1 neutrino flavor oscillation model, we study how such a type of dark matter imprints the solar neutrino fluxes, spectra, and survival probabilities of electron neutrinos. The current solar neutrino measurements allow us to define an upper limit for the ratio of the mass of a light dark matter particle $m_f$ and the Fermi constant $G_F$, such that $G_F/m_f$ must be smaller than $10^{30} G_F$ eV$^{-1}$ to be in agreement with current solar neutrino data from the Borexino, Sudbury Neutrino Observatory, and Super-Kamiokande detectors. Moreover, for models with a very small Fermi constant, the amplitude of the time variability must be lower than 3% to be consistent with current solar neutrino data. We also found that solar neutrino detectors like Darwin, able to measure neutrino fluxes in the low-energy range with high accuracy, will provide additional constraints to this class of models that complement the ones obtained from the current solar neutrino detectors.

Unified Astronomy Thesaurus concepts: The Sun (1693); Solar neutrino problem (1510); Solar neutrinos (1511); Neutrino oscillations (1104); Neutrino telescopes (1105); Neutrino astronomy (1100)

1. Introduction

The origin of dark matter has been a fundamental problem in physics for almost six decades, during which most of the proposed solutions assumed that a single massive particle that interacts weakly with baryons makes all the dark matter observed in the universe (e.g., Wang et al. 2016). Recently, research has emerged where more sophisticated solutions have been proposed to solve the dark matter problem. One of these is the possibility of the dark matter being a composite of two light particles: a light dark matter (LDM) particle $f$ and a sterile neutrino $\nu_s$.

The existence of such an LDM field can be identified with a dilution field of an extradimensional extension of the Standard Model or/and a CP-violating pseudo-Goldstone boson of a spontaneously broken global symmetry. For some of these models, $\phi$ couples to the Standard Model fields, and as such it induces periodic time variation in particle masses and couplings. In such theories the gauge invariance suggests that the $\phi$ should possess an identical coupling constant to charged leptons, in which case scalar interactions with the electrons provide a good opportunity for detection through atomic clocks (e.g., Arvanitaki et al. 2016), accelerometers (e.g., Arvanitaki et al. 2018), and gravitational wave detectors (e.g., Lopes & Silk 2014; Graham et al. 2016).

Similarly, this $\phi$ field can couple to neutrinos. Once again, these types of interactions generically result in time-varying corrections to the neutrino masses, neutrino mass differences, and mixing angles, which can be searched for in the neutrino flux signals on present and future experimental neutrino detectors (Aharmim et al. 2013; Abe et al. 2016; Borexino Collaboration et al. 2018; Aalbers et al. 2020). If $\phi$ couples weakly to the neutrinos, over a large range of masses, it can significantly modify the neutrino oscillations probabilities leading to a distorted survival electron neutrino probability function (Berlin 2016; Krnjaic et al. 2018).

The motivation for such a model comes from the possibility of this composite particle physics model resolving two observational problems:

1. The classical cold dark matter model leads to several inconsistencies with the cosmological observational data, such as the missing satellite problem and the cusp problem (e.g., Primack 2009). Light dark matter resolves such problems if dark matter is totally or partially made of light scalar particles with a mass of the order of $10^{-22}$ eV (Hu et al. 2000; Peebles 2000). In hierarchical models of structure formation, such a type of dark matter is able to explain the flatness observed on the profiles of the distribution of gas and stars in halos and filaments (Mocz et al. 2019).

2. Although the standard three-neutrino flavor model produces a reasonable good global fit to all the neutrino data (Esteban et al. 2019), there are now many hints that point out the possibility that a fourth neutrino exists. This one does not have any interaction other than gravity and for that reason it is known as a sterile neutrino (Diaz et al. 2019). It was found that a flavor oscillation model made of the three active neutrinos plus a sterile neutrino could explain some of the observed anomalies found on the short-baseline neutrino oscillation experiments (Giunti et al. 2012, 2013), Liquid Scintillator Neutrino Detector (Aguilar et al. 2001), and MiniBooNE Short-Baseline Neutrino Experiment (MiniBooNE Collaboration et al. 2018), as well as the anomalies related with GALLEX and SAGE solar neutrino detectors—the so-called gallium anomalies (Kostensalo et al. 2019). For instance, Kostensalo et al. (2019) found that the data favor a 3+1
neutrino flavor model with $m_2 = 1.1\text{eV}$ and mixing matrix element $U_{e4} = 0.11$.

One possibility to resolve both problems (neutrino anomalies and structure formation) is to consider that dark matter is made of a light scalar field that couples to a sterile neutrino (e.g., Farzan 2019). The interactions of $\phi$ and $\nu_i$ could impede the oscillations in the universe and thereby improve the agreement between the structure formation and cosmological observations (e.g., Dasgupta & Kopp 2014; Hannestad et al. 2014).

If such $\phi$ and $\nu_i$ particles exist today, they were produced abundantly in the early universe. For instance, sterile neutrinos can be produced via mixing with active neutrinos (Dodelson & Widrow 1994), in some scenarios such neutrino production is being enhanced by the oscillations between active and sterile neutrinos (Bezrukov et al. 2019, 2020; de Gouvêa et al. 2020) or by the lepton asymmetry (Shi & Fuller 1999). The production of light dark matter can take many forms, such as vector bosons by parametric resonance production (Dror et al. 2019). For instance, some models predict a sterile neutrino abundance $\Omega_{\nu_i} h^2 = 0.12(\sin^2(\theta_i)/3.5 \times 10^{-9})(m_{\nu_i}/7\text{keV})$, where $m_{\nu_i}$ and $\theta_i$ is the sterile neutrino mass and mixing angle (Kusenko 2009). For the light dark matter field some authors find $\Omega_{\phi} h^2 = 0.1(a_\phi/10^{17}\text{GeV})^2 (m_{\phi}/10^{-22}\text{eV})^{1/2}$, where $a_\phi$ is a parameter that relates to the initial misalignment of the axion, and $m_{\phi}$ is the axion mass (Hui et al. 2017; Niemeyer 2019). Conveniently, we will assume that in the present-day universe the total dark matter abundance is given by

$$\Omega_{\text{DM}} h^2 = \Omega_\phi h^2 + \Omega_{\nu_i} h^2,$$

where the $\Omega_\phi h^2$ and $\Omega_{\nu_i} h^2$ are the total $\phi$ and $\nu_i$ densities in the present universe, respectively. For future reference, we assume that the present-day total dark matter abundance $\Omega_{\text{DM}} h^2 = 0.12$ (Planck Collaboration et al. 2018), and the dark matter density in the solar neighborhood is $\rho_{\text{DM}} = 0.39\text{GeV cm}^{-3}$ (Catena & Ullio 2010).

In this paper, we study the impact that this light dark matter field has in the 3+1 neutrino flavor model. Specifically, we discuss how the light dark matter field modifies the neutrino flavor oscillations, and by using the current sets of solar neutrino data, we also put constraints in the parameters of such models and make predictions for the future neutrino experiments.

The article is organized as follows. In Section 2, we discuss how the light dark matter drives the 3+1 neutrino flavor oscillations. In Section 3, we present the neutrino flavor oscillation model in the presence of a cosmic light dark matter field. In Section 4, we compute the survival electron neutrino probabilities for the electron neutrinos produced in the proton–proton (PP) chain and carbon–nitrogen–oxygen (CNO) cycle solar nuclear reactions. In Section 5, we discuss the results in relation to current experiments and future ones. Finally, in Section 6, we present the conclusion and a summary of our results.

If not stated otherwise, we work in natural units in which $c = \hbar = 1$. In these units all quantities are measured in GeV, and we make use of the conversion rules $1\text{m} = 5.068 \times 10^{15}\text{GeV}^{-1}$, $1\text{kg} = 5.610 \times 10^{26}\text{GeV}$ and $1\text{s} = 1.519 \times 10^{24}\text{GeV}^{-1}$.

### 2. Light Dark Matter and Sterile Neutrinos in the Universe

We assume that in the present universe, the dark matter is composed of two fundamental particles: a light scalar boson $\phi$ and sterile neutrinos $\nu_s$, where $m_\phi$ and $m_{\nu_s}$ are their respective masses (Hannestad et al. 2014). The LDM field $\phi$ couples with the active neutrinos and the sterile neutrino by a Yukawa interaction $g_{\phi}\phi \nu_i \nu_s$ where $g_{\phi}$ is a dimensionless coupling (Farzan 2019). To illustrate this effect, consider an LDM scalar $\phi$ with a Yukawa coupling to active neutrinos. Then the relevant part of the Lagrangian reads

$$\mathcal{L} \supset - (m_{\nu} + g_{\phi})\nu \nu + \text{H.c.},$$

where for convenience of representation the flavor indices have been suppressed. We also assume that the dimensionless coupling is very small ($g \ll 1$). From the Euler–Lagrange equations of $\phi$ and $\nu$, it is possible to show that the effect of $\phi$ on the propagation of the neutrino is equivalent to changing the neutrino mass from $m_{\nu}$ to $m_{\nu} + \delta m_{\nu}$. As we will see later, this $\delta m_{\nu}$ perturbation will induce time variations in the mass-squared differences and mixing angles of all neutrino flavors through $\phi$. In principle the Yukawa couplings can have any structure in the neutrino flavor space. In this work, we will focus on two convenient scenarios of great interest to neutrino detectors: mass-square differences and mixing angles (e.g., Ding & Feruglio 2020). Moreover, we will also assume that $\delta m_{\nu} = g_{\phi} \phi$ (e.g., Smirnov & Xu 2019).

#### 2.1. Dark Matter Time-dependent Variation

The hypothesis that dark matter in the local universe is made of very light particles leads to the following description: the LDM field $\phi$ in the dark matter halo of the Milky Way is represented by a group of plane waves with frequency $\omega_{\phi}$, such that $\omega_{\phi} = m_{\phi}/(1 + v_s^2/2)$ where $v_s$ is the virial velocity of the particles in the dark halo. A population of such light particles will smooth inhomogeneities in the dark matter distribution on scales smaller than the de Broglie wavelength $\lambda_{\text{DB}}$ of these LDM particles. For any particle, we compute $\lambda_{\text{DM}}$ using the relation $\lambda_{\text{DM}} = 1.24 \times 10^{20}(10^{-22}\text{eV}/m_{\phi})^{1/3}/v_s \text{cm}$.

We notice that the kinetic term on $\omega_{\phi}$ is neglected once the virial velocity $v_s \sim 10^3$ is very small (e.g., Blas et al. 2017). Therefore, we dropped the corrections related to $v_s$ for the Equation (2). Accordingly, the general form of this LDM field reads

$$\phi(\mathbf{r}, t) = \phi_o \cos(m_{\phi} t + \epsilon_o) \approx \phi_o \cos(m_{\phi} t),$$

where $\phi_o$ and $\epsilon_o = v_s \cdot \mathbf{r}$ are the amplitude and phase of the wave $\phi(t)$, respectively. In this work we consider $\epsilon_o \approx 0$. Moreover, the energy momentum of a free massive oscillating field has a density given by $\rho_\phi = \rho_o m_{\phi}/2$ and a pressure given by $p_\phi = \rho_\phi \cos(2m_{\phi} t)$. Although formally $\rho_\phi$ has an oscillating part proportional to $\phi(t)$, because this component is very small we neglected its contribution in this analysis (Khmelnitsky & Rubakov 2014).

The quantity $\rho_\phi$ is a slowly varying function of the position. Conveniently, the amplitude of $\phi(t)$ can be written as $\phi_0 = \sqrt{2\rho_\phi(r)/m_{\phi}}$ where $\rho_\phi(r) = \rho_{\phi\text{DM}}(\Omega_{\phi}/3\Omega_{\text{DM}})$ is the fraction of dark matter density in $\phi$ particles at the spacetime coordinate $r$. Accordingly, the Sun immersed in this light dark matter halo will experience a periodic perturbation due to the action of the $\phi(r, t)$, which by the presence of a Yukawa coupling $g_{\phi}$ will exert a temporal variation on the propagation of all neutrinos. We estimate the dark matter density number $n_{\phi}$ in the solar neighborhood as follows: if we consider that the
main contribution arises from a single dark matter particle with mass $m_\phi$, then the relevant density in our case will take the value $n_\phi = \rho_0 / m_\phi$. If we assume that all dark matter is made of $\phi$ bosons, we have $\rho_0 = \rho_{0,DM} = 0.39$ GeV cm$^{-3}$ (Catena & Ullio 2010) and $m_\phi = 10^{-22}$ eV then $n_\phi = 3.9 \times 10^{30}$ cm$^{-3}$ (particles per centimeter cubed). This value is only 2 orders of magnitude smaller than the density of electrons in the Sun’s core, $n_e \sim 6 \times 10^{33}$ cm$^{-3}$ (Lopes & Turck-Chièze 2013). Since these particles are very light, we assume that there is no accretion of these particles in the Sun’s core during its evolution in the main sequence until the present age.

### 2.2. Neutrino Time-dependent Dark-matter-induced Oscillations

In the presence of the LDM field $\phi$, the neutrino mass $m_\nu$, according to Equation (2) (Ding & Feruglio 2020), will receive a contribution $\delta m_\nu = g_\phi$, such that from Equation (3), we obtain

$$\frac{\delta m_\nu}{m_\nu} = \epsilon_\phi \cos(m_\phi t), \quad (4)$$

where $\epsilon_\phi$ is the amplitude

$$\epsilon_\phi = \frac{g_\phi \sqrt{2\rho_0}}{m_\phi m_\nu} = \frac{g_\phi \sqrt{2\rho_{0,DM}}}{m_\phi m_\nu} \left( \frac{\Omega_\phi}{\Omega_{DM}} \right)^{1/2}. \quad (5)$$

If not stated otherwise, we will assume that all dark matter in the present universe is made of only LDM particles such that $\Omega_\phi = \Omega_{DM}$. We observe that $\epsilon_\phi$ is a relevant factor even if $\phi$ is a small fraction of the dark matter halo. In particular, $\phi$ will affect the oscillation parameters of all neutrino flavors, including the sterile neutrinos. If we only take into account the first order perturbation, thus, the neutrino mass-squared difference can be written as

$$\Delta m^2_\nu(t) = m^2_\nu - m^2_\nu \approx \Delta m^2_{\nu,0} (1 + 2\epsilon_\phi \cos(m_\phi t)), \quad (6)$$

where $\Delta m^2_{\nu,0}$ is the standard (undistorted) value and $\Delta m^2_{\nu}(t)$ evolves through $\cos(m_\phi t)$ (see Equation (3)), and a frequency $m_\phi$. The mass-squared difference $\Delta m^2_{\nu}$ between neutrinos of different flavors follows the usual convection (e.g., Lopes 2017) such that $\Delta m^2_{\nu} = m^2_\nu - m^2_\nu (i = 2, 3, 4)$. In particular for the sterile neutrino, we have $\Delta m^2_{\nu} = m^2_{\nu} - m^2_4$ where $m_4$ is the mass of the sterile neutrino. Similarly, the mixing angles variation is written as

$$\theta_{ij}(t) \approx \theta_{ij,0} + \epsilon_\phi \cos(m_\phi t), \quad (7)$$

where $\theta_{ij,0}$ is the standard (undistorted) mixing angle. The indexes $i$ and $j$ in $\theta_{ij}$ follow a convention identical but not equal for the mass-squared differences (see Lopes 2018a, and references therein). Therefore, as first suggested by Kmjaic et al. (2018), the LDM $\phi(t)$ impacts the neutrino flavor oscillations through the modified expressions for the mass-squared differences (Equation (6)) and mixing angles (Equation (7)).

### 3. Light Dark Matter and the Sterile Neutrino Model

In the following section, we consider a 3+1 neutrino flavor oscillation model to describe the propagation of active neutrinos ($\nu_e$, $\nu_\mu$, $\nu_\tau$) plus a sterile neutrino $\nu_4$ through the solar plasma. Following the usual notation ($\nu_e$, $\nu_\mu$, $\nu_\tau$, $\nu_4$) corresponds to the neutrino flavors, ($\nu_1$, $\nu_2$, $\nu_3$, $\nu_4$) are the mass neutrino eigenstates, and ($m_1$, $m_2$, $m_3$, $m_4$) are the neutrino masses. The evolution of neutrinos propagating in matter is described by the equation

$$i \frac{d\Psi}{dr} = \mathcal{H}\Psi = \frac{1}{2E} (UM^2U^\dagger + 2E\nu)\Psi, \quad (8)$$

where $\mathcal{H}$ is the Hamiltonian and $\Psi = (\nu_e, \nu_\mu, \nu_\tau, \nu_4)^T$. $M^2$ is a neutrino mass matrix, $U$ is a $(4 \times 4)$ unitary matrix describing the mixing of neutrinos and $\nu$ is the diagonal matrix of Wolfenstein potentials (Kuo & Pantalone 1989). $M^2$ is defined as $M^2 = \text{diag}(0, \Delta m^2_{31}, \Delta m^2_{32}, \Delta m^2_{33})$. The first term of the Hamiltonian describes the neutrino propagation through vacuum and the second term incorporates the matter effects or Mikheyev–Smirnov–Wolfenstein (MSW) effects (Wolfenstein 1978; Mikheyev & Smirnov 1985). In general, the Hamiltonian $\mathcal{H}$ that drives the evolution of neutrino flavor must include the Wolfenstein potentials related with $\phi(t)$ (Brdar et al. 2018).

In most studies of three-neutrino flavor models, the authors are solely interested in the modulation coming from the square mass differences $\Delta m^2_{\nu}(t)$ (by Equation (6)) and mixing angles $\theta_{ij}(t)$ (by Equation (7)). For that reason, all neutrinos are assumed to couple $\phi(t)$. As a consequence, their contribution to $\nu$ cancels out. Hence, it is correct to neglect the contribution of $\phi(t)$ to the Wolfenstein potential (Dev et al. 2020). Nevertheless, here in this 3+1 neutrino flavor model, as we will discuss later, we include the contribution of $\phi(t)$ in $\nu$.

This 3+1 neutrino flavor model with dark matter is identical to the standard (undistorted) three-neutrino flavor model (see Equation (8)). However, in this model we included a sterile neutrino, and the Wolfenstein potentials in $\nu$ are modified to take into account the new LDM field $\phi$ (Miranda et al. 2015).

### 3.1. Neutrino Matter-induced Oscillations

In the standard three-neutrino flavor model, the matter potential $\nu$ takes into account the interaction of active neutrinos ($\nu_e$, $\nu_\mu$, $\nu_\tau$) with the ordinary fermions of the solar plasma, for which the $\nu$ = diag ($V_{ee}$, $V_{em}$, $V_{en}$) where $V_{em}$ corresponds to the weak charged current (cc) that takes into the forward scattering of $\nu_e$ with electrons, and $V_{en}$ is the weak neutral current (nc) that corresponds to the scattering of the active neutrinos with the ordinary fermions of the solar plasma (e.g., Xing & Zhong 2020). $V_{en}$ can be expressed as $V_{en} = V_{nc} + V_{ip} + V_{in}$ where $V_{nc}$ with $j = e, p, n$ are the contributions coming from electrons, protons, and neutrons, respectively. However due to the electrical neutrality of the solar plasma, the contribution of $V_{nc}$ and $V_{in}$ canceled out such that $V_{en} = V_{ip}$. Accordingly, $V_{ip} = \sqrt{2}G_F \, n_p(r)$ and $V_{in} = V_{nc} = G_F / \sqrt{2} \, n_e(r)$. Here $G_F$ is the Fermi constant and $n_p(r)$ and $n_e(r)$ are the number density of electrons and neutrons inside the Sun. Nevertheless, since $V_{nc}$ is a universal term for all active neutrino flavors, and as such does not change the flavor oscillations pattern, conveniently we write $\nu$ = diag ($V_{ee} + 0$, 0, 0). Now, the inclusion of sterile neutrinos in the neutrino flavor model alters $\nu$ (from Equation (8)) by incorporating a new degree of freedom, as a consequence

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5 In this model, the intermediate particle is a heavy boson, specifically the $Z$ or $W^\pm$ bosons.
\(V = \text{diag} \{ V_{cc}, V_{nc}, 0 \}\) (e.g., Giunti & Li 2009; Maltoni & Smirnov 2016; Xing & Zhong 2020).

Finally, in our 3+1 neutrino flavor model, we include the interaction of active and sterile neutrinos with the dark matter field \(\phi\) by means of an intermediate heavy boson \(I^4\). These interactions result from the forward scattering of these neutrinos through the LDM field \(\phi\), thus \(V = \text{diag} \{ V_{cc}, V_{nc}, V_{cc,0}, V_{nc,0}\}\), where \(V_{cc}\) (with \(\nu_i = \nu_e, \nu_\mu, \nu_\tau\)) relates to the neutrino \(\nu_i\). This \(V\) corresponds to a generalization of the Wolfenstein potentials found in the literature, for which most neutrino flavor models only take into account the scattering of the sterile neutrinos on heavy dark matter (Capozi et al. 2017; Lopes 2018a; Lopes & Silk 2019).

In our model, we opt to assume that all active neutrinos experience the same interaction with the LDM field \(\phi\), such that their dark matter potentials are the same, such that \(V_{cc,0} = V_{nc,0}\) (with \(j = e, \mu, \tau\)), it follows \(V = \text{diag} \{ V_{cc}, V_{nc}, V_{cc,0}, V_{nc,0}\}\). Now, if we subtract the common term \(V_{cc} + V_{nc,0}\) to the diagonal matrix \(V\), the latter takes the simple form: \(V = \text{diag} \{ V_{cc}, 0, 0, V_{cc,0} - V_{nc,0}\}\).

The potential \(V_{cc}\) (with \(a = s, b\)) is given by \(V_{cc} = G_{c,0} m_0\) where \(G_{c,0}\) is the equivalent of the Fermi constant and \(m_0\) is the distribution of dark matter inside the Sun (Smirnov & Xu 2019). Equally, \(V_{nc}\) relates directly with the local density of dark matter \(\rho_{DM}\) by the expression: \(V_{nc} = (G_{n,0} m_0) (\rho_{DM} M/3\Omega_{DM})\), where we assume the ratio \(G_{n,0}/m_0\) is a free parameter of the LDM model. The generalized Fermi constant is defined as \(G_{n,0} = g_e g_s/m^2\) where \(g_e, g_s\) represents the coupling constant of the corresponding neutrino \(\nu_i\) and \(m\) is the mass of the intermediate boson \(I\) (Miranda et al. 2015). This expression for the potential \(V_{nc}\) is valid since we assume that \(m^{-1} \ll R_o\) where \(R_o\) is the solar radius (Smirnov & Xu 2019). In general, we could expect that the contribution of \(\phi(t)\) to \(V_{nc}\) could lead to a time-dependent relation, however, as discussed previously (in Section 2.1) and mentioned for the first time by Khmelnitsky & Rubakov (2014), this is because the oscillatory component on the local density relates with \(V_{cc}\). This term is minimal, and therefore we neglected it.

In this preliminary study, without loss of generality, we choose to simplify \(V\) further: since the term \(V_{cc,0} - V_{nc,0}\) has two Wolfenstein potentials \(V_{cc,0}\) and \(V_{nc,0}\) that effectively correspond to two new degrees of freedom, both of these have an identical impact on the neutrino flavor oscillation model. We choose to simplify the model by assuming that \(V_{cc,0}\) is much smaller than \(V_{nc,0}\). Consequently, \(V\) takes the simplified form: \(V \approx \text{diag} \{ V_{cc,0}, 0, 0, V_{nc,0} - V_{cc,0}\}\). For reference, we note that in the Sun’s core \(V_{nc,0}\) is always smaller than \(V_{cc,0}\) once \(n_o\) is more than twice as large as \(n_e\) (e.g., Lopes 2018b). This potential is identical to others found in the literature, for instance in Capozi et al. (2017) and Lopes (2018a). Therefore, the matter potential \(V_{nc,0}\) reads \(V_{nc,0} = G_{n,0} m_0\) where for convenience of analysis, we choose to define the generalized Fermi constant as \(G_{n,0} = 4\sqrt{2} G_n G_F\) where \(G_n\) is our free parameter. Since these dark matter particles have a mass much smaller than 4 GeV, the solar plasma conditions do not allow the accretion of dark matter by the Sun (e.g., Lopes & Lopes 2019), therefore we will assume that the distribution of dark matter inside the star is equal to the value measured for the solar neighborhood \(n_o\) (see Section 2.1).

### 3.2. Neutrino Flavor Oscillation Model and the Survival Probability of Electron Neutrinos

If we adopt as reference the current experimental set of parameters for the active neutrinos (e.g., Esteban et al. 2019), the propagation neutrinos in the solar interior are completely adiabatic. The same is valid for the 3+1 neutrino flavor oscillation model coupled to an LDM field \(\phi\) considered in this study. Conveniently, the propagation of neutrinos away from resonances is well represented by a two neutrino flavor oscillation model. The motivation for such approximation can be found in Lopes (2018b) and references therein. In such a case, the electron neutrino flavor oscillation is dominated by the \((\nu_1, \nu_2)\) mass eigenstates and is only slightly affected by the decoupled \((\nu_3, \nu_4)\) eigenstates, since the associated mixing angles for the latter pair are very small (Kuo & Pantaleone 1986). Moreover, \(\nu_3\) and \(\nu_4\) evolve independent of each other and are completely independent of the doublet \((\nu_1, \nu_2)\).

In this limit, as proposed by several authors (e.g., Palazzo 2011; Blennow & Smirnov 2013), the split of the 3+1 neutrino flavor model into a dominant two neutrino flavor model \((\nu_e, \nu_x)\) with additional corrections for \(\nu_x\) and \(\nu_e\) significantly simplified the calculation and allowed us to obtain an analytical solution (e.g., Kuo & Pantaleone 1989).

Among the many expressions available in the literature to compute the survival probability of electron neutrinos \(P_e\) (e.g., Lunardini & Smirnov 2000; Miranda et al. 2015) in a 3+1 neutrino flavor model developed in the approximate scenario of a two-flavor neutrino model (e.g., Kuo & Pantaleone 1989), we opted to choose the expression obtained by Capozi et al. (2017) for the case in which \(V_{ee}/\Delta m^2_{31} \ll 1\) (and \(s_{43} = 0\) which has a better numerical accuracy than others. In that case the survival probability of electron neutrinos, i.e., \(P_e \equiv P(\nu_e \rightarrow \nu_e)\), reads

\[
P_e (E, \phi) = s_{43}^2 + c_{43}^2 s_{45}^2 \left[1 + a_m + b_m\right],
\]

where \(c_{ij} = \cos \theta_{ij}\) and \(s_{ij} = \sin \theta_{ij}\). The functions \(a_m\) and \(b_m\) are dependent on the internal structure of the Sun and are given by the expressions:

\[
a_m = C_1 (s_{mm} c_{14} - c_{mm} s_{14} s_{24})^2,
\]

and

\[
b_m = C_2 (c_{mm} c_{14} + s_{mm} s_{14} s_{24})^2,
\]

where \(c_{mm} = \cos \theta_m, s_{mm} = \sin \theta_m, C_1 = c_{13}^4 (c_{14} s_{12} - c_{12} s_{14} s_{24})^2, \) and \(C_2 = s_{13}^2 (c_{12} c_{14} + s_{12} s_{14} s_{24})^2\). The angle \(\theta_m\) is obtained for the present-day Sun (i.e., the standard solar model, see details of this model in Lopes & Silk 2013) using the expression (Capozi et al. 2017): \(\cos(2\theta_m) = M_4 (M_4^2 + M_2^2)^{-1/2}\), where \(M_4 = \cos(2\theta_{12}) - \eta_V V_x \) and \(M_4 \equiv \left|\sin(2\theta_{12}) + \eta_V V_e \right|\). \(\eta_V\) is the ratio of the energy of the neutrino \(E\) in relation to \(\Delta m^2_{21}\) given by \(\eta_V(E) = 4E/\Delta m^2_{21}\). The functions \(V_e\) and \(V_x\) are given by

\[
V_e = \frac{1}{2} [V_{cc} c_{13}^2 (c_{14}^2 - s_{14}^2 s_{24}^2) + V_s (s_{14}^2 - c_{14}^2 s_{24}^2)], \]

and

\[
V_x = (V_e - V_{cc} c_{13}^2) c_{14} s_{14} s_{24}, \]
and

$$V' = V_{\nu,\phi} - V_{\nu, c}.$$  \hspace{1cm} (14)

3.3. Light Dark Matter Impact on Solar Neutrinos

The survival probability of electron neutrinos (Equation (9)) is a time-dependent function through Equations (6), (7), and (3). Conveniently we define an effective oscillation probability \( \langle P_{ee}(E, \phi) \rangle \) that corresponds to an ensemble average of all the \( P_{ee}(E, \phi) \) (Equation (9)), as such

$$\langle P_{ee}(E) \rangle = \int_0^{\tau_0} P_{ee}(E, \phi) \frac{dt}{\tau_0}.$$  \hspace{1cm} (15)

where \( \tau_0 = 2\pi/m_\phi \) is the period of the LDM field \( \phi(t) \).

The ability of a solar neutrino detector to measure the impact of the time-dependent LDM field \( \phi(t) \) on the survival probability \( P_{ee}(E, \phi) \) (Equation (9)) depends on three characteristic timescales: the neutrino flight time \( \tau_\phi \), the time between two consecutive neutrino detections \( \tau_{ev} \), and the total run time of the experiment \( \tau_{ex} \). The neutrino flight time is proportional to the Earth-Sun distance \( d_{Earth} \), such that \( \tau_\phi = d_{Earth}/c \approx 8.2 \) minutes where \( c \) is the speed of light. The number of events measured by a detector varies strongly from one to another.

The next generation of experiments will have \( \tau_{ev} \) much larger than the pioneer Homestake experiment that only detects a few events per year (Bahcall & Davis 1976). The forthcoming Jiangmen Underground Neutrino Observatory (JUNO; Adam et al. 2015) experiment expects to measure a few tens of neutrinos per day (for instance 200 events per day or \( \tau_{ev} \approx 7 \) minutes). The total experimental run time for most solar neutrino detectors is of the order of a few decades (for instance \( \tau_{ex} \approx 10 \) yr), and future experiments will also have significant running times. Hence for all models considered in this study, we assume that solar neutrino detectors will run for long periods and will collect a large number of events, therefore we assume that \( \tau_{ev} \) and \( \tau_{ex} \) have sufficient small and large values, respectively.

In such conditions, the solar neutrino spectra time modulation by \( \phi(t) \) depends on the period \( \tau_\phi \) of the LDM field in comparison to the flight time of solar neutrinos \( \tau_\phi \). Since these neutrinos have a \( \tau_\phi \approx 8.2 \) minutes, it is possible to find the value of \( m_\phi \) for which \( \tau_\phi = \tau_\phi \) which occurs for \( m_\phi = 8.3 \times 10^{-18} \) eV. Accordingly, we can define two regimes for the time modulation of survival probability of electron neutrinos:

1. For \( \tau_\phi \approx \tau_\phi \) (low-frequency regime or low LDM mass), the time modulation of \( P_{ee}(E, \phi) \) occurs when the period of \( \phi(t) \) is larger than \( \tau_\phi \). In this case a temporal variation of the neutrino signal may be observed. This corresponds to a LDM field with a mass such that \( m_\phi \leq m_{\phi,cr} \). Therefore, the LDM field can induce an observable time variation in neutrino oscillation measurements as periodicity in the solar neutrino fluxes (Berlin 2016). Obviously, if \( \tau_\phi \) becomes very large, the modulation of \( P_{ee}(E, \phi) \) becomes indistinguishable from the standard scenario (undistorted case), since the running time of the experiment is not sufficient to observe this phenomenon. Nevertheless, in our study, the LDM field has always an \( m_\phi \geq 10^{-23} \) eV or a period \( \tau_\phi \geq 13 \) yr. Therefore, it is always possible to probe such a model with current experimental running times.

2. For \( \tau_\phi \leq \tau_\phi \) (high-frequency regime or high LDM mass), the change of \( P_{ee}(E, \phi) \) due to \( \phi(t) \) is too fast to be observed as a modulating signal like in the previous case. This regime occurs for LDM fields with a mass such that \( m_\phi \geq m_{\phi,cr} \). Nevertheless, the time average of the ensemble of oscillation probability \( P_{ee}(E, \phi) \) can be distorted in such a regime, hence the effect can be detected as \( \langle P_{ee}(E) \rangle \) which will deviate from the standard scenario (Krnjaic et al. 2018). The net effect of averaging over time induces a shift in the observed values of \( P_{ee}(E, \phi) \) relative to its undistorted value.

Therefore, we can expect to study both regimes in a quite reasonable range of LDM masses using data from the present and future solar neutrino experiments. In fact, some of the current solar neutrino collaborations have already large statistics and high event rates that we can use to look for time modulations in solar neutrinos. Some of these neutrino collaborations have already searched for regular phenomena with periods varying from 10 minutes to 10 yr (e.g., Yoo et al. 2003; Aharmim et al. 2010).

In this work, we will study models that will fall in these two regimes of time modulation. Therefore to satisfy the conditions mentioned above, we decided to analyze the impact of the LDM field in solar neutrino fluxes for \( \phi(t) \) with a period \( \tau_\phi \) varying from 4 \( \mu \)s to 13 yr or equivalently with a \( m_\phi \) varying from \( 10^{-9} \) to \( 10^{-23} \) eV, which is a range possible to be scanned by future detectors like the Deep Underground Neutrino Experiment (DUNE Collaboration et al. 2015) and JUNO (An et al. 2016).

4. Light Dark Matter Impact on Electron Neutrino Spectra

Insight into the Sun, the flux variation of neutrinos with different flavors due to matter (including LDM) is strongly dependent of the local distributions of electrons and neutrons, but also on the population of dark matter particles in the solar neighborhood. This new flavor mechanism (sterile neutrinos and LDM field \( \phi \)) affects all electron neutrinos produced in the Sun’s core. A detailed discussion about the neutrino sources inside the Sun, and their specific solar properties, can be found in Lopes (2013, 2017). The average survival probability of electron neutrinos for each nuclear reaction in the solar interior, i.e., \( P_{ee,i}(E, \phi) \) is computed by

$$P_{ee,i}(E, \phi) = C_i \int_0^{R_i} P_{ee}(E, \phi, r) S_i(r) 4\pi \rho(r) r^2 dr,$$  \hspace{1cm} (16)

where \( C_i \left( \int_0^{R_i} S_i(r) 4\pi \rho(r) r^2 dr \right)^{-1} \) is a normalization constant and \( S_i(r) \) is the electron neutrino emission function for the \( i \) solar nuclear reaction. \( i \) corresponds to the following solar neutrino sources (from the PP chain and CNO cycle nuclear reactions): \( pp \), \( pep \), \( ^8\text{B} \), \( ^7\text{Be} \), \( ^{13}\text{N} \), \( ^{15}\text{O} \), and \( ^{17}\text{F} \).

Moreover, since the survival probabilities \( P_{ee,i}(E, \phi) \) (Equation (16)) are time dependent through \( \phi \), these quantities also vary with time. Therefore, the oscillation probability \( P_{ee}(E) \) (Equation (15)) is generalized for each specific nuclear reaction \( i \):

$$\langle P_{ee,i}(E) \rangle = \int_0^{\tau_0} P_{ee,i}(E, \phi) \frac{dt}{\tau_0}.$$  \hspace{1cm} (17)

The LDM field \( \phi \) can lead to different temporal imprints on the neutrino oscillation measurements. The specific impact depends on the mass of the LDM particle. In the following, we compute the spectra of neutrinos from any specific nuclear
reaction that we know to be essentially independent of the properties surrounding solar plasma. Since in the 3+1 neutrino flavor model new processes exist to change the survival probability of electron neutrinos, this will modify the solar neutrino spectra measured on Earth. These new processes will alter the conversion rates of $\nu_e$ to other flavors ($\nu_{\mu}$, $\nu_{\tau}$, and $\nu_x$) and vice versa. Accordingly, the electron neutrino spectrum of the nuclear reaction $i$ inside the core is defined as $\Phi_i$ and $\Phi_{\odot}$ is the electron neutrino spectrum arriving on Earth (Lopes 2018b) such that:

$$\Phi_{\odot}(E) = P_{\nu_e,i}(E, \phi) \Phi_i(E),$$

(18)

where $P_{\nu_e,i}(E)$ is the average survival probability of electron neutrinos for reactions in the solar interior as given by Equation (16). Equally if we take the time average of Equation (18), we obtain the following averaged spectrum for each nuclear reaction $i$:

$$\Phi_{\odot}(E) = \langle P_{\nu_e,i}(E) \rangle \Phi_i(E),$$

(19)

where $\langle P_{\nu_e,i}(E) \rangle$ is the average survival probability of electron neutrinos as given by Equation (17).

5. The Sun: Light Dark Matter and Sterile Neutrinos

Here, we will study the impact of the theoretical model presented in the previous sections, specifically we compute the survival probability of electron neutrinos (as given by Equations (9), (15), (16), and (17)) in the case of a standard solar model with low-Z (e.g., Lopes & Silk 2013; Capelo & Lopes 2020).

In the parameterization for the 3+1 neutrino flavor oscillation model, we opt to adopt the recent values obtained in the data analysis of the standard three-neutrino flavor oscillation model obtained by de Salas et al. (2020), and for the sterile neutrino additional fiducial parameters we used the values obtained by Gariazzo et al. (2015). Accordingly, for a parameterization with a normal ordering of neutrino masses, the mass-square difference and the mixing angles have the following values: $\Delta m_{31}^2 = 7.50 \pm 0.22 \times 10^{-5}$ eV$^2$, $\sin^2 \theta_{12} = 0.318 \pm 0.016$, and $\sin^2 \theta_{13} = 0.0225 \pm 0.0075$. Although $\Delta m_{31}^2 = 2.56 \pm 0.003 \times 10^{-3}$ eV$^2$ and $\sin^2 \theta_{23} = 0.50 \pm 0.025$, we mention them here for reference (de Salas et al. 2020). These new parameters are consistent with previous estimations (Esteban et al. 2019, Gonzalez-Garcia et al. 2016). For the sterile neutrino, we choose the following fiducial values for the mass-square difference and mixing angles (Gariazzo et al. 2015, 2016; Capozzi et al. 2017): $\Delta m_{31}^2 = 1.6$ eV$^2$, $\sin^2 \theta_{14} = 0.027$, $\sin^2 \theta_{24} = 0.014$ and the other mixing angle for the sterile neutrinos are fixed to zero. Moreover, we assume that all phases $\{\theta_{13,14,14a}\}$ and other angles related to the sterile neutrinos are equal to zero.

The present-day internal structure of the Sun corresponds to an up-to-date standard solar model (SSM) that has a better agreement with neutrino fluxes and helioseismic data sets. This solar model was obtained from a one-dimensional stellar evolution code allowed to evolve in time until the present-day solar age, 4.57 Gyr, having been calibrated to the values of luminosity and effective temperature of the present Sun, of $3.8418 \times 10^{33}$ erg s$^{-1}$ and 5777 K, respectively, as well as the observed abundance ratio at the Sun’s surface: $Z_i/X_i \odot = 0.0181$, where $Z_i$ and $X_i$ are the metal and hydrogen abundances at the surface of the star (Turck-Chieze & Lopes 1993; Bahcall et al. 1995, 2006). This stellar model was computed with the release version 12115 of the stellar evolution code MESA (Paxton et al. 2011, 2019). The details about the physics of this standard solar model in which we use the AGSS09 (low-Z) solar abundances (Asplund et al. 2009) are described in Lopes & Silk (2013) and Capelo & Lopes (2020).

Figures 1 and 2 show the impact of the time-dependent mass-square difference (Equation (6)) and mixing angles (Equation (7)) on the averaged electron survival probability $\langle P_{\nu_e,i}(E) \rangle$ (Equation (15)) for which the LDM field $\phi$ has a fixed amplitude (Equation (5)): $\epsilon_\phi = 0$ or $\epsilon_\phi = 1.5\%$. We also show LDM models for which the sterile neutrino couples to $\phi$ with strength $G_{\phi}$.

The overall shape of the curve $\langle P_{\nu_e,i}(E) \rangle$ depends on $G_{\phi}$ times $n_\phi$ in the potential $V_{\nu_\phi}$ or the ratio $G_{\phi}/m_\phi$ as previously mentioned. For instance, in a LDM model in which we fix $m_\phi = 10^{-9}$ eV (or $n_\phi$), an increase of $G_{\phi}$ from $10^{20}$ to $10^{22}$ G$_F$ leads $\langle P_{\nu_e,i}(E) \rangle$ to vary significantly, as shown in Figure 1. As expected this change in $\langle P_{\nu_e,i}(E) \rangle$ is more pronounced for high energy neutrinos where the MSW effect is more significant. If we choose higher values of $m_\phi$ the results will somehow be similar (see Figure 1).

Evidently, for an LDM model in which $m_\phi$ decreases by a certain amount ($n_\phi = \rho_\phi/m_\phi$), the constancy of $G_{\phi}/m_\phi$ in $V_{\nu_\phi}$ implies that $G_{\phi}$ can increase by the same order of magnitude to obtain the same MSW effect on the $\langle P_{\nu_e,i}(E) \rangle$ curve (see Figure 1). For instance, an LDM model with $m_\phi = 10^{-23}$ eV and $G_{\phi} = 10^7 G_F$ will have $\langle P_{\nu_e,i}(E) \rangle$ identical to an LDM model with $m_\phi = 10^{-9}$ eV and $G_{\phi} = 10^{21} G_F$, since in both LDM models we have the same ratio: $G_{\phi}/m_\phi \approx 10^{30}$. The same argument explains the reason why the coupling constant between sterile neutrinos and more massive dark matter particles is much smaller in those models than in the present study. For instance, this is the case for particles captured from the dark matter halo by the Sun. Since over time, the star accreted a significant amount of dark matter (e.g.,
Lopes 2018a), for these models $G_\phi$ is significantly smaller than the value found for the present study.

The most important feature of such a class of LDM models is the time dependence of the dark matter field $\phi(t)$ and its imprint in the flavor oscillation parameters’ mass-square differences (Equation (6)) and mixing angles (Equation (7)). As predicted by Equation (9), there are many $P_{ee}(E, \phi)$ with near similar behavior. Figure 1 shows an ensemble of time-dependent $P_{ee}(E, \phi)$ as a pink band. The difference between curves relates to the dependence of the oscillation parameters on time. In this LDM model it is assumed there is a negligible interaction between sterile neutrinos and $\phi$ (for which $G_\phi \approx 0$). The figure also shows $\langle P_{ee}(E) \rangle$ (red curve) the time-averaged ensemble of $P_{ee}(E, \phi)$ curves that we compute using Equation (15).

Although there are several parameters that contribute to the time variability of $P_{ee}(E, \phi)$ (Equation (9)), the main contributions come from $\theta_{23}$ and $\theta_{13}$. The variability related $\theta_{23}$ is relevant for the high energies. We notice that the contributions coming from $\Delta m_{21}^2$ and $\theta_{32}$ are much smaller than all the parameters mentioned above. The amplitude of the $P_{ee}(E, \phi)$ pink band is defined by the value of $\epsilon_\phi$ for which we adopt the fiducial value of $\epsilon_\phi = 1.5\%$. It is worth pointing out that the $P_{ee}(E, \phi)$ band is much larger for low-energy than for higher-energy values. Moreover, the averaged value of this ensemble given by $\langle P_{ee}(E) \rangle$ is identical to $\langle P_{ee} \rangle$ with $\epsilon_\phi \approx 0$. Figure 2 shows the variability of $P_{ee}(E, \phi)$ for LDM with different $G_\phi$ values. These results are identical to the model in Figure 1. Nevertheless, the LDM model with the largest $G_\phi$ has an (orange) band with a smaller amplitude around $\langle P_{ee}(E) \rangle$. Once again, the band thickness decreases for neutrinos with higher energy for all these models.

Figures 3 and 4 compare our predictions with current solar neutrino data (e.g., Aalbers et al. 2020). These figures show that LDM models with relatively low values of $\epsilon_\phi$ and $G_\phi$ are compatible overall with current solar neutrino data coming from Borexino, Super-Kamiokande, and SNO. Clearly, this analysis has also shown that the precision of our current solar neutrino experiments is not able to distinguish between some of these LDM models. Nevertheless, it is already possible to put
some constraints on these LDM models. For instance, we found that LDM models with \( G_\phi \approx 0 \) must have a \( \epsilon_\phi \) smaller than 3\% to be consistent with all data, including pp measurements of the Borexino detector (Borexino Collaboration et al. 2020) (see Figure 3); and any LDM models must have a ratio \( G_\phi/m_\phi \) smaller than \( 10^{30} \), otherwise they become inconsistent with pp and \( ^7\text{Be} \) measurements for several solar detectors (Bellini et al. 2010; Borexino Collaboration et al. 2020; Borexino Collaboration et al. 2018; Agostini et al. 2019) (see Figure 4). Figure 4 shows a LDM model with a \( m_\phi = 10^{-9} \) eV and \( G_\phi = 10^{-22} G_F \) with a ratio \( G_\phi/m_\phi \) of the order of \( 10^{31} G_F \). This ratio is one order of magnitude larger than the critical \( G_\phi/m_\phi \) value of \( 10^{30} \) discussed in the previous section. Figure 4 also shows the variability related with time dependence on \( P_{ee}(E, \phi) \) decreases for large values of \( G_\phi \).

There is another important effect that also contributes to the time variability of \( P_{ee}(E, \phi) \). The PP chain and CNO cycle nuclear reactions occur at different distances from the center of the Sun and each nuclear reaction emits neutrinos in a well-defined energy range. As a consequence, the electron neutrinos produced in each specific nuclear reaction will be affected differently by the MSW effect. As such, this effect will also contribute to the overall variability of electron neutrinos \( P_{ee}(E, \phi) \) (see Equation (16)) and their time-averaged \( \langle P_{ee}(E, \phi) \rangle \) (see Equation (17)).

The time-dependent electron neutrino survival probability will have a significant impact on the neutrino spectra of the different nuclear reactions. Accordingly, Figures 5 and 6 show the spectra correspond to two neutrino types: pp and \(^8\text{B} \) neutrinos. An essential difference between these two spectra relates to the thickness of the \( \epsilon_\phi \) band for a fixed value since thickness decreases with neutrino energy. Therefore the \( \epsilon_\phi \) band is more significant for a pp spectrum than for a \(^8\text{B} \) neutrino spectrum. This is an effect identical to the one discussed previously for the \( P_{ee}(E, \phi) \) functions. Therefore, the measurement of solar neutrino fluxes and solar neutrino spectrum in the energy range below 0.2 MeV will provide the strongest constraint for such a class of dark matter models. Figures 5 and 6 show the spectra of \(^8\text{B} \) and pp, if we assume the precision expected to be attained by the Darwin experiment (Aalbers et al. 2020). Figure 6 also shows the precision expected for the Darwin experiment.

### 6. Conclusion

This article focuses on the impact of LDM on solar neutrino fluxes, spectra, and survival probabilities of electron neutrinos, specifically a dark matter model made of two particles: a sterile neutrino and an LDM particle. In particular, we describe how the 3+1 neutrino flavor model is affected by this type of LDM particles, with an emphasis on how the LDM affects the Wolfenstein potentials. We also study how the dark matter models affect the survival probability functions of electron neutrinos related to the different nuclear reactions occurring in the solar interior, and we compute the spectra of two relevant solar neutrino sources: pp and \(^8\text{B} \) neutrino nuclear reactions.

By studying a large range of dark matter particle masses (from \( 10^{-9} \) to \( 10^{-23} \) eV) we found that depending on the mass of these LDM particles and the value of the generalized Fermi constant, the shape of electron neutrino survival probability and their spectra can vary with time. We establish that for LDM particles with low masses (low-frequency regime), the solar neutrino detectors can observe the electron neutrino survival probability changing with time. Conversely, for dark matter particles with higher masses (high-frequency regime), this impact can be determined by measuring the time-averaged electron neutrino survival probability.

It was possible to establish, using data from current solar neutrino measurements, that those models with a \( G_\phi/m_\phi \) ratio smaller than \( 10^{30} G_F \) eV\(^{-1} \) agree with current solar neutrino data from the Borexino, SNO, and Super-Kamiokande detectors. We also found that for models with a near-zero constant, the time-variability amplitude must be smaller than 3\%. Such a constraint is equivalent to the condition \( G_\phi \sqrt{2 \rho_{\text{DM}}/(m_\phi m_\nu)} (\Omega_\phi/\Omega_{\text{DM}})^{1/2} \leq 0.03 \).
Finally, we also found that the precision expected in the measurements to be made by the Darwin detector will allow us to put powerful constraints to this class of models.

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