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Embedding the texture of the neutrino mass matrix into the MaVaNs scenario

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Abstract: We have embedded the texture of the neutrino mass matrix with three families into the MaVaNs scenario. We take the power-law potential of the acceleron field and a typical texture of active neutrinos, which is derived by the D_4 symmetry and predicts the maximal mixing of the atmospheric neutrino and the vanishing U_{e3}. The effect of couplings among the dark fermion and active neutrinos are studied by putting the current cosmological data and the terrestrial neutrino experimental data. It is found that the neutrino flavor mixings evolve as well as the neutrino masses. Especially, U_{e3} develops into the non-vanishing one and $\theta_{\rm atm}$ deviates from the maximal mixing due to couplings among the dark fermion and active neutrinos.

Keywords: Neutrino Physics.

Contents

1. Introduction

One of the most challenging questions in both cosmological and particle physics is the nature of the dark energy in the Universe. At the present epoch, the energy density of the Universe is dominated by a dark energy component, whose negative pressure causes the expansion of the Universe to accelerate. In order to clarify the origin of the dark energy, one has tried to understand the connection of the dark energy with particle physics.

Recently, Fardon, Nelson and Weiner [1] proposed an idea of the mass varying neutrinos (MaVaNs), in which the neutrino couples to the dark energy. Gu, Wang and Zhang [2] also considered the coupling of the scalar, such as Quintessence, to the neutrinos. It should be also noted that the variable neutrino mass was considered at first in [3], and was discussed for neutrino clouds [4].

The renewed MaVaNs scenario [1] has tried to make a connection between neutrinos and the dark energy. In this scenario, an unknown scalar field which is called "acceleron" is introduced, and then, the neutrino mass becomes a dynamical field. The acceleron field sits at the instantaneous minimum of its potential, and the cosmic expansion only modulates this minimum through changes in the neutrino density. Therefore, the neutrino mass is given by the acceleron field and changes with the evolution of the Universe. The cosmological parameter w and the dark energy also evolve with the neutrino mass. Those evolutions depend on a model of the scalar potential strongly. Typical examples of the potential have been discussed by Peccei [5].

The MaVaNs scenario leads to interesting phenomenological results. The neutrino oscillations may be a probe of the dark energy [6, 7]. The baryogenesis [2, 8, 9], the cosmo MSW effect of neutrinos [10] and the solar neutrino [11, 12] have been studied in the context of this scenario. Cosmological discussions of the scenario are also presented [13 − 17]. The extension to the supersymmetry have been presented in ref. [18, 19].

Now one needs to construct the realistic MaVaNs scenario with three families, which is consistent with the terrestrial neutrino experimental data [20 − 24]. In order to get a realistic MaVaNs scenario, we have embedded a texture of the neutrino mass matrix with three families into the MaVaNs scenario and examined neutrino masses and flavor mixings [25]. We take a typical texture, which is derived by D_4 symmetry and predicts the maximal flavor mixing and the vanishing one in the lepton sector. The effect of couplings among the dark fermion and active neutrinos are studied by putting the current cosmological data [26] and the terrestrial neutrino experimental data [20 − 24]. It is found that neutrino flavor mixings evolve as well as neutrino masses. Especially, U_{e3} develops into the non-vanishing one and $\theta_{\rm atm}$ deviates from the maximal mixing due to the couplings among the dark fermion and active neutrinos.

In section 2, we present the formulation of the MaVaNs scenario with three families, and in section 3, we study evolutions of neutrino masses and flavor mixings. The section 4 devotes to the summary.

2. Dark energy and three active neutrino masses

In the MaVaNs scenario, one considers a dark energy sector consisting of an acceleron field, ϕ_a and a dark fermion, n. Only left-handed neutrinos are supposed to couple to the dark sector. There are two constraints on the scalar potential of the acceleron field ϕ_a. The first one comes from the observation of the Universe, which is that the present dark energy density is about $0.7\rho_c$, ρ_c being the critical density. Since the dark energy is assumed to be the sum of the energy densities of neutrinos and the scalar potential in the MaVaNs scenario,

$$\rho_{\rm dark} = \rho_\nu + V(\phi_a) \ , \tag{2.1}$$

the first constraint turns to

$$\rho_\nu^0 + V(\phi_a^0) = 0.7\rho_c \ , \tag{2.2}$$

where “0” represents a value at the present epoch, and 70% is taken for the dark energy in the Universe.

The second one comes from the fundamental assumption in this scenario, which is that $\rho_{\rm dark}$ is stationary with respect to variations in the neutrino mass. This assumption is represented by

$$\frac{\partial \rho_\nu}{\partial \sum m_{\nu i}} + \frac{\partial V(\phi_a(m_{\nu i}))}{\partial \sum m_{\nu i}} = 0 \ . \tag{2.3}$$

For our purpose it suffices to consider the neutrino mass as a function of the cosmic temperature T [5]. Since one can write down generally

$$\rho_\nu = T^4 \sum_{i=1}^{3} F(\xi_i) \ , \qquad \xi_i = \frac{m_{\nu i}}{T} \ , \tag{2.4}$$

the stationary condition eq. (2.3) turns to

$$T^4 \sum_{i=1}^{3} \frac{\partial F}{\partial \xi_i} \frac{\partial \xi_i}{\partial \phi_a} + \frac{\partial V(\phi_a)}{\partial \phi_a} = 0 \ , \tag{2.5}$$

where

$$F(\xi_i) = \frac{1}{\pi^2} \int_0^\infty \frac{dy y^2 \sqrt{y^2 + \xi_i^2}}{e^y + 1} \ . \tag{2.6}$$

We can obtain the time evolution of neutrino masses from the relation of eq. (2.5).

Since the stationary condition should be also satisfied at the present epoch, the second constraint on the scalar potential is

$$\left. \frac{\partial V(\phi_a)}{\partial \phi_a} \right|_{\phi_a = \phi_a^0} = -T^4 \sum_{i=1}^{3} \frac{\partial F(\xi_i)}{\partial \xi_i} \left. \frac{\partial \xi_i}{\partial \phi_a} \right|_{m_{\nu i}(\phi_a) = m_{\nu i}^0(\phi_a^0),\ T = T_0} \ , \tag{2.7}$$

where $T_0 = 1.69 \times 10^{-4}$eV. It is found that the gradient of the scalar potential should be negative and very small.

One can also calculate the equation of state parameter w as follows:

$$w + 1 = \frac{4 - h(T)}{3\left[1 + \frac{V(\phi_a)}{T^4 \sum F(\xi_i)}\right]} \ , \tag{2.8}$$

where

$$h(T) = \frac{\sum \xi_i \frac{\partial F(\xi_i)}{\partial \xi_i}}{\sum F(\xi_i)} \ . \tag{2.9}$$

In order to calculate the evolution of neutrino masses and w, we assume the power-law potential for the scalar potential $V(\phi_a)$ as follows:

$$V(\phi_a) = A \left(\frac{\phi_a}{\phi_a^0} \right)^k \ , \tag{2.10}$$

where A and k are fixed by the condition of eqs. (2.2) and (2.7) if the magnitude of the acceleron field at present , ϕ_a^0 is given. Then, we can calculate evolutions of neutrino masses and w.

3. Neutrinos masses and flavor mixings

Let us discuss the neutrino mass matrix. Suppose only left-handed neutrinos couple to the dark sector, which consists of an acceleron field, ϕ_a and a dark fermion, n. A lagrangian for the dark sector and the neutrino sector is given as

$$\mathcal{L} = \overline{\nu}_{L\alpha} m_D^\alpha n + \lambda \phi_a n n + \overline{\nu}_{L\alpha} M_D^{\alpha\beta} \nu_{R\beta} + \nu_{R\alpha}^T M_R^{\alpha\beta} C^{-1} \nu_{R\beta} + \text{h.c.} \ , \tag{3.1}$$

ν_L and ν_R are the left-handed and the right-handed neutrinos, respectively. Since we consider three families of active neutrinos, mass matrices m_D^α, $M_D^{\alpha\beta}$ and $M_R^{\alpha\beta}$ are 3×1,

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3×3 and 3×3 matrices, respectively. Then, the neutrino mass matrix is given as the 7×7 matrix

$$M = \begin{pmatrix} 0 & m_D & M_D \\ m_D^T & \lambda\phi_a & 0 \\ M_D^T & 0 & M_R \end{pmatrix} , \tag{3.2}$$

in the (ν_L, n, ν_R) basis. We assume the right-handed Majorana mass scale M_R to be much higher than the Dirac neutrino mass scale M_D, and $\lambda\phi_a$ to be much higher than the scale of m_D. Then, the effective neutrino mass matrix is approximately given as

$$M_\nu = M_D M_R^{-1} M_D^T + \frac{m_D m_D^T}{\lambda\phi_a} \ . \tag{3.3}$$

Furthermore, we have one sterile neutrino with the mass $\lambda\phi_a$ and heavy three right-handed Majorana neutrinos.

The first term in the right hand side of eq. (3.3) is the ordinary neutrino seesaw mass matrix, which is denoted by $\tilde{M}_\nu$, and so it depends on the flavor model of neutrinos. On the other hand, the second term originates from couplings of the left-handed neutrinos to the dark sector. The matrix m_D is parametrized as

$$m_D = D \begin{pmatrix} a \\ b \\ c \end{pmatrix} , \tag{3.4}$$

where coefficients a, b and c are introduced to indicate relative couplings to three flavors. Then, the second term is written as

$$\frac{m_D m_D^T}{\lambda\phi_a} = \frac{D^2}{\lambda\phi_a} \begin{pmatrix} a^2 & ab & ac \\ ab & b^2 & bc \\ ac & bc & c^2 \end{pmatrix} . \tag{3.5}$$

In this paper, we assume that the ordinary neutrino seesaw mass matrix $\tilde{M}_\nu$ dominates the effective neutrino mass matrix M_ν. Therefore, the contribution from the dark sector is the next-leading term. The first term depends on the model of the neutrino mass matrix. In order to find the effect of the acceleron on neutrino masses and mixings, we take the following typical texture of neutrinos in the flavor basis of the charged lepton:

$$\tilde{M}_\nu = \begin{pmatrix} X & Y & Y \\ Y & W & U \\ Y & U & W \end{pmatrix} , \tag{3.6}$$

which leads to the neutrino flavor mixing matrix:

$$\tilde{U} = \begin{pmatrix} c_{12} & s_{12} & 0 \\ -\frac{1}{\sqrt{2}} s_{12} & \frac{1}{\sqrt{2}} c_{12} & -\frac{1}{\sqrt{2}} \\ -\frac{1}{\sqrt{2}} s_{12} & \frac{1}{\sqrt{2}} c_{12} & \frac{1}{\sqrt{2}} \end{pmatrix} , \tag{3.7}$$

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with arbitrary three neutrino masses $\tilde{m}_1$, $\tilde{m}_2$ and $\tilde{m}_3$. Here, c_{12} and s_{12} denote $\cos\theta_{12}$ and $\sin\theta_{12}$, respectively. The mass matrix of eq. (3.6) was derived by the discrete symmetry D_4 and examined phenomenologically [27]. This mass matrix gives $U_{e3}=0$ and $U_{\mu 3}=-1/\sqrt{2}$ as seen in eq. (3.7). On the other hands, θ_{12} is an arbitrary mixing angle. Therefore, this neutrino mass matrix is attractive one because it guarantees the one maximal flavor mixing and one vanishing flavor mixing.

In this paper, we discuss the effect of couplings of the left-handed neutrinos to the dark fermion on the neutrino mass matrix with the flavor symmetry. Since the contribution of the dark fermion depends on couplings $a,\ b,\ c$, which denote flavor couplings, we examine typical three cases of $a=b=c=1$ (the flavor blind), $a=1, b=c=0$ (the first family coupling) and $a=b=0, c=1$ (the third family coupling). The other choices of a,b,c do not provide drastic changes of results in these three cases.

3.1 The case of $a=b=c=1$ (the flavor blind)

At first, we consider the case that couplings among the left-handed neutrinos and the dark fermion are the flavor blind. Therefore, we take $a=b=c=1$ in eq. (3.4), which leads to the effective neutrino mass matrix

$$M_\nu = \begin{pmatrix} X & Y & Y \\ Y & W & U \\ Y & U & W \end{pmatrix} + \frac{D^2}{\lambda\phi_a}\begin{pmatrix} 1 & 1 & 1 \\ 1 & 1 & 1 \\ 1 & 1 & 1 \end{pmatrix}. \tag{3.8}$$

Since the structure of the neutrino mass matrix is not changed by the dark sector as seen in eq. (3.8), the flavor mixings do not deviate from $\tilde{U}_{e3}=0$ and $\tilde{U}_{\mu 3}=-1/\sqrt{2}$. Assuming that the contribution of the dark sector in eq. (3.8) is small compared with the active neutrino mass matrix, we get the mass eigenvalues $m_i (i=1\sim 3)$ and eigenvectors u_i in the first order perturbative expansion as follows:

$$\begin{aligned} m_i &= \tilde{m}_i + M_{ii}^{\prime(1)} \ , \\ u_i &= \tilde{u}_i + u_i^{(1)} \ , \end{aligned} \tag{3.9}$$

where

$$\begin{aligned} M_{ij}^{\prime(1)} &= \frac{D^2}{\lambda\phi_a}\ (\tilde{U}_{ei}^* \quad \tilde{U}_{\mu i}^* \quad \tilde{U}_{\tau i}^*)\begin{pmatrix} 1 & 1 & 1 \\ 1 & 1 & 1 \\ 1 & 1 & 1 \end{pmatrix}\begin{pmatrix} \tilde{U}_{ej} \\ \tilde{U}_{\mu j} \\ \tilde{U}_{\tau j} \end{pmatrix} \\ &= \frac{D^2}{\lambda\phi_a}\ (\tilde{U}_{ei}^* + \tilde{U}_{\mu i}^* + \tilde{U}_{\tau i}^*)\ (\tilde{U}_{ej} + \tilde{U}_{\mu j} + \tilde{U}_{\tau j}) \ , \\ u_i^{(1)} &= C_{i1}\tilde{u}_1 + C_{i2}\tilde{u}_2 + C_{i3}\tilde{u}_3 \ , \end{aligned} \tag{3.10}$$

with

$$C_{ij} = \frac{M_{ij}^{\prime(1)}}{\tilde{m}_i - \tilde{m}_j} \ , \qquad (i \neq j) \ , \qquad\qquad C_{ii} = 0 \ , \tag{3.11}$$

and

$$\tilde{u}_1 = \begin{pmatrix} c_{12} \\ -\frac{s_{12}}{\sqrt{2}} \\ -\frac{s_{12}}{\sqrt{2}} \end{pmatrix} , \qquad \tilde{u}_2 = \begin{pmatrix} s_{12} \\ \frac{c_{12}}{\sqrt{2}} \\ \frac{c_{12}}{\sqrt{2}} \end{pmatrix} , \qquad \tilde{u}_3 = \begin{pmatrix} 0 \\ -\frac{1}{\sqrt{2}} \\ \frac{1}{\sqrt{2}} \end{pmatrix} . \tag{3.12}$$

Since we get

$$M'^{(1)} = \frac{D^2}{\lambda\phi_a} \begin{pmatrix} (c_{12}-\sqrt{2}s_{12})^2 & (c_{12}-\sqrt{2}s_{12})(\sqrt{2}c_{12}+s_{12}) & 0 \\ (c_{12}-\sqrt{2}s_{12})(\sqrt{2}c_{12}+s_{12}) & (s_{12}+\sqrt{2}c_{12})^2 & 0 \\ 0 & 0 & 0 \end{pmatrix} \tag{3.13}$$

mass eigenvalues and flavor mixings are written in the first order approximation:

$$\begin{aligned} m_1 &= \tilde{m}_1 + \frac{D^2}{\lambda\phi_a}(c_{12}-\sqrt{2}s_{12})^2 , \\ m_2 &= \tilde{m}_2 + \frac{D^2}{\lambda\phi_a}(s_{12}+\sqrt{2}c_{12})^2 , \\ m_3 &= \tilde{m}_3 , \\ U_{e2} &= s_{12} + \frac{D^2}{\lambda\phi_a}\, c_{12}\, \frac{(c_{12}-\sqrt{2}s_{12})(\sqrt{2}c_{12}+s_{12})}{\tilde{m}_2-\tilde{m}_1} , \\ U_{\mu 3} &= -\frac{1}{\sqrt{2}} , \\ U_{e3} &= 0 , \end{aligned} \tag{3.14}$$

where U_{e3} and $U_{\mu 3}$ are free from the contribution of the dark sector.

Since the parameters $\tilde{m}_i (i = 1 \sim 3)$ and s_{12} are arbitrary, these are taken to be consistent with the experimental data in the 90% CL limit [20−22, 24]:

$$\begin{aligned} 0.33 \le \tan^2\theta_{\rm sun} \le 0.49 , \quad & 7.7\times 10^{-5} \le \Delta m^2_{\rm sun} \le 8.8\times 10^{-5}\ {\rm eV}^2 , \\ 0.92 \le \sin^2 2\theta_{\rm atm} , \quad & 1.5\times 10^{-3} \le \Delta m^2_{\rm atm} \le 3.4\times 10^{-3}\ {\rm eV}^2 . \end{aligned} \tag{3.15}$$

The bound obtained by the reactor neutrinos [23] $\theta_{\rm reactor} \le 12^\circ$ is also taken in our study.

Assuming the normal hierarchy of neutrino masses, we take the following typical neutrino masses at the present epoch:

$$m_3 = 0.05\ {\rm eV} , \qquad m_2 = 0.01\ {\rm eV} , \qquad m_1 = 0.0045\ {\rm eV} . \tag{3.16}$$

These values lead to $\Delta m^2_{\rm atm} = 2.4\times 10^{-3}\ {\rm eV}^2$ and $\Delta m^2_{\rm sun} = 8.0\times 10^{-5}\ {\rm eV}^2$, which are consistent with the experimental values in eq. (3.15).[1] On the other hand, since we expect that U_{e2} does not deviate from s_{12} so much, we take a typical value $s_{12} = 0.5$, which gives the relevant value of U_{e2}. By using these values, we examine the effect of the dark sector on neutrino masses and flavor mixings. Evolutions of masses and mixings depend on the

[1] There are other choices for m_2 and m_1. For example, the set of $m_2 = 0.009$ eV and $m_1 = 0.001$ eV also gives $\Delta m^2_{\rm sun} = 8.0\times 10^{-5}$ eV. However, the numerical results in table 1 do not so change.

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Couplings	$\frac{D^2}{\lambda\phi_a}$ (eV)	$\tilde{m}_3$(eV)	$\tilde{m}_2$(eV)	$\tilde{m}_1$(eV)	k	U_{e2}	$U_{\mu 2}$	U_{e3}
$a=b=c=1$	10^{-3}	0.05	0.0070	0.0045	9.47×10^{-5}	0.59	$-1/\sqrt{2}$	0
	5×10^{-4}	0.05	0.0085	0.0045	4.73×10^{-5}	0.53	$-1/\sqrt{2}$	0
	10^{-4}	0.05	0.0097	0.0045	9.47×10^{-6}	0.51	$-1/\sqrt{2}$	0
$a=1$	10^{-3}	0.05	0.0098	0.0037	3.16×10^{-5}	0.56	$-1/\sqrt{2}$	0
$b=c=0$	5×10^{-4}	0.05	0.0099	0.0041	1.58×10^{-5}	0.53	$-1/\sqrt{2}$	0
	10^{-4}	0.05	0.0100	0.0044	3.16×10^{-6}	0.51	$-1/\sqrt{2}$	0
$a=b=0$	10^{-3}	0.0495	0.0096	0.0044	3.16×10^{-5}	0.46	-0.698	0.0006
$c=1$	5×10^{-4}	0.0498	0.0098	0.0044	1.58×10^{-5}	0.48	-0.703	0.0003
	10^{-4}	0.0499	0.0099	0.0045	3.16×10^{-6}	0.49	-0.706	0.0001

Table 1: The summary of output parameters and predictions of mixings. The input parameters are taken as $m_3 = 0.05$ eV, $m_2 = 0.01$ eV, $m_1 = 0.0045$ eV and $s_{12} = 0.5$

parameter of the dark sector $D^2/\lambda\phi_a^0$, which is unknown one. We show the numerical results for three cases of

$$\frac{D^2}{\lambda\phi_a^0} = 10^{-3}\ \text{eV}\ , \quad 5\times10^{-4}\ \text{eV}\ , \quad 10^{-4}\ \text{eV}\ , \tag{3.17}$$

because the effect of the dark sector is tiny in the case of $D^2/\lambda\phi_a^0 \leq 10^{-4}$eV. On the other hand, the speed of sound becomes imaginary in the case of $D^2/\lambda\phi_a^0 \geq 10^{-3}$eV, which will be discussed in the section 4.

For each case of eq. (3.17), we solve equations in eq. (3.14) at the present epoch, and then we can get $\tilde{m}_1$, $\tilde{m}_2$, $\tilde{m}_3$, U_{e2}, $U_{\mu 3}$ and U_{e3}. The numerical results are presented in table 1. It is noticed that m_3, $U_{\mu 3}$ and U_{e3} are free from the effect of the dark sector, while m_2 and U_{e2} have the seizable contribution from the dark sector, and m_1 has the tiny one. Therefore, Δm^2_{atm}, Δm^2_{sun} and U_{e2} are time dependent in the universe.

On the other hand, the sterile neutrino mass is predicted for each case of eq. (3.17) as $1.0D^2$keV, $2.0D^2$keV and $10D^2$keV, respectively, which depends on the unknown parameter D in the eV unit.

Since we know ϕ_a dependence of neutrino masses, we get the parameter A and k in the scalar potential $V(\phi_a)$ of eq. (2.10). The conditions of eqs. (2.2) and (2.7) turn to be

$$\begin{aligned} A &= 0.7\rho_c - \rho_\nu^0 = 2.79\times10^{-11}\text{eV}^4\ , \\ k &= T^3 \sum_{i=1}^{3} \frac{\partial F(\xi_i)}{\partial \xi_i}\bigg|_{m_{\nu i}(\phi_a)=m^0_{\nu i}(\phi^0_a),\ T=T^0} \times \frac{D^2}{A\lambda\phi_a^0}[(c_{12}-\sqrt{2}s_{12})^2 + (s_{12}+\sqrt{2}c_{12})^2] \\ &= 3n_\nu^0 \frac{D^2}{A\lambda\phi_a^0}\ , \end{aligned} \tag{3.18}$$

where $n_\nu^0 = 8.82\times10^{-13}\text{eV}^3$ is taken for the neutrino number density at the present epoch . The values of k are also summarized in table 1.

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3.2 The case of $a = 1, b = c = 0$ (the first family coupling)

Let us consider the case that flavor dependent couplings of the left-handed neutrinos to the dark fermion. The case of $a = 1,\ b = c = 0$ in eq. (3.4) leads to the effective neutrino mass matrix

$$M_{ij}^{\prime(1)} = \frac{D^2}{\lambda\phi_a}\begin{pmatrix} c_{12}^2 & c_{12}s_{12} & 0 \\ c_{12}s_{12} & s_{12}^2 & 0 \\ 0 & 0 & 0 \end{pmatrix}, \tag{3.19}$$

which gives mass eigenvalues and flavor mixings in the first order approximation:

$$\begin{aligned} m_1 &= \tilde{m}_1 + \frac{D^2}{\lambda\phi_a}c_{12}^2\ , \\ m_2 &= \tilde{m}_2 + \frac{D^2}{\lambda\phi_a}s_{12}^2\ , \\ m_3 &= \tilde{m}_3\ , \\ U_{e2} &= s_{12} + \frac{D^2}{\lambda\phi_a}\frac{c_{12}^2 s_{12}}{\tilde{m}_2 - \tilde{m}_1}\ , \\ U_{\mu 3} &= -\frac{1}{\sqrt{2}}\ , \\ U_{e3} &= 0\ . \end{aligned} \tag{3.20}$$

Taking same input parameters in the case of $a = b = c = 1$, we get numerical results as shown in table 1. In this case, m_3, $U_{\mu 3}$ and U_{e3} are free from the effect of the dark sector as well as the case of $a = b = c = 1$. On the other hand, m_1, m_2 and U_{e2} have the seizable contribution from the dark sector. Furthermore, we get

$$A = 2.79 \times 10^{-11}\text{eV}^4\ , \tag{3.21}$$

which is the same one in eq. (3.18), and k is summarized in table 1.

3.3 The case of $a = b = 0, c = 1$ (the third family coupling)

We discuss another case of couplings. Since the case of $a = b = 0,\ c = 1$ in eq. (3.4) leads to the effective neutrino mass matrix

$$M_{ij}^{\prime(1)} = \frac{1}{2}\frac{D^2}{\lambda\phi_a}\begin{pmatrix} s_{12}^2 & -c_{12}s_{12} & -s_{12} \\ -c_{12}s_{12} & c_{12}^2 & c_{12} \\ -s_{12} & c_{12} & 1 \end{pmatrix}, \tag{3.22}$$

we get mass eigenvalues and flavor mixings in the first order approximation:

$$\begin{aligned} m_1 &= \tilde{m}_1 + \frac{1}{2}\frac{D^2}{\lambda\phi_a}s_{12}^2\ , \\ m_2 &= \tilde{m}_2 + \frac{1}{2}\frac{D^2}{\lambda\phi_a}c_{12}^2\ , \\ m_3 &= \tilde{m}_3 + \frac{1}{2}\frac{D^2}{\lambda\phi_a}\ , \end{aligned} \tag{3.23}$$

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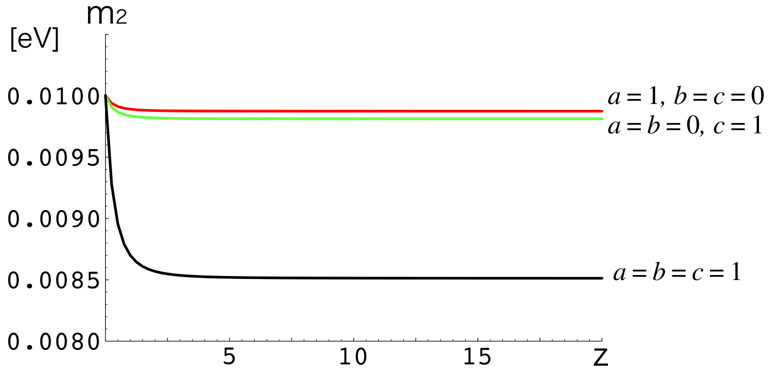

Figure 1: Plot of m_2 versus the redshift z for three cases.

$$
\begin{aligned}
U_{e2} &= s_{12} - \frac{1}{2}\frac{D^2}{\lambda\phi_a}\ \frac{c_{12}^2 s_{12}}{\tilde{m}_2 - \tilde{m}_1}\ , \\
U_{\mu 3} &= -\frac{1}{\sqrt{2}}\ + \frac{1}{2\sqrt{2}}\frac{D^2}{\lambda\phi_a}\left(\frac{c_{12}^2}{\tilde{m}_3 - \tilde{m}_2} + \frac{s_{12}^2}{\tilde{m}_3 - \tilde{m}_1}\right), \\
U_{e3} &= \frac{1}{2}\frac{D^2}{\lambda\phi_a}\left(\frac{c_{12}s_{12}}{\tilde{m}_3 - \tilde{m}_2} - \frac{c_{12}s_{12}}{\tilde{m}_3 - \tilde{m}_1}\right).
\end{aligned}
$$

Taking the same input parameters in the case of $a = b = c = 1$, we get numerical results as shown in table 1. In this case, m_3, $U_{\mu 3}$ and U_{e3} also catch the effect of the dark sector as well as m_1, m_2 and U_{e2}. It is remarked that the non-vanishing U_{e3} is obtained as well as the deviation from the maximal mixing of $U_{\mu 3}$. The parameter A and k are the same one in the subsection 3.2.

4. Evolutions of neutrino masses and w

If ϕ_a^0 is fixed, in other words, D^2/λ is given, we can calculate the evolution of neutrino masses and w by using eqs. (2.5) and (2.8). In order to see these evolutions, we take $D^2/\lambda = 1\text{eV}^2$ in eq. (3.17) in this section.

In the cases of $a = b = c = 1$ and $a = 1,\ b = c = 0$, m_3 does not evolve because the dark sector does not contribute to the third family, while it does evolve a little in the case of $a = b = 0,\ c = 1$. Therefore, we show the evolution of neutrino mass m_2 versus the redshift $z = T/T_0 - 1$ for three cases in figure 1.

As seen in figure 1, the remarkable evolution of m_2 is found in the region of $z = 0 \sim 1$ for the case of $a = b = c = 1$. Evolutions of cases of $a = 1$, $b = c = 0$ and $a = b = 0$, $c = 1$ are small compared with the one in the case of $a = b = c = 1$.

Evolutions of w versus z are shown in figure 2. These behaviors of w are similar in three cases. The value of w turns to positive around $z = 10$, which is somewhat different from the result by Peccei [5], in which the power-law potential was also taken in the one family model and then, w becomes positive near $z = 20$.

It was remarked [15] that the speed of sound, c_s becomes imaginary at the non-relativistic limit at the present and then the Universe cease to accelerate. However, our prediction of c_s^2 is positive and around 0.1 at the present epoch if the condition

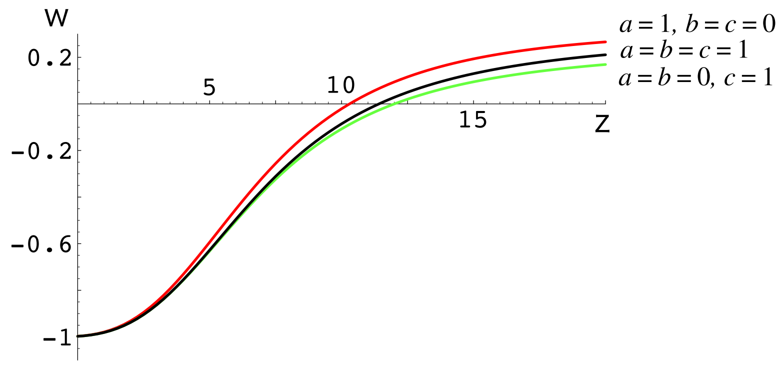

Figure 2: Plot of the equation of state parameter w versus z for three cases.

$D^2/(\lambda\phi_a) \leq 5 \times 10^{-4}$ is satisfied. It means that one needs careful study of the next-leading term at the non-relativistic limit of neutrinos.

In order to get the next-leading term at the non-relativistic limit, the function $F(\xi_i)$ in eq. (2.6) is expanded as

$$F(\xi_i) = \frac{\xi_{\nu i}}{\pi^2}\left(\int_0^\infty \frac{dy y^2}{e^y+1} + \frac{1}{2\xi_{\nu i}^2}\int_0^\infty \frac{dy y^4}{e^y+1} + \cdots\right) . \tag{4.1}$$

Then, the speed of sound is given as[2]

$$c_s^2 = w + \frac{\dot{w}}{\dot{\rho}_{\rm dark}}\,\rho_{\rm dark} \tag{4.2}$$

$$= \frac{\sum_{i=1}^3\left(\frac{\partial m_{\nu i}}{\partial z}\hat{n}_\nu\right)}{\sum_{i=1}^3 m_{\nu i}\frac{\partial \hat{n}_\nu}{\partial z}} + \frac{5}{3}f\hat{n}_\nu\frac{\left[4T_0\left(\sum_{i=1}^3\frac{1}{\xi_i}\right) - T\left(\sum_{i=1}^3\frac{1}{\xi_i^2}\frac{\partial\xi_i}{\partial z}\right)\right]}{\sum_{i=1}^3 m_{\nu i}\frac{\partial\hat{n}_\nu}{\partial z}} , \tag{4.3}$$

where

$$\hat{n}_\nu = \frac{T^3}{\pi^2}\int_0^\infty \frac{dy y^2}{e^y+1} , \qquad f = \frac{1}{2}\frac{\int_0^\infty \frac{dy y^4}{e^y+1}}{\int_0^\infty \frac{dy y^2}{e^y+1}} \simeq 6.47 . \tag{4.4}$$

In eq. (4.3), the time derivative is replaced with z derivative. The first term in the right hand side of eq. (4.3) is the leading term at the non–relativistic limit and is negative definite because of $\partial m_{\nu i}/\partial z < 0$ and $\partial n_\nu^{(0)}/\partial z > 0$. Since the numerator of the second term, which is the next-leading term, is reduced to

$$4T_0\left(\sum_{i=1}^3\frac{1}{\xi_i}\right) - T\left(\sum_{i=1}^3\frac{1}{\xi_i^2}\frac{\partial\xi_i}{\partial z}\right) = 5T_0\left(\sum_{i=1}^3\frac{1}{\xi_i}\right) - \left(\sum_{i=1}^3\frac{1}{\xi_i^2}\frac{\partial m_{\nu i}}{\partial z}\right) , \tag{4.5}$$

the second term is positive definite. The magnitude of $\partial m_{\nu i}/\partial z$ is model dependent. In our power-law potential for $V(\phi_a)$, the magnitude of the first term becomes very small.

[2]A paper including the detail derivation of c_s^2 and its physical implication will appear soon.

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Then, the second term is non-negligible one in the case of $m_i \leq \mathcal{O}(10^{-2})$eV. If we take the condition $D^2/(\lambda\phi_a) \leq 5 \times 10^{-4}$, the leading term is very small and then, c_s^2 becomes positive due to the second term. Actually, we checked numerically relative magnitude of each term in the right hand side in eq. (4.3). Thus, the positive c_s^2 is understandable in our model.

5. Summary

We have embedded the texture of the neutrino mass matrix with three families into the MaVaNs scenario. We take a typical texture of active neutrinos, which is derived by the D_4 symmetry and predicts the maximal mixing of the atmospheric neutrino and the vanishing U_{e3}. The effect of couplings among the dark sector and active neutrinos are discussed by putting the current cosmological data and the terrestrial neutrino experimental data.

It is found that flavor mixings evolve as well as neutrino masses. Especially, U_{e3} develops into the non-vanishing one and θ_{atm} deviates from the maximal mixing due to the dark sector in the case of coupling $a = b = 0,\ c = 1$. It is also remarked that the speed of sound, c_s^2 could be positive in our model.

Thus, the three family neutrino texture can be embedded into the MaVaNs scenario, where the power-law potential of the acceleron field, ϕ_a is taken. It is noticed that our results do not change so much even if we take the exponential potential of ϕ_a. The related phenomena of our flavor mixing varying scenario will be discussed elsewhere.

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