Landskepticism or Why Effective Potentials Don’t Count String Models

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Abstract

This paper is a synthesis of talks I gave at the Cargese Workshop in June 2004 and the Munich Conference on Superstring Vacua in November 2004. I present arguments which show that the landscape of string theory is not a well established feature of the theory, as well as a brief discussion of the phenomenological prospects of the landscape and the use of the anthropic principle.
1 Introduction

This paper is a synthesis of my summary remarks at the Cargese School, and my talk at the Munich conference. Its main message is that the landscape of string theory is far from an established fact. I point out that a hypothetical rigorous formulation of the landscape will require us to construct mathematical models of both the Big Bang and Eternal inflation. I emphasize that without such a rigorous formulation we are unlikely to be able to calculate the properties of the landscape with sufficient accuracy to make comparisons with experiment. I briefly review the phenomenological difficulties of the landscape idea, and make some comments about the anthropic principle.

The reader should be warned in advance that my personal bottom line on this subject is that the Landscape probably does not exist. However, I have tried to give it the maximal benefit of the doubt. Thus in many sections of this paper I will be trying as hard as possible to think of what the landscape might mean in a real theory of quantum gravity, which has the sort of solidity we associate with perturbative string theory, M-theory on Calabi-Yau manifolds times a torus, AdS/CFT, or Type 0 string theory in 1 + 1 dimensional linear dilaton space-times. I hope that the encouraging tone of some of these sections will not distract the reader from my main message.

One other caveat. I use the term FRW cosmology as a short-hand for a Friedman-Robertson-Walker cosmology which does not undergo eternal acceleration.

2 Effective actions in string theory

Part of the confusion in the debate about the landscape is the conflation of two notions of effective action, both of which occur in quantum field theory. The first is the 1PI effective action, which is an exact summary of the entire content of a field theory. Knowledge of it enables us to construct all of the correlation functions of the theory, in any of its vacuum states. In perturbation theory, we can compute the 1PI action around any vacuum state, and get the same result. It is important to realize that there is NO known analog of the 1PI action in string theory, which applies to all solutions of the theory, and enables us to pass from one to the other by “shifting the field”.

The other concept of effective action in field theory is the low energy or Wilsonian action.
This is defined, either in a single vacuum, or in a set of quasi-degenerate vacua whose energy density differences (as well as the heights of the barriers between them) are small compared to some cutoff scale. The Wilsonian action only contains degrees of freedom whose fluctuations are significant at these low energy scales. It is important to avoid using it when the conditions of its validity are violated. For example, in minimally SUSic QCD with \( N_F < N_C - 1 \), the low energy degrees of freedom consist of a meson superfield \( M \) and one can compute the exact low energy superpotential of \( M \). The low energy Kahler potential is canonical. If one uses this low energy Lagrangian to compute the energy density of states with a given expectation value of \( M \), this Lagrangian gives a divergent answer at the symmetric point of vanishing eigenvalue. This is not the correct physical answer. The true Kahler potential (in the 1PI sense) of \( M \) is modified at the origin of moduli space, and the energy density there is really of order the QCD scale. We should stop paying attention to the predictions of the Wilsonian action when they are outside its range of validity.

In string theory, the effective actions we compute are analogous to Wilsonian actions, but their range of validity is even more constrained. In particular, the stringy derivation of the effective action views it as a tool for calculating boundary correlation functions in a \textit{fixed asymptotic space-time background}. We tend to forget this because, particularly in situations with a lot of SUSY, the leading low energy term in the effective action is independent of the background. This fosters the illusion that different backgrounds can be viewed as \textit{vacuum states of the same theory}. In fact, as emphasized in [2], the italicized phrase is borrowed from quantum field theory, and refers to concepts that depend entirely on the separation between IR and UV physics of that formalism. In string theory/quantum gravity, UV and IR physics are much more intimately entangled, and the concept of different vacuum states of the same underlying string theory Hamiltonian is much more circumscribed.

When we have a continuous moduli space of super-Poincare invariant S-matrices, we can do experiments[1] at one value of the moduli, which are sensitive to the S-matrix at other values of the moduli\(^1\). Note however, that the only Hamiltonian form we have for such models is in light cone frame, where different values of the moduli correspond to different Hamiltonians. Similarly, a moduli space of correlation functions on the boundary of Anti-deSitter space, corresponds to a one parameter set of different Hamiltonians, rather than different superselection sectors of the same Hamiltonian. The notion of different vacua of the same theory, in any of the senses that this is meant in quantum field theory, is simply not applicable to theories of quantum gravity,

\(^1\)These experiments are much more difficult than they would be in a SUSic quantum field theory, without gravity.
beyond the very limited context of continuous moduli spaces of Super-Poincare invariant S-matrices.

The recognition that changes in background asymptotics correspond to changes in the Hamiltonian, rather than changes of superselection sectors for a given Hamiltonian goes back to [3], and has become commonplace with the advent of AdS/CFT. Changes in non-normalizable modes of bulk fields\(^2\) add relevant terms to the Hamiltonian. Changes in the (negative) cosmological constant correspond to changes in the fixed point which defines the boundary CFT, and thus to a completely different set of high energy degrees of freedom.

These facts lead us\(^3\) to be suspicious of attempts to find new string theory models by patching together an effective potential using the degrees of freedom of e.g. Type II string theory in flat space-time. Indeed, once we include gravitational effects, the low energy action itself gives us reason to be suspicious of meta-stable de Sitter vacua constructed in this way.

In my opinion, the concept of an effective potential on moduli space as a tool for finding string models of gravity, is a snare and a delusion, fostered by wishful thinking, and without regard to the actual evidence in front of us. There is no evidence for this concept in solid string theory calculations, and lots of evidence against it\(^4\). As I will explain below, the utility of the effective potential concept in low energy effective non-gravitational physics is not an argument for its validity in string theory. Gravitational low energy effective field theory shows us the limitations of the effective potential concept, and is a good tool for demonstrating that different solutions of a low energy effective action do not define different states of the same quantum theory unless they have the same space-time asymptotics. Rigorous quantum string theory constructions, including \(c \leq 1\) matrix models, Matrix Theory, and AdS/CFT, confirm these conclusions.

The fact that the effective actions we compute in string theory are not analogs of a 1PI action is easily confirmed by calculations we can do in well defined string theories. We can consider, among many examples, the Type IIB supergravity solutions \(M^{1,d} \times K^{9-d}\) and \(\text{AdS}_5 \times S^5\), where \(K\) is a SUSY preserving compact manifold. The effective actions appropriate to these different examples, depend on different kinds of variables. Different \(K\) give rise to different

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\(^2\)Note that if we want to make finite rather than infinitesimal changes we must restrict our attention to Breitenlohner-Freedman allowed tachyon fields.

\(^3\)well, at least they lead me

\(^4\)Those who think these judgements too harsh, should recall that I’ve spent a considerable fraction of my career writing about effective potentials on moduli space. I’ve also spent the last four years trying to draw the community’s attention to what I believe is the true situation - with frustratingly little effect. Perhaps some over the top rhetoric is in order.
moduli fields. The constant value of the dilaton is a variable in the Minkowski actions, while it is fixed and corresponds to a parameter in the quantum Hamiltonian in AdS. The AdS space has singleton fields, which have no Minkowski analog. Higher order corrections to these actions involve virtual contributions from states that only exist for some choices of $K$. Even the terms involving fields common to all the different non-compact space-times will, in general, have different coefficients in the different backgrounds (except for terms which are completely determined by symmetries). There is no way to change $K$ from e.g. a torus to a Calabi-Yau manifold, or to change Minkowski into AdS, by “shifting a field in the quantum effective action”. And, indeed, these actions are computed, by calculating boundary correlation functions with a fixed asymptotic boundary geometry, and finding an action functional whose tree level solutions match the low energy expansion of these boundary correlators.

2.1 The effective potential with and without gravity

Consider a simple scalar field theory model with a potential that has one positive energy minimum and one at exactly zero energy. We can imagine the zero energy minimum to be related to an exactly supersymmetric string theory vacuum state. Let the scale of all parts of the potential in the region of the two minima be small compared to both the string and Planck scales. It is then a correct deduction from low energy effective field theory, that a subclass of solutions of this Lagrangian, including some in which the field is homogeneous and equal to its false vacuum value over a large region of space, and asymptotes to the true vacuum at infinity, can be studied without worrying about gravity. However, *it is not true that all solutions whose energy density is small compared to the Planck scale can be studied in this way*. We can certainly excite a bubble of the meta-stable minimum, and watch it decay, but if the size of the bubble gets too large, we run into problems.

In order to exhibit a meta-stable dS space, we have to excite a region of order the putative Hubble radius into the metastable minimum. Old results of Guth and Farhi[4] show that the external observer can never verify the existence of the inflating region. A black hole forms around it. Any observer in the asymptotically flat region, who tries to jump into the black hole to find inflation, first encounters a singularity. So effective field theory tells us that, given a well defined, Poincare invariant S-matrix for gravity, we can find and explore meta-stable positive energy density minima of an effective potential, but not the meta-stable dS spaces that these minima have been thought to imply. That is, the Guth-Farhi results suggest that *the stable asymptotically flat vacuum state does not have excitations which correspond to metastable dS*
vacua, even when it has an effective potential with positive energy meta-stable minima. Rather, it has excitations in which fields are excited into metastable minima only over regions small compared to the Hubble radius at those minima. The attempt to create larger regions succeeds only in creating black holes. Similar remarks apply to stable AdS vacua, whether they are supersymmetric or not.

We see that in gravitational theories, the criterion for the validity of non-gravitational effective field theory reasoning depends on more than just the value of the energy density in Planck units. When the Schwarzschild radius of a region exceeds its physical size (in the approximation in which gravity is neglected), a black hole forms. Effective field theory remains valid outside the black hole horizon (if it is large enough), but not inside. In the above example, no external observer can probe the putative dS region, without first encountering a singularity.

We also see another example of the principle that the solutions of the same effective equations of motion may not reside in the same quantum theory. This is in stark contrast to non-gravitational quantum field theory with the same potential. There, the meta-stable minimum of the potential represents an unstable Lorentz invariant state of the Hamiltonian whose ground state is the zero vacuum energy state. We can create arbitrarily large regions where the field lies in its meta-stable minimum.

The results of Coleman and De Lucia[5], on vacuum tunneling in the presence of gravity, give us a sort of converse to these observations. Given the same effective Lagrangian we used in the previous paragraph, we can assume the existence of the meta-stable dS space, and ask what it decays into. Here there is a surprise, particularly for those who constantly repeat the mantra “dS space decays into flat space”. In fact, the analytic continuation of the CDL instanton is a negatively curved Friedmann Robertson Walker cosmology, which (if the potential has a stable zero cosmological constant minimum), is asymptotically matter dominated. Although it locally resembles flat space on slices of large cosmological time, its global structure is completely different. In particular, an attempt to set up an asymptotically Minkowskian coordinate system, starting from the local Minkowski frame of some late time observer, inevitably penetrates into regions where the energy density is high (the energy density is constant on slices of constant negative spatial curvature, and is of order the energy density of the false vacuum at early FRW times). Another indication that the local point of view is misleading is that if we make the vacuum energy even slightly negative, the asymptotic state is radically altered: we have a Big Crunch rather than a small perturbation of Minkowski space. Models of quantum gravity with AdS asymptotics cannot have meta-stable de Sitter (or Minkowski) excitations.
Thus, both analysis of creation and decay of meta-stable dS states, suggests that if a potential has a stable minimum with vanishing cosmological constant, and another with positive energy density, the Minkowski solution and the meta-stable dS solution are simply not part of the same theory. There remains a possibility of the existence of a theory with a stable, matter dominated, FRW cosmological solution with a meta-stable dS excitation. The problem with this is that it is very unlikely that we can make a reliable exploration of this scenario within the realm of low energy effective field theory.

The CDL instanton solution is non-singular. The $a = 0$ point of the FRW coordinate system is just a coordinate singularity marking the boundary between the FRW region and a region of the space-time which continues to inflate. However, arbitrary homogeneous perturbations of the CDL solution have curvature singularities. Further, there is a large class of localized perturbations which evolve to Big Crunch singularities, rather than passing smoothly through the $a = 0$ point. For example, consider a localized perturbation on some hyperbolic time slice a finite proper time prior to $a = 0$. Let it be homogeneous in a large enough region, so that signals from its inhomogeneous tail cannot propagate to the $a = 0$ point. Then we will have a singularity. Below we will see an argument that singularities are associated with generic initial\(^5\) and final data in this geometry.

If we re-examine the Guth-Farhi argument in the FRW context, we see that it continues to hold until we go back in time to a point where the cosmic energy density is of order the barrier to the meta-stable minimum. At high enough energy density, there are FRW solutions in which the field classically evolves into the meta-stable minimum. It will then decay by tunneling, into the FRW continuation of the CDL instanton. Generic FRW solutions (including arbitrary homogeneous perturbations of the CDL instanton) have curvature singularities at a finite cosmic time in the past. To establish their existence as genuine theories of quantum gravity one must go beyond effective field theory, and probably beyond perturbative string theory.

Freivogel and Susskind\(^6\) have suggested a scattering theory in which asymptotic states are associated with incoming and outgoing wave perturbations of the nonsingular, time symmetric Lorentzian continuations of the, CDL instantons for the various meta-stable vacua of string theory. They claim that in this framework, the breakdown of effective field theory is avoided, as long as the effective potential is everywhere smaller than Planck scale. I find this suggestion interesting, but it is not based on reliable calculations, and there is a strong indication that it does not work. If one considers black holes with radius larger than the dS radius, formed in the

\(^5\)In mentioning initial data, we are referring to the time symmetric Lorentzian CDL instanton.
remote past of the FRW part of the time symmetric Lorentzian CDL geometry, it is hard to see how these solutions asymptote to the future CDL geometry without encountering a singularity. Near the \( a = 0 \) point where the FRW coordinates on the CDL solution become singular, most of the space-time is isometric to dS space at its minimal radius. If we have formed a black hole in the past, with radius much larger than this, then the entire spacetime must end in a singularity.

Thus, I would claim that unlike asymptotically flat or AdS space-times, we have no reliable effective field theory argument that there are an infinite number of states in space-times that asymptote to the time symmetric Lorentzian CDL instanton. An infinite number of states is a minimal requirement for the existence of a quantum theory that can make precise mathematical predictions, which can be self consistently measured in the theory. We must understand the nature of the singularities in these perturbations of the CDL instanton geometry before we can conclude that the framework makes sense.

I think it is more likely, that meta-stable dS vacua exist\(^6\) only in the context of a Big Bang cosmology. If we consider the problem of accessing a meta-stable dS minimum of an effective potential, at times when the cosmic energy density is of order the barrier height of the potential, then the Guth-Farhi problem does not appear to exist. Starting from a Big Bang singularity, one can find homogeneous solutions where a scalar field wanders over its potential surface at a time when the energy is higher than the barriers between minima, and then settles in to a meta-stable minimum with positive cosmological constant. One must understand the Big Bang to make a reliable theory of such a situation, but apart from that the solution is non-singular. In particular, the problem of large black holes in the initial state, is not present for this situation. The difficulties are all associated with understanding the Big Bang singularity, and with Eternal Inflation, which I will discuss below. However, I will also point out below that generic final states in the CDL instanton geometry cannot have evolved from a meta-stable dS tunneling event - only a finite number of states are compatible with this history.

To summarize, it is clear from semi-classical calculations alone, that the concept of a vacuum state associated with a point in scalar field space is not a valid one in theories of quantum gravity. A given low energy effective field theory may have different solutions which do not have anything to do with each other in the quantum theory. One solution may be a classical approximation to a well defined quantum theory, while the other is not. It seems likely that the context in which we will have to investigate the existence or non-existence of

\(^6\)if they exist at all!
the landscape is Big Bang cosmology. This is the only situation in which we can “reliably” construct a universe which gets stuck in a meta-stable dS minimum.

Thus, I claim that if string theory really has a multitude of meta-stable dS states, then the exact theory into which they fit is a theory of a Big Bang universe which temporarily gets stuck, with some probability, in each of these states. This is ultimately followed by decay to negatively curved FRW universes. These FRW universes have four infinite dimensions and 6 or 7 large compact dimensions, which are expanding to infinity. It is clear that the probability for finding a particular dS vacuum is partly determined by the density matrix at the Big Bang and not just by counting arguments. In the next subsection I will describe existing proposals for the cosmological distribution of vacua. My main point here is that the nature of the Big Bang will have to be addressed before we can hope to understand the correct statistics of what are called stringy vacua.

2.2 Eternal inflation

The string landscape seems to fit in well with older ideas which go under the name of eternal inflation, the self reproducing inflationary universe, etc. The simplest model which exhibits this sort of behavior is one with a single scalar field with two minima, one with positive vacuum energy and the other with vanishing energy. One considers the expanding branch of the meta-stable dS universe with the field in the false minimum. In the quantum field theory approximation this seems to produce an ever-expanding region of space and one allows the dS space to decay by CDL bubble formation, independently in each horizon volume. In the eternal inflation picture, one tries to interpret the result as a single classical space-time. One obtains a Penrose diagram with a future space-like boundary. The future space-like boundary is fractal, with regions corresponding to singularities\textsuperscript{7} as well as FRW asymptotics, interspersed in a causally disconnected way. In the landscape there will be many different singular regions of the boundary, as well as many different FRW regions. Advocates of the landscape/eternal inflation picture then make the analogy to maps of the observable universe, with different causally disconnected regions being the analogs of different planets. Physics it is said, depends on “where you live” and organisms like ourselves can only live in certain regions of the map. A key difference between different minima of the effective potential in eternal inflation, and different planets is that we cannot, even in principle, communicate with causally disconnected regions of the universe.

\textsuperscript{7}The singular regions correspond to decays to parts of the potential where the vacuum energy is negative. They would exist in the landscape context, but not in the simple model we are discussing.
How is one to interpret such a picture in terms of conventional quantum mechanics? I believe that the fundamental clue comes from the principle of Black Hole Complementarity[7]. Black holes also present us with two regions of space-time which are causally disconnected. Hawking showed long ago that this was an artifact of the semi-classical approximation, and that black holes return their energy to the external space-time in which they are embedded. If we assume that the region behind the horizon has independent degrees of freedom, commuting with those in the external space, then we are confronted by the information loss paradox.

String theorists have believed for some years now, that this is not the case. The principle of Black Hole Complementarity is the statement that the observables behind the horizon do not commute with those in the external space-time. For a large black hole, (and for a long but finite time as measured by the infalling observer), these two sets of observables are both individually well described by semiclassical approximations, but the two descriptions are not compatible with each other.

Fischler and I tried to relate this principle to the Problem of Time[8]. In the semi-classical quantization of gravity one attempts to solve the Wheeler-DeWitt equation

\[ \mathcal{H}\Psi = 0 \]

with an ansatz

\[ \Psi = e^{iS}\chi(t,\phi), \]

where \( S \) is the action of some classical space-time background solution, and \( \chi \) is the wave functional of a quantum field theory in this space-time

\[ i\partial_t\chi = H(t)\chi. \]

The (generally time dependent) Hamiltonian \( H(t) \) depends both on the choice of classical background, and on the particular time slicing chosen for that background. For example, for the Schwarzschild background we could choose \( H_{Sch} \), the Schwarzschild Hamiltonian, or some time dependent \( H(t) \) where \( t \) is the proper time of a family of in-falling observers. Even ignoring subtle questions of whether these two Hamiltonian evolutions act on the same Hilbert space, it is clear that they are different and that \([H(t), H_{Sch}] \neq 0 \) at any time \( t \). It is therefore not surprising that the semi-classical observables of different observers do not commute with each other.

Given this description of black holes, it is natural to conjecture that a similar phenomenon occurs for any space-time with horizons. In [8] this was called Cosmological Complementarity.
for asymptotically dS spaces. E. Verlinde has suggested the name Observer Complementarity for the general case.

If we apply this logic to Eternal Inflation, we obtain a picture quite different from the original description of these space-times. We simply associate a single Hilbert space and many different (generally time dependent) Hamiltonians to the fractal Penrose diagram. Each Hamiltonian is associated with the causal patch of a given observer. Mathematically, the situation can be equally well described by saying we have a collection of different theories of the universe. There is a philosophical cachet, associated with the phrase, “physics depends on where you live in the multiverse”, which is absent from this alternative way of describing the physics.

In the formalism described in [6] this is almost precisely what is conjectured. For each meta-stable dS point \( L \) in the landscape, which can decay into the Dine-Seiberg region of moduli space, and for each (typically 10 or 11 dimensional) Super-Poincare invariant solution \( V \) of string theory “into which \( L \) can decay” there is a different unitary S-matrix \( S_{L,V} \). It is claimed that each of these S-matrices contains all the physics of the landscape, because there is a canonical way to compute the unitary equivalence \( U \) in the formula

\[
S_{L,V} = U_{L,V,L',V'} S_{L',V'} U_{L',V,L}^\dagger.
\]

The statement that there is a theory of eternal inflation in which all of the points \( L \) are meta-stable states is really the statement that the theory contains a canonical algorithm for computing the “gauge transformations” \( U \). The authors of [6] claim that all of the meta-stable dS states will show up as resonances in every S-matrix, \( S_{V,L} \). This claim is plausible if the S-matrices are indeed related by unitary conjugation. The spectrum of the S-matrix is then gauge invariant. In ordinary scattering theory, time delays, which are related to resonance lifetimes, are related to the spectrum of the S-matrix. In the eternal inflation context, there is no universal notion of time for the different asymptotic states, so more work is necessary to understand these concepts.

If indeed the information about each meta-stable dS vacuum can be extracted from the spectral density of a given S-matrix \( S_{V,L} \), and if we can find a reliable framework, for defining and calculating these S-matrices, then the landscape will have a mathematical definition. From a practical point of view however, we just say that string theory gives us an algorithm for constructing models of the world (in the landscape context this means choosing particular stable or meta-stable minima) which is not unique and that we are trying to use data to constrain

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8. We can also consider initial and final states corresponding to different CDL instantons. In [6] these are claimed to be different gauge copies of the same information in the S matrices. I will mention a different interpretation below.

9. For the moment, no approximate statement of what this algorithm is has been proposed. I would conjecture that, if the formalism makes sense, the transition amplitudes between two different FRW spacetimes, mentioned in the previous footnote, provide the algorithm for calculating the \( U \) mappings.
which model we choose. In the penultimate section, where I summarize the phenomenological difficulties of the landscape, I will describe the situation in the language of the approximate Hamiltonian of a given metastable dS observer, rather than that of eternal inflation.

I have emphasized the problems with this S-matrix point of view and suggested the alternative notion that the landscape could only make sense in the context of Big Bang Cosmology. One must thus understand how to describe the initial states. There are two possibilities, either there is some principle which picks out a fixed initial state$^{10}$ at the Big Bang, or there is a generalized S-matrix in which we relate a particular state at the Big Bang to a particular linear combination of final scattering states in one of the FRW backgrounds defined by a decaying meta-stable dS space. In a manner analogous to Freivogel and Susskind, one would conjecture that the descriptions in terms of different future FRW backgrounds are unitarily equivalent to each other, by unitary transformations which do not respect locality. Since we are unlikely to have much control over the initial conditions at the Big Bang, one should choose the in-state at the Big Bang to be a high entropy density matrix. Thus, the practical difference between the Hartle-Hawking (S-vector) and S-matrix proposals is the entropy of the initial state.

I know of two proposals for the initial density matrix at the Big Bang, which might lead to a set of cosmological selection rules for meta-stable points in the landscape. The first, modular cosmology$^{[9]}$ postulates an early era in which the universe can be described semi-classically, but the potential on moduli space is smaller than the total energy density. The metric on moduli space then provides a finite volume measure. Furthermore, the motion of the moduli is chaotic. These facts suggest that the probability of finding the universe in a given meta-stable minimum of the potential is the volume of the basin of attraction of that minimum, divided by the volume of moduli space.

Holographic cosmology$^{[10]}$ gives an alternative view of the initial state of the universe, as a “dense black hole fluid” where standard notions of local field theory do not apply. The model contains two phenomenological parameters, which govern the transition between this phase of the universe and a normal phase in which the field theory description is valid. There are indications that the transition occurs at an energy density well below the unification scale of standard model couplings. We might then expect a transition directly into a state with most of the moduli frozen. In order for the model to provide an adequate account of the fluctuations in the CMB, one must have at least one “active” modulus at these low energies, which can provide for a modest number of e-folds of inflation. In such a model, minima of the potential on moduli

$^{10}$e.g. the Hartle Hawking Wave Function of the Universe.
space, with energy higher than the scale at which a field theoretic description of the universe is possible, cannot make any sense. At best a small class of low energy minima could be compatible with holographic cosmology\textsuperscript{11}. If holographic cosmology is compatible with landscape ideas, the probability of accessing a particular minimum will be determined by quantum gravitational considerations, far removed from effective field theory. In this framework, the dense black hole fluid is stable and is the most probable state of the universe. A normal universe like our own is determined by a somewhat improbable initial condition, but one expects the maximum entropy initial state that does not collapse to a dense black hole fluid. The survival probability of a given normal state depends on both the properties of the black hole fluid and the low energy physics of the normal state, so the determination of the most probable meta-stable minimum would be a complicated quantum gravitational calculation.

In both of these classes of models, simple enumeration of meta-stable minimum is not a good account of the physical probability distributions.

3 Lessons from de Sitter’s anti

For completeness, I want to record here a number of results about AdS space-times, which bear on the questions at hand. In my opinion, the most important result of the AdS/CFT correspondence is that \textit{the negative cosmological constant is a discrete input parameter, which controls the high energy behavior of the theory. It cannot be affected by low energy renormalization.} In $\text{AdS}_5 \times \text{S}^5$ the cosmological constant in Planck units is $N^{1/4}$. It is determined by the rank of the gauge group. This is not an accident. The coincidence of the AdS black hole entropy formula and that of CFT, shows us that the c.c. in Planck units is always a power of the c-coefficient in the CFT entropy formula. Note that this is completely consistent with low energy bulk effective field theory. In that formalism, the c.c. is an undetermined parameter, which we have to fix by a renormalization condition. AdS/CFT tells us that, at least for negative c.c., this normalization can only be understood in the high energy theory. The effective field theory must be tuned to reproduce that high energy input\textsuperscript{12}.

Everything we know about semi-classical AdS physics is consistent with this. There is no such thing as a “meta-stable conformally invariant state” in CFT. Correspondingly, there are

\textsuperscript{11}The potentials calculated in the landscape have no indication of a cut-off at energy scales far below the unification scale. This suggests that the two theoretical frameworks are not compatible, but I am trying to avoid jumping to conclusions.

\textsuperscript{12}In 1 + 2 dimensional AdS, we can even describe this as forcing the low energy effective theory to preserve the correct Virasoro symmetry algebra of the high energy theory.
no tunneling solutions from metastable points with any vacuum energy, into AdS space. The view of the c.c. in an AdS “vacuum” as the minimum of a low energy effective potential finds confirmation neither in AdS/CFT, nor in semiclassical tunneling calculations.

AdS/CFT also gives us some clues about how the minimization of effective potentials might be related to more robust concepts in quantum gravity[?]. The key idea is the holographic renormalization group[19], which reinterprets the space-time equations for domain wall solutions to the renormalization group flow in the boundary field theory. Certain minima and Breitenlohner-Freedman allowed maxima of the effective potential of a supergravity theory, correspond to fixed points of the RG flow. Thus, the correct quantum gravitational notion of defining a boundary conformal field theory can, in certain limits, be related to finding stationary points of an effective potential.

There are a large number of puzzling features about this correspondence, as well as features which indicate that conventional ideas about the effective potential are unfounded. I list some of them, in random order:

- There are often stationary points of the potential which do not appear to have any meaning in boundary field theory. These include points with tachyons that violate the BF bound. These could only be non-unitary CFTs, but they appear in the flow initiated by relevant perturbations of a perfectly unitary theory. Although there are no examples with meta-stable, positive energy, minima, their possible role in the holographic RG is completely obscure.

- The holographic RG makes it clear that different stationary points of the same potential do not represent different states of the same theory. The theory defined at the IR end of the flow has fewer fundamental degrees of freedom than the UV theory. It is defined by decoupling. In global dS space, the flow is cut off by the compact sphere, and we never really lose the UV degrees of freedom. The perturbed theory has the same asymptotic value of the c.c. as the unperturbed one. The global space-time picture of these domain walls is that of stable defects imbedded in a given AdS space-time, rather than walls between two different asymptotic space-times.

- There are domain walls in the low energy effective theory, between positive curvature AdS stationary points, which have no RG flow interpretation, because the CFTs have no relevant operators.

- The RG interpretation of the effective potential suggests that only the behavior near some
of its stationary points has physical relevance. It is well known that the RG has a vast *scheme dependence* which makes most aspects of the RG flow simply conventions. It is only in the presence of large symmetry groups (integrable flows) that there is physical relevance to the behavior of the RG flow far from fixed points. Note that the effective potential, away from the stationary points determining the ends of the flow, would then have a similar scheme dependence.

There is a very specific puzzle posed by the attempt to make the landscape compatible with AdS/CFT. In Type IIB string compactifications with fluxes up to order $N$ on manifolds with third Betti number $b_3$, one claims to find stable AdS minima with c.c. as small as $N^{-b_3}$. This poses three problems: one must find CFTs corresponding to the AdS space with radius much larger than both the Planck and the string scale. Furthermore, it is claimed that all of these vacua are in the region where “string perturbation theory” is a good approximation.

The AdS radius in Planck units is a measure of the $c$ coefficient in the $2 + 1$ CFT entropy formula $S(E) = c(ER)^{2/3}$. In $2 + 1$ dimensions we would expect any CFT to be in the IR basin of attraction of a RG flow from a Gaussian fixed point, and we would expect to see at least as much entropy at the Gaussian point as we do at the fixed point representing the putative landscape state. Silverstein[18] has suggested that the relevant Gaussian fixed points are the world volume theories of branes carrying the fluxes in the string compactification. The entropy of these theories appears to scale like $b_3N^2$ rather than $N^{b_3}$. Silverstein suggested that string junctions should (for reasons that are not immediately apparent) be included among the UV degrees of freedom in the world volume gauge theory. These transform in the $\otimes b_3N_i$ under the $\prod U(N_i)$ gauge group of the branes, and could provide the requisite entropy.

Gauge theories with large representations do have a large number of fixed points in $2 + 1$ dimensions, which are easily accessible. The classical running of the gauge and quartic couplings can be balanced by the large quantum correction coming from loops of the field in the large representation. The usual vector large $N$ resummation gives a reliable calculation of the properties of these fixed points. The problem is that, since the anomalous dimensions at these points are close to those of the Gaussian model, a large radius interpretation of the CFT spectrum is not possible.

This is a general point which is worth emphasizing. The number of primary operators of fixed dimension grows exponentially with a power of the dimension in any CFT which is an order one deformation of a Gaussian model. This is much too rapid to be explained by Kaluza
Klein states in a large radius $AdS_k \times K_p$ compactification. In $\mathcal{N} = 4$ Super Yang-Mills theory with $g^2 N$ of order 1, we recognize this as an indication that “the AdS radius is of order the string scale”, and attribute the growth of the spectrum to excited string states. Here this is not possible. We are not in a regime where the planar diagram expansion makes sense. Indeed, the effect of the huge representation we invoked to explain the entropy is to make $g^2 N \ll 1$ at the fixed point.

The landscape predicts a plethora of AdS flux vacua at weak string coupling, with AdS radius ranging from $\sim 10$ to $N^{b_3}$ times the string scale. The CFT duals to these states have not yet been found. The search for them leads to a large number of $2 + 1$ CFTs which have no dual interpretation in terms of large radius SUGRA, or weakly coupled string theory on manifolds whose size is of order string scale. It seems to me that the continued search for the CFT duals of AdS landscape states is the most likely avenue for finding a rigorous justification for part of the landscape picture or for falsifying it.

4 Phenomenology of the landscape

For practical purposes, the landscape gives us a large set of alternative effective Lagrangians for describing the physics we have observed or will observe in our universe. These are parametrized by a collection of numbers, which include the dimension of space-time, the name, rank and representation content of the low energy gauge theory, the value of the cosmological constant, and the values of all the coupling constants and masses of fields in the Lagrangian\(^{13}\). These numbers can be collected together and viewed as a multidimensional probability space. In the supergravity approximation, we have a way of calculating an \textit{a priori} distribution for these numbers. Proponents of the landscape would claim that this is an approximation to some more exact distribution, though no-one has suggested a procedure for calculating the corrections. In the previous section I have suggested that early universe cosmology may make important modifications of the distribution of metastable minima. If it turned out that the distribution predicted the Lagrangian we observe with high probability, it would be a great triumph for string theory.

The value of the cosmological constant tells us that this is not the case. Weinberg’s bound\([11]\), which constrains cosmological parameters by insisting on the existence of galaxies,

\(^{13}\)In principle we could also have non-trivial conformal field theories in the low energy world, at least in some approximation.
has the form (in Planck units).

$$\Lambda \leq K \rho_0 Q^3,$$

where $\rho_0$ is the dark matter density at the beginning of the matter dominated era, $Q$ is the amplitude of primordial density perturbations at horizon crossing, and $K$ is a pure number of order 1.. For any reasonable values of the other parameters, this means $\Lambda$ is much smaller than the typical value found in the landscape.

It is clear then that we must supply additional data from experiment in order to fix our description of the world. The landscape framework supplies some theoretical guidance - it tells us that there are a finite number of possibilities, of order $10^{10^2}$ (to order of magnitude accuracy in the logarithm). Various authors [12] have begun to investigate the a priori distribution of properties like the gauge group and number of generations, the scale of supersymmetry breaking, the existence of large warp factors which give rise to large hierarchies of energy scales etc., assuming a uniform distribution on the space of minima. The hope is that correlations will become evident which will tell us that a small number of inputs is enough to extract the Lagrangian of the world we live in from the ensemble of Lagrangians the landscape presents us with.

The anthropic principle has also been invoked as an input datum to impose on the ensemble of Lagrangians. For its proponents, the attraction of this principle is that the answer to a single yes/no question, “Is there carbon based life?” puts strong constraints on a collection of parameters in the Lagrangian (assuming all others fixed at their real world values). This attraction may be an illusion. In a probability space, the characteristic function of any subset of data points is a single yes/no question. So physicists must ask if there is any special merit to the particular characteristic function chosen by the anthropic principle. This is a hard question to answer, because we do not have much theoretical understanding of life and intelligence, and we have no experimental evidence about other forms of life in the universe we inhabit.

If the typical life form resembles the great red spot on Jupiter rather than us, then this life form would think that the criterion that is most appropriate to apply in our universe is the Redspotthropic principle. To put this in a more positive manner: if considerations of carbon based life lead to explanations of the values of the fundamental parameters, then we are making a prediction about the typical form of life that our descendants will find when they explore the universe. It should look just like us, or at least be sufficiently similar that the criteria for its
existence are close to the criteria for ours. If instead, our descendants’ explorations show that the typical life form could tolerate much larger variations in the fundamental constants, then our so-called explanations would really be a fine-tuning puzzle. The Red Spot People could calculate and understand that we wouldn’t be there if the up quark mass were a little bigger, but they might reasonably ask “Who ordered them?”.

Of all soi disant anthropic arguments, Weinberg’s bound on cosmological parameters is the least susceptible to this kind of criticism. If there are no galaxies, there are no planets, no Red Spots, no Black Clouds, perhaps no conceivable form of life. Many physicists cite the numerical success of this bound as evidence that anthropic reasoning may have relevance to the real world. It is important to realize that this numerical success depends on keeping all other parameters fixed. Arkani-Hamed reported on unpublished work with Dimopoulos, which showed that if both $\Lambda$ and $M_P$ (really the ratio of these parameters to particle physics scales, which are held fixed) are allowed to vary subject only to anthropic constraints, then the preferred value of $\Lambda$ is larger than experimental bounds by many orders of magnitude. Similarly the authors of [13] following [14] argued that if both $\Lambda$ and $Q$ are allowed to vary then the probability of finding a universe like our own is of order $10^{-4}$. A contrary result was reported in [15], but only by assuming an a priori probability distribution that favored small values of $Q$.

In inflationary models of primordial fluctuations, the value of $Q$ depends on details of the inflaton potential at the end of slow roll. We would certainly expect this parameter to vary as we jump around the landscape. Similar remarks apply to $\rho_0$. Allowing $\rho_0$ to vary would further reduce the probability that the anthropic distribution favors the real world. Thus, at least with our current knowledge of the landscape, it seems likely that the numerical success of Weinberg’s bound in the landscape context is not terribly impressive\(^\text{14}\).

The greatest challenge to all methods of dealing with the landscape is the large number of parameters in the standard model which have to be finely tuned to satisfy experimental constraints. These include the strength of baryon, lepton and flavor violating couplings, $\theta_{QCD}$ and the values of many quark and lepton masses. Anthropic reasoning helps with some of these parameters, but not all, and is insufficient to explain the lifetime of the proton, the value of $\theta_{QCD}$ and many parameters involving the second and third generation quarks and leptons. From the landscape point of view, the best way to deal with this (in my opinion) is to find classes of vacua in which all of these fine tuning problems are solved, perhaps by symmetries. One can then ask whether there are enough vacua left to solve the cosmological constant problem. This

\(^\text{14}\)I cannot resist remarking that in the context of Cosmological Supersymmetry Breaking[16], only $\Lambda$ varies, and Weinberg’s bound retains its original numerical status.
might be a relatively easy task. One could then go on to see whether other features of this class
of vacua are in concordance with the real world.

It is clear that at a certain point in this process, if we don’t falsify the landscape easily,\(^{15}\)
we will run into the problem that current technology does not allow one to calculate the low
energy parameters with any degree of precision. Indeed, the error estimates are only guesses
because we don’t even know in principle how to calculate the next term in the expansion in
large fluxes. A more fundamental framework for the discussion of the landscape is a practical
necessity as well as a question of principle. I have suggested that if a rigorous framework for the
landscape exists, it is probably to be found in the context of a theory of a Big Bang universe
with Eternal Inflation, and future FRW asymptotics for any given observer. It is likely that
we will be unable to define more precise calculations of the properties of the landscape without
finding a rigorous mathematical definition of such a space-time.

The fundamental object in such a space-time would be a scattering matrix\[^{8}\] relating a
complete set of states at the Big Bang to states in Lorentzian CDL bubble space-times corre-
sponding to decays of meta-stable dS landscape states into the Dine-Seiberg region of moduli
space. The final states, in addition to particle labels, would carry indices \((V, L)\) describing
a particular dS minimum \(L\) and a particular “asymptotic vacuum”, \(V\), into which it decays.
Thus, we would have matrix elements

\[
S(I|V, L, p_i),
\]

where \(I\) labels an initial state at the Big Bang and \(p_i\) a set of “particle” labels for localized
scattering states in a given CDL bubble. An important unanswered question is whether the
S-matrix is unitary for each \((L, V)\)\(^{16}\) or only when all \((L, V)\) sectors are taken into account.
The first alternative is analogous to the proposal of [6].

The following argument has a bearing on this question. I have stated above that there
was no problem with an infinite number of final states for fixed \(L\). This is not necessarily the
case. If I extrapolate scattering data on \(I\) backwards, using the classical equations of motion,
and assuming a minimal finite energy for each particle, then all but a finite number of states
will encounter a space-like singularity before transition to the metastable dS regime. This is
the time reverse of the argument about black holes in the initial state of the time symmetric
CDL bubble. This suggests that for fixed \(L\), the matrix \(S(I|V, L, p_i)\) has finite rank: only a
\(^{15}\)e.g. by showing that the number of vacua left after all the other fine tuning problems are solved is too small
to solve the cosmological constant problem.
\(^{16}\)One would then invoke the existence of unitary mappings taking the different unitary S-matrices into each
other.
finite subspace of the space of all out states on the CDL geometry labeled by \( L, V \) would be allowed. This leads to a modification of the proposal of [6] in which only the S-matrix for fixed \( V \), keeping all possible values of \( L \), is unitary. However, there is no clear reason now to assume that \( V \) should be fixed, so perhaps only the full S-matrix is unitary.

A more disturbing conclusion is reached if one combines the claim of [17] that the number of \( L \) sectors is finite, with the above argument. One then concludes that the whole subsector of the S-matrix that has meta-stable dS resonances, also has finite rank, and that the entire landscape fits into a Hilbert space with a finite number of states. There will be an infinite rank matrix of direct transitions between the Big Bang and a final FRW state, which do not go through meta-stable dS resonances. One of the supposed virtues of the landscape picture of metastable dS was that, unlike a stable dS space, the landscape was part of a system with an infinite number of states, which could make infinitely precise quantum measurements on itself. If both Douglas’ claim, and that of the last paragraph are true, this is no longer obvious. All but a finite number of the final states in a given CDL instanton geometry, would not connect to a tunneling process from a meta-stable dS vacuum, but instead would evolve directly from a Big Bang. The whole issue of a rigorous framework for the landscape remains as murky as ever.

It seems to me that the obvious conclusion to draw from the conjunction of the difficulties of principle from which the landscape suffers, with its phenomenological problems, is that the landscape does not exist. This is not meant to be a firm scientific conclusion, but a warning and a guide to future research. In my opinion, much of the justification for the landscape relies on “ancient string theory folklore”, which has been discredited by the insights of the duality revolution, most particularly the AdS/CFT correspondence. Careful consideration of classical black hole physics also casts doubt on the world view of the landscape gardener.

At times the landscape is also hailed as the only way to understand de Sitter space-time in string theory. A (formerly? )famous footnote in the field theory textbook by Bjorken and Drell, warns the reader to be careful of making the “what else could it be?” argument, because, “if you’re not careful, someone else will tell you what else it could be”. In this case, perhaps, someone already has [16].
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