INTRODUCTION

The cosmological constant [1] and inflation [2–6] are important issues in modern cosmology. In speaking about the former issue within quantum field theory, one would expect that the scale of vacuum energy would be determined by a value that is characteristic of the interactions between elementary particles. Therefore, the value of the cosmological constant is extremely large. This value is defined, if not by the Planck mass, at least by a scale like a gauge symmetry breaking scale in particle physics, i.e., the grand unification scale, the electroweak symmetry breaking scale, or quark–gluon condensates. However, in each of these cases, the cosmological constant is significantly above the limits ensuing from observations of the anisotropy of the cosmic microwave background radiation (CMBR) [7–9], the large scale structure (LSS) of the universe [10], and the dependence of the brightness of type Ia supernovae (SNIa) on redshift [11–14]. According to these observations, the allowable density of vacuum energy is on the order of $10^{-3}$ eV, which does not fit into the concept of characteristic energies inherent in field interactions. Therefore, the literature discusses the mechanisms of transformation of the original large cosmological constant into a reduced value that is close to the one observed.

Among the models, the renormalization group running of the cosmological constant, which evolves depending on the expansion rate of the universe, i.e., the Hubble parameter [15–23], is noteworthy. In this case, the minimum of the renormalization-invariant effective potential, i.e., the vacuum energy density, depends on the coupling constants specified by the values of fields that are related to the Hubble scale, which leads to a slow logarithmic evolution of the cosmological constant according to the renormalization group equations.

Within another approach, a special class of asymptotically safe theories [24] is singled out among nonrenormalizable theories, including Einstein’s gravity. This class of theories has the following properties: dimensional factors in the form of powers of a certain scale are determined for a countable number of local operators of a nonrenormalizable theory with arbitrary coefficients so that, for the remaining dimensionless constants—charges, there are renormalization-group equations with fixed points, the set of which is finite when the theory asymptotically gains predictive power near the points of attraction. Asymptotically safe gravity may also include the evolution of the cosmological term, say, to its zero value, and inflation [25].

The possibility of using the quantum–mechanical seesaw mechanism with different energy density values characteristic of particle physics is also considered. This leads to a stationary vacuum whose cosmological constant has a reduced small scale [26–33].

At the same time, the inflation model, which solves many problems of observational cosmology, leads to an inflaton (scalar field) potential that is characterized by a mass of $m \approx 1.5 \times 10^{13}$ GeV, a large vacuum expectation value (much larger than the Planck energy scale), a very low self-interaction of the field (the constant of self-interaction of the fourth order is $\lambda \sim 10^{-13}$), and a flat plateau at the level of energy density $\Lambda^4$ for $\Lambda \sim 10^{16}$ GeV. This raises the question of whether the parameters of such an exotic potential are natural (see an extensive review on the relationship of inflation to the particle
physics and the mechanism of heating and thermalization of the universe after inflation in [34]).

In order to answer this question, the Higgs boson model is now being intensely investigated within the standard model of elementary particles. The scalar Higgs boson ϕ has a nonminimal coupling to gravity (scalar curvature R) in the form of a term of the Lagrangian \( L_{\text{int}} = \xi R \phi^4 \) with a coupling constant on the order of \( \xi \sim 10^4 \), which does lead to a sufficiently flat potential with a necessary plateau during a transition to an effective inflaton field minimally associated with gravity by means of a conformal transformation \([38–47]\). Consideration of the renormalization group corrections within this approach leads to a stringent constraint on the Higgs boson mass: \( 135.6 < m_H < 184.5 \text{ GeV} \).

The Higgs boson of the standard model may also play a significant role in the dynamics of the early universe given the minimum connection of this scalar field with gravity (\( \phi \rightarrow 0 \)) \([48]\). Namely, there is a critical mass of the Higgs boson which, considering the two-loop corrections \([49]\), is equal to \( m_H^{\text{crit}} = (153 \pm 3) \text{ GeV} \). Hence, a Higgs scalar with a supercritical mass cannot cause the inflation of the universe; if the Higgs particle is the only scalar field in the theory up to the Planck energy scale, then Higgs boson with a subcritical mass is forbidden because the inflation would have caused would have engendered a universe with a large-scale structure of matter that is completely different from the observed distribution. In this case, however, given a supercritical mass of the Higgs boson, the inhomogeneous matter distribution would be explained only by a fine tuning of the initial data, which can be avoided in the inflation theory. Consequently, subcritical masses of the Higgs boson require the introduction of an additional (in terms of the standard model) inflaton scalar field, which would ensure the dynamic formation of the necessary properties of the universe’s large-scale structure \([50]\).

The cosmological constraints on the Higgs boson mass can also be derived in other approaches, e.g., when this scalar is considered together with another field, as was done in \([51]\), where the constraint \( m_H < 134 \text{ GeV} \) was derived.

The work \([52]\) proposed a supergravity inflation model by choosing an appropriate type of the Kahler potential with an additional symmetry, i.e., the independence of the Kahler potential from the imaginary part of the scalar field. On a deeper level, the supergravity theory cannot find a way leading to models such as those in \([52]\). The work \([53]\) proposed a realistic model of inflation within the superstring theory that overcomes the problem of instability arising in the compactification of extra dimensions and passes from the anti-de Sitter vacuum of the initial state to a minus minimum inflaton potential with positive energy (de Sitter vacuum). However, in our view, the development of inflation may occur in the domain of such energy densities where supergravity is obviously violated. Hence, a self-consistent theory of the sub-Planck energy field should be based on the principle of renormalizability, whereas the potentials obtained within a complete and accurate supergravity theory include contributions from higher powers according to the Newtonian gravitational coupling constant \( G \), which are known to be nonrenormalizable. Therefore, these contributions are likely to be effectively reduced after supergravity is violated; hence, an effective Lagrangian should contain only those terms that do not contradict the renormalizability requirement, i.e., are controlled by quantum loop corrections. Otherwise, the field theory would have no predictive power at energies much lower than the Planck level. In this sense, the consideration of supergravity corrections to the potential should be limited to the leading terms that are linear with respect to the Newtonian constant \( G \), as was done by S. Weinberg \([54]\), and a requirement should be added for the order of the scalar field self-interaction, which should be four or less.

The problems of the cosmological constant and inflationary dynamics within generalizations of general relativity, e.g., conformal relativity theory, are a topic in and of themselves \([55]\).

In this paper, we propose to look at the problem of relaxation of the cosmological constant and naturalness of the parameters of the scalar field potential (which causes inflation) from a unified perspective based on the phenomenological introduction of a dynamic field whose potential is constant \( (V = \text{const}) \), i.e., flat, within the accuracy of contributions no greater than a certain power with respect to the inverse Planck mass. Hence, this value of the potential defines the cosmological constant. This field has not only a zero mass, but zero self-interaction. Fields of this kind appear naturally in superstring theory and are called modules; they determine the flat direction of the effective potential.

We introduce this field as a scalar component of the chiral superfield in supersymmetry theory. Next, we set the task of determining the parameters of the field superpotential so that, after taking into account the leading corrections due to supergravity in weak gravitational fields, i.e., corrections that are linear with respect to the gravitational coupling constant, the resulting renormalized field potential would become flat, which means that the field would indeed be a module. This procedure for determining the parame-

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1 Realistic models of low-energy inflation considering the restrictions ensuing from primordial nucleosynthesis, CMB anisotropy, and the inhomogeneous distribution of matter in the large-scale structure of the universe are presented in the works \([35–37]\), which investigate the supersymmetric version of the standard model of interactions between elementary particles with flat directions in the superpotential.
ters of the original superpotential is called a cosmological bootstrap, because the original mass and self-interaction of the field are bootstrapped according to the requirement to reduce their contributions by means of the leading corrections due to supergravity.

Here we assume that the scale of the potential plateau is much smaller than the Planck scale, but large enough to fit into the concept of interactions between elementary particles. For example, we consider it to be close in order of magnitude to the scale of the grand unification of gauge interactions. It turns out then that the bootstrap is possible and the initial parameters of the field are close to the mass and self-interaction of the inflation that ensures the large-scale structure of the universe.

The next step is a small violation of the bootstrap relations for the mass and self-interaction of the field, which may be caused by taking into account higher order corrections in the inverse Planck mass such as loop corrections taking into account heavy particles with Planck scale masses. This violation leads to a dynamical instability of the flat potential and the relaxation of the original cosmological constant during inflation caused by a field that was originally a modulus. Then the parameters of inflation are naturally defined by the small violation of the bootstrap and are consistent in order of magnitude to the observed parameters. It is important to note that the mass, vacuum expectation value, and self-interaction of the field are essentially determined by the introduction of a single parameter, i.e., the scale of the original cosmological constant, whose value is characteristic in the particle physics.

1. BOOTSTRAP

In the theory of gravity with a dimensional coupling constant $G$, it would be natural to expect that the vacuum energy $\rho_G$, which specifies the cosmological constant, is determined by the Planck scale $\tilde{m}_p = 1/\sqrt{8\pi G} \approx 2.4 \times 10^{18}$ GeV, so that

$$\rho_G = M^4, \quad (1)$$

where $M \sim \tilde{m}_p$. However, supersymmetry can be dynamically connected with another energy scale $\Lambda$, which is assumed to be much smaller than the Planck scale:

$$\Lambda \ll M. \quad (2)$$

From the phenomenological perspective, i.e., without going into the mechanism for introducing the scale $\Lambda$ along with $\tilde{m}_p$, we assume that, in theory, there are a couple of dimensional quantities with a definite hierarchy (2). It is natural to assume that local supersymmetry, i.e., supergravity, leads to a vacuum energy

$$\rho_S = \Lambda^4. \quad (3)$$

This finite renormalization from (1) to (3) is quite possible in supergravity, since the introduction of a superpotential in the form of a constant

$$W_0 = i\omega_0^3, \quad (4)$$

leads to an additional contribution to the vacuum energy density, so that [54]

$$V_0 = \rho_G - 24\pi G|W_0|^2 = M^4 - 24\pi G\omega_0^6, \quad (5)$$

and condition $V_0 = \rho_S$ is the determination of the scale $\omega_0$. The procedure described above is step “zero” of the bootstrap: the successive determination of the theory parameters through the initially specified values, in particular, the constant term in the superpotential through the primary scales $M$ and $\Lambda$.

At the same time, our aim is to describe the cosmological constant as a flat potential.

Vacuum energy in supersymmetry is specified by the potential

$$V_S = \left| \frac{\partial W}{\partial \Phi} \right|^2. \quad (6)$$

Where, in the case of the cosmological constant, i.e., a potential that is independent of the field $\Phi$, the superpotential should be written as

$$W \mapsto W_1 = i\omega_0 + \Lambda^2 \Phi, \quad (7)$$

so that

$$V_S = \Lambda^4. \quad (8)$$

The consideration of supergravity with linear contributions with respect to the Newtonian constant $G$ leads, in general, to the potential [54]

$$V = \rho_G + \left| \frac{\partial W}{\partial \Phi} \right|^2 - 24\pi G\left| W - \frac{1}{3} \Phi \frac{\partial W}{\partial \Phi} \right|^2 + \frac{16\pi G}{3} |\Phi|^2 \left| \frac{\partial W}{\partial \Phi} \right|^2. \quad (9)$$

Then, the substitution of a superpotential of the form (7) into (9) yields

$$V \mapsto V_1 = (M^4 - 24\pi G\omega_0^6) + \Lambda^4 \frac{4}{3} \Phi^2 \superscript{16\pi G}{\Lambda^4}, \quad (10)$$

3 In this paper we do not consider the issues of regularization of infinities and the renormalization group.

4 Following S. Weinberg, we confine ourselves to corrections to the potential that are linear with respect to the gravitational coupling constant $G$.

5 Here, of course, it is assumed that all values that differ only by a factor on the order of one are equivalent.
if the field $\phi = \Phi \sqrt{2}$ is assumed to be real (the general case is discussed in Appendix A). Moreover, we assume below for simplicity that the potential with supergravity corrections (9) has a symmetry with respect to the discrete operation of reflection $\phi \mapsto -\phi$, i.e., does not contain odd powers of the field $\phi$. In this step of the bootstrap, the cosmological constant is still specified by the energy density $V_S$ in (8) if the parameter $\omega_0$ is somewhat adjusted, so

$$M^4 - 24\pi G \omega_0^5 = 0. \quad (11)$$

However, potential (10) contains, in addition to the constant energy density, a field-dependent contribution; moreover, at this stage, supergravity corrections cause instability in the cosmological constant. This instability indicates that supergravity renormalizes the field-square term of the potential; hence, the field-square interaction should be introduced already at the stage of the superpotential $W$. Furthermore, we see that the module field in the theory with global supersymmetry loses this property in the theory with local supersymmetry, i.e., supergravity. Therefore, we set the task of determining the nontrivial field superpotential, which corresponds to the module field, provided that leading supergravity corrections (9) are considered.

Following the bootstrap procedure, we can write the superpotential as

$$W = i\omega_0^3 + \Lambda^2 \Phi + \frac{1}{2} \mu_0 \Phi^2 + \frac{g_0^2}{3} \Lambda^4. \quad (12)$$

Then, within an accuracy of the fourth order with respect to the field, i.e., taking into consideration the renormalized terms only, the potential $V = \hat{V} \Lambda^4$ takes the form

$$\hat{V} = 1 - \frac{1}{2} \mu^2 \phi^2 + \frac{\lambda^2}{4} \phi^4, \quad (13)$$

which considers the reduction based on (11) at the zero step of the bootstrap and introduces the notation for the dimensionless values with “caps”:

$$\hat{\phi}^2 = \frac{32\pi G}{3} \phi^2, \quad \hat{g}_0 = -\frac{3}{16\pi G \Lambda^2} g_0, \quad (14)$$

$$\hat{\omega}_0 = \frac{16\pi G}{3\Lambda^4} \omega_0^6, \quad \hat{\mu}_0 = \frac{3}{16\pi G \Lambda^2} \mu_0^2,$$

so that

$$\hat{\mu}^2 = 1 - \hat{\mu}_0^2 + 2\hat{g}_0 + \frac{3}{2} \hat{\omega}_0 \hat{\mu}_0, \quad (15)$$

Potential (13) is a constant energy density if we have the bootstrap relation

$$\hat{\mu}^2 = 0, \quad \hat{\lambda} = 0. \quad (16)$$

Provided that $\hat{\omega}_0^6 \gg 1$, it is easy to find solutions of Eqs. (16) for the parameters $\hat{\mu}_0$ and $\hat{g}_0$, since there are two sets corresponding to the solutions of the quadratic equation for $\hat{g}_0$:

$$\hat{g}_0 = 1 \pm \frac{1 - 7\hat{\mu}_0^2}{8\hat{\omega}_0^6}. \quad (17)$$

Namely,

$$\hat{\mu}_0^1 \approx \frac{-2}{3\hat{\omega}_0^3}, \quad \hat{\mu}_0^2 \approx \frac{-10}{3} \hat{\omega}_0^3, \quad (18)$$

$$\hat{g}_0^1 \approx \frac{7}{16} \hat{\mu}_0^1, \quad \hat{g}_0^2 \approx 2 - \frac{7}{16} \hat{\mu}_0^2.$$

The subsequent corrections to (18) can be written as an expansion in even powers of the relationship $\Lambda/\hat{m}_{Pl}$, and, since the coefficients of this expansion are strictly defined, one can speak, first, about the fine tuning of the bootstrap parameters and, second, about the fact that a mismatch in this fine tuning would violate the cosmological bootstrap. It is this violation pattern that we study in this paper: the introduction of corrections with respect to the powers of $\Lambda/\hat{m}_{Pl}$ to the “bare” superpotential is the reason for the dynamic instability of the cosmological constant, leading to the inflationary expansion of the universe.

Sets (18) not only specify different (in terms of hierarchy) values for the original constant of field self-interaction $\hat{g}_0$, since $\hat{g}_0^1 \gg \hat{g}_0^1$, but also, according to (6), lead to initial field potentials in supersymmetry that are significantly different in terms of physics. Indeed, the initial potential $V_S = \hat{V}_S \Lambda^4$ can be written as

$$\hat{V}_S = 1 + \frac{\hat{\mu}_0^6}{2} \phi^2 + \frac{\hat{\lambda}_0^4}{4} \phi^4, \quad (19)$$

where

$$\hat{\mu}_S = \hat{\mu}_0 - 2\hat{g}_0, \quad \hat{\lambda}_S = \hat{g}_0^2. \quad (20)$$

For solutions (18), the initial potential is stable: $\hat{\omega}_0^6 > 0$.

Moreover, considering the approximation $\hat{\omega}_0^6 \gg 1$, we can write explicitly

$$\hat{V}_S^I \approx 1 + \left( \frac{\hat{\mu}_0^6}{4} \right) \phi^2 + \left( \frac{7}{32} \hat{\mu}_0^2 \right) \phi^4, \quad (21)$$

$$\hat{V}_S^{II} \approx (1 - \phi^2) + \frac{1}{2} \hat{\mu}_0^2 \phi^2.$$
The pattern of the dependence of the initial potential \( \tilde{V}_S \) for sets I and II of the bootstrap solutions is shown in Fig. 1. As can be seen in the figure, set I represents a situation whereby the supersymmetry violation in the potential corresponds to the vacuum energy density \( \Lambda^4 \) and the vacuum expectation value of the field is zero, whereas set II assumes a reduced contribution of the field in the supersymmetry violation, which goes to zero, whereas set II assumes a reduced contribution of the field in the supersymmetry violation, which goes well with the violation of the above-suggested symmetry with respect to the discrete transformation \( \phi \mapsto -\phi \). However, set II is characterized by “natural” scaling values both for the mass parameter and self-interaction constant, whereas these parameters of set I have reduced values.

The violation of bootstrap relations (16) due to quantum loop corrections leads to an instability of the cosmological constant \( \Lambda = 1 \) and the inflationary expansion of the universe if the minimum of the potential corresponds to an energy density that is negligibly small compared with unity. In this case, the cosmological constant relaxes from an extremely large value to a value which we assume to be zero.

2. PHENOMENOLOGICAL ANALYSIS

Cosmological bootstrap relations are violated due to a variation in the parameters \( \omega_0, \mu_0, \) and \( g_0 \). We do not specify the nature of this variation in this section, assuming only, in a most general form, that quantum corrections can be expanded in a series with respect to the minor parameter of the ratio \( \Lambda / \tilde{m}_{pl} \). This raises a number of natural requirements.

First, corrections to \( \omega_0 \) contribute to the vacuum energy, so it receives, in the leading approximation, a contribution on the order of \( \delta V_0 \sim \Lambda^4 \) and, hence, \( \delta \omega_0 \sim O(1) \), since, according to the introduced definitions \( \delta V_0 \sim \Lambda^4 \omega_0 \). According to the construction,

\[
\hat{\omega}_0 \sim \left( \frac{\tilde{m}_{pl}}{\Lambda} \right)^2 \Rightarrow \hat{\omega}_0^3 \sim \left( \frac{\Lambda}{\tilde{m}_{pl}} \right)^4 . \tag{22}
\]

This correction may also lead to a dependent variation in the parameter \( \hat{\mu}_0 \sim \hat{\omega}_0^3 \), so that

\[
\delta \hat{\mu}_0 \sim \left( \frac{\Lambda}{\tilde{m}_{pl}} \right)^4 , \tag{23}
\]

and a consistent variation in the self-interaction constant \( \delta \hat{g}_0 \); however, such a transfer of variations

6 Otherwise we would have to introduce a dynamic field, which again would cause instability of the residual nonzero energy density and inflation if the residual energy density is positive. A negative value of the residual energy density would lead to the collapse of the universe, which is not observed. See also the discussion in the next section.

between the relations derived from the bootstrap does not generally lead to a violation of the bootstrap. Therefore, we believe that these parameters have independent sources of corrections.

Second, the bootstrap parameters have an order of magnitude of

\[
\hat{\mu}_0 \sim \left( \frac{\Lambda}{\tilde{m}_{pl}} \right)^2, \quad \hat{\lambda}_0 \sim \mathcal{O}(1) + \left( \frac{\Lambda}{\tilde{m}_{pl}} \right)^4 , \tag{24}
\]

and, consequently, the leading corrections can be written as

\[
\delta \hat{\mu}_0 \sim \left( \frac{\Lambda}{\tilde{m}_{pl}} \right)^{2+\tilde{q}}, \quad \delta \hat{\lambda}_0 \sim \left( \frac{\Lambda}{\tilde{m}_{pl}} \right)^{4+\tilde{q}} , \tag{25}
\]

with integer degrees of \( \tilde{q}, \tilde{\lambda} \geq 0 \). Then the bootstrap violation according to (15) is explained by nonzero values of

\[
\hat{\mu}^2 = -2\hat{\mu}_0 \delta \hat{\mu}_0 + 2\hat{\lambda}_0 \delta \hat{\lambda}_0 + \frac{3}{2} \hat{\lambda}_0 \delta \hat{\mu}_0 , \tag{26}
\]

so that

\[
\hat{\mu}^2 \sim \left( \frac{\Lambda}{\tilde{m}_{pl}} \right)^q + \mathcal{O}(1) \left( \frac{\Lambda}{\tilde{m}_{pl}} \right)^{4+\tilde{q}} , \tag{27}
\]

Then, the minimum of the potential \( \tilde{V} \) at the point \( \hat{\phi}^2 = \hat{\mu}^2 / \hat{\lambda} \) has a value

\[
\tilde{V}_{\text{min}} = 1 - \frac{\hat{\mu}^4}{4\hat{\lambda}} . \tag{28}
\]
and the relaxation of the cosmological constant, i.e., a
decrease in the contributions to the vacuum energy
density on the order of $\Lambda^4$, is possible only if
\begin{equation}
\frac{\hat{\mu}^4}{4\lambda} \sim \mathcal{O}(1). \tag{29}
\end{equation}
Condition (29) means that, in the case $q \geq \tilde{q}$,
\begin{equation}
2q = 4 + \tilde{q}, \tag{30}
\end{equation}
hence, there is a finite set of values for the degrees of the
corrections
\begin{equation}
\{ q, \tilde{q} \} \mapsto \{ 2, 0 \}, \{ 3, 2 \}, \{ 4, 4 \}. \tag{31}
\end{equation}
If $q < \tilde{q}$, there is a solution $q = 4$ with an arbitrary value
of $\tilde{q} \geq 5$. In all these cases, the scaling behavior of the
parameters of the potential can be reduced to
\begin{equation}
\hat{\mu}^4 \sim \left( \frac{\Lambda}{m_{\text{Pl}}} \right)^q, \quad \hat{\lambda} \sim \left( \frac{\Lambda}{m_{\text{Pl}}} \right)^{2q}, \tag{32}
\end{equation}
where $q = \{ 2, 3, 4 \}$.

As a result, we have a general phenomenological
description of the corrections violating the cosmological
bootstrap. We recall here that the corrections to
the parameter $\sigma^2_{\delta}$ can be reduced to the variation of
$\delta V_{\text{min}} \sim \mathcal{O}(1)$ so that the contributions to the vacuum
density on the order of $\Lambda^4$ are reduced completely. In other words, the condition $\hat{V}_{\text{min}} = 0$ is fulfilled with the necessary accuracy. This requirement is
quite natural from the standpoint of bootstrap design,
because a “survival” of the contributions to the vacuum
density on the order of $\Lambda^4$ would lead to the need for the introduction of a secondary module field. Therefore, if we consider the final physical field of the module, the reduction of the cosmological term
in the bootstrap violation is the very determination of this field.

Thus, we assume that a violation of the cosmological
bootstrap for the module field leads to a scaling potential
\begin{equation}
\hat{V} = \left( 1 - \frac{\hat{\mu}^2}{\hat{\lambda} \phi^2} \right)^2 . \tag{33}
\end{equation}
Hence, for a real physical field $\phi = \sqrt{2}|\Phi|$, the mass
and the self-interaction constant are
\begin{equation}
m^2 = \frac{32\pi G}{3} \Lambda^4 \mu^2, \quad \lambda = \left( \frac{8\pi G}{3} \Lambda^2 \right)^2 \mu^4, \tag{34}
\end{equation}
and in the order of magnitude
\begin{equation}
m \sim \bar{m}_{\text{Pl}} \left( \frac{\Lambda}{m_{\text{Pl}}} \right)^2 q/2, \quad \lambda \sim \left( \frac{\Lambda}{m_{\text{Pl}}} \right)^4 q^2, \tag{35}
\end{equation}
while the vacuum expectation value of the field
\begin{equation}
\nu = \frac{m}{\sqrt{2} \bar{m}_{\text{Pl}}} \sim \bar{m}_{\text{Pl}} \left( \frac{m_{\text{Pl}}}{\bar{m}_{\text{Pl}}} \right)^{q/2}. \tag{36}
\end{equation}
The inflaton mass is rather rigidly fixed by observational data; hence, assuming that $m/m_{\text{Pl}} \sim 10^{-5}$, we
find typical values of the model parameters by the
order of magnitude (see the table).

An analysis of the WMAP observational data over five years [7] in the inflaton model with a Higgs-type potential [56] also leads to a constraint on the vacuum expectation of the field:

(i) in a hilltop inflation scenario (the field slides
down to a minimum from a state that is close to an
unstable maximum at $\phi = 0$):
\begin{equation}
\frac{\nu}{m_{\text{Pl}}} \geq 10, \tag{37}
\end{equation}
(ii) in a chaotic inflation scenario (the field evolves
to a minimum from $|\phi| > \nu$):
\begin{equation}
\frac{\nu}{m_{\text{Pl}}} \geq 100, \tag{38}
\end{equation}
so that $q = 2$ is considered preferable among the
admissible values of the parameter characterizing the
contribution of the corrections, since large vacuum
expectation values correspond to a situation whereby
the potential degenerates within the limit $V \sim \phi^2$,
which is actually located on the edge of the confidence
domain with an interval of $1\sigma$. For $q = 2$, the scale of
the original cosmological constant is $\Lambda \sim 5 \times 10^{16}$ GeV,
which is essentially consistent with the hypothesis
about a scale corresponding to the grand unification
of gauge interactions.

This observation is further confirmed by the
WMAP observational data over seven years [9], which
already exclude the inflation scenario whereby the
field slides to a minimum from a state with a field value
that is larger than its vacuum expectation value $|\phi| > \nu$
at a confidence level of $1\sigma$, and establishes that
\begin{table}
\begin{tabular}{|c|c|c|c|}
\hline
$q$ & $\Lambda/m_{\text{Pl}}, 10^{-2}$ & $\lambda$ & $\nu/m_{\text{Pl}}$ \\
\hline
2 & 2 & $4 \times 10^{-14}$ & 50 \\
3 & 4 & $5 \times 10^{-14}$ & 150 \\
4 & 6 & $10^{-15}$ & 300 \\
\hline
\end{tabular}
\end{table}

\footnote{We proceed from an analysis that we carried out by the method
of attractor drift in the phase plane, which is used to describe the
inflationary dynamics [56–58].}
(i) the mass of the inflaton field is

$$m_{\text{inf}} \approx (1.30 - 1.74) \times 10^{13} \text{ GeV},$$  \hspace{1cm} (39)

(ii) the vacuum expectation value of the inflaton $$\langle \phi \rangle = v$$ is

$$2.5 m_{\text{Pl}} < v < 54 m_{\text{Pl}},$$  \hspace{1cm} (40)

where $$m_{\text{Pl}} = \sqrt{8 \pi m_{\text{Pl}}}$$ is the Planck mass;

(iii) a new scenario of chaotic inflation—hilltop inflation (the field slides down from a plateau on the hill of the potential near the zero field value in the direction of the potential minimum)—is achieved, with the plateau height being given by the parameter $$V(0) = V_{\text{hill}} = \Lambda_{\text{hill}}^4$$.

$$\Lambda_{\text{hill}} = (1.2 - 6.0) \times 10^{16} \text{ GeV}. \hspace{1cm} (41)$$

An analysis of the parameters is presented in Fig. 2, whence it follows that the preferred value of $$q = 2$$.

Thus, in a viable model for corrections, the parameters of the original superpotential are expected to shift according to

$$\delta \mu_0^2 \sim \mu_0^2 \left( \frac{\Lambda}{m_{\text{Pl}}} \right)^2, \hspace{0.5cm} \delta g_0 \sim \pm \left( \frac{\Lambda}{m_{\text{Pl}}} \right)^6,$$  \hspace{1cm} (42)

where the sign for $$\delta g_0$$ corresponds to a stable potential $$(\lambda > 0)$$ for sets I and II, respectively.

### 3. ONE-LOOP STRUCTURE

In the simplest case, after the supersymmetry violation, the model has one light scalar real field, with the term light meaning that the field mass is significantly smaller than the scale of the original cosmological constant:

$$m \ll \Lambda.$$  

Moreover, light fields comprise the massless graviton and the gravitino whose mass is specified in the leading approximation by the formula [54]

$$m_2 = \frac{8\pi G}{3} \Lambda^4 \ll \Lambda^2.$$  \hspace{1cm} (43)

In the low-energy approximation, the heavy fields of inflatino and the imaginary part of the scalar field do not propagate. The propagators and the interaction vertices of the inflaton are given in Appendix C. They are determined by the nontrivial superpotential and the chiral superfield supercurrent.

Then the one-loop contributions to the effective potential of the inflaton field are a result of the loops of (1) the field itself; (2) the gravitino of the inflatino supercurrent, whose propagator is reduced to a constant at low energies; and (3) the gravitino.

In this section, we analyze a model with a chiral superfield and regularization in Euclidean space by introducing a loop momentum cutoff.

#### 3.1. Generation of $$\mu_0$$

It is interesting to note that the contraction of the inflatino propagator to a point at low energies due to the Planck scale of the inflatino mass $$m'$$ leads to the natural introduction of a mass parameter in the superpotential $$\mu_0$$, which is initially equal to zero. In fact, the vertices of interaction of a real scalar field with inflatino due to the self-interaction constant $$g_0$$ and that of inflatino with gravitino due to vacuum energy with a scale of $$\Lambda$$ can be effectively reduced to the vertex of interaction of a scalar field with inflatino and gravitino (Fig. 3), i.e., to the generation of $$\mu_0$$ in the form

$$\mu_0 = 2 g_0 \Lambda \frac{\Lambda}{m' \cdot m}.$$  \hspace{1cm} (44)

Considering the scale of the inflatino mass $$m' \sim m_{\text{Pl}}$$, we see that a situation that corresponds to the cosmological bootstrap is achieved, with the potential near the zero field value in the direction of the potential minimum—is achieved, with the plateau height being given by the parameter $$V(0) = V_{\text{hill}} = \Lambda_{\text{hill}}^4$$.

$$\Lambda_{\text{hill}} = (1.2 - 6.0) \times 10^{16} \text{ GeV}. \hspace{1cm} (41)$$

An analysis of the parameters is presented in Fig. 2, whence it follows that the preferred value of $$q = 2$$.

Thus, in a viable model for corrections, the parameters of the original superpotential are expected to shift according to

$$\delta \mu_0^2 \sim \mu_0^2 \left( \frac{\Lambda}{m_{\text{Pl}}} \right)^2, \hspace{0.5cm} \delta g_0 \sim \pm \left( \frac{\Lambda}{m_{\text{Pl}}} \right)^6,$$  \hspace{1cm} (42)

where the sign for $$\delta g_0$$ corresponds to a stable potential $$(\lambda > 0)$$ for sets I and II, respectively.

### 3. ONE-LOOP STRUCTURE

In the simplest case, after the supersymmetry violation, the model has one light scalar real field, with the term light meaning that the field mass is significantly smaller than the scale of the original cosmological constant:

$$m \ll \Lambda.$$  

Moreover, light fields comprise the massless graviton and the gravitino whose mass is specified in the leading approximation by the formula [54]

$$m_2 = \frac{8\pi G}{3} \Lambda^4 \ll \Lambda^2.$$  \hspace{1cm} (43)

In the low-energy approximation, the heavy fields of inflatino and the imaginary part of the scalar field do not propagate. The propagators and the interaction vertices of the inflaton are given in Appendix C. They are determined by the nontrivial superpotential and the chiral superfield supercurrent.

Then the one-loop contributions to the effective potential of the inflaton field are a result of the loops of (1) the field itself; (2) the gravitino of the inflatino supercurrent, whose propagator is reduced to a constant at low energies; and (3) the gravitino.

In this section, we analyze a model with a chiral superfield and regularization in Euclidean space by introducing a loop momentum cutoff.

#### 3.1. Generation of $$\mu_0$$

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Thus, in a viable model for corrections, the parameters of the original superpotential are expected to shift according to

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where the sign for $$\delta g_0$$ corresponds to a stable potential $$(\lambda > 0)$$ for sets I and II, respectively.
logical bootstrap is one whereby \( g_0 \sim (\Lambda/\tilde{m}_{Pl})^2 \), and consequently, \( \mu_0 \sim \tilde{m}_{Pl}(\Lambda/\tilde{m}_{Pl})^4 \), as it should be in a solution with parameter set II.

But then the original bare superpotential with a zero value of \( \mu_0 \) leads to the potential of \( V_\beta = \Lambda^4 (\phi^2 - 1)^2 \) (see (21)), i.e., corresponds to a zero vacuum energy. It is clear that the generation of an effective mass parameter \( \delta \mu_0 \) at low energies, as described above, occurs at the tree level and, therefore, does not change the vacuum energy, which is equal to zero.

As a result of the introduction of corrections, the variation of the bare value (44) has two sources, \( \delta g_0 \) and \( \delta \Lambda^2 \), so

\[
\delta \mu_0 = \mu_0 \left( \frac{\delta g_0}{g_0} + \frac{\delta \Lambda^2}{\Lambda^2} \right) \approx \mu_0 \frac{\delta \Lambda^2}{\Lambda^2},
\]

where the approximation is made in accordance with scheme II, when the relative contribution of the variation in the self-interaction constant \( g_0 \) is suppressed. Hence, from (42) and (45), we should expect that the correction will have the form

\[
\frac{\delta \Lambda^2}{\Lambda^2} \sim -\left( \frac{\Lambda}{m_{Pl}} \right)^2.
\]

We emphasize that, in the case of the cosmological bootstrap, relation (44) indicates a connection between the inflatino mass and the primary Planck scale vacuum energy density, since the parameter \( \mu_0 \) with set II is specified by the value of the parameter \( \omega_0 \) (see (18)) and its value in the first step of the bootstrap procedure (see (11)), i.e., by the reduction of Planck contributions to the vacuum energy. As a result, we find within an accuracy of small corrections that

\[
m' = \frac{16}{5\sqrt{6}} M^2 \sqrt{\pi G} \sim \tilde{m}_{Pl}.
\]

This relationship meets the requirement of the existence of a flat potential of the module field, and we consider it not a condition of fine tuning of the model parameters, but rather a definition of the module field as such, i.e., the starting point of our model.

### 3.2. Field Zero-Point Modes and Regularization

A free scalar real field has a canonical energy-momentum tensor

\[
T_{\mu\nu} = \partial_\mu \phi \partial_\nu \phi - \frac{1}{2} g_{\mu\nu} (\partial_\alpha \phi \partial^\alpha \phi - m^2 \phi^2),
\]

and the zero-point modes of the field contribute to the imaginary part of the propagator in the loop (Fig. 4):

\[
\langle T_{\mu\nu} \rangle_0 = \int \frac{d^4 p}{(2\pi)^4} \left( p_\mu p_\nu - \frac{1}{2} g_{\mu\nu} (p^2 - m^2) \right) i^2 \times \operatorname{Im} \frac{1}{p^2 - m^2 + i0}.
\]

Considering that

\[
\operatorname{Im} \frac{1}{p^2 - m^2 + i0} = -\pi \delta (p^2 - m^2),
\]

\[(p^2 - m^2) \delta (p^2 - m^2) = 0,
\]

we find that the contribution of the energy-momentum tensor to the vacuum expectation value due to the field’s Lagrangian with a metric factor vanishes and the expression

\[
\langle T_{\mu\nu} \rangle_0 = \int \frac{d^4 p}{(2\pi)^4} p_\mu p_\nu \pi \delta (p^2 - m^2),
\]

in the Minkowski space is reduced to the standard contribution of field zero-point modes after the

9 In general, of course, this contribution gives a classical expression, i.e., the value of the potential at the minimum, which is zero for a free field.
removal of the delta function by integrating over the zero component of the momentum

\[ \langle T_{\mu\nu}\rangle_0 = \pi \int \frac{d^3p}{(2\pi)^3} \frac{p_\mu p_\nu}{|p|^4}; \] 

(51)

This formula makes sense after the introduction of regularization, e.g., by cutting the integral over the modulus of a three-dimensional momentum on the scale \( \Lambda_M \). Moreover, by virtue of spherical symmetry, the resulting averaged energy-momentum tensor corresponds to ultrarelativistic matter if the cutoff is significantly larger than the field mass: \( m \ll \Lambda_M \), or non-relativistic matter (or “dust”) if the mass is substantially larger than the cutoff: \( m \gg \Lambda_M \). Indeed,

\[ \langle T_{00}\rangle_0 = \int \frac{d^3p}{(2\pi)^3} \frac{1}{2} |p|^2, \quad \langle T_{0a}\rangle_0 = \langle T_{a0}\rangle_0 = 0, \]
\[ \langle T_{ab}\rangle_0 = \int \frac{d^3p}{(2\pi)^3} \frac{p_a p_b}{|p|^2} = \frac{1}{3} \delta_{ab} \int \frac{d^3p}{(2\pi)^3} \frac{p^2}{|p|^2}. \]

However, this result appears unphysical, since zero-point modes have an energy-momentum tensor that is different from the vacuum tensor, i.e., proportional to the metric tensor. This indicates that the regularization should satisfy the physical requirements (see [59]). A trivial condition is the introduction of a normal ordering of operators into the definition of energy-momentum tensor, which simply reduces the contribution of zero-point modes to zero. We follow a different path, introducing regularization in Euclidean space after the Wick rotation \( p_0 = ip_\gamma \); as a result, the expression for the contribution of zero-point modes to the averaged energy-momentum tensor is

\[ \langle T_{\mu\nu}\rangle_0^E = \frac{i d^4p}{(2\pi)^4} \frac{-ip_\mu p_\nu}{p^2 + m^2}, \]
\[ = \frac{1}{4} g_{\mu\nu} \int d^4p \frac{p^2}{(2\pi)^4} \frac{1}{p^2 + m^2}; \] 

(52)

which considers the spherical symmetry of Euclidean space. As a result, the contribution of one real scalar vacuum mode yields

\[ \langle T_{\mu\nu}\rangle_0^E = -g_{\mu\nu} \frac{1}{(16\pi)^2} \left[ \frac{1}{2} p_\mu^2 - m^2 p_\mu^2 + m^4 \ln \frac{p_\mu^2 + m^2}{m^2} \right] \Lambda_u^2, \]
\[ \Lambda_d^2 \] 

(53)

where \( \Lambda_u, d \) denote the upper and lower boundaries of the cutoff with respect to the absolute value of the Euclidean momentum, respectively.

For a Majorana fermion with an energy-momentum tensor

\[ T_{\mu\nu}^E = \frac{1}{2} \bar{\psi} \gamma_\mu \gamma_\nu \psi - \frac{1}{2} g_{\mu\nu} \bar{\psi} (\not{p} - m') \psi, \]

(54)

a similar procedure considering the elementary calculation of the trace of Dirac gamma matrices \( \text{tr}[\gamma_i (\not{p} + m')] = 4p_\gamma \) and the negative sign for the fermion loop leads to the following expression for the contribution of zero-point modes:

\[ \langle T_{\mu\nu}\rangle_0^E = g_{\mu\nu} \frac{2}{(16\pi)^2} \times \left\{ \frac{1}{2} p_\mu^2 - m^2 p_\mu^2 + m^4 \ln \frac{p_\mu^2 + m^2}{m^2} \right\} \Lambda_u^2, \]
\[ \Lambda_d^2 \] 

(55)

where factor 2 (when compared with the scalar field contribution) corresponds to the consideration of two—left and right—modes of the Majorana particle.

The summation of the contributions of the chiral superfield components to the vacuum energy-momentum tensor yields zero in the case of exact supersymmetry (all masses in the supermultiplet are equal). If supersymmetry is violated, considering the rules of sum

\[ \sum (-1)^F = 0, \quad \text{and} \quad \sum (-1)^F m^2 = 0, \]

10 results in the expression

\[ \langle T_{\mu\nu}\rangle_0^E = g_{\mu\nu} \frac{1}{(16\pi)^2} \times \left\{ m^2 m^2 \ln \frac{p_\mu^2 + m^2}{m^2} - m^4 \ln \frac{p_\mu^2 + m^2}{m^2} \right\} \Lambda_u^2, \]
\[ \Lambda_d^2 \] 

(56)

which we neglect the contribution of the real scalar field, assuming its mass to be zero, but then it follows from the rules of sum that \( \tilde{m}^2 = 2m^2 \) and, finally,

\[ \langle T_{\mu\nu}\rangle_0^E = g_{\mu\nu} \frac{1}{(16\pi)^2} \left\{ m^2 m^2 \ln \frac{p_\mu^2 + m^2}{m^2} - m^4 \ln \frac{p_\mu^2 + m^2}{m^2} \right\} \Lambda_u^2, \]
\[ \Lambda_d^2 \] 

(57)

\[ = g_{\mu\nu} \frac{2 m^4}{(16\pi)^2} \ln \frac{1 + \frac{p_\mu^2}{m^2}}{2 m^2} \Lambda_u^2. \]

10Here, \( \tilde{m} \) is the mass of the imaginary part of the scalar superfield.
Since the logarithm in (57) takes negative values, the following can be inferred:

(i) If the upper limit of integration takes values on the order of the Planck scale, \( \Lambda_u \sim \tilde{m}_{p1} \), then it makes a negative contribution to the vacuum energy density, which corresponds to the reduction of the initial bare value of vacuum energy \( M^0 \) by means of introducing the parameter \( \omega_0 \) in the superpotential;

(ii) If the lower limit of integration is \( \Lambda_d \ll m' \), then we have a limit

\[
\langle T^{(5)}_{\mu\nu} \rangle_u = g_{\mu\nu} \frac{1}{(16\pi)^2} \frac{1}{2} \Lambda_d^4 \sim g_{\mu\nu} \Lambda^4, \tag{58}
\]

which corresponds to the introduction of vacuum energy on the order of \( \Lambda^4 \) at \( \Lambda_d \sim \Lambda \).

Hence, the integration over the square of the Euclidean momentum within \( [\Lambda_E^2, \tilde{m}_{p1}^2] \) at \( \Lambda_E \sim \Lambda \) and \( \tilde{m}_{p1} \sim m \) corresponds to the formation of the initial conditions for the introduction of a chiral superfield for the module field. The integration interval \( [0, \Lambda_E] \), i.e., the regularization in Euclidean space with a cutoff \( \Lambda_E \), corresponds to the low-energy contribution of light fields. It is these integrals over loops for the calculation of the effective potential that will be considered in our further calculations.

Then, the correction to the vacuum energy-momentum tensor due to zero-point modes takes the form

\[
\delta T^E_{\mu\nu} = -g_{\mu\nu} \frac{1}{(16\pi)^2} \frac{1}{2} \Lambda_E^4. \tag{59}
\]

It is important to note that the cutoff with splitting of the Planck integration interval that we carried out by introducing an intermediate point \( \Lambda_E \) indicates an exact reduction of vacuum energy on the order of \( \Lambda^4 \)!

Furthermore, it is clear that the inclusion of the interaction of fields with constants in the form of power corrections with respect to \( \Lambda/\tilde{m}_{p1} \) does not change the statements that we have made.

### 3.3. Generation of \( g_0 \)

If the origin of the mass parameter \( \mu_0 \) is completely naturally explained by introducing an effective constant given a compression of the propagator in the presence of field self-interaction, the origin of the self-interaction with the constant \( g_0 \), which corresponds to the cosmological bootstrap, is not trivial, since it would mean that the module field has been in existence from the beginning. However, this question can be somewhat clarified by emphasizing that, provided that the interaction with graviton is considered, the modification of the canonical energy-momentum tensor of a free scalar real field according to the formula

\[
T^\text{mod}_{\mu\nu} = \phi_p \phi_p \phi - \frac{1}{2} g_{\mu\nu} \phi (p^2 - m^2) \phi, \tag{60}
\]

leads to a quadratic term of the interaction of the scalar field in the effective potential (diagram in Fig. 5).\(^{12}\)

Indeed, the calculation of the loop in Euclidean space with leading accuracy\(^{13}\) yields

\[
-i 2 g_0 \Lambda^2 = \int \frac{i d^d p_E}{(2\pi)^d} 4\pi G \Rightarrow g_0 = -\frac{G c_g \Lambda^4_E}{16\pi} \Lambda^2, \tag{61}
\]

where we introduced a dimensionless constant \( c_g \) to parameterize the arbitrariness in the choice of the ultraviolet cutoff according to the renormalization procedure and we assume here that \( c_g \) is on the order of unity. Consequently, in the first place, the resulting initial value of \( g_0 \) has a correct order of smallness with respect to \( \Lambda/\tilde{m}_{p1} \) and, second, it assumes, in the leading approximation, a value that is necessary for the bootstrap if the cutoff is put equal to

\[
c_g \Lambda^4_E = \frac{2}{3} (16\pi \Lambda^2)^2 \Rightarrow g_0 = -\frac{32\pi G \Lambda^2}{3}. \tag{62}
\]

In this case, the correction to vacuum energy due to the loop of the zero-point modes takes the form

\[
\delta T^E_{\mu\nu} = \frac{-1}{3 c_g} g_{\mu\nu} \Lambda^4, \tag{63}
\]

which indicates the need for a contribution of other fields in the vacuum energy on the order of \( \Lambda^4 \) if \( c_g \neq 1/3 \). It should be noted, however, that we proceeded from the assumption that the loop contributes only to the constant of the quadratic self-interaction of the field, leaving the field normalization unchanged.

---

\(^{11}\)This correction, of course, corresponds to the leading contribution of the scalar real field, which is the only chiral superfield component whose mass is much smaller than the cutoff; given this cutoff, “heavy” fields result in minor corrections to \( \delta T^E_{\mu\nu} \) on the order of \( (\Lambda/\tilde{m}_{p1})^2 \).

\(^{12}\)The momentum operator acts on one of the scalar fields, e.g., the right one. Of course, this can be rewritten in terms of partial derivatives, which act symmetrically twice, both on the left and right fields.

\(^{13}\)We neglect here the contribution of the mass term.
although, in the general case, one should consider the possibility of this renormalization by the power corrections $\Lambda/\tilde{m}_{\text{Pl}}$.

It follows from our study that the initial model parameters can be generated by the introduction of one-loop corrections involving gravitons and, hence, gravitino for a free scalar field if the cutoff is considered comparable in order of magnitude to the scale of the primary cosmological constant $\Lambda$. However, this implies that the loops involving the graviton and gravitino should not be considered when analyzing violations of the cosmological bootstrap due to loops with this cutoff, or this cutoff in the loops with the graviton and gravitino should be considered suppressed by the (even) degrees of $\Lambda/\tilde{m}_{\text{Pl}}$. Otherwise, we cannot draw a line between the concepts of the leading contribution and corrections. The same conclusion can be reached if we consider the contributions of the zero-point corrections. The same conclusion can be reached when one-loop contributions for corrections involving the cosmological bootstrap is meaningful precisely when one-loop contributions for corrections involving the graviton and gravitino have a cutoff with the above suppression factor with respect to $\Lambda/\tilde{m}_{\text{Pl}}$.

### 3.4. Inflaton Loops

The one-loop corrections involving the inflaton scalar real field only correspond to the contributions to the effective potential according to the diagrams shown in Fig. 6, i.e., to the corrections to the field self-interaction of the second and fourth orders. They lead to the following expressions for the amplitudes

$$-iL_1 = -\frac{3g_0}{16\pi^2}ie_2\Lambda_E^2,$$

$$-iL_2 = \frac{54g_0^4}{16\pi^2}\mu_0\frac{c_4\Lambda_E^2}{\Lambda_{\text{reg}}}$$

which correspond to corrections in the Lagrangian

$$\delta\mathcal{L} = \frac{1}{2}L_1\phi^2 + \frac{1}{4!}L_2\phi^4,$$

Note that expressions (64) introduce the normalization scale of the logarithmic corrections $\Lambda_{\text{reg}}$ and dimensionless constants $c_{2,4}$ that correspond to the variation in the cutoff parameter for various physical quantities in the regularization and renormalization.

Then $L_2$ ensures a standard renormalization of the constant $g_0$:

$$\delta g_0 = -\frac{9g_0^3}{32\pi^2}\ln\frac{c_4\Lambda_E^2}{\Lambda_{\text{reg}}} \Rightarrow \frac{dg_0}{d\ln\Lambda_{\text{reg}}} = \frac{9g_0^3}{16\pi^2},$$

which at $\Lambda_{\text{reg}} = \Lambda$ and $g_0 < 0$ leads to the change

$$\delta g_0 \sim \left(\frac{\Lambda}{\tilde{m}_{\text{Pl}}}\right)^6 > 0,$$

as might be expected in scenario II of the bootstrap model (see (42)).

Then, $L_1$ is naturally interpreted as a correction to the original value $2g_0\Lambda^2$, which determines the quadratic self-interaction of the field, and, since the variation of $g_0$, as we have just established, has a greater degree of smallness with respect to $\Lambda/\tilde{m}_{\text{Pl}}$, we should assume that

$$L_1 = 2g_0\delta\Lambda^2,$$

whence it follows that

$$\delta\Lambda^2 = \frac{3}{2}\frac{g_0}{16\pi^2}e_2\Lambda_E^2 < 0.$$

According to (45), we find

$$\delta\mu_0 \sim \frac{3}{2}\frac{g_0}{16\pi^2}c_2\Lambda_E^2 \sim -\left(\frac{\Lambda}{\tilde{m}_{\text{Pl}}}\right)^2 < 0.$$ (67)

Now we redefine the dimensionless values with “caps” in (14) by substituting $\Lambda^2 \mapsto \Lambda^2 \pm \delta\Lambda^2$. This procedure will give us the scalar field potential in the form

$$V = (\Lambda^2 \pm \delta\Lambda^2)\left(1 - \frac{\Lambda^2}{2\phi} + \frac{\Lambda^2}{4!}\phi^4\right),$$ (68)
where

\[
\mu^2 = -\frac{5}{2} \frac{\delta \mu_0}{\mu_0} = \frac{15}{16 \pi^2} \frac{g_0 \sigma_E^2 \Lambda^2}{\Lambda^2},
\]

\[
\lambda = 4 \frac{\delta g_0}{g_0} = -9 \frac{g_0^2}{4 \pi} \ln \frac{\Lambda E \sqrt{\Lambda_4}}{\Lambda_{\text{reg}}},
\]

(69)

Apparently, it is easy to select a quite reasonable value of the normalization point \( \Lambda_{\text{reg}} \) that would ensure a zero vacuum energy density: \( \mu^4 = 4 \lambda \). Under this condition, the normalization point is close to the vacuum value of the field:

\[
\langle \phi \rangle^2 = \frac{3}{16 \pi G} \frac{\langle \phi \rangle^2}{4 \pi G \mu^2} \sim \Lambda_{\text{reg}}^2.
\]

Thus, the contributions that we have just analyzed satisfy the natural condition for the cosmological bootstrap.

However, the scalar field loops also lead to corrections to the vertices of the contact interaction of the inflaton with gravitino and inflatino given the introduction of the supercurrent, with these loops being quite similar to those discussed above. However, it should be noted that the inflatino with a Planck scale mass is “frozen” at low energies and, hence, the external inflatino fields are zero; therefore, the scalar field loops at the aforementioned supercurrent vertices are not relevant to the contributions being studied, because the outer ends of inflatino in such diagrams lead to zero values of these corrections.

### 3.5. Graviton and Gravitino Loops

In order to find a correction that is quadratic with respect to the inflaton field, we consider the loop diagram involving a gravitino and an inflatino (Fig. 7) according to the Feynman rules (see Appendix 3).14

The diagram, the input and output momenta of the inflaton are assumed to be zero and the momentum circulating through the loop is assumed to be \( k \). Then, the contribution of the inflatino loop to the quadratic order self-interaction of the inflaton:

\[
\delta \mu^2 = \frac{\mu_0^2 (8 \pi G)}{4} \int \frac{d^4 k}{(2 \pi)^4} \text{Tr} \left\{ \frac{P(k)^{\mu \nu}}{k^2 + m_g^2} \frac{\gamma_\mu \gamma_5}{k^2 + m^2} \gamma_\nu \right\}.
\]

Here we introduce the constant \( c_m \) as a parameter of the ultraviolet cutoff. Since a gravitino propagates in the loop, \( c_m \ll 1 \); therefore, considering the expression for the gravitino mass, we find that this contribution is suppressed as

\[
\delta \mu^2 \sim c_m \left( \frac{\Lambda}{m_{pl}} \right)^2,
\]

and it is insignificant in the case of a violation of the cosmological bootstrap.

The correction to \( g_0 \) is given by two diagrams with zero input momenta and a circulating momentum of \( k \) (Fig. 8).

According to the first diagram in Fig. 8, we have

\[
g_0 \delta g_0 = \frac{3}{12} \frac{g_0^2 (8 \pi G)}{2} \int \frac{d^4 k}{(2 \pi)^4} \text{Tr} \left\{ \frac{P(k)^{\mu \nu}}{k^2 + m_g^2} \frac{\gamma_\mu \gamma_5}{k^2 + m^2} \gamma_\nu \right\}.
\]

(72)
From the second diagram in Fig. 8, we derive the following analytical expression:

$$g_0\delta g_0 = \frac{6}{12} \frac{\mu_0^4 (8\pi G)^2}{16} \int \frac{d^4 k}{(2\pi)^4} \text{Tr} \left\{ \frac{P(k)_{\mu\nu}}{k^2 + m_g^2} \gamma^5 \left( \frac{-i\gamma^\mu k + m^*}{k^2 + m^2} \right)^5 \right\} \times \gamma^\nu \delta \Lambda^2$$

(73)

$$\approx \frac{\mu_0^4 (8\pi G)^2 c_g^2 \Lambda^2}{8m^2} + \Theta\left( \frac{1}{m^3} \right) \sim -g_0 c_g^2 \left( \frac{\Lambda}{m_{Pl}} \right)^{12}.$$  

It follows from (72) that, in the presence of a gravitino in the loop, the cutoff constant should be suppressed as $c_g \sim (\Lambda/m_{Pl})^2$. Then the cosmological bootstrap model is self-consistent, but the inflatino contributes significantly to the variation in the field self-interaction, which is comparable to a logarithmic renormalization and might be playing a decisive role. This, however, does not change the order of magnitude of the previous numerical estimating calculations. It follows from (73) that the role of the constant $c_g$ is insignificant.

The corrections as a result of the interaction of the canonical energy-momentum tensor with the graviton are presented in the diagrams in Fig. 9.

The first diagram in Fig. 9 gives a logarithmic correction to the mass of the form

$$\delta \mu^2 = -(8\pi G)(\mu_0^2 + 2g_0\Lambda^2) \eta_{\mu\nu} \times (\eta^{\mu\nu} \eta^{\mu\nu} - \eta^{\mu\nu} \eta^{\mu\nu}) \eta_{\mu\nu}$$

$$\times \int \frac{d^4 k}{(2\pi)^4} \frac{1}{k^2 + m^2} = \frac{(8\pi G)(\mu_0^2 + 2g_0\Lambda^2)^2}{4\pi^2} \times \ln \frac{c_g^4 \Lambda^2}{\Lambda_{reg}^2} \sim \Lambda^2 \left( \frac{\Lambda}{m_{Pl}} \right)^6 \ln \frac{c_g^4 \Lambda^2}{\Lambda_{reg}^2},$$

and the second diagram gives a correction to the self-interaction constant

$$g_0\delta g_0 = \frac{6}{12} \frac{\mu_0^4 (8\pi G)^2}{16} \int \frac{d^4 k}{(2\pi)^4} \text{Tr} \left\{ \frac{P(k)_{\mu\nu}}{k^2 + m_g^2} \gamma^5 \left( \frac{-i\gamma^\mu k + m^*}{k^2 + m^2} \right)^5 \right\} \times \gamma^\nu \delta \Lambda^2$$

(74)

$$\times \ln \frac{c_g^4 \Lambda^2}{\Lambda_{reg}^2} \sim -g_0 \left( \frac{\Lambda}{m_{Pl}} \right)^{10}.$$  

The contribution in (74) can be attributed both to the suppressed logarithmic correction to the scale $\Lambda^2$ of the form $g_0\delta \Lambda^2$, which is insignificant, and to a similar variation in $\delta g_0$, which is comparable to the previously calculated inflaton-loop logarithmic correction. In this case, considering that $c_m \sim \Lambda/m_{Pl}$, we find that the logarithm argument has an additional degree of the minor parameter; nonetheless, it is still of the same order of magnitude as the renormalization group correction. In this case, the correction in (74) has the opposite sign, and the contributions may offset one another. This suggests that the main contribution appears to be specified in (72).

We infer from (75) that even at $c_g \sim (\Lambda/m_{Pl})^2$ this contribution is negligible for our consideration.

As a result, we have shown that the one-loop structure of the theory is consistent with the relations necessary for a reasonable violation of the cosmological bootstrap.

**CONCLUSIONS**

We have established the possibility of constructing a realistic model in which

(i) after a consideration of the leading corrections with respect to the gravitational constant in supergravity, the primary cosmological constant with a characteristic scale on the order of the grand unification energy corresponds to the flat potential of a real scalar module field that has a nontrivial superpotential;

(ii) the fine tuning of the superpotential parameters is violated due to quantum loop corrections, which leads to the instability of the original cosmological constant;
(iii) the primary cosmological constant is relaxed as a result of inflation caused by the instability of the potential of the module field, which acts as an inflaton;

(iv) after the consideration of the quantum loop corrections, the parameters of the potential are naturally consistent in order of magnitude with the values that are phenomenologically observed in the description of the large-scale structure of our universe.

The existence of a nontrivial primary superpotential of the module field is a result of a connection between its parameters, i.e., the cosmological bootstrap.

The constructed model has only two energy scales: the Planck mass \( \tilde{m}_{\text{Pl}} \) and the scale of the primary cosmological constant on the order of the grand unification \( \Lambda \). These energy scales are in a strict hierarchy: \( \Lambda/\tilde{m}_{\text{Pl}} \gg 1 \). This is sufficient for the realization of the cosmological bootstrap and a natural explanation for the real parameters of the inflaton.

In Section 3 we calculated the one-loop structure of the theory and formulated the conditions for the derivation of the necessary hierarchy of loop corrections for the cosmological bootstrap. These corrections violate the fine tuning of the initial superpotential: the parameters for the momentum cutoff in a loop with an inflaton must have a scale of \( \Lambda \), whereas the cutoff in the loops with a graviton, gravitino, or inflatino must have a diagram suppression factor on the order of \( (\Lambda/\tilde{m}_{\text{Pl}})^2 \).

We should point out some technical issues. First, we considered the corrections to the potential due to supergravity in the leading order with respect to the gravitational coupling constant. Second, our analysis was confined to the contributions to the renormalizable self-interaction terms of the scalar field, i.e., self-interaction of the fourth order or less with respect to the field. At the same time, we assumed that self-interaction of highest orders with respect to the inflaton is out of control, since these terms of the Lagrangian allow for an arbitrary finite renormalization. Therefore, we believed the latter to be negligible in our model. Third, the very procedure of fine tuning of the potential parameters—the cosmological bootstrap—also includes a fixed expansion of these parameters in the highest powers of the small ratio between the scale of the cosmological constant and the Planck mass \( \Lambda/\tilde{m}_{\text{Pl}} \). Therefore, the introduction of loop corrections with respect to \( \Lambda/\tilde{m}_{\text{Pl}} \) would destroy the structure of the cosmological bootstrap and create conditions for inflation. We can therefore conclude that, in general, the proposed pattern for the relationship between the primary cosmological constant and inflation seems quite plausible.

APPENDIX A

IMAGINARY PART OF THE SCALAR FIELD AND GENERAL FORM OF THE POTENTIAL

Without any loss in generality, we can assume that the parameter \( \Lambda \) in superpotential (12) is real, while \( \mu_0 \) and \( g_0 \) should, generally speaking, be treated as complex. Then the potential has the form

\[
V_S = \left[ \frac{2W}{\partial \Phi} \right]^2 = \Lambda^2 + g_0 \mu_0 (\Phi^* \Phi)^2 + \mu_0 \mu_0^* \Phi \Phi^* + \Lambda^2 \left( g_0^* (\Phi^* \Phi)^2 + g_0 \mu_0^2 \right) + i \left\{ \mu_0 \Phi - \mu_0^* \Phi^* \right\} \tag{A1.1}
\]

+ \( i \Lambda^2 \Phi^* \Phi \{ \mu_0 g_0 \Phi^* - \mu_0^* g_0 \Phi \} \).

Therefore, a discrete symmetry under the inversion \( \Phi \leftrightarrow -\Phi^* \) leads immediately to the conditions of reality:

\[
\mu_0 = \mu_0^*, \quad g_0 = g_0^* \tag{A1.2}
\]

and the potential

\[
V_S = \Lambda^2 + \frac{1}{2} \{ \mu_0^2 + 2g_0^2 \Lambda^2 \} \phi^2 + \frac{1}{4} g_0^2 \phi^4 + \Delta V_S, \tag{A1.3}
\]

where the addition

\[
\Delta V_S = \mu_0 \{ g_0^2 \phi^2 - 2\Lambda^2 \} \phi^2 - 4g_0^2 \Lambda^2 \phi \tag{A1.4}
\]

in terms of the real fields

\[
\phi = \sqrt{2} |\Phi|, \quad \tilde{\phi} = \text{Im} \Phi, \tag{A1.5}
\]

and hence, there is a constraint

\[
|\tilde{\phi}| \leq \frac{1}{\sqrt{2}} \phi. \tag{A1.6}
\]

In the model under study, \( g_0 < 0 \); consequently, the addition \( \Delta V_S \) has a minimum with respect to \( \tilde{\phi} \) at a fixed absolute value of the field \( \Phi \), with the zero value of the imaginary part of the field being unstable.

The minimum with respect to \( \tilde{\phi} \) is attained at

\[
\tilde{\phi}_* = \mu_0 g_0 \phi^2 - 2\Lambda^2 g_0^2 \Lambda^2, \tag{A1.7}
\]

and, consequently, this field can be approximated by the constant

\[
\tilde{\phi}_* \approx -\frac{\mu_0}{4g_0}, \tag{A1.8}
\]

in the domain

\[
\phi^2 \leq \frac{\Lambda^2}{g_0}. \tag{A1.9}
\]

if \( \phi^2 > \mu_0^2/g_0^2 \). In the realistic model, \( g_0 \sim -\left( \Lambda/\tilde{m}_{\text{Pl}} \right)^6 \) and constraint (A1.9) can be reduced to \( \phi \ll
\[ \tilde{m}_p (\tilde{m}_p / \Lambda)^2, \] i.e., to the condition that the field is necessarily less than its vacuum expectation value.

The substitution of (A1.7) yields a potential for the field \( \phi \) of the form

\[ V_S \mapsto \frac{\lambda_0}{4} (\phi^2 - \phi_0^2)^2, \quad (A1.10) \]

where

\[ \phi_0 = \frac{2 \Lambda^2}{g_0}, \quad \lambda_0 = g_0 + \frac{\mu_0^2}{4 \lambda^2}, \quad (A1.11) \]

This potential has a zero value of vacuum energy, as should be in the case of one chiral superfield when there is always a complex-valued solution of the quadratic equation \( |\Phi|^2 = \text{const} \) for the field \( \Phi \), which is the same, of course, as \( |\Phi| = |\phi| \).

As a result, we inferred that, in the case of a real cosmological role of the scalar field, the phenomenological requirement of “elimination” of its imaginary part from the dynamics means the introduction of a supersymmetry-violating term that would ensure that \( \bar{\phi} \to 0 \). This can be achieved, for example, by adding a “mass term” of the form

\[ \Delta \tilde{V} = \tilde{m}^2 (\bar{\phi} + C \mu_0 \Lambda^2/m^2)^2 + \text{const}, \quad (A1.12) \]

where the mass takes values on the order of the Planck mass \( \tilde{m} \sim \tilde{m}_p \). The first variant of \( \Delta \tilde{V} \) with \( C = 0 \) contains no fine-tuning parameters and yields a negligible value of the imaginary part of the scalar field \( \phi_* \sim \mu_0 \Lambda^2/m^2 \to 0 \), which is weakly dependent on the field \( \phi \). However, the second variant with \( C = 1 \) leads to a reduction in the potential term that is linear with respect to the field \( \phi \) and, thus, independent of the real part of the field and to an even greater suppression of the imaginary part. Thus, the second variant significantly expands the scope of applicability of the approximation with a zero value of the field \( \phi \). Finally, in the third variant with \( C = 2 + \mu_0/(2g_0 \Lambda^2) \), the imaginary part of the field is zero in a vacuum, i.e., at the minimum of the potential for the real part of the field at \( \phi_* = -2 \Lambda^2/g_0 - \mu_0^2/4 \lambda^2 > 0 \); therefore, in this phenomenologically relevant case, the vacuum is invariant under the operation of complex conjugation of the field \( \Phi \). The dynamics of the imaginary part of the field becomes completely insignificant at energies much below the Planck scale, i.e., in the classical description of gravity.

Thus, the inclusion of the scalar field in the cosmological model truly allows us to consider this field as real.

Finally, we note that the introduction of mass for the imaginary part of the scalar field with a supersymmetry-violating term leads to the standard rules of sum for the squared masses of the chiral superfield components with the fermion number \( F = \{0,1\} \):

\[ \sum (-1)^F m^2 = 0 \Rightarrow m^2 + \tilde{m}^2 = 2 m^2, \quad (A1.13) \]

where \( m' \) is the mass of the scalar field’s superpartner, the Majorana field of inflatino. As a result, inflatino necessarily acquires a mass of the order of the Planck mass.

**APPENDIX B**

**CHIRAL ROTATION**

The superpartner of the scalar field—inflatino—is described by a Majorana spinor \( \psi \), i.e., a charge-self-adjoint Dirac spinor, whose left and right components in the chiral representation are

\[ \psi = \begin{pmatrix} \psi_L \cr \psi_R \end{pmatrix} = \begin{pmatrix} \chi \cr \bar{\chi} \end{pmatrix}, \quad (B2.1) \]

where the two-component spinors are charge conjugates: \( \bar{\psi} = i \sigma_2 \bar{\chi} \); hence, superpotential (12) leads to a Lagrangian

\[ \mathcal{L}_2 = -\frac{1}{2} (i \mu_0 \chi \bar{\chi} - i \mu_0 \bar{\chi} \chi + 2 g_0 \Phi \chi \bar{\chi} + 2 g_0 \Phi^* \bar{\chi} \bar{\chi}), \quad (B2.2) \]

with quadratic terms with respect to inflatino.

This Lagrangian can be reduced to the standard form with a real mass of the Majorana field using a chiral rotation:

\[ \psi_u = e^{\gamma_5 u} \psi \Rightarrow \chi_u = e^{iu} \chi, \quad \bar{\chi}_u = e^{-iu} \bar{\chi}, \quad (B2.3) \]

with the parameter \( u = \pi/4 \), so that

\[ \mathcal{L}_2 \rightarrow -\frac{1}{2} (\mu_0 \chi_u \bar{\chi}_u + \mu_0 \bar{\chi}_u \chi_u - 2 i g_0 \Phi \chi_u \bar{\chi}_u + 2 i g_0 \Phi^* \bar{\chi}_u \chi_u), \]

\[ = -\frac{1}{2} (\mu_0 \bar{\psi}_R \psi_L + \mu_0 \bar{\psi}_L \psi_R - 2 i g_0 \Phi \bar{\psi}_R \psi_L \]

\[ + 2 i g_0 \Phi^* \bar{\psi}_L \psi_R), \quad (B2.4) \]

\[ = -\frac{1}{2} (\mu_0 \bar{\psi} \psi - \sqrt{2} i g_0 \Phi \bar{\psi} \gamma_5 \psi + 2 g_0 \Phi \bar{\psi} \gamma_5 \psi), \quad (B2.6) \]

where we omitted for brevity the bispinors’ index \( u \), which is irrelevant for what follows.

Thus, after the chiral rotation, we have a Majorana field with well-defined peaks of the Yukawa interaction with real and imaginary scalar fields of the chiral superfield. It is important to note that the kinetic terms for the left and right spinor components are
invariant under the chiral rotation. Moreover, the contact terms of the interaction of the inflatino with the gravitino and the scalar field are also invariant under the chiral rotation. This is explained by the fact that gravitino is also described by a Majorana field and the contact interaction terms retain their chirality.

APPENDIX C
FEYNMAN RULES

Here we describe the rules of the diagrammatic technique. We use the definition of a metric and the relevant Dirac matrices according to Weinberg [54]:

\[ \eta_{\mu \nu} = \text{diag}(-1, +1, +1, +1), \]  

(C3.1)

We denote the inflaton propagator with a solid line and compare the expression

\[ (-i) \frac{1}{-p^2 + m^2}, \]  

(C3.2)

where \( m \) is the mass of the inflaton, which is assumed to be zero.

The inflatino is denoted with a double line. The analytical expression for the propagator has the form

\[ (-i) \frac{-i\gamma p + m'}{p^2 + m'^2}, \]  

(C3.3)

where \( m' \) is the mass of the inflatino.

The gravitino is denoted with a dashed line, and the propagator is

\[ (-i) \frac{\gamma^0}{p^2 + m^2} \left( \frac{1}{m^2} \right) \left( \eta^{\mu \nu} + i\frac{p^\mu p^\nu}{m^2} \right) \left( -i\gamma p + m_\pi \right), \]  

(C3.4)

\[ \quad \frac{1}{3} \left( \gamma^\mu - i\frac{p^\mu}{m_\pi} \right) \left( i\gamma p + m_\pi \right) \left( \gamma^\nu - i\frac{p^\nu}{m_\pi} \right), \]

where \( m_\pi = \sqrt{\frac{(8\pi G)}{12}} \Lambda^4 \) is the mass of the gravitino.

The expressions for the vertices appear in the consideration of supersymmetry: the gravitino interacts with the supercurrent through the Lagrangian term:

\[ \sqrt{8\pi G} \int d^4x \frac{1}{2} \bar{\psi}_\mu \psi_\mu, \]  

(C3.5)

where \( \psi_\mu \) is the gravitino field, and the supercurrent

\[ S^\mu = \sqrt{2} \left[ \gamma^\nu \partial_\nu \phi \psi_R + \bar{\gamma}^\nu \partial_\nu \phi^* \gamma^\mu \psi_L \right. \]

\[ \left. + \left( \frac{\partial W}{\partial \phi} \right)^* \gamma^\mu \psi_L + \left( \frac{\partial W}{\partial \phi} \right)^* \gamma^\mu \psi_R \right]. \]

with \( \phi \) standing for the inflaton, \( \psi \) for the inflatino, and \( W \) for the superpotential.

The Lagrangian also has potential terms responsible for the self-interaction of the inflaton field and the interaction between the inflaton and the inflatino; these terms have the form

\[ \int dx \frac{1}{2} \frac{\partial^2 W}{\partial \phi^2} (\bar{\psi}_L \psi_L) + \frac{1}{2} \left( \frac{\partial^2 W}{\partial \phi^2} \right)^* (\bar{\psi}_L \psi_L)^* \]  

(C3.6)

\[ + \left( \frac{\partial W}{\partial \phi} \right) \left( \frac{\partial W}{\partial \phi} \right)^*. \]

In theory, there are three types of vertices ensuing from the supercurrent:

(i) The first is the point of convergence of three lines—the inflaton, the inflatino, and the gravitino, respectively; the vertex and the analytical expression are as follows:

\[ (i) (i) \frac{1}{2} \mu_0 \sqrt{8\pi G} \gamma^5 = \frac{1}{2} \mu_0 \sqrt{8\pi G} \gamma^5, \]  

(C3.7)

where \( \gamma^5 = i\gamma^0 \gamma^1 \gamma^2 \gamma^3 \).

(ii) The second type is possible as a result of the consideration of a self-interacting field; four lines converge here: inflatino, gravitino and two inflatons; the vertex and the analytical expression are of the form

\[ (ii) \frac{1}{\sqrt{2}} \mu_0 \sqrt{8\pi G} \gamma^5, \]  

(C3.8)

(iii) In the third type, the inflatino transforms into a gravitino:

\[ (iii) \frac{1}{\sqrt{2}} \sqrt{8\pi G} \Lambda^4 \gamma^5. \]  

(C3.9)

The graviton propagator has the form

\[ G^{\mu\nu, \mu' \nu'} \]  

\[ = (-i) \frac{1}{2 \rho^2} (\eta^{\mu \nu} \eta^{\mu' \nu'} + \eta^{\mu \nu'} \eta^{\mu' \nu} - \eta^{\mu' \nu} \eta^{\mu \nu'}). \]  

(C3.10)
The vertex of interaction of the graviton with the scalar particle follows from the Lagrangian term
\[ \sqrt{8\pi G} \int d^4x T^{\mu\nu} h_{\mu\nu}, \]
where \( T^{\mu\nu} \) is the energy-momentum tensor, and the correction \( h_{\mu\nu} \) is specified as
\[ g_{\mu\nu} = \eta_{\mu\nu} + 2\sqrt{8\pi G} h_{\mu\nu} \]
and is equal to
\[ \eta_{\mu\nu}^2 + 2g_0^2\Lambda^2. \] (C3.11)

As a result of the field self-interaction of the fourth order with respect to the field, we have a vertex
\[ -i\lambda_0 = -6i\mu_0^2, \] (C3.12)
and the vertex of interaction of the inflaton field with the inflatino is derived from the Lagrangian term after the chiral rotation (Appendix B):
\[ -\sqrt{2g_0} \gamma^5. \] (C3.13)

The law of momentum conservation must hold at each vertex: the input momentum is equal to the output momentum. The momenta of the internal lines, which remain nonfixed after the consideration of all conservation laws, are used to perform integration by \( d^4p/(2\pi)^4 \). Each fermion loop adds the factor \((-1)\). The Wick rotation corresponds to the substitution \( p_0 \rightarrow ip_4 \) in the transition to Euclidean coordinates.

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