

Chaotic scalar fields as models for dark energy

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We consider stochastically quantized self-interacting scalar fields as suitable models to generate dark energy in the Universe. Second quantization effects lead to new and unexpected phenomena if the self-interaction strength is strong. The stochastically quantized dynamics can degenerate to a chaotic dynamics conjugated to a Bernoulli shift in fictitious time, and the right amount of vacuum energy density can be generated without fine-tuning. It is numerically observed that the scalar field dynamics distinguishes fundamental parameters such as the electroweak and strong coupling constants as corresponding to local minima in the dark energy landscape. Chaotic fields can offer possible solutions to the cosmological coincidence problem, as well as to the problem of uniqueness of vacua.

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I. INTRODUCTION

There is by now convincing observational evidence that the Universe is currently in a phase of accelerated expansion [1,2]. The favored explanation for this behavior is the existence of vacuum energy or, in a more general setting, of dark energy. The observations suggest that the Universe currently consist of approximately 73% dark energy, 23% dark matter, and 4% ordinary matter [3]. The nature and origin of the dominating dark energy component are not understood, and many different models co-exist. The simplest models associate dark energy with the vacuum energy of some unknown self-interacting scalar field, whose potential energy yields a cosmological constant [4]. In quintessence models slowly evolving scalar fields with a nontrivial equation of state are considered [5]. String theory also yields possible candidates of scalar fields that might generate dark energy, in the form of runaway dilatons and moduli fields [6]. Various exotic forms of matter such as phantom matter [7] and Born-Infeld quantum condensates [8] are currently being discussed. For some superstring cosmology ideas related to small cosmological constants, see also [9].

When trying to formulate a suitable model for dark energy, at least two unsolved fundamental problems arise.

(1) *The cosmological constant problem.* Why is the observed vacuum energy density so small, as compared to typical predictions of particle physics models? From electroweak symmetry breaking via the Higgs mechanism one obtains a vacuum energy density prediction that is too large by a factor 10^{55} as compared to the currently observed value. Spontaneous symmetry breaking in grand unified theory models is even worse, it yields a discrepancy by a factor of 10^{11} .

(2) *The cosmological coincidence problem.* Why is the order of magnitude of the currently observed vacuum energy density the same as that of the matter density? A true cosmological constant stays constant during the expansion of the

Universe, whereas the matter energy density decreases with a^{-3} , where $a(t)$ is the scale factor in the Robertson-Walker metric. It looks like a very strange coincidence that right now we live at an epoch where the vacuum energy density and matter density have the same order of magnitude, if during the evolution of the Universe one is constant and the other one decreases as $a(t)^{-3}$.

To this list one may add yet another fundamental problem which we may call the uniqueness problem.

(3) *The uniqueness problem.* String theory allows for an enormous amount of possible vacua after compactification. In each of these states the fundamental constants of nature can take on different possible values. But what is the mechanism that selects out of these infinitely many possibilities the physically relevant vacuum state, with its associated fundamental constants that give rise to a Universe of the type we know (that ultimately even enabled the development of life)? Relating the answer purely to an anthropic principle seems unsatisfactory.

In this paper we consider a new model for dark energy which, as compared to other models, is rather conservative. It just associates dark energy with self-interacting scalar fields corresponding to a φ^4 theory, which is second quantized. However, the fundamental difference from previous approaches is that these fields are very strongly (rather than weakly) self-interacting, and that 2nd quantization effects play an important role. We will use as the relevant method to quantize the scalar fields the stochastic quantization method introduced by Parisi and Wu [10]. In the fictitious time variable of this approach, the fields will turn out to perform rapid deterministic chaotic oscillations, due to the fact that we consider not a weakly but a very strongly self-interacting field. This chaotic behavior is a new effect not present in any classical treatment. It is generally well known that chaos plays an important role in general relativity [11], quantum field theories [12–14], and string theories [15]. The main result of our consideration is that the chaotic field theories considered naturally generate a small cosmological constant and have the scope to offer simultaneous solutions to the cosmological coincidence and uniqueness problems.

Our physical interpretation is to associate the chaotic behavior of the scalar fields with tiny vacuum fluctuations

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which are allowed within the bounds set by the uncertainty relation, due to the finite age of the Universe. This interpretation naturally leads to the right amount of dark energy density being generated, and fine-tuning can be avoided. The chaotic fields (presently) have a classical equation of state close to $w = -1$, and can thus account for the accelerated expansion of the Universe. However, during the early evolution of the Universe they behave in a different way: They effectively track radiation and matter. This property will help to avoid the cosmological coincidence problem.

The chaotic model also contains an interesting symmetry between gravitational and gauge couplings. In our model the role of a metric for the 5th coordinate (the fictitious time) is taken over by dimensionless coupling constants which are given by the ratio of the fictitious time lattice constant and physical time lattice constant squared (both lattice constants can still go to zero, just their ratio is fixed). These coupling constants do not occur in any classical treatment but are entirely a consequence of our second quantized treatment. The vacuum energy generated depends on these couplings in a non-trivial way. The physical significance of our model is illustrated by the fact that we numerically observe the vacuum energy to have local minima for coupling constants that numerically coincide with running electroweak coupling strengths, evaluated at the known fermionic mass scales, as well as running strong coupling constants evaluated at the known bosonic mass scales. This numerical observation, previously reported in [14], is now embedded into a cosmological context. The role of the chaotic fields in the Universe can be understood in the sense that they are responsible for fixing and stabilizing fundamental parameters as local minima in the dark energy landscape. This is somewhat similar to the role the dilaton field plays in string theory after supersymmetry breaking.

Our numerical discovery of local minima that coincide with known standard model coupling constants makes it very unlikely that there are different Universes with different fundamental parameters. In fact, the numerical results provide strong evidence that there is a unique vacuum state of the Universe that possesses minimum vacuum energy precisely for the known set of standard model parameters.

This paper is organized as follows. In Sec. II we show how a second-quantized scalar field dynamics can degenerate to a chaotic dynamics in fictitious time. Our main example is a chaotic φ^4 theory leading to 3rd order Tschebyscheff maps, which is dealt with in Sec. III. In Sec. IV we present a physical interpretation of the chaotic dynamics using the uncertainty relation, which in a natural way fixes the order of magnitude of the vacuum energy density to be generated. Section V deals with the energy, pressure and classical equation of state of the chaotic fields. In Sec. VI we consider the Einstein equations associated with our model and discuss a possible way to avoid the cosmological coincidence problem. Section VII yields a prediction for the current ratio of matter energy density and dark energy density to the critical energy density. In Sec. VIII we describe how local minima of the dark energy landscape generated by the chaotic fields can fix the fundamental parameters. Finally, in Sec. IX we dis-

cuss spontaneous symmetry breaking phenomena for the chaotic fields.

II. STOCHASTIC QUANTIZATION OF STRONGLY SELF-INTERACTING SCALAR FIELDS

Let us consider a self-interacting scalar field φ in the Robertson-Walker metric. For a complete theory describing all quantum mechanical fluctuations we need to second-quantize it. This can be done via stochastic quantization. In the Parisi-Wu approach of stochastic quantization one considers a stochastic differential equation evolving in a fictitious time variable s , the drift term being given by the classical field equation [10]. Quantum mechanical expectations correspond to expectations with respect to the generated stochastic processes in the limit $s \rightarrow \infty$. The fictitious time s is different from the physical time t ; it is just a helpful fifth coordinate to do the 2nd quantization. Neglecting spatial gradients the field φ is a function of physical time t and fictitious time s . The 2nd quantized equation of motion is

$$\frac{\partial}{\partial s} \varphi = \ddot{\varphi} + 3H\dot{\varphi} + V'(\varphi) + L(s,t), \quad (1)$$

where H is the Hubble parameter, V is the potential under consideration and $L(s,t)$ is Gaussian white noise, δ correlated in both s and t . For e.g. a numerical simulation we may discretize Eq. (1) using

$$s = n\tau, \quad (2)$$

$$t = i\delta, \quad (3)$$

where n and i are integers, τ is a fictitious time lattice constant, and δ is a physical time lattice constant. The continuum limit requires $\tau \rightarrow 0$, $\delta \rightarrow 0$, but we will later argue that it makes physical sense to keep small but finite lattice constants of the order of the Planck length. We obtain

$$\begin{aligned} \frac{\varphi_{n+1}^i - \varphi_n^i}{\tau} &= \frac{1}{\delta^2} (\varphi_n^{i+1} - 2\varphi_n^i + \varphi_n^{i-1}) + 3\frac{H}{\delta} (\varphi_n^i - \varphi_n^{i-1}) \\ &+ V'(\varphi_n^i) + \text{noise}. \end{aligned} \quad (4)$$

This can be written as the following recurrence relation for the field φ_n^i :

$$\begin{aligned} \varphi_{n+1}^i &= (1 - \alpha) \left\{ \varphi_n^i + \frac{\tau}{1 - \alpha} V'(\varphi_n^i) \right\} + 3\frac{H\tau}{\delta} (\varphi_n^i - \varphi_n^{i-1}) \\ &+ \frac{\alpha}{2} (\varphi_n^{i+1} + \varphi_n^{i-1}) + \tau \cdot \text{noise}, \end{aligned} \quad (5)$$

where a dimensionless coupling constant α is introduced as

$$\alpha := \frac{2\tau}{\delta^2}. \quad (6)$$

We also introduce a dimensionless field variable Φ_n^i by writing $\varphi_n^i = \Phi_n^i p_{max}$, where p_{max} is some (so far) arbitrary en-

ergy scale. The above scalar field dynamics is equivalent to a spatially extended dynamical system (a coupled map lattice [16]) of the form

$$\begin{aligned} \Phi_{n+1}^i &= (1-\alpha)T(\Phi_n^i) + \frac{3}{2}H\delta\alpha(\Phi_n^i - \Phi_n^{i-1}) \\ &+ \frac{\alpha}{2}(\Phi_n^{i+1} + \Phi_n^{i-1}) + \tau \cdot \text{noise}, \end{aligned} \quad (7)$$

where the local map T is given by

$$T(\Phi) = \Phi + \frac{\tau}{p_{max}(1-\alpha)}V'(p_{max}\Phi). \quad (8)$$

Here the prime means

$$\frac{\partial}{\partial\varphi} = \frac{1}{p_{max}} \frac{\partial}{\partial\Phi}. \quad (9)$$

Note that a symmetric diffusively coupled map lattice of the form

$$\Phi_{n+1}^i = (1-\alpha)T(\Phi_n^i) + \frac{\alpha}{2}(\Phi_n^{i+1} + \Phi_n^{i-1}) + \tau \cdot \text{noise} \quad (10)$$

is obtained if $H\delta \ll 1$, equivalent to

$$\delta \ll H^{-1}, \quad (11)$$

meaning that the physical time lattice constant δ is much smaller than the age of the Universe. In this case the term proportional to H in Eq. (7) can be neglected. The local map T depends on the potential under consideration. Since we restrict ourselves to real scalar fields φ , T is a 1-dimensional map.

The main result of our consideration is that iteration of a coupled map lattice of the form (10) with a given map T has physical meaning: It means that one is considering the second-quantized dynamics of a self-interacting real scalar field φ with a force V' given by

$$V'(\varphi) = \frac{1-\alpha}{\tau} \left\{ -\varphi + p_{max} T\left(\frac{\varphi}{p_{max}}\right) \right\}. \quad (12)$$

Integration yields

$$V(\varphi) = \frac{1-\alpha}{\tau} \left\{ -\frac{1}{2}\varphi^2 + p_{max} \int d\varphi T\left(\frac{\varphi}{p_{max}}\right) \right\} + \text{const.} \quad (13)$$

In terms of the dimensionless field Φ this can be written as

$$V(\varphi) = \frac{1-\alpha}{\tau} p_{max}^2 \left\{ -\frac{1}{2}\Phi^2 + \int d\Phi T(\Phi) \right\} + \text{const.} \quad (14)$$

III. CHAOTIC φ^4 THEORY

An interesting observation is the following one. The lattice constant τ of fictitious time should be small, in order to approximate the continuum theory, which is ordinary quantum field theory. If τ is small one naively expects the map T given by Eq. (8) to be close to the identity for finite forces V' , since $\tau V'/p_{max}$ is small. What about, however, very strong forces V' due to very strongly self-interacting fields? If p_{max}/τ is of the same order of magnitude as V' then a nontrivial map T can arise. In particular, this map may even exhibit chaotic behavior.

As a distinguished example of a φ^4 theory generating strongest possible chaotic behavior, let us consider the map

$$\Phi_{n+1} = T_{-3}(\Phi_n) = -4\Phi_n^3 + 3\Phi_n \quad (15)$$

on the interval $\Phi \in [-1, 1]$. T_{-3} is the negative third-order Tschebyscheff map, a standard example of a map exhibiting strongly chaotic behavior. It is conjugated to a Bernoulli shift, thus generating the strongest possible stochastic behavior possible for a smooth low-dimensional deterministic dynamical system. The corresponding potential is given by

$$V_{-3}(\varphi) = \frac{1-\alpha}{\tau} \left\{ \varphi^2 - \frac{1}{p_{max}^2} \varphi^4 \right\} + \text{const.}, \quad (16)$$

or, in terms of the dimensionless field Φ ,

$$V_{-3}(\varphi) = \frac{1-\alpha}{\tau} p_{max}^2 (\Phi^2 - \Phi^4) + \text{const.} \quad (17)$$

Apparently, starting from this potential we obtain by second quantization a field φ that rapidly fluctuates in fictitious time on some finite interval, provided that initially $\varphi_0 \in [-p_{max}, p_{max}]$. The small noise term in Eq. (10) can be neglected as compared to the deterministic chaotic fluctuations of the field.

Of physical relevance are the expectations of suitable observables with respect to the ergodic chaotic dynamics. For example, the expectation $\langle V_{-3}(\varphi) \rangle$ of the potential is a possible candidate for vacuum energy in our Universe. One obtains

$$\langle V_{-3}(\varphi) \rangle = \frac{1-\alpha}{\tau} p_{max}^2 (\langle \Phi^2 \rangle - \langle \Phi^4 \rangle) + \text{const.} \quad (18)$$

For uncoupled Tschebyscheff maps ($\alpha=0$), expectations of any observable A can be evaluated as the ergodic average

$$\langle A \rangle = \int_{-1}^{+1} A(\Phi) d\mu(\Phi), \quad (19)$$

with the natural invariant measure being given by

$$d\mu(\Phi) = \frac{d\Phi}{\pi\sqrt{1-\Phi^2}} \quad (20)$$

(see any textbook on chaotic dynamics, e.g. [17]). This measure describes the probability distribution of the iterates under long-term iteration. From Eq. (20) one obtains $\langle \Phi^2 \rangle = \frac{1}{2}$ and $\langle \Phi^4 \rangle = \frac{3}{8}$; thus

$$\langle V_{-3}(\varphi) \rangle = \frac{1}{8} \frac{p_{max}^2}{\tau} + \text{const.} \quad (21)$$

Alternatively, we may consider the positive Tschebyscheff map $T_3(\Phi) = 4\Phi^3 - 3\Phi$. This basically exhibits the same dynamics as T_{-3} , up to a sign. Repeating the same calculation we obtain

$$V_3(\varphi) = \frac{1-\alpha}{\tau} \left\{ -2\varphi^2 + \frac{1}{p_{max}^2} \varphi^4 \right\} + \text{const} \quad (22)$$

and

$$V_3(\varphi) = \frac{1-\alpha}{\tau} p_{max}^2 (-2\Phi^2 + \Phi^4) + \text{const.} \quad (23)$$

For the expectation of the vacuum energy one gets

$$\langle V_3(\varphi) \rangle = \frac{1-\alpha}{\tau} p_{max}^2 (-2\langle \Phi^2 \rangle + \langle \Phi^4 \rangle) + \text{const.} \quad (24)$$

which for $\alpha=0$ reduces to

$$\langle V_3(\varphi) \rangle = -\frac{5}{8} \frac{p_{max}^2}{\tau} + \text{const.} \quad (25)$$

Symmetry considerations between T_{-3} and T_3 suggest to take the additive constant const as

$$\text{const} = + \frac{1-\alpha}{\tau} p_{max}^2 \frac{1}{2} \langle \Phi^2 \rangle. \quad (26)$$

In this case one obtains the fully symmetric equation

$$\langle V_{\pm 3}(\varphi) \rangle = \pm \frac{1-\alpha}{\tau} p_{max}^2 \left\{ -\frac{3}{2} \langle \Phi^2 \rangle + \langle \Phi^4 \rangle \right\}, \quad (27)$$

which for $\alpha \rightarrow 0$ reduces to

$$\langle V_{\pm 3}(\varphi) \rangle = \pm \frac{p_{max}^2}{\tau} \left(-\frac{3}{8} \right). \quad (28)$$

The simplest model for dark energy in the Universe, as generated by a chaotic φ^4 theory, would be to identify $\frac{3}{8} p_{max}^2 / \tau = \rho_\Lambda$, the constant vacuum energy density corresponding to a classical cosmological constant Λ , which stays constant during the expansion of the Universe. This is certainly a possible simple model. On the other hand, such an approach would solve neither the cosmological constant nor the cosmological coincidence problem. For this reason we will turn to a more advanced model in the following sections, which will naturally produce the right amount of vacuum energy density in the Universe.

Before proceeding to this model, let us provide some general comments on the physical meaning of the parameter τ .

The vacuum energy generated by the chaotic fields is inversely proportional to τ [see Eq. (28)]. Strict equivalence of the stochastic quantization method with quantum field theory requires the continuum limit $\tau \rightarrow 0$. In this limit Eq. (28) can still generate a finite amount of vacuum energy, provided both $\tau \rightarrow 0$ and $p_{max} \rightarrow 0$ such that the ratio p_{max}^2 / τ stays finite. In fact, many of the results presented in this paper depend only on the ratio p_{max}^2 / τ and not on the individual values attributed to τ and p_{max} .

From the viewpoint of ordinary quantum field theory, the limit $\tau \rightarrow 0$ and $p_{max} \rightarrow 0$ means that one formally considers a potential $V(\varphi) = \mu^2 \varphi^2 + \lambda \varphi^4$ for which both potential parameters $\mu^2 \sim \tau^{-1}$ and $\lambda \sim \tau^{-1} p_{max}^{-2}$ diverge [see Eq. (22)]; moreover the field φ lives on an infinitely small support $[-p_{max}, p_{max}]$. Clearly, this is a very singular type of quantum field theory, which in principle can also be studied by other means than stochastic quantization, although perturbation theory will not be applicable. The advantage of our formulation in terms of stochastic quantization is that for the dimensionless chaotically evolving field Φ the potential remains finite [see, e.g., Eq. (23)]. If there is no fictitious time, then the parameter τ enters in the form of the (singular) potential parameters $\mu^2 \sim \tau^{-1}$ and $\lambda \sim \tau^{-1} p_{max}^{-2}$. In the next section we will argue that on physical grounds it makes sense to consider very small but finite τ .

IV. REPRODUCING THE CURRENTLY MEASURED DARK ENERGY DENSITY

To obtain quantitative statements on the dark energy density as generated by some chaotically evolving field φ , let us fix the free parameters τ and p_{max} by some physical arguments. Let us start with the parameter τ . It is the lattice constant of fictitious time s and has dimension GeV^{-2} . Ordinary stochastic quantization based on Gaussian white noise requires the continuum limit $\tau \rightarrow 0$. But since quantum field theory runs into difficulties at the Planck scale m_{pl} and is expected to be replaced by a more advanced theory at this scale, it is most reasonable to take the small but finite value

$$\tau \sim m_{pl}^{-2}. \quad (29)$$

Next, consider the parameter p_{max} . It has dimension GeV and describes the maximum energy scale of our rapidly fluctuating scalar fields φ , which take on values on the finite interval $[-p_{max}, p_{max}]$. A natural value of p_{max} follows if one associates the rapidly fluctuating chaotic fields φ_n^i with vacuum fluctuations that are allowed due to the uncertainty relation

$$\Delta E \Delta t = O(\hbar). \quad (30)$$

Taking $\Delta t \sim t$ of the order of the age of the Universe, a corresponding energy uncertainty ΔE arises. This ΔE is very large for a very young Universe, and then decreases to extremely small values for the current age of the Universe of about 13.7 Gyr. Any finite age t of the Universe implies that spontaneous vacuum fluctuations with energies of order $\Delta E \sim t^{-1}$ can occur. It is physically plausible to identify these

energy fluctuations ΔE with the rapidly fluctuating chaotic fields $\varphi = p_{max} \Phi_n^i$, since both ΔE and φ live on a finite interval, and both fluctuate in an unpredictable way. The uncertainty relation (30) together with $\Delta t \sim t$ implies (in units where $\hbar = c = 1$)

$$p_{max} \sim \frac{1}{t}. \quad (31)$$

In this way the energy scale p_{max} occurring in the chaotic field theories is most naturally identified with the inverse age of the Universe. However, note that this quantum mechanical interpretation in terms of vacuum fluctuations requires a *chaotic* map T , since some regularly evolving Φ_n^i cannot be associated with fluctuations at all.

It is remarkable that by taking Eq. (29) and Eq. (31) together, the right amount of vacuum energy follows without any fine-tuning. One has for generic chaotic maps T

$$\langle V(\varphi) \rangle \sim \frac{p_{max}^2}{\tau} \sim H^2 m_{pl}^2, \quad (32)$$

since $t^{-1} \sim H$. Moreover,

$$H^2 = \frac{8\pi G}{3} \rho_c \sim \frac{1}{m_{pl}^2} \rho_c \quad (33)$$

where ρ_c denotes the critical density of a flat Universe and $G = m_{pl}^{-2}$ is the gravitational constant. Combining Eqs. (32) and (33) one obtains

$$\langle V(\varphi) \rangle \sim \rho_c, \quad (34)$$

as required and confirmed by current astronomical observations. Our simple physical interpretation, namely to interpret the chaotic fluctuations as vacuum fluctuations allowed due to the finite age of the Universe, thus yields the right order of magnitude of dark energy density.

In general, Eq. (34) yields only order of magnitude estimates. Nevertheless, it is instructive to work out some concrete numbers, based on simple model assumptions. For example, we may assume that the entire vacuum energy of the Universe is due to one chaotic field described by V_{-3} . The current age of the Universe is $t_0 = (13.7 \pm 0.2)$ Gyr $= (4.32 \pm 0.06) \times 10^{17}$ s [3]. Using an uncertainty relation of the form $\Delta E \Delta t = \hbar/2$ we get $p_{max} = 1/(2t_0) = (7.62 \pm 0.08) \times 10^{-43}$ GeV. Choosing $\tau = \kappa m_{pl}^2$, where κ is some dimensionless number of $O(1)$, we get

$$\langle V(\varphi) \rangle = \frac{3}{8} \frac{p_{max}^2}{\tau} = (3.19 \pm 0.05) \times 10^{-47} \kappa^{-1} \text{ GeV}^4. \quad (35)$$

The current observational estimate of dark energy density in the Universe is [3]

$$\rho_\varphi^{obs} = (2.9 \pm 0.3) \times 10^{-47} \text{ GeV}^4, \quad (36)$$

which is consistent with $\kappa \approx 1$. If the observed dark energy in the Universe is produced by our chaotic theory, then the measured data imply

$$\kappa = 1.10 \pm 0.10. \quad (37)$$

Once again, this estimate is based on concrete model assumptions. In general, we do not know the precise values of the proportionality constants in Eqs. (29) and (31); moreover we do not know how many different chaotic fields may contribute to the dark energy of the Universe.

Of course, the actual properties of the chaotic dark energy component depend on the classical equation of state of the chaotic fields, which will be dealt with in the next section.

V. ENERGY DENSITY, PRESSURE, AND EQUATION OF STATE

The kinetic energy term of our chaotic fields is given by

$$T_{kin} = \frac{1}{2} \left(\frac{\partial}{\partial t} \varphi \right)^2. \quad (38)$$

Discretized with lattice constant δ we obtain for the expectation of T_{kin}

$$\begin{aligned} \langle T_{kin} \rangle &= \frac{1}{2} \frac{p_{max}^2}{\delta^2} \langle (\Phi_n^i - \Phi_n^{i-1})^2 \rangle \\ &= \frac{p_{max}^2}{\tau} \frac{1}{2} \alpha (\langle \Phi^2 \rangle - \langle \Phi_n^i \Phi_n^{i-1} \rangle). \end{aligned} \quad (39)$$

In particular, for $\alpha \rightarrow 0$ the expectation of kinetic energy vanishes, and a Universe mainly filled with such a field is vacuum energy dominated.

In general, the expectation of the total energy density $\langle \rho \rangle$ is given by

$$\langle \rho \rangle = \langle T_{kin} \rangle + \langle V \rangle \quad (40)$$

and the expectation of the pressure by

$$\langle p \rangle = \langle T_{kin} \rangle - \langle V \rangle. \quad (41)$$

For the map T_{-3} one obtains

$$\begin{aligned} \langle \rho \rangle &= \frac{p_{max}^2}{\tau} \left\{ \frac{\alpha}{2} (\langle \Phi^2 \rangle - \langle \Phi_n^i \Phi_n^{i-1} \rangle) + (1 - \alpha) \right. \\ &\quad \left. \times \left(\frac{3}{2} \langle \Phi^2 \rangle - \langle \Phi^4 \rangle \right) \right\}, \end{aligned} \quad (42)$$

$$\begin{aligned} \langle p \rangle &= \frac{p_{max}^2}{\tau} \left\{ \frac{\alpha}{2} (\langle \Phi^2 \rangle - \langle \Phi_n^i \Phi_n^{i-1} \rangle) - (1 - \alpha) \right. \\ &\quad \left. \times \left(\frac{3}{2} \langle \Phi^2 \rangle - \langle \Phi^4 \rangle \right) \right\}, \end{aligned} \quad (43)$$

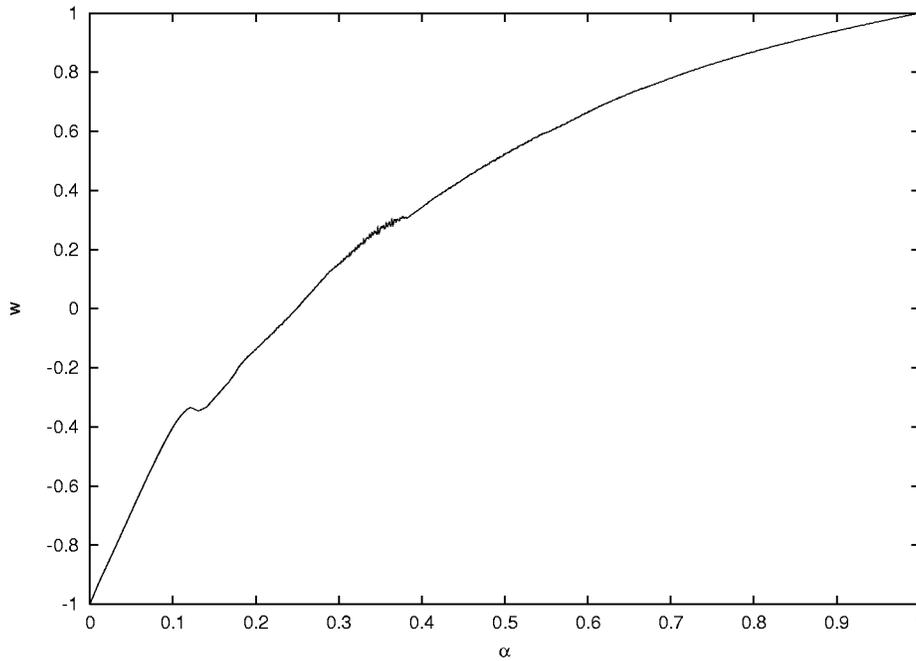


FIG. 1. Classical equation of state $w = \langle p \rangle / \langle \rho \rangle$ of the chaotic field φ as a function of the coupling α .

where the additive constant of the self-interacting potential is fixed by the symmetry consideration of Sec. III. The above equations yield the equation of state

$$w = \frac{\langle p \rangle}{\langle \rho \rangle} \tag{44}$$

which varies as a function of the coupling α in a nontrivial way.

For $\alpha = 0$, the equation of state of our fields is $w = -1$, since the expectation of kinetic energy vanishes (all expectations should be interpreted as quantum mechanical expectations with respect to second quantization). For small α , w is close to -1 . It should be clear that although our fields

fluctuate rapidly in both physical and fictitious time, these fluctuations are averaged away when doing the quantum mechanical expectations. Thus the classical picture that arises out of this 2nd quantized rapidly fluctuating model is a very homogeneous field.

The expectations in Eqs. (42), (43) are easily numerically calculated by long-term iterating the coupled map lattice for random initial conditions and averaging over all i and n . We used lattices of size 10000 with periodic boundary conditions. The result for the equation of state $w(\alpha)$ is displayed in Fig. 1. For small α , w grows approximately in a linear way. It monotonically increases from $w = -1$ for $\alpha = 0$ to $w = +1$ for $\alpha = 1$, up to a wiggle at $\alpha \approx 0.12$. Figure 2 shows the corresponding (classical) energy density and pressure of

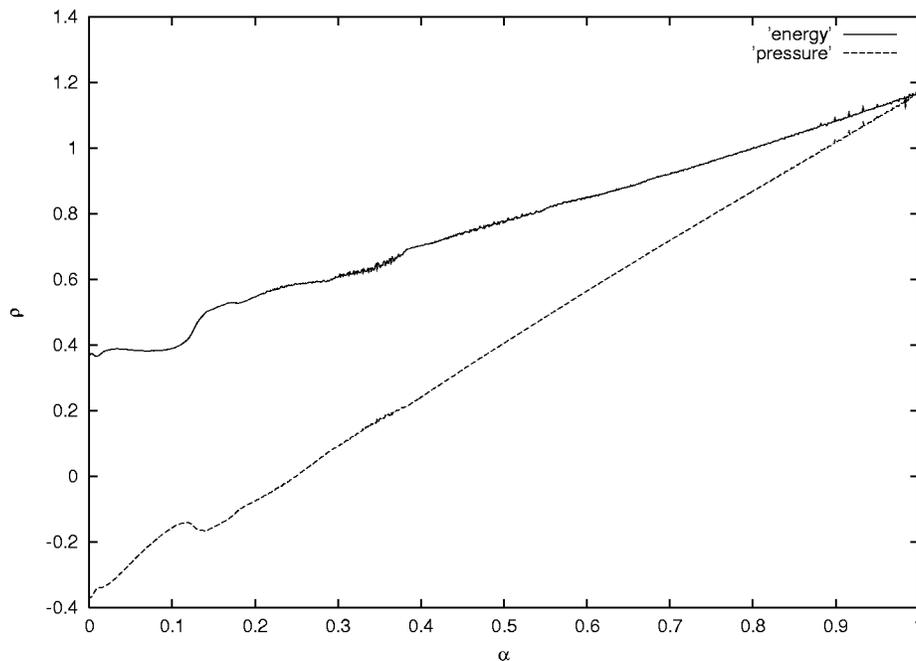


FIG. 2. Expectation of energy and pressure of the chaotic field as a function of the coupling α .

the field. To account for the currently observed dark energy in the Universe, most chaotic fields must have a coupling α smaller than about 0.04. Larger α are ruled out by the observations providing evidence for $w < -0.78$ [3].

VI. EINSTEIN EQUATIONS AND DYNAMICAL EVOLUTION

Let us now consider the Einstein equations for a homogeneous and isotropic Universe that consists of three different components: matter, radiation, and chaotic fields. These three components are labeled by the indices m, r, φ , respectively. One has

$$H^2 = \left(\frac{\dot{a}}{a}\right)^2 = \frac{8}{3} \pi G (\rho_\varphi + \rho_m + \rho_r), \quad (45)$$

$$\frac{\ddot{a}}{a} = -\frac{4}{3} \pi G (\rho_\varphi + 3p_\varphi + \rho_m + \rho_r + 3p_r). \quad (46)$$

Here ρ_j denotes the (classical) energy density of component j , and p_j the pressure. For simpler notation we omit the expectation values $\langle \dots \rangle$. The equation of state of each component is $w_j = p_j / \rho_j$. As is well known, one has for matter $w_m = 0$, for radiation $w_r = 1/3$, whereas the equation of state w_φ of the chaotic fields φ depends on the coupling α (see Fig. 1).

The Einstein (or Friedmann) equations are usually supplemented by the assumption of conservation of energy for each species j ,

$$\dot{\rho}_j = -3H(\rho_j + p_j). \quad (47)$$

These equations can be derived from the Einstein equations under the assumption of adiabatic expansion, i.e. one assumes that no entropy is produced.

For a Universe dominated by a species j with constant equation of state $w_j = p_j / \rho_j$ Eq. (47) leads to

$$\rho_j \sim a^{-3(1+w_j)}. \quad (48)$$

We obtain the well-known result that for matter $\rho_m \sim a^{-3}$, for radiation $\rho_r \sim a^{-4}$, whereas for true classical vacuum energy (a cosmological constant Λ) with $w = -1$ one has no dependence on a at all, $\rho_\Lambda = \text{const}$.

For the chaotic fields, energy conservation is a non-trivial issue, for the following reasons: (i) These fields model vacuum fluctuations, and vacuum fluctuations by definition do not conserve energy. (ii) The chaotic field dynamics (like any chaotic dynamics) constantly produces entropy in fictitious time, whereas the Friedmann equations (containing no fictitious time) describe an adiabatic expansion of the Universe. (iii) The coupling constant α and hence the equation of state of the chaotic field may change in time. (iv) There may be an entire spectrum of chaotic fields with different couplings α that can interact with each other. (v) The chaotic fields may interact with dark matter.

In the following, we want to restrict ourselves to a simple scenario where the coupling α is small and where there is energy conservation of quantum mechanical expectations in full agreement with the Friedmann equations. Two interesting possibilities arise in this context:

(a) A *symmetric phase* of the Universe, where both the negative and positive Tschebyscheff dynamics are active. In this case the vacuum energies $\rho_\varphi^- := \langle V_{-3}(\varphi) \rangle$ and $\rho_\varphi^+ := \langle V_{+3}(\varphi) \rangle$ cancel to zero:

$$\rho_\varphi = \rho_\varphi^- + \rho_\varphi^+ = 0 \quad (49)$$

[a consequence of Eq. (27)]. Since the total vacuum energy ρ_φ is zero, there is no statement from the Einstein (or Friedmann) equations on the time evolution of the absolute value $|\rho_\varphi^\pm|$. We are free to postulate the validity of Eq. (32)

$$|\rho_\varphi^\pm| = |\langle V_{\pm 3}(\varphi) \rangle| \sim \frac{p_{max}^2}{\tau} \sim \frac{m_{Pl}^2}{t^2} \quad (50)$$

for arbitrary couplings α .

(b) An *asymmetric phase* of the Universe, where only the negative Tschebyscheff dynamics T_{-3} is active (or where it dominates as compared to T_{+3}). In this case one has

$$\rho_\varphi > 0. \quad (51)$$

If there is a chaotic field with coupling constant α that does not interact with other fields then energy conservation implies

$$\rho_\varphi \sim a^{-3(1+w_\varphi)}, \quad (52)$$

where $w_\varphi = w_\varphi(\alpha)$ is the equation of state as displayed in Fig. 1. Note that chaotic fields with $w_\varphi \approx 0$ ($\alpha \approx 0.25$) behave similar as dark matter, whereas chaotic fields with $w_\varphi \approx -1$ ($\alpha \approx 0$) act like a cosmological constant (see also [18]).

We now show that Eq. (50), derived from the uncertainty relation, naturally leads to tracking behavior of the vacuum energy density. Suppose the Universe surrounding the chaotic fields is dominated by a species with equation of state $w > -1$ (typically matter or radiation), then

$$a(t) \sim t^{2/[3(1+w)]} \Leftrightarrow t^{-2} \sim a^{-3(1+w)}. \quad (53)$$

Putting this into Eq. (50) we obtain

$$|\rho_\varphi^\pm| = |\langle V_{\pm 3}(\varphi) \rangle| \sim a^{-3(1+w)}, \quad (54)$$

i.e. the vacuum energy density associated with the chaotic fields decays *in the same way* with a as the density of the dominating species.

Equation (54) can help to naturally avoid the cosmological coincidence problem. Consider e.g. the following scenario. Initially (say, shortly after inflation) we may have a symmetric phase where $\rho_r \sim |\rho_\varphi^\pm| \gg \rho_m$. Here ρ_φ^- denotes the vacuum energy density of the negative Tschebyscheff map, which is canceled by the vacuum energy density $\rho_\varphi^+ = -\rho_\varphi^-$ of the positive Tschebyscheff map. Then $|\rho_\varphi^\pm|$ first decays approximately as a^{-4} , since the Universe is radiation domi-

nated. At some stage we arrive at $\rho_r \sim |\rho_\varphi^\pm| \sim \rho_m$, and from then on matter dominates over radiation, so that from then on $|\rho_\varphi^\pm|$ decays approximately as a^{-3} . During the late-time evolution of the Universe, $|\rho_\varphi^\pm|$ will always stay of the same order of magnitude as ρ_m , since both $|\rho_\varphi^\pm|$ and ρ_m decay as a^{-3} . In spite of this, for small enough couplings α the chaotic fields have a classical equation of state close to $w_\varphi = -1$, and can thus produce the accelerated expansion of the Universe via Eq. (46), provided there is symmetry breaking between T_{+3} and T_{-3} at some late stage of the evolution of the Universe. A concrete mechanism for this will be discussed in Sec. IX.

What is our physical interpretation of the chaotic fields in the Universe? For $\alpha=0$, it can be rigorously proved that rescaled deterministic chaotic Tschebyscheff maps can be used to generate spatio-temporal Gaussian white noise on a larger scale [12,13]. In other words, on fictitious time scales $\tau' \gg \tau$ and physical time scales $\delta' \gg \delta$ the chaotic noise just looks like ordinary Gaussian white noise. We may thus couple the chaotic fields φ to ordinary standard model fields in order to second quantize the standard model fields, i.e. use the chaotic fields as a source of quantization noise. This is the basic idea of the so-called chaotic quantization approach [12]. The chaotic fields are well embedded in this way and since they are just playing the role of quantization noise, we do not expect them to have any disturbing influence on, say, baryogenesis and similar processes in the early Universe. In this interpretation vacuum energy just arises out of the expectation of a classical potential that generates quantization noise.

VII. PREDICTION OF Ω_φ AND Ω_m

Our approach allows for the prediction of the order of magnitude of cosmological parameters such as $\Omega_\varphi = \rho_\varphi / \rho_c$ and $\Omega_m = \rho_m / \rho_c$ at the present time. Let us start from the uncertainty relation in the form

$$\Delta E \Delta t = \frac{\hbar}{2}, \quad (55)$$

which implies

$$p_{max} = \frac{1}{2t}. \quad (56)$$

Choosing the time scale $\tau = \kappa m_{pl}^{-2}$ we get for $\alpha \approx 0$

$$\rho_\varphi^- = \langle V_{-3}(\varphi) \rangle = \frac{3}{8} \frac{p_{max}^2}{\tau} = \frac{3}{32} \frac{m_{pl}^2}{\kappa} \frac{1}{t^2}. \quad (57)$$

During the radiation dominated period of the Universe one has for the energy density of radiation

$$\rho_r = \frac{\pi^2}{30} N(T) T^4, \quad (58)$$

where T is the temperature and $N(T)$ is the number of relativistic particle degrees of freedom. There is also a relation between time and temperature, namely

$$t = \sqrt{\frac{90}{32\pi^3 N(T)}} \frac{m_{pl}}{T^2}. \quad (59)$$

Putting Eq. (59) into Eq. (57) one obtains

$$\rho_\varphi^- = \frac{1}{30} \pi^3 \frac{1}{\kappa} N(T) T^4 = \frac{\pi}{\kappa} \rho_r. \quad (60)$$

This equation once again shows that it is reasonable to assume that ρ_r and ρ_φ^- have the same order of magnitude. Since ρ_φ^- decays in the same way as ρ_r , Eq. (60) is valid during the entire radiation dominated epoch. Finally ρ_r falls below ρ_m and from then on we have

$$\rho_\varphi^- \approx \frac{\pi}{\kappa} \rho_m. \quad (61)$$

After symmetry breaking we have $\rho_\varphi^- = \rho_\varphi \approx \text{const}$. This implies a prediction for $\Omega_\varphi := \rho_\varphi / \rho_c$ at the present time, namely

$$\Omega_\varphi = \frac{\rho_\varphi}{\rho_\varphi + \rho_m + \rho_r} \approx \frac{\pi/\kappa}{1 + \pi/\kappa}, \quad (62)$$

neglecting ρ_r at the present epoch and assuming that the symmetry breaking took place rather recently. In Sec. IV we saw that the currently observed dark energy density is best fitted by the value $\kappa = 1.10 \pm 0.10$. Equation (62) yields with this value the prediction

$$\Omega_\varphi \approx 0.74, \quad \Omega_m \approx 0.26, \quad (63)$$

which is consistent with observations [3]. Again, due to the reasons that were already mentioned at the end of Sec. IV, Eq. (63) should be regarded only as an order-of-magnitude estimate.

VIII. FIXING FUNDAMENTAL PARAMETERS

We have seen that chaotic fields can generate the right amount of vacuum energy and have the scope to avoid the cosmological constant and coincidence problem. We now show that they also offer solutions to the problem of uniqueness of vacua.

First, let us slightly generalize the chaotic field dynamics (10) to

$$\Phi_{n+1}^i = (1 - \alpha) T(\Phi_n^i) + \sigma \frac{\alpha}{2} [T^b(\Phi_n^{i-1}) + T^b(\Phi_n^{i+1})] \quad (64)$$

(we neglect the small noise term). The case $\sigma = +1$ is called ‘‘diffusive coupling,’’ the case $\sigma = -1$ ‘‘anti-diffusive coupling.’’ Chaotic fields with $b=1$ are said to be of type A [$T^1(\Phi) =: T(\Phi)$], chaotic fields with $b=0$ to be of type B [$T^0(\Phi) =: \Phi$]. In [14] the chaotic fields were called ‘‘chaotic strings,’’ but this is only a different name for the same dy-

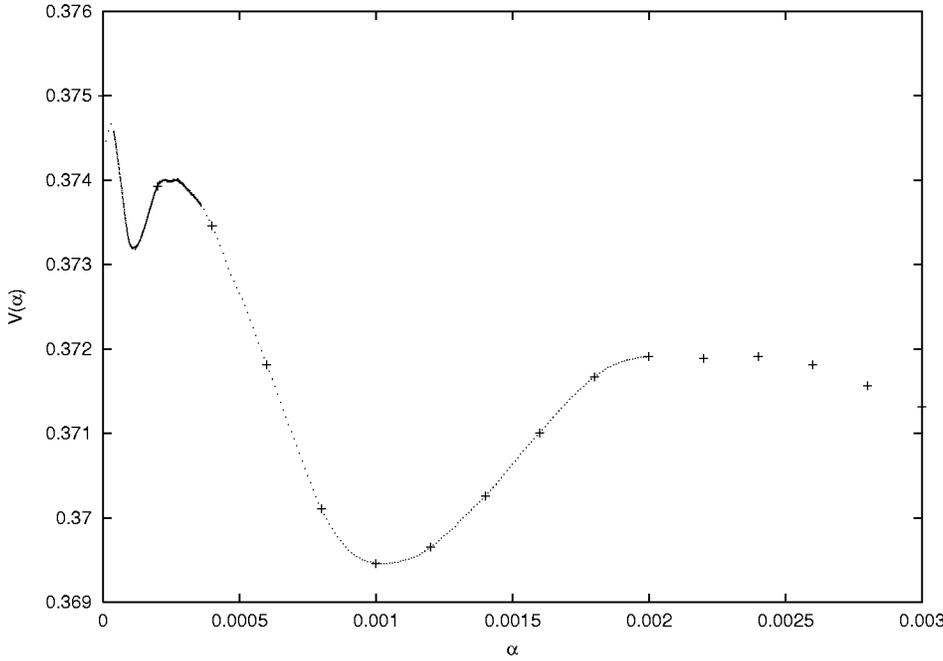


FIG. 3. Self-energy $V(\alpha)$ of the type-A chaotic field in the low-coupling region. There are local minima at couplings a_i that coincide with the weak coupling constants of right-handed fermions in the standard model.

namics. Our derivation in Sec. II leads to chaotic fields of B type with diffusive coupling, but from a dynamical systems point of view all 4 degrees of freedom ($b=0,1$, $\sigma=\pm 1$) exist and are of relevance. As shown in detail in [14], there are two different types of vacuum energies for the chaotic fields, namely

(1) the self-energy

$$V(\alpha) := \frac{P_{max}^2}{\tau} \left(\frac{3}{2} \langle \Phi^2 \rangle - \langle \Phi^4 \rangle \right) \quad (65)$$

and

(2) the interaction energy

$$W(\alpha) := \frac{P_{max}^2}{2\tau} \langle \Phi_n^i \Phi_n^{i+1} \rangle. \quad (66)$$

Equation (65) actually represents the self-energy of the map T_{-3} ; the self-energy of the map T_{+3} has opposite sign and cancels the self-energy of T_{-3} in the symmetric phase.

Basically, the self-energy is the expectation of the potential that generates the chaotic dynamics in fictitious time, and the interaction energy is the expectation of the potential that generates the diffusive coupling in physical time. One may also define a total vacuum energy as $H^\pm(\alpha) := V(\alpha) \pm \alpha W(\alpha)$, where the $-$ sign corresponds to diffusive and the $+$ sign to anti-diffusive coupling. All additive constants are fixed by the postulate of invariance of the theory under global and local Z_2 transformations [13]. In other words, we allow for the existence of a symmetric phase. For small α the interaction energy can be neglected as compared to the self-energy; moreover, the type-A and type-B forms are observed to have the same self-energy in this limit.

The central hypothesis of this paper is a symmetry between standard model coupling constants and the chaotic field couplings α . We assume that for any dimensionless

coupling constant α that appears in the standard model of electroweak and strong interactions, there is a corresponding chaotic field that is just coupled with this α . The Universe then tries to reach a state of minimum vacuum energy by adjusting its free parameters in such a way that the chaotic fields reach a state of minimum vacuum energy.

While at first sight this may look like a purely theoretical concept, there is numerical evidence that this principle is indeed physically realized. As an example, Fig. 3 shows the self-energy $V(\alpha)$ of our chaotic fields of type A with diffusive coupling in the low-coupling region. We observe that $V(\alpha)$ has local minima at

$$a_1 = 0.000246(2), \quad (67)$$

$$a_2 = 0.00102(1), \quad (68)$$

$$a_3 = 0.00220(1) \quad (69)$$

(a_1 and a_3 are actually small local minima on top of the hill).

On the other hand, in the standard model of electroweak interactions the weak coupling constant is given by

$$\alpha_{weak} = \alpha_{el} \frac{(T_3 - Q \sin^2 \theta_W)^2}{\sin^2 \theta_W \cos^2 \theta_W}. \quad (70)$$

Here Q is the electric charge of the particle ($Q = -1$ for electrons, $Q = 2/3$ for u -like quarks, $Q = -1/3$ for d -like quarks), and T_3 is the third component of the isospin ($T_3 = 0$ for right-handed particles, $T_3 = -\frac{1}{2}$ for e_L and d_L , $T_3 = +\frac{1}{2}$ for ν_L and u_L). Consider right-handed fermions f_R . With $\sin^2 \theta_W = \bar{s}_W^2 = 0.2318$ (as experimentally measured [19]) and the running electric coupling $\alpha_{el}(E)$ taken at energy scale $E = 3m_f$ we obtain from Eq. (70) the numerical values

$$\alpha_{weak}^{d_R}(3m_d)=0.000246, \quad (71)$$

$$\alpha_{weak}^{c_R}(3m_c)=0.001013, \quad (72)$$

$$\alpha_{weak}^{e_R}(3m_e)=0.00220. \quad (73)$$

There is an amazing numerical coincidence between the local minima a_1, a_2, a_3 of $V(\alpha)$ and the weak coupling constants of $f_R = u_R, c_R, e_R$, respectively.

Now regard the fine structure constant α_{el} and the Weinberg angle $\sin^2\theta_W$ as *a priori* free parameters. Suppose these parameters change to slightly different values. Then immediately this would produce larger vacuum energy $V(\alpha)$, since we get out of the local minima. The system is expected to be driven back to the local minima, and the fundamental parameters are stabilized in this way, provided the Universe is in an asymmetric phase.

The above example is only one example of a large number of numerical coincidences observed. In [13,14] an extensive numerical investigation of self-energies, interaction energies, and total vacuum energies was performed for the above chaotic field theories. A large number of amazing numerical coincidences was found.¹ These results are described in detail in [13,14]; here we summarize only the main results.

(1) The smallest (stable) zeros of the interaction energy $W(\alpha)$ coincide with running electroweak coupling constants, evaluated at energies given by the smallest fermionic mass scales. Type A describes d quarks and electrons interacting electrically, type B u quarks and neutrinos interacting weakly.

(2) Local minima of the self-energy $V(\alpha)$ coincide with running weak coupling constants of right-handed fermions, evaluated at the lightest fermionic mass scales.

(3) Local minima of the total vacuum energy $H^+(\alpha)$ occur at running strong coupling constants evaluated at the lightest baryonic energy scales.

(4) Local minima of the total vacuum energy $H^-(\alpha)$ occur at running strong couplings evaluated at the lightest mesonic energy scales.

In [13,14] chaotic fields corresponding to 2nd order Tschebyscheff maps were also investigated (see the Appendix), and the following numerical coincidences were found:

(1) The smallest (stable) zeros of $W(\alpha)$ coincide with running strong coupling constants evaluated at the smallest bosonic mass scales. Type A describes the W boson, type B the Higgs boson.

(2) Local minima of the self-energy $V(\alpha)$ coincide with Yukawa and gravitational couplings evaluated at the fermionic mass scales.

For more details, see [13,14].

All these numerically observed coincidences are not explainable as a random coincidence. Rather, they suggest in-

terpreting the coupling constant α of our second-quantized chaotic fields φ as a running gauge coupling. We are free to identify $\alpha = 2\pi/\delta^2$ with a gauge coupling, since the occurrence of a ratio of lattice constants τ and δ^2 is a new effect in our 2nd quantized discretized theory, and there is no theory of this dimensionless number so far, which represents a kind of metric for the 5th coordinate (the fictitious time). So we are indeed free to make the hypothesis that α coincides with a running gauge coupling. By doing so, we implicitly construct a symmetry between gauge couplings and gravitational couplings, since usually the strength of the kinetic term in the action of a field is determined by the metric, i.e. gravitational effects, whereas here it is fixed by standard model coupling strengths. The chaotic fields appear to select out of the infinitely many vacua allowed by string theory the unique ground state that corresponds to the known coupling constants of the Universe. All free parameters are fixed in the sense that if the fundamental parameters (masses, coupling constants, and mixing angles) had different values, larger vacuum energy would arise.

IX. SPONTANEOUS SYMMETRY BREAKING AND CANCELLATION OF UNWANTED VACUUM ENERGY

Chaotic scalar fields not only allow for a simple mechanism to produce dark energy, they also yield a simple mechanism to cancel unwanted dark energy. If we assume that both the positive and negative Tschebyscheff dynamics are physically realized, the corresponding vacuum energies precisely cancel for symmetry reasons [see Eq. (28)]. This symmetry is a Z_2 symmetry which is not there for ordinary smoothly evolving scalar fields (where opposite potentials lead to unstable or ill-defined behavior). On the other hand, if only the negative Tschebyscheff field dynamics is active, or if it dominates, then positive dark energy arises. This positive dark energy can drive inflation, fix standard model parameters as local minima in the dark energy landscape, and generate late-time acceleration. It is therefore desirable to construct a theory that allows for Z_2 symmetry breaking between positive and negative Tschebyscheff maps.

It is clear that in order to fix fundamental parameters with the methods described in the previous section, we must have a broken Z_2 symmetry at some stage of the evolution of the Universe. Indeed, the minimum requirement we need is at least one very early stage of broken symmetry, in order to *first-time fix* the fundamental parameters to the values which make the Universe work, and another late-time asymptotic state of broken symmetry, in order to *stabilize* the parameters to their known values so that they cannot drift away to other values. It is natural to identify the first phase of broken symmetry with the inflationary phase [21], and the other phase of broken symmetry with the late-time state of the Universe. In between, we may allow for a symmetric state, which has the advantage that nucleosynthesis is not spoilt.²

¹For the φ^4 theory the relevant energy scale is always $E = 3m_f$. The factor of 3 can be related to the index of the Tschebyscheff polynomial considered [13].

²The abundance of light elements is correctly predicted by standard big bang nucleosynthesis but is spoilt if there is too much dark energy [22]. The measured cosmic microwave background also

As a concrete simple model, consider a scalar field σ which takes on the value $\sigma=0$ in the symmetric phase and the values $\sigma=\pm 1$ in the phase where the Z_2 symmetry is spontaneously broken. The total potential describing the chaotic field dynamics is given by

$$V(\sigma, \Phi_-, \Phi_+) = \frac{1-\sigma}{2} V_{-3}(\Phi_-) + \frac{1+\sigma}{2} V_{+3}(\Phi_+), \quad (74)$$

where

$$V_{-3}(\Phi_-) = \frac{p_{max}^2}{\tau} \left(\Phi_-^2 - \Phi_-^4 + \frac{1}{2} \langle \Phi_-^2 \rangle \right) \quad (75)$$

is the potential generating the negative Tschebyscheff field dynamics and

$$V_{+3}(\Phi_+) = \frac{p_{max}^2}{\tau} \left(-2\Phi_+^2 + \Phi_+^4 + \frac{1}{2} \langle \Phi_+^2 \rangle \right) \quad (76)$$

the one generating the positive Tschebyscheff field dynamics in fictitious time ($\alpha \approx 0$). In the symmetric phase ($\sigma=0$) we obtain from Eq. (74)

$$\langle V(\sigma, \Phi_-, \Phi_+) \rangle = 0, \quad (77)$$

whereas in a broken phase with $\sigma=-1$ we obtain

$$\langle V(\sigma, \Phi_-, \Phi_+) \rangle = \langle V_{-3}(\Phi_-) \rangle = \frac{p_{max}^2}{\tau} \left(\frac{3}{2} \langle \Phi_-^2 \rangle - \langle \Phi_-^4 \rangle \right) > 0, \quad (78)$$

where we have relabeled $\Phi_- = \Phi$.

We assume that the symmetry is first spontaneously broken to $\sigma=-1$ at the onset of inflation. A large amount of positive vacuum energy is generated via Eq. (78), since at this stage the Universe is very young and $p_{max} \sim t^{-1} \sim H$. The chaotic fields can help to drive inflation, and fundamental parameters are pre-fixed as local minima in the dark energy landscape.

Then, there is a symmetric phase with $\sigma=0$. The consideration of Sec. VII applies but the dark energy is suppressed for symmetry reasons. Big bang nucleosynthesis and galaxy formation can go ahead without any problems. Note that during the symmetric epoch the fundamental parameters are no longer stabilized as local minima in the dark energy landscape. They can drift to slightly different values. This is consistent with the experimental findings of a varying fine structure constant [25].

Finally, there is late-time symmetry breaking to $\sigma=-1$. This phase is necessary because otherwise the fundamental parameters would keep on drifting to different values. By the late-time symmetry breaking, the parameters are finally forced back and stabilized at their equilibrium values, which

were already pre-fixed during inflation. This gives physical sense to the role of late-time dark energy in the Universe.

Late-time symmetry breaking could be physically understood as follows. Suppose the negative Tschebyscheff dynamics (which generates positive vacuum energy) is always uniformly distributed in space. It is a property of empty space-time. At a very early stage, the corresponding vacuum energy was strong and may have driven inflation. Then matter and radiation are created. Assume that the matter and radiation particles are second quantized by chaotic noise generated by the positive Tschebyscheff map (which generates negative vacuum energy). Then, as soon as sufficiently much matter and radiation have been created and quantized, the two vacuum energies $\langle V_{+3} \rangle$ and $\langle V_{-3} \rangle$ can precisely (or almost) cancel, as long as matter and radiation are uniformly distributed in space. Inflation may stop in this way and we obtain a symmetric phase of the Universe after inflation. But at a late stage of the evolution of the Universe matter clumps into galaxies. Since the negative vacuum energy is generated by chaotic quantization noise for each particle it follows the spatial distribution of matter. The positive vacuum energy remains uniformly distributed. Hence after structure formation there is no longer any spatially uniform cancellation of vacuum energy. Empty space has an excess of positive vacuum energy, galaxies are spatial regions with an excess of negative vacuum energy due to quantization noise (the negative vacuum energy in the galaxy can be partially compensated by positive kinetic terms that arise out of spatial inhomogeneities in the galaxy). In this physical interpretation the late-time symmetry breaking is related to structure formation.

X. CONCLUSION

We have presented a new model for dark energy in the Universe. This model is based on a rather conservative approach, the assumption of the existence of second quantized self-interacting scalar fields described by a ϕ^4 theory. However, the main difference is that these fields are strongly self-interacting, rather than weakly. When doing 2nd quantization using the Parisi-Wu approach, rapidly fluctuating chaotic fields arise. The expectation of the underlying potentials yields the currently observed dark energy density.

The advantage of this new chaotic model is that many of the questions raised in the introduction seem to have natural solutions. The cosmological constant problem is avoided; in our model the right order of magnitude of vacuum energy is naturally produced if we interpret the chaotic dynamics in terms of vacuum fluctuations allowed by the uncertainty relation, for a given finite age of the Universe. The cosmological coincidence problem is also avoided, since in our model the generated dark energy is no longer constant, but thins out with the expansion of the Universe in the same way as the energy density of the dominating species (matter or radiation). In spite of that, the (classical) equation of state of the chaotic component is close to $w=-1$, and can account for the accelerated expansion of the Universe, provided there is late-time symmetry breaking. The chaotic fields are physically interpreted in terms of vacuum fluctuations. As such

seems to indicate little or no dark energy at the time of last scattering [23]. Galaxy formation is disturbed as well if there is too much dark energy [24].

they can temporarily violate energy conservation, but quantum mechanical expectations are fully compatible with the Friedmann equations.

The physical relevance of our model is emphasized by the observation of a large number of numerical coincidences between local minima in the dark energy landscape and running standard model coupling constants evaluated at the known fermionic and bosonic mass scales. It thus appears that chaotic fields have the potential to fix and stabilize fundamental parameters and to select the physically relevant vacuum state out of infinitely many possibilities.

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APPENDIX: GENERAL TSCHEBYSCHIEFF MAPS

Our approach can be easily generalized to Tschebyscheff maps of arbitrary order N . One has $T_1(\Phi) = \Phi$, $T_2(\Phi) = 2\Phi^2 - 1$, $T_3(\Phi) = 4\Phi^3 - 3\Phi$, generally $T_N(\Phi) = \cos(N \arccos \Phi)$ with $\Phi \in [-1, 1]$. A Tschebyscheff map of order N is conjugated to a Bernoulli shift of N symbols, it is ergodic and mixing for $N \geq 2$. It exhibits the strongest possible chaotic behavior that is possible for a 1D smooth map, characterized by a minimum skeleton of higher-order correlations [20].

It is useful to consider both positive and negative Tschebyscheff maps and to define

$$T_{-N}(\Phi) := -T_N(\Phi). \quad (\text{A1})$$

The behavior of T_{-N} under iteration is identical to that of T_N up to a sign, the trajectory of T_{-N} differs by a constant sign (N even) or an alternating sign (N odd) from that of T_N .

Equation (13) implies that the maps T_N correspond to potentials V_N given by

$$V_N(\varphi) = \frac{1-\alpha}{\tau} \left\{ -\frac{1}{2}\varphi^2 + p_{max} \int d\varphi T_N\left(\frac{\varphi}{p_{max}}\right) \right\} + \text{const.} \quad (\text{A2})$$

In particular, one obtains for $N = \pm 1, \pm 2, \pm 3$

$$V_{\pm 1}(\varphi) = \frac{1-\alpha}{\tau} \left(-\frac{1}{2}\varphi^2 \pm \frac{1}{2}\varphi^2 \right) + \text{const.}, \quad (\text{A3})$$

$$V_{\pm 2}(\varphi) = \frac{1-\alpha}{\tau} \left[-\frac{1}{2}\varphi^2 \pm \left(\frac{2}{3p_{max}}\varphi^3 - p_{max}\varphi \right) \right] + \text{const.}, \quad (\text{A4})$$

$$V_{\pm 3}(\varphi) = \frac{1-\alpha}{\tau} \left[-\frac{1}{2}\varphi^2 \pm \left(-\frac{3}{2}\varphi^2 + \frac{1}{p_{max}^2}\varphi^4 \right) \right] + \text{const.} \quad (\text{A5})$$

Of physical relevance are the expectations of these potentials, formed with respect to the ergodic dynamics. Since negative and positive Tschebyscheff maps generate essentially the same dynamics, up to a sign, any physically relevant expectation should also be the same for T_N and T_{-N} , up to a possible sign. For all N , this symmetry condition fixes the additive constant to be

$$\text{const} = + \frac{1-\alpha}{\tau} \frac{1}{2} \langle \varphi^2 \rangle. \quad (\text{A6})$$

With this choice one obtains the following formulas for the self-energy which are fully symmetric under the transformation $N \rightarrow -N$:

$$\langle V_{\pm 1}(\varphi) \rangle = \pm \frac{1-\alpha}{\tau} \frac{1}{2} \langle \varphi^2 \rangle, \quad (\text{A7})$$

$$\langle V_{\pm 2}(\varphi) \rangle = \frac{1-\alpha}{\tau} \left(\frac{2}{3p_{max}} \langle \varphi^3 \rangle - p_{max} \langle \varphi \rangle \right), \quad (\text{A8})$$

$$\langle V_{\pm 3}(\varphi) \rangle = \pm \frac{1-\alpha}{\tau} \left(-\frac{3}{2} \langle \varphi^2 \rangle + \frac{1}{p_{max}^2} \langle \varphi^4 \rangle \right). \quad (\text{A9})$$

Written in terms of the dimensionless field variable $\Phi = \varphi/p_{max}$ this is

$$\langle V_{\pm 1}(\varphi) \rangle = \pm \frac{1-\alpha}{\tau} p_{max}^2 \frac{1}{2} \langle \Phi^2 \rangle, \quad (\text{A10})$$

$$\langle V_{\pm 2}(\varphi) \rangle = \frac{1-\alpha}{\tau} p_{max}^2 \left(\frac{2}{3} \langle \Phi^3 \rangle - \langle \Phi \rangle \right), \quad (\text{A11})$$

$$\langle V_{\pm 3}(\varphi) \rangle = \pm \frac{1-\alpha}{\tau} p_{max}^2 \left(-\frac{3}{2} \langle \Phi^2 \rangle + \langle \Phi^4 \rangle \right). \quad (\text{A12})$$

For Tschebyscheff maps of arbitrary order N one obtains

$$V_{\pm N}(\varphi) = \frac{1-\alpha}{\tau} p_{max}^2 \left\{ -\frac{1}{2}\Phi^2 \pm \int \cos(N \arccos \Phi) d\Phi \right\} \quad (\text{A13})$$

$$= \frac{1-\alpha}{2\tau} p_{max}^2 \left\{ -\Phi^2 \pm \left(\frac{1}{N+1} T_{N+1}(\Phi) - \frac{1}{N-1} T_{N-1}(\Phi) \right) \right\} + \text{const} \quad (\text{A14})$$

$$\langle V_{\pm N}(\varphi) \rangle = (\pm 1)^N \frac{1-a}{2\tau} p_{max}^2 \left\{ \frac{1}{N+1} \langle T_{N+1}(\Phi) \rangle - \frac{1}{N-1} \langle T_{N-1}(\Phi) \rangle + C \right\}. \quad (\text{A15})$$

For uncoupled Tschebyscheff maps with $|N| \geq 2$, any expectation of an observable $A(\Phi)$ is given by Eqs. (19) and (20). For $\alpha \neq 0$ the invariant density changes in a nontrivial way, but expectations can still be easily calculated numerically by long-time iteration of the coupled map lattice.

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