False vacuum in the supersymmetric mass varying neutrinos model

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We present detailed analyses of the vacuum structure of the scalar potential in a supersymmetric mass varying neutrinos model. The observed dark energy density is identified with false vacuum energy and the dark energy scale of order $(10^{-3}\ \mathrm{eV})^4$ is understood by the gravitationally suppressed supersymmetry breaking scale, $F(\mathrm{TeV})^2/M_{\mathrm{Pl}}$, in the model. The vacuum expectation values of sneutrinos should be tiny in order that the model works. Some decay processes of superparticles into an acceleron and sterile neutrino are also discussed in the model.

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I. INTRODUCTION

In recent years, many cosmological observations suggest that the Universe is flat and its expansion is accelerating [1]. The most common explanation for the question of the origin of such an acceleration is to assume the existence of an unknown dark energy component.

A dynamical model, which relates a dark energy scalar with neutrinos, so-called mass varying neutrinos (MaVaNs) scenario, was proposed as a candidate for the dark energy [2,3]. In the context of this scenario there are a lot of works [4–6]. Among these works, we consider a MaVaNs model in the supersymmetric theory.

The MaVaNs scenario assumes an unknown scalar field so-called "acceleron" and a sterile neutrino. They are gauge singlets under the gauge group of the standard model. In the extension of the scenario to the supersymmetry, these fields are embedded into chiral multiplets [5,6]. Typical time evolutions of the neutrino mass and the equation of state parameter of the dark energy are presented in Ref. [6], where the acceleron and a sterile neutrino are embedded into the single chiral multiplet. In this model, the observed dark energy density is identified with the false vacuum energy and the dark energy scale of order $(10^{-3}\ \mathrm{eV})^4$ is understood by the gravitationally suppressed supersymmetry breaking scale, $F(\mathrm{TeV})^2/M_{\mathrm{Pl}}$. These works are based on assumptions that the imaginary part of the acceleron and vacuum expectation values of sneutrinos vanish.

In this paper, we present a modified supersymmetric MaVaNs model of Ref. [6] and detailed analyses of the vacuum structure of the scalar potential taking account of the finite imaginary part of the acceleron. Since the vanishing imaginary part of the acceleron should be dynamically realized, discussions based on the assumption in Ref. [6] are inadequate to analyze the vacuum structure of the potential. Actually, we find the global minimum is different from one in the case of the vanishing imaginary part of the acceleron. The dark energy is identified with the false vacuum energy and the dark energy scale is given by $\rho_{\mathrm{DE}}^{1/4} \sim F(\mathrm{TeV})^2/M_{\mathrm{Pl}}$ in Ref. [6]. We show that the identification and the observed dark energy scale are also realized in the modified model where the acceleron is appropriately stabilized. Furthermore, we present cosmological discussions for the magnitude of the Yukawa coupling among the acceleron and sterile neutrinos which is assumed to be of order one in Ref. [6]. In this framework, the model with nonvanishing vacuum expectation values of sneutrinos are also discussed. We find that the magnitude of VEVs of sneutrinos are strictly constrained in the supersymmetric MaVaNs model. Then, two body decay processes of the superparticles are also studied.

In Sec. II, we present a supersymmetric MaVaNs model and detailed analyses of the vacuum structure of the scalar potential with the finite imaginary part of the acceleron and vanishing VEVs of sneutrinos. In Sec. III, we discuss the case of nonvanishing VEVs of sneutrinos in the model and show constraints on the magnitude of VEVs. In Sec. IV, two body decays of the superparticles into the acceleron and sterile neutrino are investigated. Section V is devoted to the summary.

II. VACUUM STRUCTURE OF THE SCALAR POTENTIAL

A. Scalar potential for the acceleron

Let us discuss vacuum structure in a supersymmetric mass varying neutrinos model. We consider the following superpotential W in the MaVaNs model,

$$W = m_D LA + M_D LR + \frac{\lambda_1}{6}A^3 + \frac{\lambda_2}{2}AAR + \frac{M_A}{2}AA + \frac{M_R}{2}RR, \quad (1)$$

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where L and R are chiral superfields corresponding to the left-handed lepton doublet and the right-handed neutrino, respectively [7]. The A is a gauge singlet chiral superfield. The scalar and spinor components of A are represented by ϕ and ψ, respectively. The scalar component corresponds to the acceleron which causes the present cosmic acceleration. The spinor component is a sterile neutrino. Coupling constants and mass parameters are represented by $\lambda_i (i = 1, 2)$ and M_A, M_D, M_R, m_D, respectively. The first and second terms of the right-hand side in Eq. (1) are derived from Yukawa interactions such as $y_1 HLA$ and $y_2 HLR$ with $y_1\langle H\rangle = m_D$ and $y_2\langle H\rangle = M_D$, respectively, where H is the usual Higgs doublet. In general, other interactions among gauge singlet fields such as $\lambda_1' R^3$ and $\lambda_2' AR^2$ are allowed but contributions from these interactions to the acceleron potential and neutrino masses are negligibly small. The terms appearing in Eq. (1) are important to realize the MaVaNs scenario. In this model, there is no assignment to conserve the R-parity due to the emergence of a trilinear coupling. In the limit of $y_1 \to 0$, a dark energy sector composed of A is decoupled from the minimal supersymmetric standard model (MSSM) (visible) sector and then R-parity conservation is restored as in the MSSM. The lightest superparticle (LSP) in the MSSM sector is absolutely stable in the case of conserved R-parity, however some decays of sparticles into the acceleron and a sterile neutrino are generally allowed in the model given by Eq. (1). We will return to this point later.

Taking supersymmetry breaking effects into account, the scalar potential for the acceleron is given by

$$\begin{aligned} V(\phi) = {} & \frac{\lambda_1^2+\lambda_2^2}{4}|\phi|^4 + (M_A^2 + m_D^2 - m^2)|\phi|^2 + \left(\frac{\lambda_1}{2}M_A|\phi|^2\phi - \frac{\kappa_1}{3}\phi^3 + \text{H.c.}\right) + \frac{\lambda_1}{2}m_D\tilde{\nu}_L\phi^{\dagger 2} + m_D M_A\tilde{\nu}_L\phi^\dagger \\ & + m_D M_D\tilde{\nu}_R\phi^\dagger + \frac{\lambda_2}{2}m_D\phi^\dagger\tilde{\nu}_L\tilde{\nu}_R^\dagger + \frac{\lambda_1\lambda_2}{4}|\phi|^2\phi\tilde{\nu}_R^\dagger + \frac{\lambda_2}{2}M_A|\phi|^2\tilde{\nu}_R + \frac{\lambda_2}{2}M_D\tilde{\nu}_L\phi^{\dagger 2} + \frac{\lambda_2}{2}M_R\tilde{\nu}_R\phi^{\dagger 2} \\ & - \kappa_2 H\tilde{\nu}_L\phi - \kappa_3 H\tilde{\nu}_R\phi + \text{H.c.} + V_0, \end{aligned} \tag{2}$$

where m and $\kappa_i (i = 1 \sim 3)$ are supersymmetry breaking parameters and V_0 is a constant determined by imposing the condition that the cosmological constant is vanishing at the true minimum of the acceleron potential.

In the mass varying neutrinos scenario, the dark energy is supposed to be composed of the neutrinos and the scalar potential for the acceleron:

$$\rho_{\text{DE}} = \rho_\nu + V(\phi), \tag{3}$$

where ρ_{DE} and ρ_ν correspond to the energy densities of the dark energy and the neutrinos, respectively. For convenience later, we rewrite the scalar potential (2) in terms of two real scalar fields ϕ_i $(i = 1, 2)$ instead of ϕ,

$$\begin{aligned} V(\phi_i) = {} & \frac{\lambda_1^2+\lambda_2^2}{4}(\phi_1^2+\phi_2^2)^2 + \left(\frac{\lambda_1}{2}M_A - \frac{2}{3}\kappa_1\right)\phi_1^3 + \left(\frac{\lambda_1}{2}M_A + 2\kappa_1\right)\phi_1\phi_2^2 + \frac{m_\phi^2}{2}(\phi_1^2+\phi_2^2) \\ & + (2M_A\tilde{\nu}_L + 2M_D\tilde{\nu}_R + \lambda_2\tilde{\nu}_L\tilde{\nu}_R)m_D\phi_1 + \left(\frac{\lambda_1\lambda_2}{2}\phi_1 + \lambda_2 M_A\right)\tilde{\nu}_R(\phi_1^2+\phi_2^2) \\ & + \{\tilde{\nu}_L(\lambda_1 m_D + \lambda_2 M_D) + \lambda_2 M_R\tilde{\nu}_R\}(\phi_1^2-\phi_2^2) - 2(\kappa_2\tilde{\nu}_L + \kappa_3\tilde{\nu}_R)H\phi_1 + V_0, \end{aligned} \tag{4}$$

where we define as $\phi \equiv \phi_1 + i\phi_2$ and $m_\phi^2 \equiv 2(M_A^2 + m_D^2 - m^2)$.

First we discuss the vacuum structure of this scalar potential in the case of the vanishing vacuum expectation values of sneutrinos,

$$\begin{aligned} V(\phi_i) = {} & \frac{\lambda_1^2+\lambda_2^2}{4}(\phi_1^2+\phi_2^2)^2 + \left(\frac{\lambda_1}{2}M_A - \frac{2}{3}\kappa_1\right)\phi_1^3 \\ & + \left(\frac{\lambda_1}{2}M_A + 2\kappa_1\right)\phi_1\phi_2^2 + \frac{m_\phi^2}{2}(\phi_1^2+\phi_2^2) + V_0. \end{aligned} \tag{5}$$

If $(\lambda_1 M_A/2 - 2\kappa_1/3) < 0$, the scalar potential has a local minimum at the origin in field space. Moreover, when parameters satisfy a relation,

$$m_\phi < -\sqrt{\frac{2}{\lambda_1^2+\lambda_2^2}}\left(\frac{\lambda_1}{2}M_A - \frac{2}{3}\kappa_1\right), \tag{6}$$

the potential has a second local minimum at $\phi_1 > 0$ and $\phi_2 = 0$, V_2, which is lower than the local minimum at the origin, V_1. Then, V_2 is given as follows:

$$V_2 \simeq -\frac{27}{4(\lambda_1^2+\lambda_2^2)^3}\left(\frac{\lambda_1}{2}M_A - \frac{2}{3}\kappa_1\right)^4. \tag{7}$$

In a region of

$$-\frac{(\lambda_1 M_A + 4\kappa_1) + M}{\lambda_1^2+\lambda_2^2} < \phi_1 < -\frac{(\lambda_1 M_A + 4\kappa_1) - M}{\lambda_1^2+\lambda_2^2}, \tag{8}$$

ϕ_2 has nonvanishing VEVs,

$$\langle\phi_2\rangle = \pm\sqrt{-\frac{(\lambda_1^2+\lambda_2^2)\phi_1^2+(\lambda_1 M_A+4\kappa)\phi_1+m_\phi^2}{\lambda_1^2+\lambda_2^2}}, \tag{9}$$

where $M \equiv \sqrt{(\lambda_1 M_A + 4\kappa)^2 - 4(\lambda_1^2+\lambda_2^2)m_\phi^2}$. The potential energy of the local minimum in this region, V_3, is given by

$$V_3 \simeq -\frac{(\lambda_1 M_A + 4\kappa_1)^6}{3\cdot 32^2\kappa_1^2(\lambda_1^2+\lambda_2^2)}. \tag{10}$$

Since V_3 is smaller than V_2 in the case of $\lambda_1 M_A \ll \kappa_1$, V_3 is the value of global minimum of this scalar potential. Therefore, we find

$$V_0 = |V_3| \simeq \frac{(\lambda_1 M_A + 4\kappa_1)^6}{3\cdot 32^2\kappa_1^2(\lambda_1^2+\lambda_2^2)}. \tag{11}$$

In summary, there are four local minima in the scalar potential (5). They are shown in Fig. 1. The first one is the origin in field space and the second one is situated at $\phi_1 > 0$ and $\phi_2 = 0$. The remaining two local minima are degenerated and global ones of the scalar potential,

$$V_3 \simeq V(\phi_2 = \pm\langle\phi_2\rangle_g), \tag{12}$$

where

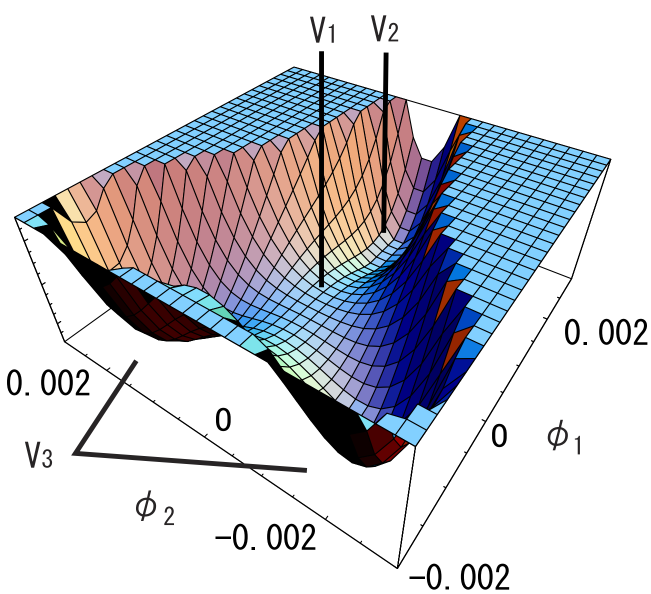

FIG. 1 (color online). The acceleron potential given in Eq. (5). There are four local minima, the first one is the origin in field space, V_1, the second one is situated at $\phi_1 > 0$ and $\phi_2 = 0$, V_2, and the remaining two local minima are degenerated and global ones, V_3.

$$\langle\phi_2\rangle_g = \sqrt{\frac{-(\lambda_1^2+\lambda_2^2)\phi_1^2-(\lambda_1 M_A+4\kappa)\phi_1-m_\phi^2}{\lambda_1^2+\lambda_2^2}}. \tag{13}$$

We find that a constant V_0 in Eq. (2), which cancels the cosmological constant, is tuned such as $V_0 = |V_3|$.

B. Numerical analysis

Now, let us consider a possible mass varying neutrinos model with this scalar potential. Since the dark energy in the MaVaNs scenario is the sum of the scalar potential and the energy density of neutrinos, the stationary condition is given by

$$\frac{\partial\rho_{\rm DE}}{\partial\phi_i} = \frac{\partial\rho_\nu}{\partial\phi_i} + \frac{\partial V(\phi_i)}{\partial\phi_i} = 0. \tag{14}$$

We find the local minimum at the origin is slightly deviated due to the contribution of the neutrino energy density. And a numerical calculation shows that the model works near the origin in field space: $\langle\phi_1\rangle^0 \simeq 0$ eV and $\langle\phi_2\rangle^0 = 0$ eV, where superscript zero means the value at the present epoch. Since the ϕ field is trapped in the local minimum near the origin, one observes the potential energy $|V_3|$ as the dark energy at the present epoch. The ϕ field trapped near the origin is plausible because it is likely that the Universe starts at the origin due to the thermal effects. Typical values of parameters are given as

$$\begin{aligned} \kappa_1 &\simeq 2.17\times 10^{-3}\ \text{eV},\\ M_A &\simeq 1.003\times 10^{-2}\ \text{eV},\\ |M_D^2/M_R| &\simeq 4.998\times 10^{-2}\ \text{eV},\\ m &\simeq 1.008\times 10^{-2}\ \text{eV}, \end{aligned} \tag{15}$$

by numerical calculation when we take

$$\begin{aligned} m_{\nu_L}^0 &= 5\times 10^{-2}\ \text{eV}, & m_\psi^0 &= 10^{-2}\ \text{eV},\\ m_D &= 10^{-3}\ \text{eV}, & m_\phi &= 10^{-4}\ \text{eV},\\ \lambda_1 &= 10^{-9}, & \lambda_2 &= 1, \end{aligned} \tag{16}$$

as inputs masses and couplings. Now, we assume that supersymmetry is broken at the TeV scale in a hidden sector and its effects are transmitted to a dark energy sector, which includes a chiral superfield A, only through gravitational interaction. Then, the scale of soft terms of order (TeV) $2/M_{\rm pl} \sim \mathcal{O}(10^{-2}\text{–}10^{-3}\ \text{eV})$ is expected. Soft parameters κ_2 and κ_3 are three-scalar couplings such as $\kappa_2\tilde{\nu}_L H\phi$ and $\kappa_3\tilde{\nu}_R H\phi$. If the dark energy sector is communicated with the visible sector, which includes the right-handed sneutrino, only through the gravitational interaction except for a peculiar Yukawa interaction of the MaVaNs scenario such as $y_1 H\nu_L\psi$, $\kappa_{2,3}$ are roughly estimated of order $m_{\rm soft}^2/M_{\rm pl}$ by dimensional analysis, where $m_{\rm soft}$ means the scale of soft terms in the visible sector. The magnitude of soft parameters in the visible sector depends on a mediation mechanism between the visible and a hidden sector. If we take $m_{\rm soft} \sim \mathcal{O}$ (TeV), the magnitude of $\kappa_{2,3}$ becomes $\mathcal{O}(10^{-3}$ eV). A larger scale of $\kappa_{2,3}$, which corresponds to $m_{\rm soft}$, prefers tiny or vanishing VEVs of sneutrinos in the MaVaNs model. We will return to this point later.

The parameter M_A affects the neutrino masses. The effective neutrino mass matrix is given by

$$\mathcal{M} \simeq \begin{pmatrix} c + \epsilon_1 & m_D + \epsilon_2 \\ m_D + \epsilon_2 & M_A + \lambda_1\langle\phi_1\rangle + \epsilon_3 \end{pmatrix}, \tag{17}$$

in the basis of (ν_L, ψ), where ν_L and ψ are the left-handed and a sterile neutrino, respectively. The c in the (1, 1) element of this matrix corresponds to the usual term given by the seesaw mechanism in the absence of the acceleron, $c \equiv -M_D^2/M_R$, where $\lambda_1\langle\phi_1\rangle \ll M_D \ll M_R$ is assumed. The parameters $\epsilon_i (i = 1 \sim 3)$ are functions of ϕ_1, which are suppressed by large scales such as the right-handed neutrino mass scale M_R or the Higgsino mass $m_{\tilde{h}}$. If the VEV of ϕ_1 and the mixing between the active and a sterile neutrino are small, the mass of a sterile neutrino is almost determined by M_A. We assume that the magnitude of a sterile neutrino mass is the same order as the left-handed neutrino one, $m_\psi^0 = 10^{-2}$ eV, and notice $m_\psi^0 \simeq M_A \simeq 1.003 \times 10^{-2}$ eV as shown in Eqs. (15) and (16).

The parameter m_D determines the mixing between the active and sterile neutrinos. When we take $m_D = 10^{-3}$ eV, the mixing angle, $\sin^2\theta \simeq 4 \times 10^{-4}$, is expected. The active and sterile neutrinos mix maximally in the case of 10^{-3} eV $\ll m_D$, which is ruled out by the current astrophysical, cosmological, and laboratory bounds [8]. Since m_D is defined as $y_1\langle H\rangle$ in our model, $m_D \lesssim 10^{-3}$ eV and $\langle H\rangle \simeq 10^2$ GeV require $y_1 \lesssim 10^{-14}$. In general, MaVaNs models need such a tiny Yukawa coupling among the left-handed neutrino, the Higgs boson, and the sterile neutrino. In the limit of $y_1 \to 0$, a dark energy sector is decoupled from the visible sector and the mixing between the left-handed and sterile neutrinos vanishes. Such a small Yukawa coupling may be explained in a brane model, however the discussion is beyond the scope of this paper. We use the value $m_D = 10^{-3}$ eV, which corresponds to $y_1 \simeq 10^{-14}$, in numerical analyses.

In the MaVaNs scenario, the mass of the scalar field should be less than $\mathcal{O}(10^{-4}$ eV) in order that the scalar field does not vary significantly on the distance of interneutrino spacing, $n_\nu^{-1/3}$. Therefore, $m_\phi = 10^{-4}$ eV is taken in our calculation.

Next, we comment on couplings λ_i. The size of λ_1 is constrained by the successful big bang nucleosynthesis (BBN) theory and supernova data. In the early universe, the acceleron and the sterile neutrino have to decouple in order that the scenario of BBN is not radically changed. If the acceleron production rates through acceleron-strahlung and pair annihilation of neutrinos are less than the expansion rate at the era of BBN, these particles are not in thermal equilibrium at the epoch [3]. These considerations give us

$$\left(\frac{\partial m_\nu}{\partial\phi}\right)^2 \leq 1, \tag{18}$$

$$\left(\frac{\partial m_\nu}{\partial\phi}\right)^4 < 10^{-22}, \qquad \left(\frac{\partial^2 m_\nu}{\partial\phi^2}\right)^2 < 10^{-34}\ \text{eV}^{-2}. \tag{19}$$

Equations (18) and (19) show constraints that ϕ-production rates through ϕ-strahlung and neutrino pair annihilation are smaller than the expansion rate at the BBN epoch, respectively.

Similar constraints come from consideration in supernovae. We require that ϕ-emission does not cool a protoneutron star too quickly to change the observed neutrino spectra emitted in the first ten seconds. This requirement is realized in the following conditions:

$$\left(\frac{\partial m_\nu}{\partial\phi}\right)^2 \leq 10^{-12}, \tag{20}$$

$$\left(\frac{\partial m_\nu}{\partial\phi}\right)^4 < 10^{-23}, \qquad \left(\frac{\partial^2 m_\nu}{\partial\phi^2}\right)^2 < 10^{-37}\ \text{eV}^{-2}. \tag{21}$$

We find that constraints from supernova data, Eqs. (20) and (21), are more severe than the ones from the BBN scenario, Eqs. (18) and (19).

In our model, the acceleron dependences of masses of the left-handed and sterile neutrinos are given by diagonalizing the effective neutrino mass matrix (17),

$$m_{\nu_L}(\phi_1) = \frac{c + M_A + \lambda_1\langle\phi_1\rangle}{2} + \frac{\sqrt{[c - (M_A + \lambda_1\langle\phi_1\rangle)]^2 + 4m_D^2}}{2}, \tag{22}$$

$$m_\psi(\phi_1) = \frac{c + M_A + \lambda_1\langle\phi_1\rangle}{2} - \frac{\sqrt{[c - (M_A + \lambda_1\langle\phi_1\rangle)]^2 + 4m_D^2}}{2}, \tag{23}$$

where we ignore ϵ_i. Figure 2 shows the acceleron dependences of neutrino masses, $(\partial m_\nu/\partial\phi_1)^2$ and $(\partial^2 m_\nu/\partial\phi_1^2)^2$, and the constraint on λ_1 from supernovae. Since the neutrino masses depend on only ϕ_1, $(\partial/\partial\phi)$ is replaced with $(\partial/\partial\phi_1)$. The left and right figures in Fig. 2 correspond to the constraint in Eq. (20) and the second one in Eq. (21), $(\partial^2 m_\nu/\partial\phi^2)^2 < 10^{-37}$ eV^{-2}. We find that both constraints are satisfied if $\lambda_1 \lesssim 4 \times 10^{-9}$.

Next, we consider the coupling λ_2. In numerical calculations, we take the same number density of sterile neutrinos as the one of the left-handed neutrinos, $n_\nu \propto T^3$,

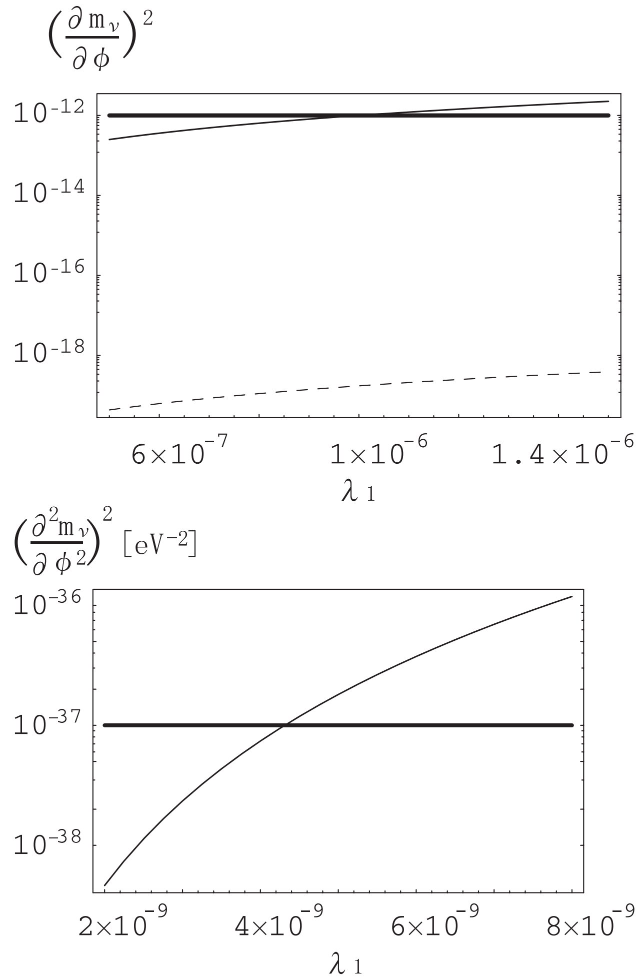

FIG. 2. The acceleron dependences of neutrino masses and the constraint on λ_1 from supernova data. The top figure shows the relation between $(\partial m_\nu/\partial\phi)^2$ and λ_1. The horizontal line shows the constraint of Eq. (20). The dashed and thin curves correspond to $(\partial m_{\nu_L}/\partial\phi_1)^2$ and $(\partial m_\psi/\partial\phi_1)^2$, respectively. The lower region than the horizontal line, $(\partial m_\nu/\partial\phi_1)^2 < 10^{-12}$, is allowed by the supernova data. In the bottom figure, the horizontal line shows the constraint of Eq. (21). The thin curve corresponds to $(\partial^2 m_{\nu_L}/\partial\phi_1^2)^2$, which is exactly equal to $(\partial^2 m_\psi/\partial\phi_1^2)^2$. The lower region than the horizontal line, $(\partial^2 m_\nu/\partial\phi^2)^2 < 10^{-37}$ eV^{-2}, is allowed by the supernova data. These figures mean that if λ_1 is smaller than about 4×10^{-9}, all constraints are satisfied.

where T is the neutrino temperature. This means that sterile neutrinos should be thermalized in the early universe. The interaction rate of a sterile neutrino through the interaction $(\lambda_2/2)ARR$ is typically given by $\Gamma \sim \lambda_2^2 T$. The thermalization of sterile neutrinos, $\Gamma > H(\sim T^2/M_{\rm Pl})$, requires $\lambda_2^2 M_{\rm Pl} > T$. For instance, if we assume that sterile neutrinos thermalized after the inflation $\lambda_2^2 M_{\rm Pl} > T_{\rm RH}$ should be realized, where $T_{\rm RH}$ is the reheating temperature. The thermalization of sterile neutrinos in this MaVaNs model with the coupling λ_2 of order one is easily realized by many inflation models.

C. False vacuum

The presented MaVaNs model works at the false vacuum. We estimate the decay probability of the false vacuum. The decay probability is given by e^{-B}, where B is the action for the classical solution in Euclidean space. We use the result of [9],

$$B = 2\pi^2 \frac{(\Delta\phi)^4}{[(\Delta V_-)^{1/3} - (\Delta V_+)^{1/3}]^3}, \tag{24}$$

where $\Delta\phi \equiv \phi_- - \phi_+$, ϕ_+, and ϕ_- are the vacuum expectation values at the false and the true vacuum, respectively, and

$$\Delta V_\pm \equiv V_T - V(\phi_\pm), \tag{25}$$

where V_T is the potential energy of the top of the barrier between the false and the true vacuum. After calculating the dark energy density numerically in our model, we find

$$\phi_+ \simeq 3.61\times10^{-6}\ \text{eV}, \qquad \phi_- \simeq 2.83\times10^{-3}\ \text{eV},$$
$$\Delta V_+ \simeq 2.3\times10^{-20}\ \text{eV}^4, \qquad \Delta V_- \simeq 2.95\times10^{-11}\ \text{eV}^4, \tag{26}$$

and thus $B \simeq 42.9 \gg 1$. Therefore, the decay probability is sufficiently small. This probability is of the decay from the false vacuum near the origin in the field space, V_1, into one of two degenerated true vacua, V_3. There is another possibility of the false vacuum decay, which is the one into another false vacuum, V_2. In this case, we get $B \simeq 41.1$, where we take $\phi_- \simeq 2.80\times10^{-3}$ eV. In this numerical calculation, we take the square potential approximation in Ref. [9]. In the limit of $\Delta V_+ \to 0$, the result of this approximation equation (24) describes the sheer drop which is the extreme case of "tunnelling without barriers" discussed in Ref. [10],

$$B = 2\pi^2 \frac{(\Delta\phi)^4}{\Delta V_-}. \tag{27}$$

Since $\Delta V_+ \ll \Delta V_-$ in our model, this approximation is appropriate rather than the "triangle" [9] or "thin-wall" approximation [11].

III. VACUUM EXPECTATION VALUES OF SNEUTRINOS

In the previous section, we consider the vanishing vacuum expectation values of sneutrinos. Let us consider effects of nonvanishing VEVs of sneutrinos. VEVs of sneutrinos interacting with the acceleron are strictly constrained in the mass varying neutrinos model. The most severe one comes from the stationary condition equation (14). We rewrite the condition equation (14) as

$$\frac{\partial V}{\partial \phi_i} = -T^3 \sum_{I=\nu_L,\psi} \frac{\partial m_I}{\partial \phi_i}\frac{\partial F(\xi_I)}{\partial \xi_I}, \tag{28}$$

where $\rho_\nu = T^4 \sum F(\xi_I)$, $\xi_I \equiv m_I/T$ and

$$F(\xi_I) \equiv \frac{1}{\pi^2}\int_0^\infty \frac{dyy^2\sqrt{y^2+\xi_I^2}}{e^y+1}. \tag{29}$$

Since the right-hand side in Eq. (28) is approximately proportional to the number density of neutrinos and both active and sterile neutrino masses derived from the mass matrix Eq. (17) do not depend on ϕ_2, the stationary condition at the present epoch is approximated by

$$\frac{\partial V}{\partial \phi_1}\Bigg|_{\phi_1=\langle\phi_1\rangle^0} \simeq -n_\nu^0 \sum_{I=\nu_L,\psi}\frac{\partial m_I}{\partial \phi_1}\Bigg|_{\phi_1=\langle\phi_1\rangle^0}. \tag{30}$$

A viable MaVaNs model should satisfy this constraint. Of course, the model with the vanishing VEVs of sneutrinos satisfies it. When we take $\lambda_1 = 10^{-9}$, we have

$$\left(\frac{\partial m_{\nu_L}}{\partial \phi_1}\right)^2 \sim \mathcal{O}(10^{-25}), \tag{31}$$

$$\left(\frac{\partial m_{\psi}}{\partial \phi_1}\right)^2 \sim \mathcal{O}(10^{-18}). \tag{32}$$

Therefore, the magnitude of the right-hand side in Eq. (30) is $\mathcal{O}(10^{-22}\ \mathrm{eV}^3)$, where we take $n_\nu^0 \sim \mathcal{O}(10^{-13}\ \mathrm{eV}^3)$. Since the term $m_\phi^2\phi_1$ is dominant in the left-hand side of Eq. (30), we have $\langle\phi_1\rangle^0 \sim \mathcal{O}(10^{-14}\ \mathrm{eV})$, where we take $m_\phi = 10^{-4}$ eV.

When the vacuum expectation values of sneutrinos are nonvanishing, the following terms are strictly constrained unless miraculous cancellation among each term is realized,

$$\begin{aligned}\frac{\partial V(\phi_1)}{\partial \phi_1} \supset\ & (2M_A\tilde{\nu}_L + 2M_D\tilde{\nu}_R + \lambda_2\tilde{\nu}_L\tilde{\nu}_R)m_D + \lambda_1\lambda_2\tilde{\nu}_R\phi_1^2 \\ & + 2\lambda_2 M_A\tilde{\nu}_R\phi_1 + 2\{\tilde{\nu}_L(\lambda_1 m_D + \lambda_2 M_D) \\ & + \lambda_2 M_R\tilde{\nu}_R\}\phi_1 - 2(\kappa_2\tilde{\nu}_L + \kappa_3\tilde{\nu}_R)H. \end{aligned} \tag{33}$$

Especially, restrictions on the VEVs of sneutrinos are obtained from $\kappa_2 H\tilde{\nu}_L$ and $\kappa_3 H\tilde{\nu}_R$. In order that these terms do not spoil the model, $\langle\tilde{\nu}_{L,R}\rangle < 10^{-39}$ eV must be realized when we take $\kappa_{2,3} = 10^{-3}$ eV and $\langle H\rangle = 10^2$ GeV. The larger values of $\kappa_{2,3}$ gives the smaller VEVs of sneutrinos. The vanishing or tiny VEVs of sneutrinos are favored in this supersymmetric mass varying neutrinos model.

IV. SPARTICLE DECAY

The mass varying neutrinos scenario assumes some light particles such as the acceleron and the sterile neutrino. In the supersymmetric realization as we discussed, these particles are embedded into a single chiral multiplet. Since Eq. (1) leads to R-parity violation in order that a given scalar potential and the neutrinos mass are consistent with the current cosmological observations, some heavy sparticle such as the sneutrino and the Higgsino decay into light particles due to such R-parity violating interactions.

We consider two body decay processes of the sneutrino and the Higgsino in this section. A peculiar decay process of the sneutrino in the MaVaNs scenario is $\tilde{\nu}_L \to \phi\phi$, whose decay rate is given by

$$\Gamma_{\tilde{\nu}_L\to\phi\phi} \simeq \frac{\lambda_1^2 m_D^2}{16\pi m_{\tilde{\nu}_L}}. \tag{34}$$

When we take $m_{\tilde{\nu}_L} = 1$ TeV, $\lambda_1 = 10^{-9}$, and $m_D = 10^{-3}$ eV, we have $1/\Gamma_{\tilde{\nu}_L\to\phi\phi} \simeq 3\times 10^{22}$ sec. A larger left-handed sneutrino mass or smaller λ_1 and m_D suppress this process. The dependence of $1/\Gamma_{\tilde{\nu}_L\to\phi\phi}$ on the left-handed sneutrino mass and λ_1 is shown in Fig. 3.

The Higgsino decay is another peculiar one in this scenario. There are two important processes, $\tilde{h}\to\nu_L\phi$ and $\tilde{\nu}_L\psi$. The decay rates of these processes are given by

$$\Gamma_{\tilde{h}\to\nu_L\phi} \simeq \frac{y_1^2 m_{\tilde{h}}}{32\pi}, \qquad \Gamma_{\tilde{h}\to\tilde{\nu}_L\psi} \simeq \frac{y_1^2(m_{\tilde{h}}^2 - m_{\tilde{\nu}_L}^2)^2}{32\pi m_{\tilde{h}}^3}. \tag{35}$$

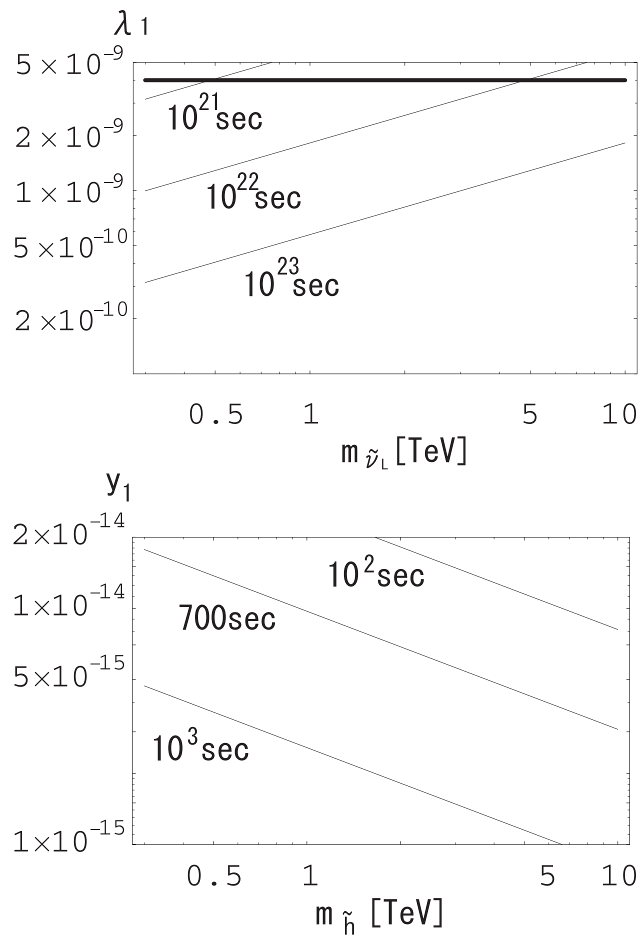

FIG. 3. The top figure shows the dependence of $1/\Gamma_{\tilde{\nu}_L\to\phi\phi}$ on the left-handed sneutrino mass and λ_1, where $m_D = 10^{-3}$ eV is taken. The region lower than the bold horizontal line, which corresponds $\lambda_1 \simeq 4\times 10^{-9}$, is allowed. The bottom figure shows $1/\Gamma_{\tilde{h}\to\tilde{\nu}_L\psi} \simeq 1/\Gamma_{\tilde{h}\to\nu_L\phi} \simeq y_1^2 m_{\tilde{h}}/32\pi$.

When we take $y_1 = 10^{-14}$ and $m_{\tilde{h}} = 1$ TeV, $1/\Gamma_{\tilde{h}\to\nu_L\phi} \simeq 700$ sec. If we assume that the sneutrino mass is less than the Higgsino mass but the same size, $1/\Gamma_{\tilde{h}\to\tilde{\nu}_L\psi}$ is the same order as $1/\Gamma_{\tilde{h}\to\nu_L\phi}$. Larger $m_{\tilde{h}}$ or y_1 enhances these processes as shown in Fig. 3. The above processes as we discussed may affect many LSP cold dark matter models. Such effects of the mass varying neutrinos scenario will be discussed elsewhere.

V. SUMMARY

We presented a supersymmetric mass varying neutrinos model and detailed analyses of the vacuum structure of the scalar potential taking account the finite imaginary part of the acceleron. The potential has four local minima and the dark energy is identified with the false vacuum energy. A metastable vacuum is realized near the origin in field space, $\phi_1 \simeq 0$ and $\phi_2 = 0$, due to the supersymmetry breaking effect. This metastable vacuum is enough stable in the age of the Universe. In the model, the observed dark energy scale of order $(10^{-3}\ \mathrm{eV})^4$ is understood by gravitationally suppressed supersymmetry breaking scale, $F(\mathrm{TeV})^2/M_{\mathrm{Pl}}$.

Following these considerations, the case with nonvanishing vacuum expectation values of sneutrinos has been discussed. We have found that the VEVs of sneutrinos are strictly constrained such as $\langle\tilde{\nu}_{L,R}\rangle < 10^{-39}$ eV in order that terms with respect to sneutrinos in the acceleron potential do not spoil the model. This means that the vanishing or tiny VEVs of sneutrinos are favored in the proposed supersymmetric MaVaNs model. Finally, two body decay processes $\tilde{\nu}_L \to \phi\phi$, $\tilde{h} \to \nu_L\phi$, and $\tilde{h} \to \tilde{\nu}_L\psi$ are discussed. The obtained decay rates are $1/\Gamma_{\tilde{\nu}_L\to\phi\phi} \simeq 3 \times 10^{22}$ sec. and $1/\Gamma_{\tilde{h}\to\nu_L\phi} \simeq 1/\Gamma_{\tilde{h}\to\tilde{\nu}_L\psi} \simeq 700$ sec. Effects on many LSP cold dark matter models from these processes is expected to provide interesting physics.

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