Properties of potential modelling three benchmarks: the cosmological constant, inflation and three generations

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I. INTRODUCTION

Recently the seesaw mechanism for the mixing of two virtual vacuum-levels due to fluctuations described by thin domain walls, has been explored in order to derive the natural scale of cosmological constant \cite{1} in terms of supersymmetry (SUSY) breaking scale $\mu_X$ and Planck mass defined by the Newton constant $G$ in gravitation, $m_{Pl} = 1/\sqrt{G}$, so that the vacuum state of de Sitter spacetime (dS) constitutes the stationary level composed by the superposition of flat and Anti-de Sitter (AdS) vacua, and it has got the constant energy density $\rho = \mu_X$ at $\mu_X \sim \mu^2_{X}/m_{Pl}$. Then, the modern value of cosmological constant $\Lambda \sim \mu^2_{X}/m_{Pl}^2$ gives $\mu_{X} \approx 0.25 \cdot 10^{-2}$ eV, hence, $\mu_{X} \sim 10^4$ GeV, i.e. the low scale of SUSY breaking down. The virtual flat vacuum corresponds to a scalar field positioned in the local minimum of its potential with zero energy provided by the exact SUSY. The virtual AdS vacuum-state is described by the field positioned at the local minimum of primary potential with positive energy, which breaks down SUSY, while the supergravity contribution linear in $G$ \cite{11} gives the term providing the negative overall sign of cosmological constant due to the dominant energy density of zero-point quantum-field modes with masses $m \sim \mu_X$. The virtual flat vacuum does not decay to the AdS one\footnote{1} due to stabilization effect of gravity \cite{13, 14}, but it suffers from fluctuations in the form of spherically symmetric AdS-bubbles surrounded by domain walls. Then, due to such the specific seesaw mechanism, one of two true stationary states is the dS-vacuum. Note, that the low energy contributions due to such phenomena like the electroweak symmetry breaking down or condensates in QCD can modify the vacuum energy of initial AdS-state with broken SUSY, only, while zero energy of flat state is preserved by exact SUSY. However, the numerical value of low energy condensates is negligible in comparison with the dominant term coming from the SUSY breaking itself due to the appropriate hierarchy of relevant scales. Similarly, any dynamics at energies higher than the scale of SUSY breaking down cannot disturb the vacuum energy, since such the dynamics is supersymmetric.

The quite general idea of incorporating the seesaw mechanism for the derivation of naturally small cosmological constant from the Planck mass and SUSY breaking scale is not originally new itself. The point of view was presented in scientific e-folklore as discussion in blogs\footnote{2}, for instance. In addition, in the framework of supergravity G. Chalmers argued for the relevant suppression of cosmological constant in \cite{13}. However, the idea becomes more actual, when it is realized in terms of reasonable model. M. McGuigan has modified the Wheeler-DeWitt equation in order to switch on a coupling between two sectors characterized by the Planck scale and SUSY, correspondingly. So, the seesaw mechanism has been involved into the gravity, and the cosmological constant of natural scale has been generated \cite{16}. The other way has been formulated in our approach invented in \cite{1}.

In \cite{17} we have constructed a model of potential, which has allowed us to investigate the scale parameters in the problem within the suggested approach. So, we have established the following general features:

\footnote{1 If the gravity is switched off, the false vacuum decay, see \cite{12}.}

\footnote{2 See the following blog sites \url{http://cosmicvariance.com/2005/12/05/duff-on-susskind/#comment-8629} (on Dec 6, 2005) and \url{http://motls.blogspot.com/2005/12/cosmological-constant-seesaw.html} (on Dec 19, 2005).}
• thin domain walls correspond to the low-scale SUSY breaking down due to the gauge mediation, when the distance between the extremal positions of scalar field takes sub-Planckian values, while

• thick domain walls are related with the gravity-mediated SUSY breaking down at high energies $\mu_{\chi} \sim 10^{12-13}$ GeV and super-Planckian values of field increment between the extremals.

This potential is suitable also for demonstrating the origin of fermion generations observed in the Standard model. Indeed, the non-trivial vacuum structure is described by a superposition of initial states with definite masses of fermions. The superposition can be represented by two-dimensional (2D) column, say. Therefore, it is natural that a 2×2 mass-matrix is involved for the fermion states, i.e. two generations appear. One can easily introduce three generations by considering the superposition of three initial vacuum-states: two flat levels and single AdS-vacuum in the model discussed. However, such the vacuum structure does not answer the question, why we are living in the vacuum we have got, since all of three stationary superpositions can be occupied, but one of them, at least, is the AdS-state with a huge negative energy density irrelevant to the astronomical observations. Moreover, it is the flat vacuum suffering from the domain-wall fluctuations. Therefore, evolving the scalar field to the position of flat vacuum will incorporate the mixing with the AdS-vacuum only, and the evolution will not see the mixing with another flat vacuum even through the AdS-states, since beyond the domain wall the boundary condition at spatial infinity remains unchanged, i.e. ascribed to the first flat vacuum. This fact does not influence the analysis as concerns for scaling properties of potential, of course. However, the potential needs a modification compatible with the scalar-field evolution during the Universe expansion in order to get the natural reason for the living in the dS-vacuum. In addition, the potential of [17] suggests a kind of fine tuning, because of its two flat vacua with the coinciding zero vacuum-energy at different values of scalar field.

The direction of modification is clear: one should get a potential with a single flat vacuum and a couple of AdS-vacua, which energy densities $\rho = -\rho^2_s$ do not coincide, in general, but $\rho^2_s$ take values of the same order of magnitude. Then, the flat vacuum will fluctuate due to bubbles of both AdS-vacua, and the Universe will get the observed cosmological constant in the dS-state, if the evolution will drive it to the field-position in the flat vacuum.

In Section II we present the model of potential satisfying all requirements listed above. The potential is composed of several contributions. The first term is the bare quadratic potential, which generates the second contribution being the quartic term due to the supergravity correction linear in Newton constant $G$. The third term modifies the potential at low energies due to the modelled contribution by zero-point modes, so that it has got the form with the flat local minimum and two AdS minima. The barrier between the flat and AdS minima is tuned in order to produce the thin domain wall with the thickness given by the inverse bare mass in the first term mentioned above. Moreover, the behavior of potential at $\phi \to 0$ is well approximated by the quartic term with the coupling constant of the same order of magnitude at large fields. Then, we can evaluate the bare mass $m_{\text{bare}} \sim 10^{12}$ GeV and quartic coupling $\lambda \sim 10^{-14}$.

The vacuum structure formed due to the mixing between the initial flat and AdS vacua because of domain wall fluctuations, is described in Section III. Then, the stationary dS-vacuum state appears due to the fluctuations in vicinity of flat vacuum.

In Section IV we study the inflation $^3$ $^{18-21}$ governed by the scalar field with the quadratic and quartic self-couplings. The field evolution corresponds to the dynamical system possessing the properties of parametric attractor: the kinetic and potential terms rapidly reach stable critical points, which have got a slow driftage with the growth of e-folding in the scale factor of accelerated expansion. Then, the quasi-attractor provides us with the tool to quantify the inhomogeneity generated by the quantum fluctuations of scalar field during the inflation. The data signalize for the dominance of quadratic term in the potential. We give a numerical constraint for the dominance and discuss the conditions of its realization in consistency with the observations. After the inflation, the scalar field enters the stage of preheating due to the tachyonic mechanism: the potential barrier generates a negative square of effective mass, that results in the scalar field decay to massless quanta in vicinity of flat vacuum. The preheating should take place at the temperature $T_{\text{preh.}} \sim 10^9$ GeV determined by the barrier height. Thus, the universe evolution drives to the dS-vacuum state.

Note that the review on the relation of inflation to the particle physics and on the mechanism of preheating and thermalization of Universe after the inflation $^4$ can be found in [23].

Section V is devoted to the analysis of textures in the mass matrices for three generations of charged fermions as caused by the vacuum structure. We show that the hierarchies in masses and mixings of charged weak currents as well as a fine violation of combined inversion of charge and space can be natural for the fermions of observable sector, while the hidden sector responsible for the SUSY breaking down can remain heavy. The picture similar

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$^3$ See modern review in [23].

$^4$ Realistic models of low-energy inflation taking into account of constraints following from the primordial nucleosynthesis (Big Bang nucleosynthesis, anisotropy of cosmic background radiation and inhomogeneity of matter density in the large scale structure of Universe, are presented in [24-26]), wherein a supersymmetric version of Standard Model for the particle interactions is studied with the use of flat directions in superpotentials.
II. THE POTENTIAL MODEL

For the sake of simplicity, we ascribe the energy density of initial vacuum-state \( \rho = -\rho_X \) to the effective contribution of single fermionic zero-point quantum-field mode of observable fermions. During the inflation, the supersymmetry is suggested to be broken down, so that the imaginary part of scalar field \( \phi \) representing the component of chiral superfield acquires a mass of Planckian scale, that makes its dynamics irrelevant (or simply frozen) at the inflationary stage under study. In that case, the shift symmetry of superpotential by setting \( \mu \to \mu_X \) modifies the quadric bare potential by the quartic term with the constant

\[
\lambda_{\text{bare}} = \frac{28\pi}{3} G m^2_{\text{bare}}.
\]

Sure, the bare values run in accordance with both the renormalization group and redefinitions in the effective action\(^6\). Thus, they can depend on the field value at low energies, at least, i.e. when \( \phi \) is close to zero.

We accept the nonperturbative or low-energy term\(^7\) of superpotential by setting

\[
f^2_{\text{LR}} (\phi) = \frac{1}{24\pi G} \hat{\rho}(M),
\]

Consider the ansatz

\[
\left( \frac{M}{\mu_X} \right)^3 = 1 - \left( 1 - \exp \left\{ -\frac{\phi^2}{m^2} \left[ 1 + \mathcal{C}(\phi) \right] \right\} \right)^{\nu},
\]

where \( \tilde{m} \) introduces the mass parameter, while \( \mathcal{C}(\phi) \) is a polynomial function, describing corrections to the quadratic dependence of the exponent argument versus the filed.

Then, at \( \phi \to 0 \) we get \( M \to \mu_X \), and the vacuum density of energy nullifies, so that at \( \nu = 3 \) the superpotential behaves like

\[
f_{\text{LR}} \sim m_{\text{Pl}}^2 \mu_X^2 \cdot \sqrt{1 - \frac{M}{\mu_X}} \sim \frac{m_{\text{Pl}}^2 \mu_X^2}{\tilde{m}^3} \phi^3.
\]

\(^5\) See, for instance, [11].

\(^6\) Moreover, the bare mass squared can change its sign, that can lead to the appearance of local minima in the potential.

\(^7\) The notion of “low energy” means the region of potential values close to zero in comparison with its values during inflation, say.
and
\[
\frac{\partial f_{\text{LE}}}{\partial \phi} \sim \frac{m_{\text{pl}} \mu_X^2}{m^3} \phi^2,
\]
which gives the low-energy correction to the bare quartic potential
\[
V_{\text{LE}}(\phi) = \frac{\lambda_{\text{LE}}}{4} \phi^4
\tag{11}
\]
with
\[
\lambda_{\text{LE}} \sim \frac{m^2 \mu_X^4}{m^6}.
\tag{12}
\]
Setting
\[
\lambda_{\text{LE}} \sim \lambda_{\text{bare}}, \quad \text{and} \quad \tilde{m} \sim m_{\text{bare}},
\tag{13}
\]
we find
\[
\lambda_{\text{LE}} \sim \frac{\mu_X}{m_{\text{pl}}}, \quad \tilde{m} \sim \sqrt{\mu_X m_{\text{pl}}}.
\tag{14}
\]

We consider the situation with thin domain walls corresponding to the gauge-mediated breaking down SUSY. It suggests that the correction function $C(\phi)$ could look as the expansion in inverse $\phi_g \sim \mu_X$ determined by a strong-field interaction in the gauge sector, so that to the leading order one could expect
\[
C(\phi) \to \frac{\phi^2}{\phi_g^2}.
\tag{15}
\]

Hence, at $\phi^2 \gg \phi_g^2 \sim \tilde{m}\phi_g \sim \tilde{m}\mu_X$ we arrive to $M \to 0$, which gives $U_{\text{LE}}(\phi) \to -\rho_X$ and $V_{\text{LE}}(\phi) \to 0$.

We suppose that the actual superpotential is well approximated by its low-energy term at $\phi^2 < \phi_g^2$, while at $\phi^2 \gg \phi_g^2$ the superpotential tends to the bare form. It means the followings:

- The first condition in (13) should naturally take place.
- The bare mass in the quadratic term of potential is substituted by its running value $m_{\text{bare}} \to m(\phi)$, which tends to zero at $\phi^2 < \phi_g^2$, so that, at least,
\[
m^2(\phi_*) \phi_*^2 \ll \lambda_{\text{LE}} \phi_*^4,
\tag{16}
\]
equivalent to
\[
m^4(\phi_*) \ll \frac{\mu_X^5}{m^4_{\text{pl}}} \sim \lambda_{\text{LE}} \mu_X^4,
\tag{17}
\]
that means $m(\phi_*) \ll \mu_X$.

Then, the actual potential
\[
V_{\text{act}} \approx V_{\text{LE}}(\phi) + \frac{\lambda_{\text{bare}}}{4} \phi^4
\]
acquires a positively-valued extremum at $\phi_*^2 \sim \tilde{m}\mu_X$, that means the breaking down SUSY, while
\[
U_{\text{act}} \approx U_{\text{LE}}(\phi) + \frac{\lambda_{\text{bare}}}{4} \phi^4
\]
takes a negative value at the extremum, that guarantees its AdS-position.

The domain-wall thickness is of the order of
\[
\delta r \sim \frac{(\delta \phi)^2}{m_{\text{pl}} \mu_X^2},
\tag{18}
\]
where $\delta \phi$ is the field increment between the fields corresponding to the flat and AdS-states, i.e. $\delta \phi \sim \phi_*$, hence,
\[
\delta r \sim \frac{\tilde{m}}{m_{\text{pl}} \mu_X},
\tag{19}
\]
The second condition of (14) gives
\[
\delta r \sim \frac{1}{\tilde{m}},
\tag{20}
\]
that means that the thickness of domain wall is fixed by the bare mass of scalar field in the theory. Then, at $\mu_X \sim 10^4$ GeV we numerically get
\[
\lambda \sim 10^{-14}, \quad \tilde{m} \sim 10^{12} \text{ GeV}.
\tag{21}
\]

Fig. 1 represents the qualitative behavior of actual potential $U_{\text{act}}(\phi)$ in comparison with the case, when the bare potential is set to zero.

\[\text{FIG. 1: The actual potential } U_{\text{act}}(\phi) \text{ with account of quartic bare contribution (solid line), and the low-energy term } U_{\text{LE}}(\phi) \text{ alone (dashed line). The ball and arrows symbolize the field evolution to the flat extremum during the inflation and after it (see Section IV).} \]

The local peak of actual potential corresponds to
\[
U_0 \sim \frac{f_{\text{LE}}^2}{\phi_*^4} \sim \sqrt{m_{\text{pl}}^2 \mu_X^2}.
\tag{22}
\]
Note, that without the supergravity correction linear in $G$, the potential remains positive.

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Note, that the simplest tadpole diagram due to the bare quartic interaction generates the mass determined by $m^4 \sim \lambda_{\text{LE}}^2 \mu_X^2$, if one puts the cut off about $\mu_X$. Such the mass will run to the bare value, when the cut off tends to the Planck scale.
III. THE VACUUM STRUCTURE

The offered ansatz for the potential gives two AdS-vacuum states $|\Phi_\pm\rangle$ positioned in the extremal points connected by the $\phi \leftrightarrow -\phi$ symmetry. These states possess equal values of energy density $\rho = -\rho_\pm$. In general, this condition can be perturbed under constraint of $\rho_\pm \sim \rho_X$, which does not essentially break the vacuum texture. The single flat vacuum-state $|\Phi_0\rangle$ is positioned at $\phi = 0$ with exact SUSY and $\rho = 0$.

The domain-wall fluctuations cause the mixing of such initial vacuum-states. Indeed, the bubbles of both AdS-vacua can appear in the flat state. Then, the mixing matrix of vacuum-states takes the form

$$H = \begin{pmatrix} -\rho_+ & 0 & \bar{\rho}_+ \\ 0 & -\rho_+ & \bar{\rho}_- \\ \bar{\rho}_+ & \bar{\rho}_- & 0 \end{pmatrix}, \quad (23)$$

where the mixing elements can be taken positive, and by construction $\bar{\rho}_\pm \ll \rho_\pm$, that represents the seesaw mechanism usually applied in the phenomenology of quarks [27].

Let us put $\rho_\pm = \rho_X$ and $\rho_\pm = \rho_X(1 + u)$ at $u \ll 1$. Therefore, we can separate the leading approximation suggesting

$$H_0 = \begin{pmatrix} -\rho_+ & 0 & \rho_+ \\ 0 & -\rho_+ & \rho_- \\ \rho_+ & \rho_- & 0 \end{pmatrix}, \quad (24)$$

and the correction in the form

$$V = -u \cdot \rho_+ \begin{pmatrix} 0 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & 0 \end{pmatrix}, \quad (25)$$

so that $H = H_0 + V$. Then, the eigenstate problem can be solved perturbatively in $u \rightarrow 0$.

The eigensystem of $H_0$ is easily found by the transformation $H_0 \rightarrow H_0' = U_0^\dagger H_0 U_0$ at

$$U_0 = \begin{pmatrix} \cos \varphi & \sin \varphi & 0 \\ -\sin \varphi & \cos \varphi & 0 \\ 0 & 0 & 1 \end{pmatrix}, \quad (26)$$

yielding

$$H_0' = \begin{pmatrix} -\rho_+ & 0 & \bar{\rho} \\ 0 & -\rho_+ & \bar{\rho} \\ \bar{\rho} & \bar{\rho} & 0 \end{pmatrix}, \quad (27)$$

with

$$\bar{\rho} = \sqrt{\bar{\rho}_+^2 + \rho_+^2}, \quad (28)$$

if

$$\tan \varphi = \frac{\bar{\rho}_-}{\rho_-}. \quad (29)$$

Matrix $H_0'$ in (27) has got the isolated 2×2-block, which can be further transformed to the diagonal form by matrix

$$U = \begin{pmatrix} 1 & 0 & 0 \\ 0 & \cos \theta & \sin \theta \\ 0 & -\sin \theta & \cos \theta \end{pmatrix}, \quad (30)$$

at

$$\tan 2\theta = \frac{2\rho_+}{\rho_X} \ll 1. \quad (31)$$

Then,

$$U^\dagger H_0' U = \text{diag} [-\rho_X, -\rho_{\text{AdS}}, \rho_{\text{AdS}}], \quad (32)$$

where

$$\rho_{\text{AdS}} = -\rho_X \cos^2 \theta - \bar{\rho}_+ \sin 2\theta,$$

$$\rho_{\text{AdS}} = -\rho_X \sin^2 \theta + \bar{\rho}_+ \sin 2\theta, \quad (33)$$

that can be further expanded in $\theta \rightarrow 0$ due to (31),

$$\rho_{\text{AdS}} \approx -\rho_X \frac{\bar{\rho}_+^2}{\rho_X}, \quad (34)$$

The eigenstates are determined by the product of rotations $U_0'u$ as follows:

$$|\Phi_{\text{AdS}}'\rangle = \cos \varphi |\Phi_+^\dagger \rangle - \sin \varphi |\Phi_-^\dagger \rangle,$$

$$|\Phi_{\text{AdS}}\rangle = \cos \theta \left[ \sin \varphi |\Phi_+^\dagger \rangle + \cos \varphi |\Phi_-^\dagger \rangle \right] - \sin \theta |\Phi_0\rangle, \quad (35)$$

$$|\Phi_{\text{AdS}}\rangle = \sin \theta \left[ \sin \varphi |\Phi_+^\dagger \rangle + \cos \varphi |\Phi_-^\dagger \rangle \right] + \cos \theta |\Phi_0\rangle.$$}

In the simplest realization of potential with the incorporation of $\phi \leftrightarrow -\phi$ symmetry, we get $\bar{\rho}_- = \bar{\rho}_+$ in $H_0$, so that $\varphi = \pi/4$. Remember, that for thin domain walls we evaluate the mixing elements in [28] by

$$\bar{\rho} \sim \frac{\mu_X^6}{m_{F_4}^2},$$

so that $\mu_X \sim 10^4 \text{ GeV}$ is consistent with the observed scale of cosmological constant. Such the kind of relation between the scales of SUSY breaking down and cosmological constant was derived by T.Banks [28] in other way of arguments.

The corrections to the energy densities linear in $u$ are straightforwardly determined by appropriate diagonal matrix elements of perturbation $V$, so that

$$\delta \rho_{\text{AdS}} = -u \cdot \sin^2 \varphi \cdot \rho_X,$$

$$\delta \rho_{\text{AdS}} = -u \cdot \cos^2 \varphi \cdot \cos^2 \theta \cdot \rho_X \approx -u \cdot \cos^2 \varphi \cdot \rho_X, \quad (36)$$

$$\delta \rho_{\text{AdS}} = -u \cdot \cos^2 \varphi \cdot \sin^2 \theta \cdot \rho_X \approx -u \cdot \cos^2 \varphi \cdot \frac{\bar{\rho}_+^2}{\rho_X}.$$
Non-diagonal matrix elements of perturbation result in the mixing of states defined in (35). These elements take hermitian values,

\[
\langle \Phi_{\text{AdS}} | \Phi \prime_{\text{AdS}} \rangle = \frac{u}{2} \rho_\chi \sin 2 \varphi \cos \theta, \\
\langle \Phi_{\text{dS}} | \Phi \prime_{\text{AdS}} \rangle = \frac{u}{2} \rho_\chi \sin 2 \varphi \sin \theta, \\
- \langle \Phi_{\text{dS}} | \Phi \prime_{\text{dS}} \rangle = \frac{u}{2} \rho_\chi \cos^2 \varphi \sin 2 \theta,
\]

so that at \( \theta \ll 1 \) we approximately get

\[
\delta \langle \Phi_{\text{AdS}} \rangle \approx \frac{u \sin 2 \varphi}{2 \sin^2 \theta} \left\{ -|\Phi_{\text{AdS}}^\prime| + \cot \varphi \sin^2 \theta |\Phi_{\text{dS}}\rangle \right\},
\]

\[
\delta \langle \Phi_{\text{dS}} \rangle \approx \frac{u \sin 2 \varphi}{2 \sin^2 \theta} \left\{ -|\Phi_{\text{dS}}^\prime| + \sin \theta |\Phi_{\text{dS}}\rangle \right\}.
\]

wherein one could further expand in \( \theta \), replacing \( \sin \theta \) by \( \theta \) itself. Moreover, we can give the above corrections to the states in terms of initial basis \( \{ |\Phi_{\text{dS}}^\pm \rangle, |\Phi_{\text{dS}}^\mp \rangle, |\Phi_{\text{dS}} \rangle \} \), so that

\[
\delta \langle \Phi_{\text{AdS}} \rangle \approx \frac{u \sin 2 \varphi}{2 \theta^2} \left\{ \sin \varphi |\Phi_{\text{dS}}^\mp \rangle + \cos \varphi |\Phi_{\text{dS}}^\mp \rangle - \theta |\Phi_{\text{dS}} \rangle \right\},
\]

\[
\delta \langle \Phi_{\text{dS}} \rangle \approx \frac{u \sin 2 \varphi}{2 \theta^2} \left\{ - \cos \varphi |\Phi_{\text{dS}}^\mp \rangle + \sin \varphi |\Phi_{\text{dS}}^\mp \rangle - \theta^3 \cot \varphi |\Phi_{\text{dS}} \rangle \right\},
\]

\[
\delta \langle \Phi_{\text{dS}} \rangle \approx u \theta \cos \varphi \left\{ \theta^2 \sin 2 \varphi |\Phi_{\text{dS}}^\mp \rangle - |\Phi_{\text{dS}}^\mp \rangle + \theta \cos \varphi |\Phi_{\text{dS}} \rangle \right\}.
\]

Therefore, the mixing remains under the control of perturbative theory, if

\[
u \ll \sin^2 \theta,
\]

i.e., when the splitting between \( \rho_{\text{AdS}}^\pm \) and \( \rho_{\text{dS}} \) is much less than \( \rho_{\text{dS}} \). Otherwise, the texture of \( \text{(35)} \) is essentially unchanged. Nevertheless, we can generically expect that the stationary AdS-levels include the mixtures of initial \( |\Phi_{\text{dS}}^\pm \rangle \)-states with amplitudes of the order of unit, while the contribution of flat state \( |\Phi_{\text{dS}} \rangle \) is suppressed. The dS-level is dominantly represented by the flat state, while the SUSY breaking amplitudes give a suppressed admixture.

Some other approaches to the cosmological constant problem in the framework of seesaw mechanism can be found in [24, 30].

IV. INCORPORATING THE INFLATION

The standard scenario of inflation with the single scalar field \( \phi \) possessing the quadratic and quartic self-couplings suggests that the inflaton takes values about the Planck scale. As we have mentioned above, at such the field values the running mass can be close to its bare value, that leads to approximate equality of quadratic and quartic terms of potential. Then, we get well elaborated picture of inflation consistent with observed features of matter inhomogeneity in the Universe: the magnitude of fluctuations of matter density, spectral index of scalar fluctuations and fraction of tensor perturbations, if one takes the couplings in agreement with \( \text{(21)} \). Thus, the parameters of inflaton potential are inherently related with the scale of SUSY breaking down.

However, a numerical analysis should be performed with more caution, since the relation between the scales in \( \text{(7)} \) due to the simplest connection of bare quadratic potential with the quartic coupling suggests a too steep growth of quartic term, which starts to dominate during the inflation. But the quartic-potential inflation is almost inconsistent with the data obtained by WMAP [3–5] and experiments on baryonic acoustic oscillations and spacial distribution of galaxies (BAO) [6] as well as on the supernovae Ia (SN) [7–10], and it can be marginally accepted, only. Then, one should consider a mechanism for the dominance of quadratic term alone, that is still consistent with the data.

Let us consider the problem in more detail. So, at Planckian values of scalar field we get the potential in the form

\[
V = \frac{m^2}{2} \phi^2 + \frac{\lambda}{4} \phi^4,
\]

with positive mass squared \( m^2 > 0 \). In the homogeneous Friedmann–Robertson–Walker metric

\[
ds^2 = dt^2 - a^2(t)dr^2
\]

with the scale factor of expansion \( a(t) \), the standard evolution equations with respect to time read off as followings:

- The field runs according to

\[
\dot{\phi} + 3H \phi + m^2 \phi + \lambda \phi^3 = 0,
\]

where the over-dot denotes the derivative with respect to time \( \dot{\phi} = \frac{d\phi}{dt} \), and the Hubble rate is defined by \( H = \dot{a}/a \).

- The Friedmann equation determines the rate of expansion

\[
H^2 = \frac{\kappa^2}{3} \left\{ \frac{1}{2} \dot{\phi}^2 + \frac{m^2}{2} \phi^2 + \frac{\lambda}{4} \phi^4 \right\},
\]

with \( \kappa^2 = 8\pi G \).

- The acceleration of expansion is given by

\[
\dot{H} = -\frac{1}{2} \kappa^2 \dot{\phi}^2.
\]

During the inflation the kinetic energy is suppressed with respect to the potential term, hence, the Hubble rate
slowly changes in accordance with (44), so it is almost a constant.

It is spectacular that the homogeneous evolution demonstrates the behavior of parametric attractor: the kinetic and potential energies of inflaton rapidly tend to definite critical values independent of initial data, while the critical points gain a driftage with the slowly changing Hubble rate\(^9\). In order to show this fact, we follow the method developed in [36, 37] and introduce appropriate scaling variables

\[
x = \frac{k}{\sqrt{6}} \frac{\dot{\phi}}{H}, \quad y^2 = \frac{k^2}{12H^2} \left(2m^2\phi^2 + \lambda \phi^4\right),
\]

as well as the control parameter of driftage

\[
z^4 = \frac{3\lambda}{k^2H^2}.
\]

Defining the amount of e-folding to the end of inflation by \(N = \ln a_{\text{end}} - \ln a\) and denoting the derivative with respect to \(N\) by prime \(\frac{d}{dN} \equiv \cdot\), we get the evolution equations of autonomous system with the parameter \(z\)

\[
x' = -3x^3 + 3x + 2z \xi(y, z) \zeta(y, z),
\]

\[
yy' = -3x^2y^2 - 2xyz \xi(y, z) \zeta(y, z),
\]

where

\[
\xi^2(y, z) = y^2 + u^4z^4,
\]

\[
\zeta^2(y, z) = \xi(y, z) - u^2z^2,
\]

while

\[
z' = -\frac{3}{2} x^2 z,
\]

and \(u\) is the constant parameter defined by

\[
u^2 = \frac{k^2m^2}{6\lambda}.
\]

Hence, the field takes the values

\[
\phi = \pm \frac{m\xi(y, z)}{uz\sqrt{\lambda}}.
\]

The Friedmann condition of (43) is transformed into

\[
x^2 + y^2 = 1,
\]

that is conserved by the dynamical system of (47)–(50), of course.

The acceleration takes place at \(\ddot{a} > 0\), that gives

\[
\frac{\ddot{a}}{a} = \dot{H} + H^2 > 0,
\]

equivalent to \(-\dot{H}/H^2 < 1\) yielding the condition of inflation end

\[
x^2 < \frac{1}{3}.
\]

The nonzero critical points \(\{x_c, y_c\}\) put \(x' = y' = 0\), and they are positioned on the Friedmann circle of (51) and related by

\[
3x_c y_c^2 = -2z \xi(y_c, z) \zeta(y_c, z).
\]

Linear perturbations \(x = x_c + \bar{x}\) and \(y = y_c + \bar{y}\) in (47) give

\[
\begin{pmatrix}
\bar{x}' \\
\bar{y}'
\end{pmatrix} = \begin{pmatrix}
3 - 9x_c^2 & -\frac{2z^2}{3x_c} (3\xi_c - 2u^2z^2) \\
-3x_c y_c & \frac{2z^2}{3x_c y_c} (3\xi_c - 2u^2z^2) - 6x_c^2
\end{pmatrix} \begin{pmatrix}
\bar{x} \\
\bar{y}
\end{pmatrix}
\]

so that under the constraint of (51) resulting in the relation \(x_c \bar{x} + y_c \bar{y} = 0\), we find

\[
\bar{x}' = \gamma_c \cdot \bar{x},
\]

at

\[
\gamma_c = 3 - 9x_c^2 - \frac{x_c^2}{\xi_c} (2u^2 + \xi_c),
\]

where we put \(\xi_c = \xi(y_c, z)\) and \(\zeta_c = \zeta(y_c, z)\). Then, perturbations exponentially decline with the expansion as

\[
\bar{x} = \bar{x}_{\text{tot}} e^{-\gamma_c (N_{\text{tot}} - N)},
\]

if

\[
\gamma_c > 0,
\]

yielding

\[
x_c^2 \left(1 - \frac{y_c^4}{2} \frac{2u^2 + \xi_c}{\xi_c^2} \frac{\gamma_c}{\gamma_c^2}ight) < \frac{1}{3},
\]

consistent with the condition of inflationary expansion [52]. Therefore, the critical point is stable during the inflation, and we have got the quasi-attractor at

\[
x_c^2 \ll 1, \quad y_c^2 \approx 1,
\]

because the control parameter has got a slow driftage as

\[
z \approx z_c \left\{1 - \frac{3}{2} x_c^2 (N - N_c)\right\},
\]

effective at large intervals

\[
|N - N_c| \ll \frac{1}{x_c^2}.
\]

We are interested in two limits:

\[\text{footnote 9: The dependence of inflation on initial data were originally studied in [33, 35].}\]
\[ u^2 z^2 \ll 1 \quad \text{quartic term dominance}, \]
\[ u^2 z^2 \gg 1 \quad \text{quadratic term dominance}. \]

So, at \( y \sim 1 \)
\[
\xi = \begin{cases} 
  y, & u^2 z^2 \ll 1, \\
  u^2 z^2, & u^2 z^2 \gg 1,
\end{cases}
\]
\[
(61)
\]
and
\[
\zeta = \begin{cases} 
  \sqrt{y}, & u^2 z^2 \ll 1, \\
  \frac{y}{u z \sqrt{2}}, & u^2 z^2 \gg 1.
\end{cases}
\]
\[
(62)
\]
Then, the attractor stability takes place at
\[
x_c^2 < \frac{2}{3}, \quad u^2 z^2 \ll 1, \\
x_c^2 < \frac{1}{2}, \quad u^2 z^2 \gg 1.
\]
\[
(63)
\]
The amount of e-foldings is accurately approximated by
\[
N \approx \frac{2}{3} \int \frac{dz}{x_c^2 z},
\]
\[
(64)
\]
with \( [53] \), so that
\[
N \approx \begin{cases} 
  \frac{3}{4 u^2 z^4}, & u^2 z^2 \ll 1, \\
  \frac{3}{u^2 z^2}, & u^2 z^2 \gg 1.
\end{cases}
\]
\[
(65)
\]
The inhomogeneities are approximated by the scalar and tensor densities of spectra versus the wave-vector at \( k = a(t) H \) as follows:
\[
\mathcal{P}_S(k) = \left( \frac{H}{2\pi} \right)^2 \left( \frac{H}{\phi} \right)^2, \\
\mathcal{P}_T(k) = 8\kappa^2 \left( \frac{H}{2\pi} \right)^2,
\]
\[
(66)
\]
which can be accurately evaluated in terms of quasi-attractor dynamics by
\[
\mathcal{P}_S(k) = \frac{\lambda}{8\pi^2} \frac{1}{z^4 x_c^2}, \\
\mathcal{P}_T(k) = \frac{6\lambda}{\pi^2} \frac{1}{z^4},
\]
\[
(67)
\]
while the ratio
\[
r = \frac{\mathcal{P}_T}{\mathcal{P}_S} = 48 x_c^2 \ll 1,
\]
\[
(68)
\]
and it determines the relative contribution of tensor spectrum.

The spectral index of scalar spectrum is defined by
\[
n_S - 1 \equiv \frac{d \ln \mathcal{P}_S}{d \ln k}.
\]
\[
(69)
\]
It can be calculated under the condition
\[
\ln \frac{k}{k_{\text{end}}} = -N - 2 \ln \frac{z}{z_{\text{end}}},
\]
\[
(70)
\]
which gives
\[
\frac{d \ln k}{d N} \approx -1,
\]
\[
(71)
\]
to the leading order in \( 1/N \).

Then, following \([59], [61], [62], [65]\), we get the limits,
\[
\begin{array}{|c|c|c|c|}
\hline
\text{limit} & \mathcal{P}_S(k) & n_S - 1 & r \\
\hline
u^2 z^2 \ll 1 & \frac{2\lambda}{3\pi^2} N^3 & -\frac{3}{N} & 16 \frac{N}{N} \\
\hline
u^2 z^2 \gg 1 & \frac{\lambda u^2}{\pi^2} N^2 & -\frac{2}{\pi} & 8 \frac{N}{N} \\
\hline
\end{array}
\]
\[
(72)
\]
The data on the correlation in the plain of \( \{n_S, r\} \) \([1, 3–10]\) prefer for the case of quadric coupling at \( N \approx 60 \) \([38]\), while the quartic coupling alone is marginally consistent with the data under \( N \approx 80 \), which is rather unrealistic \([37]\).

Therefore, we conclude that realistic scenario suggests \( u^2 z^2 \gg 1 \), i.e. \( u^4 z^4 \gg 1 \), which results in
\[
u^2 \gg \frac{4}{3} N,
\]
\[
(73)
\]
due to \([53]\). From the relation between the bare values in \([7]\), we get
\[
u_{\text{bare}}^2 = \frac{1}{3},
\]
\[
(74)
\]
which is inconsistent with the referable dominance of quadratic term. However, we have to take into account the running of potential parameters mentioned above. In this respect one could put the quartic constant approximately equal to its bare value, since this value is extremely small, so that one could expect no significant renormalization of the constant even at large logarithmic increment of scale. In contrast, one can put
\[
m(m_{\text{bare}}) \sim m_{\text{bare}}, \quad m(mp_1) \sim K \cdot m_{\text{bare}},
\]
\[
(75)
\]
at \( K \gg 1 \). Then,
\[
u^2 \sim \frac{K^2}{7},
\]
\[
(76)
\]
or \( K \gg 25 \), hence, the actual mass of inflaton should be about \( 10^{13} \) GeV in agreement with the phenomenological analysis of observed data. Such the situation is not in any contradiction with quite general properties of the potential, as concerns its scales, since factors of the form \( 4 \pi^2 \) could be responsible for the finite rescaling used above. Thus, the dominance of quadratic term can be actual.

One could simply require \((73)\) from the form of potential at \( \phi \approx m_{p1} \) by setting \( m^2 \gg \lambda m_{p1}^2 \sim \mu_X m_{p1} \sim m_{\text{bare}}^2 \), of course. However, the estimate of \((76)\) can be derived from the analysis presented above only.

Another aspect of inflation is related with the potential behavior in vicinity of \( \phi = 0 \). First, the AdS minimum is not essential for the inflation, since its scale of energy density \( \rho_X \sim \mu_X^2 \) is essentially less than the energy density to the end of inflation \( \rho_{\text{end}} \sim \lambda m_{p1}^2 \sim \mu_X m_{p1}^2 \), i.e. there is the hierarchy \( \rho_X \ll \rho_{\text{end}} \). Second, the potential barrier between the flat and AdS vacua has the height about

\[ U_0 \sim \sqrt{m_{p1}^3 \mu_X^2} \sim \frac{\mu_X^3}{m_{p1}} \rho_{\text{end}} \ll \rho_{\text{end}}. \]

Therefore, the local peak of potential is also inessential for the inflation, too. However, both these items can be involved into the mechanism of reheating.

In this respect, the most essential effect is related with the potential barrier, since near the peak the effective inflation-mass squared is negative, that causes the tachyonic instability resulting in the preheating mechanism: a rapid decay of the inflaton to quanta at the moment, when the energy density becomes comparable with the peak height \([39]\). Then, we can evaluate the preheating temperature \( T_{\text{preh}} \) by

\[ T_{\text{preh}}^2 \sim U_0, \quad \text{if} \quad m_{\text{eff}} \ll T_{\text{preh}}, \quad (77) \]

where \( m_{\text{eff}} \) is an effective mass of quanta. Numerically, \((77)\) gives

\[ T_{\text{preh}} \sim 10^9 \text{ GeV}. \quad (78) \]

In the model under consideration, the inflaton quanta with respect to the flat vacuum possess the effective mass equal to zero,

\[ m_{\text{eff}} = 0, \]

while the quanta with respect to the AdS vacuum have nonzero mass, that can be roughly evaluated by

\[ \left( m_{\text{eff}}^{\text{AdS}} \right)^2 \sim \frac{\partial^2 U}{\partial \phi^2} \approx \frac{U_0}{\phi^2}, \quad \Rightarrow \quad m_{\text{eff}}^{\text{AdS}} \sim m_{\text{bare}}, \]

though the mass can have got a much less value. Nevertheless, the decay to massless quanta would be kinematically preferable, that leads to the relaxation of inflaton in vicinity of flat vacuum, as pictured in Fig. 1.

Thus, the scalar field could inflationary evolve and relax in the flat vacuum after the tachyonic preheating in agreement with current experimental constraints from the observational data. Such the evolution could explain why we are living in the vacuum we have got.

A discussion of some other aspects of inflation in supergravity can be found in review \([22]\). For instance, the mechanism with an effectively real inflaton was considered by Kawasaki, Yamaguchi and Yanagidain in \([40]\), while the mechanism for the vacuum stabilization was offered by Kachru, Kallosh, Linde and Trevedi in \([41]\). Further developments include also the problem of irreversible vacuum decays that serve as sinks for the probability flow \([42]\).

V. THREE GENERATIONS

The vacuum structure considered in Section \([11]\) suggests that quantum field vibrations in vicinity of initial vacuum-states can mix. We can easily investigate the main features of such the mixing in the simplest case of potential symmetry versus \( \phi \leftrightarrow -\phi \). Then, the mass matrix for fermions takes the form

\[ M = \begin{pmatrix} m_X & \mu_\Lambda & \mu_B \nu \r \mu_\Lambda \nu & m_X & \mu_B \nu \r \nu^* \mu_B & \nu^* m_S & \nu^* \mu_B \nu \r \end{pmatrix}, \quad (79) \]

where \( m_X \) stands for the fermion mass in the AdS vacuum, when SUSY is broken down, while \( m_S \) denotes the fermion mass in the flat supersymmetric vacuum. Elements \( \mu_{A,B} \) introduce the mixing. In \((79)\) all of elements except \( \mu_0 \) are real due to the freedom in the definition of complex phases for the initial states, while \( |\mu_0| = \mu_B \) because of symmetry.

Let us introduce the complex phase \( \gamma \) by setting

\[ \mu_0 = \mu_B e^{i\gamma}. \quad (80) \]

Transforming the matrix to \( M_U = U \cdot M \cdot U^\dagger \) with

\[ U = \begin{pmatrix} c_0 & s_0^+ & 0 \\ -s_0^- & c_0 & 0 \\ 0 & 0 & 1 \end{pmatrix} \quad (81) \]

at

\[ c_0 = -s_0 = \frac{1}{\sqrt{2}}, \quad s_0^+ = s_0 e^{i\gamma}, \quad s_0^- = s_0 e^{-i\gamma}, \quad (82) \]

we get

\[ M_U = \begin{pmatrix} m_X - \mu_\Lambda x \cos \gamma - i \mu_\Lambda x \sin \gamma e^{i\gamma} & 0 \\ i \mu_\Lambda x \sin \gamma e^{-i\gamma} & m_X + \mu_\Lambda x \cos \gamma \sqrt{2} \mu_B \r \nu \mu_B & \nu^* m_S \nu \r \end{pmatrix}. \quad (83) \]

From \([83]\) we see that the analysis is essentially simplified at \( \gamma = \{0, \pm \pi\} \), when one can neglect effects caused by violation of combined invariance with respect to the
charge conjugation $\mathbb{C}$ and mirror inversion of space $\mathbb{P}$, and 
the matrix takes the symmetric form, so that at $\gamma = \pm \pi$

$$
\mathcal{M}^{(0)}_\mu = \begin{pmatrix}
m_X + \mu_A & 0 & 0 \\
0 & m_X - \mu_A & \sqrt{2}\mu_B \\
0 & \sqrt{2}\mu_B & m_s
\end{pmatrix},
$$

(84)

while the case of $\gamma = 0$ can be obtained from (84)
by changing the sign of $\mu_A$.

A small complex phase $\varepsilon \to 0$ of $\gamma = \pi - \varepsilon$ produces
the perturbation to (84), so that to the linear order in $\varepsilon$
it is equal to

$$
\mathcal{V} = \begin{pmatrix}
0 & i\mu_A\varepsilon & 0 \\
-i\mu_A\varepsilon & 0 & 0 \\
0 & 0 & 0
\end{pmatrix}.
$$

(85)

Matrix (84) can cause the hierarchy in both the masses
and mixings of fermion generations. Indeed, its eigenvalues
are given by

$$
m_{1,2} = \frac{1}{2} \left( \tilde{\mu}_A + m_s \pm \sqrt{(\tilde{\mu}_A + m_s)^2 + 8\mu_B^2} \right),
$$

$$
m_3 = m_X + \mu_A,
$$

(86)

where $\tilde{\mu}_A = m_X - \mu_A$. Setting

$$
\mu_B^2 \ll (\tilde{\mu}_A + m_s)^2
$$

(87)

at $\tilde{\mu}_A > 0$ and $m_s > 0$, we obtain

$$
m_2 \approx \tilde{\mu}_A + m_s,
$$

$$
m_1 \approx -\frac{2\mu_B^2}{\mu_A + m_s}.
$$

(88)

Therefore, conditions $m_X + \mu_A \gg \tilde{\mu}_A \gg m_s$ lead to the
hierarchy of fermion masses

$$
m_3 \gg m_2 \gg |m_1|.
$$

(89)

Two lighter generations are formed by superposition
of two initial states defining matrix (84). The superposition
is simply the rotation with angle $\theta_C$

$$
\tan 2\theta_C = \frac{2\sqrt{2}\mu_B}{\mu_A + m_s};
$$

(90)

which is the analogue of Cabibbo angle, since the
electroweak partners of fermion fields could have got
the mass matrix with the same texture, that leads to
the similar mixing of initial states, so that initially diagonal
electroweak charge currents acquire the mixing of two lighter
generations with the angle given by the difference of $\theta_C$
parameters for two kinds of fields.

Thus, the realistic scenario suggests the texture with

$$
m_X \sim \mu_A \sim \mu_X \gg |m_X - \mu_A| \gg m_s,
$$

$$
|m_X - \mu_A| \gg \mu_B,
$$

$$
\mu_B^2 \gg |m_s - \mu_A| m_s.
$$

(91)

Such the hierarchy can be natural, since parameters $m_X$
and $\mu_A$ are determined by the vacuum state with SUSY
broken down at the scale $\mu_X$, while $m_s$ could be equal
to zero in the case of exact SUSY. Then, the only condition
required is a small mixing between the ordinary fermionic
fields in the sectors with broken and exact SUSY, that
could serve as the definition of ordinary matter fields
in contrast to the hidden sector, wherein one could expect
$\mu_3 \sim \mu_X$, which breaks the hierarchy down, and it leads
to hidden fermions with masses of the order of $\mu_X$.

At the limit of $\mu_B \to 0$, one can expect the observation
of almost massless generation of ordinary fermions with
respect to the scale of SUSY breaking, of course. This
fact is in agreement with the experimental data.

The correction to the symmetric case produces a small
mixing of heavy generation of ordinary matter with
two lighter generations as well as the violation of CP-
invariance due to (91). This correction can be treated
perturbatively, if appropriate matrix elements of $\mathcal{V}$
is much less than the splitting between the levels, i.e. at
$|\mathcal{V}| \ll \Delta E$, that simply gives $\epsilon \ll 1$.

Similar features of generation structure can be exported
to the sector of fermion superpartners, i.e. scalar
fields of sfermions. Then, in the action quadratic versus
the sfermions, the mass matrix is composed by both
squares of initial masses positioned at the diagonal and
non-diagonal mixing parameters. Since the eigenvalues
could turn out to be negative, the negative sign would
indicate the generation of sfermion condensates.

Anyway, the observational situation suggests that sfermions
have no light states analogous to the almost massless generation.
Therefore, sfermions should imitate the texture of hidden sector.

The question about gauge vector fields is more specific,
since such the fields acquire masses due to the higgs ef-
fect. We can suppose that the observed mediators of
gauge interactions couple to the lightest generation of
higgs field, while two more heavy hidden generations of
higgs scalars as well as, probably, gauge fields can be
discovered at the energy scale of SUSY breaking down.

VI. DISCUSSION AND CONCLUSION

In the present paper we have offered the ansatz for the
low energy modification of bare quadratic potential for
the phenomenological real scalar field, so that the cor-
rection parameterizes the energy density given by zero-
point modes of quantum fields due to the supergravity
relation between the superpotential and energy density
linear in the Newton constant $G$. So, the low energy su-

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10 We do not consider the neutrino masses and mixing, here, since
this problem requires a more fine treatment because of an ex-
tremely low value of neutrino mass scale caused by physics be-
yond the Standard Model.
perpotential generates three local minima: one minimum corresponds to supersymmetric flat vacuum-state, while two other minima give the SUSY breaking down. To the same order in $G$, the bare superpotential induces the bare quartic term of potential, so that the actual potential taking into account all of bare and low energy terms has the characteristic form with the barrier separating the flat vacuum from two AdS-states. These initial vacuum-states are not stationary, since the bubbles of AdS-vacua in the flat vacuum generate the fluctuations inducing the mixing. Putting the domain wall thickness equal to the inverse bare mass by the order of magnitude and the low energy quartic coupling equal to the bare value, we have determined the bare mass scale $m \sim 10^{12}$ GeV and quartic coupling $\lambda \sim 10^{-14}$.

Having took into account the mixing described phenomenologically, we have found the stationary vacua represented by superpositions of initial flat and AdS-states. So, we have got two AdS-states with vacuum energy scale about the scale of SUSY breaking down, while the single dS-state acquires the cosmological constant consistent with the experimental data. Then, statically the offered potential naturally gives the vacuum with the desired small value of cosmological constant determined via the seesaw mechanism of mixing in terms of tuned potential barrier.

Further, we have shown, that the same scalar field can serve as the inflaton. The observations prefer for the quadratic term dominance in the region of Planckian fields. This constraint has been analyzed by means of quasi-attractor approach. Then, the inhomogeneity of matter in the Universe is in agreement with the quantum fluctuations of field during the inflation. The inflaton decays into massless quanta in vicinity of potential barrier, the height of which determines the temperature of preheating. This decay can take place on the background of flat vacuum, only. Then, the domain wall fluctuations transform the flat vacuum into the stationary dS-state, while the stationary AdS-states remain beyond the play.

The vacuum structure causes three fermion generations. We have analyzed the textures of mass matrices. So, the hierarchies for masses and mixing of the ordinary matter in the observable sector can be consistently constructed, or otherwise heavy states in the hidden sector and in the sector of superpartners for the ordinary fields can be introduced.

Thus, the offered model of potential is suitable for the description of three benchmarks: the naturally small cosmological constant, inflation with further preheating stage driving to the dS-state, and three fermion generations with appropriate hierarchies and mixing.

In this scheme all of SUSY breaking effects as well as consequent low energy condensates and phase transitions are accumulated in the energy density of initial AdS-vacua, while the flat vacuum is preserved from its influence due to the exact SUSY. The mixing of flat and AdS-states is described phenomenologically in terms of real scalar field. We have argued for the dynamical evolution to the dS-vacuum. So, we have presented the natural mechanism of driving to the observed cosmological constant, that is alternative to the renormalization group arguments developed in [43–49].

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