

Life at the interface of particle physics and string theory

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If the results of the first LHC run are not betraying us, many decades of particle physics are culminating in a complete and consistent theory for all nongravitational physics: the standard model. But despite this monumental achievement there is a clear sense of disappointment: many questions remain unanswered. Remarkably, most unanswered questions could just be environmental, and disturbingly to some the existence of life may depend on that environment. Meanwhile there has been increasing evidence that the seemingly ideal candidate for answering these questions, string theory, gives an answer few people initially expected: a large “landscape” of possibilities that can be realized in a multiverse and populated by eternal inflation. At the interface of “bottom-up” and “top-down” physics, a discussion of anthropic arguments becomes unavoidable. Developments in this area are reviewed, focusing especially on the last decade.

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I. INTRODUCTION

In popular accounts, our Universe is usually described as unimaginably large. Indeed, during the last centuries we have seen our horizon expand many orders of magnitude beyond any scale humans can relate to.

But the earliest light we can see has traveled a mere 13.8×10^9 years, just about 3 times the age of our planet. We might be able to look a bit further than that using intermediaries other than light, but soon we inevitably reach a horizon beyond which we cannot see.

We cannot rule out the possibility that beyond that horizon there is just more of the same, or even nothing at all, but widely accepted theories suggest something else. In the theory of inflation, our Universe emerged from a piece of a larger “space” that expanded by at least 60 e-folds. Furthermore, in most theories of inflation our Universe is not a one-off event. It is much more plausible that the mechanism that gave rise to our Universe was repeated a large, even infinite, number of times. Our Universe could just be an insignificant bubble in a gigantic cosmological ensemble, a “multiverse.” There are several classes of ideas that lead to such a picture, but there is no need to be specific here. The main point is that other universes than our own may exist, at least in a mathematical sense. The Universe we see is really just *our* Universe. Well, not just ours, presumably.

The existence of a multiverse may sound like speculation, but one may as well ask how we can possibly be certain that this is *not* true. Opponents and advocates of the multiverse idea are both limited by the same horizon. On whom rests the burden of proof? What is the most extraordinary statement: that what we can see is precisely all that is possible or that other possibilities might exist?

If we accept the logical possibility of a multiverse, the question arises in which respects other universes might be different. This obviously includes quantities that vary even within our own Universe, such as the distribution of matter and the fluctuations in the cosmic microwave background (CMB). But the cosmological parameters themselves, and not just their fluctuations, might vary as well. And there may be more that varies: the laws of physics could be different.

Since we observe only one set of laws of physics, it is a bit precarious to contemplate others. Could there exist alternatives to quantum mechanics or could gravity ever be repulsive rather than attractive? None of that makes sense in any way we know, and hence it seems unlikely that anything useful can be learned by speculating about this. If we want to consider variations in the laws of physics, we should focus on laws for which we have a solid underlying theoretical description.

The most solid theoretical framework we know is that of quantum field theory, the language in which the standard model of particle physics is written. Quantum field theory provides a large number of theoretical possibilities, distinguished by some discrete and some continuous choices. The discrete choices are a small set of allowed Lorentz group

representations, a choice of gauge symmetries (such as the strong and electroweak interactions), and a choice of gauge-invariant couplings of the remaining matter. The continuous choices are the low-energy parameters that are not yet fixed by the aforementioned symmetries. In our Universe we observe a certain choice among all of these options, called the standard model, discussed in Sec. II. But the quantum field theory we observe is just a single point in a discretely and continuously infinite space. Infinitely many other choices are mathematically equally consistent.

Therefore the space of all quantum field theories provides the solid underlying description we need if we wish to consider alternatives to the laws of physics in our own Universe. This does not mean that nothing else could vary, just that we cannot discuss other variations with the same degree of confidence. But we can certainly theorize in a meaningful way about universes where the gauge group or the fermion masses are different or where the matter does not even consist of quarks and leptons.

We have no experimental evidence of the existence of such universes, although there are speculations about possible observations in the cosmic microwave background (see Sec. III.E.2). We may get lucky, but our working hypothesis is the pessimistic one that all we can observe is our own Universe. But even then, the claim that the only quantum field theory we can observe in principle, the standard model of particle physics, is also the only one that can exist mathematically would be truly extraordinary.

Why should we even care about alternatives to our Universe? One could adopt the point of view that the only reality is what we can observe, and that talking about anything else amounts to leaving the realm of science. But even then there is an important consequence. If other sets of laws of physics are possible, even just mathematically, this implies that our laws of physics cannot be derived from first principles. They would be (at least partly) environmental, and deducing them would require some experimental or observational input. Certainly this is not what many leading physicists have been hoping for in the last decades. Consider, for example, Feynman’s question about the value of the fine-structure constant α : “Immediately you would like to know where this number for a coupling comes from: is it related to pi or perhaps to the base of natural logarithms?” Indeed, there exist several fairly successful attempts to express α in terms of pure numbers. But if α varies in the multiverse, such a computation would be impossible, and any successes would be mere numerology.

There is a more common “phenomenological” objection, stating that even if a multiverse exists, still the only universe of phenomenological interest is our own. The latter attitude denies the main theme of particle physics in the last three decades. Most activity has focused on the “why questions” and on the problem of “naturalness.” This concerns the discrete structure of the standard model, its gauge group, the couplings of quarks and leptons, and the questions of why they come in three families and why certain parameters have strangely small values. The least one can say is that if these features could be different in other universes, this might be part of the answer to those questions.

But there is a more important aspect to the latter discussion that is difficult to ignore in a multiverse. If other environments

are possible, one cannot avoid questions about the existence of life. It is not hard to imagine entire universes where nothing of interest can exist, for example, because the only stable elements are hydrogen and helium. In those universes there would be no observers. Clearly, the only universes in the multiverse that can be observed are those that allow the existence of observers. This introduces a bias: what we observe is not a typical sample out of the set of possible universes, unless all universes that (can) exist contain entities one might plausibly call “observers.” If the standard model features we are trying to understand vary over the multiverse, this is already crucial information. If there is furthermore a possibility that our own existence depends on the values of these parameters, it is downright irresponsible to ignore this when trying to understand them. Arguments of this kind are called “anthropic,” and tend to stir up strong emotions. These are the kind of emotions that always seem to arise when our own place in the cosmos and its history is at stake. One is reminded of the resistance against heliocentricity and evolution. But history is not a useful guide to the right answer; it only serves as a reminder that arguments should be based on facts, not on emotions. We discuss some general objections in Sec. III.

The fact that at present the existence of other universes and laws of physics cannot be demonstrated experimentally does not mean that we will never know. One can hope that one day we will find a complete theory of all interactions by logical deduction, starting from a principle of physics. For more than half a century, it has been completely acceptable to speculate about such theories provided the aim was a unique answer. But it is equally reasonable to pursue such a theory even if it leads to a large number of possible realizations of quantum field theories. This is not about giving up on the decade long quest for a unique theory of all interactions. It is simply pointing out a glaring fallacy in that quest. Nothing we know, and nothing we argue for here, excludes the possibility that the traditional path of particle physics toward shorter distances or higher energies will lead to a unique theory. The fallacy is to expect that there should be a unique way back: that starting with such a theory we might derive our Universe uniquely using pure mathematics.

A theoretical construction exists that may have a chance to fulfill the hope of finding the underlying theory: string theory. It is the third main ingredient of the story and is introduced in Sec. IV. It describes both gravitational and gauge interactions as well as matter. Initially it seemed to deliver the unique outcome many were hoping for as the strong constraints it has to satisfy appeared to allow only very few solutions.

But within two years, this changed drastically. The “very few solutions” grew exponentially to astronomically large numbers. One sometimes hears claims that string theorists were promising a unique outcome. But this is simply incorrect. In several papers from around 1986 one can find strong statements about large numbers of possibilities, starting with Narain (1986) and Strominger (1986), shortly thereafter followed by Kawai, Lewellen, and Tye (1986), Antoniadis, Bachas, and Kounnas (1987), and Lerche, Lust, and Schellekens (1987). Large numbers of solutions had already been found earlier in the context of Kaluza-Klein supergravity, reviewed by Duff, Nilsson, and Pope (1986), but the demise of uniqueness of string theory had a much bigger impact.

The attitudes toward these results differed. Some blamed the large number of solutions on our limited knowledge of string theory and speculated about a dynamical principle that would determine the true ground state; see, for example, Strominger (1986). Others accepted it as a fact and adopted the phenomenological point of view that the right vacuum would have to be selected by confrontation with experiment as stated by Kawai, Lewellen, and Tye (1987). In a contribution to the European Physical Society (EPS) conference in 1987 the hope for a unique answer was described as “unreasonable and unnecessary wishful thinking” (Schellekens, 1987).

It began to become clear to some people that string theory was not providing evidence *against* anthropic reasoning, but *in favor* of it. But the only person to state this explicitly at that time was Linde (1986b), who simply remarked that “the emergent plenitude of solutions should not be seen as a difficulty but as a virtue.” It took ten more years for a string theorist to put this point of view into writing (Schellekens, 1998), and 15 years before the message was advertised loud and clear by Susskind (2003).

In the intervening 15 years a lot had changed. An essential role in the story is played by *moduli*, continuous parameters of string theory. String theorists like to emphasize that “string theory has no free parameters,” and indeed this is true, since the moduli can be understood in terms of vacuum expectation values (VEVs) of scalar fields, and hence are not really parameters. All parameters of quantum field theory, the masses and couplings of particles, depend on these scalar VEVs. The number of moduli is 1 or 2 orders of magnitude larger than the number of standard model parameters. This makes those parameters “environmental” by definition and opens the possibility that they could vary over an ensemble of universes.

The scalar potential governing the moduli is flat in the supersymmetric limit. Supersymmetry (SUSY) is a symmetry between boson and fermions, which is, at best, an approximate symmetry in our Universe, but also a nearly indispensable tool in the formulation of string theory. If supersymmetry is broken, there is no reason why the potential should be flat. But this potential could very well have a disastrous runaway behavior toward large scalar VEVs or have computationally inaccessible local minima (Dine and Seiberg, 1985). Indeed this potential catastrophe was looming over string theory until the beginning of this century, when a new ingredient known as “fluxes” was discovered by Bousso and Polchinski (2000). This gave good reason to believe that the potential can indeed have controllable local minima, and that the number of minima (often referred to as “string vacua”) is large: an estimate of 10^{500} given by Douglas (2004a) is leading a life of its own in the literature. These minima are not expected to be absolutely stable; a lifetime of about 14×10^9 years is sufficient.

This ensemble has been given the suggestive name “the landscape of string theory.” Our Universe would correspond to one of the minima of the potential. The minima are sampled by means of tunneling processes from an eternally inflating de Sitter (dS) space (Linde, 1986a). If this process continues eternally, if all vacua are sampled, and if our Universe is one of them (three big if’s that require more discussion), then this provides a concrete setting in which anthropic reasoning is not only meaningful, but inevitable.

This marks a complete reversal of the initial expectations of string theory and is still far from being universally accepted or formally established. Perhaps it will just turn out to be a concept that forced us to rethink our expectations about the fundamental theory. But a more optimistic attitude is that we have in fact reached the initial phase of the discovery of that theory.

The landscape also provided a concrete realization of an old idea regarding the value of the cosmological constant Λ , which is smaller by more than 120 orders of magnitude than its naive size in Planckian units. If Λ varies over the multiverse, then its smallness is explained at least in part by the fact that for most of its values life would not exist. The latter statement is not debatable. What can be debated is if Λ does indeed vary, what the allowed values are and if anthropic arguments can be made sufficiently precise to determine its value. The anthropic argument, already noted by many, was sharpened by Weinberg (1987). It got little attention for more than a decade, because Λ was believed to be exactly zero and because a physical mechanism allowing the required variation of Λ was missing. In the string theory landscape the allowed values of Λ form a “discretuum” that is sufficiently dense to accommodate the observed small value.

This gave a boost to the landscape hypothesis in the beginning of this millennium and led to an explosion of papers in a remarkably broad range of scientific areas: string theory, particle physics, nuclear physics, astrophysics, cosmology, chemistry, biology, and geology, numerous areas in mathematics, even history and philosophy, not to mention theology. It is impossible to cover all of this in this review. It is not easy to draw a line, but on the rapidly inflating publication landscape we use a measure that has its peak at the interface of the standard model and string theory.

II. THE STANDARD MODEL

Despite its modest name, the standard model is one of the greatest successes in the history of science. It provides an amazingly accurate description of the three nongravitational interactions we know: the strong, the electromagnetic, and the weak interactions. Its success ranges from the almost 10-digit accuracy of the anomalous magnetic moment of the electron to the stunningly precise description of a large number of high-energy processes currently being measured at the LHC at CERN, and prior to that at the Tevatron at Fermilab, and many other accelerators around the world. Its success was crowned on July 4, 2012, with the announcement of the discovery of the Higgs boson at CERN, the last particle that was still missing. But this success has generated somewhat mixed reactions. In addition to the understandable euphoria, there are clear overtones of disappointment. Many particle physicists hoped to see the first signs of failure of the standard model. A few would even have preferred *not* finding the Higgs boson.

This desire for failure on the brink of success can be explained in part by the hope of simply discovering something new and exciting, something that requires new theories and justifies further experiments. But there is another reason. Most particle physicists are not satisfied with the standard model because it is based on a large number of seemingly *ad hoc* choices. Next we enumerate them.

We start with the “classic” standard model, the version without neutrino masses and right-handed neutrinos. In its most basic form it fits on a T shirt, a very popular item in the CERN gift shop these days. Its Lagrangian density is given by

$$\begin{aligned}\mathcal{L} = & -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} \\ & + i\bar{\psi}\not{D}\psi + \text{conjugate} \\ & + \bar{\psi}_i Y_{ij} \psi_j \phi + \text{conjugate} \\ & + |D_\mu \phi|^2 - V(\phi).\end{aligned}\quad (2.1)$$

In this form it looks barely simple enough to be called “elegant,” and furthermore many details are hidden by the notation.

A. The gauge sector

The first two lines are nearly completely fixed by symmetries and depend only on the discrete choices of gauge group and representations, plus the numerical value of the three real coupling constants of the gauge group $SU(3) \times SU(2) \times U(1)$. The left-handed fermions couple to this gauge group according to the following representations:

$$(3, 2, \tfrac{1}{6}) + (\bar{3}, 1, -\tfrac{2}{3}) + (\bar{3}, 1, \tfrac{1}{3}) + (1, 2, -\tfrac{1}{2}) + (1, 1, 1).$$

This repeats 3 times for no known reason. There is no theoretical reason why this particular combination of representations is the one we observe, although there is an important restriction on four cubic traces and one linear trace of the representation matrices from a condition called “anomaly cancellation.”

B. Yukawa couplings

The third line introduces a new field ϕ , a complex Lorentz scalar coupled to the gauge group as $(1, 2, \tfrac{1}{2})$, another choice dictated by observation and not by fundamental physics. This line consists of all terms allowed by the gauge symmetry, with an arbitrary complex coefficient Y_{ij} , the Yukawa coupling, for each term. The allowed couplings constitute three complex 3×3 matrices, for a total of 54 parameters (not all of which are observable, see below).

C. Scalar bosons

The last line specifies the kinetic terms of the scalar boson, with a minimal coupling to the gauge bosons. The last term is a potential, a function of ϕ . This potential has the form

$$V(\phi) = \tfrac{1}{2}\mu^2 \phi^* \phi + \tfrac{1}{4}\lambda(\phi^* \phi)^2. \quad (2.2)$$

This introduces two more real parameters. By means of the Higgs mechanism this sector of the theory gives masses to the W and Z bosons and all quarks and leptons, and to four weak mixing angles [the Cabibbo-Kobayashi-Maskawa (CKM) matrix].

D. The CKM matrix

The CKM matrix is obtained by diagonalizing two complex matrices, the up-quark mass matrix M_u and the down-quark

mass matrix M_d , which are the products of the corresponding Yukawa coupling matrices and the Higgs VEV v :

$$D_u = U_L^\dagger M_u U_R, \quad D_d = V_L^\dagger M_d V_R, \quad U_{\text{CKM}} = U_L^\dagger V_L, \quad (2.3)$$

where D_u and D_d are real, positive diagonal matrices. For three families, U_{CKM} can be parametrized by three angles and a phase. It turns out to be nearly diagonal, which presumably is an important clue. An often used approximate parametrization is

$$U_{\text{CKM}} \approx \begin{pmatrix} 1 - \lambda^2/2 & \lambda & A\lambda^3(\rho - i\eta) \\ -\lambda & 1 - \lambda^2/2 & A\lambda^2 \\ A\lambda^3(1 - \rho - i\eta) & -A\lambda^2 & 1 \end{pmatrix},$$

where $\lambda = 0.226$, and corrections of order λ^4 have been ignored. For values of the other parameters, see [Beringer et al. \(2012\)](#). They will not matter in the remainder of this review, because the current state of the art does not go beyond getting the leading terms up to factors of the order of 1, especially the hierarchy of the three mixing angles $\theta_{12} = \lambda$, $\theta_{23} \propto \lambda^2$, and $\theta_{13} \propto \lambda^3$. The degree of nonreality of the matrix can be expressed in terms of the Jarlskog invariant J , which is defined as

$$\text{Im}[V_{ij}V_{kl}V_{il}^*V_{kj}^*] = J \sum_{m,n} \epsilon_{ikm} \epsilon_{jln}. \quad (2.4)$$

This is a very small number: $J \approx 3 \times 10^{-5}$.

E. Quark and lepton masses

The values of the quark and lepton masses in GeV are listed below. See [Beringer et al. \(2012\)](#) for errors and definitions.

u, c, t	d, s, b	e, μ, τ
0.0023	0.0048	0.000511
1.275	0.095	0.105
173.5	4.5	1.777

The masses and hierarchies are not explained within the standard model; they are simply put in by means of the Yukawa coupling matrices.

F. The number of parameters

We now have a total of 18 observable parameters, which have now finally all been measured. From the measured values of the $W^\pm$ and Z masses and the electromagnetic coupling constant e we can compute $g_1 = (M_Z/M_W)e$, $g_2 = M_Z/(\sqrt{M_Z^2 - M_W^2})$, and the vacuum expectation value v of the scalar ϕ , using $M_W = \frac{1}{2}g_2 v$. This vacuum expectation value is related to the parameters in the potential as $v = 2\sqrt{-\mu^2/\lambda}$ and has a value of about 246 GeV. The Higgs mass determines μ^2 , and hence now we also know λ .

G. CP violating terms

There is, however, one more dimensionless parameter that does not appear on the T shirt. One can consistently add a term of the form

$$\theta \frac{g_3^2}{32\pi^2} \sum_{a=1}^8 F_{\mu\nu}^a F_{\rho\sigma}^a \epsilon^{\mu\nu\rho\sigma}, \quad (2.5)$$

where the sum is over the eight generators of SU(3). This term is not forbidden by any symmetries. The parameter $\theta \in [0, 2\pi)$ is shifted by the quark mass diagonalization. The physical combination $\bar{\theta} = \theta - \arg\det(M_u M_d)$ is observable in dipole moments of the neutron and nuclei. Nothing has been seen so far, which implies that $\bar{\theta} < 10^{-10}$. Note that one could also introduce a similar term for the SU(2) and U(1) gauge groups, with parameters θ_2 and θ_1 . However, θ parameters of Abelian theories are not observable, and θ_2 can be rotated to zero using baryon number phase rotations. Therefore we get only one extra parameter $\bar{\theta}$ bringing the total to 19.

H. Renormalizability

The 19 parameters were obtained by writing down all interactions allowed by the symmetry with a mass dimension less than or equal to 4. Without this restriction, infinitely many terms could be added to Eq. (2.1), such as four-fermion interactions or polynomials in $(\phi^* \phi)$. Any such term defines a new mass scale, and we can consistently “decouple” these terms by sending these mass scales to infinity.

In theories like the standard model, all unknown (and unknowable) virtual short-distance contributions are lumped together in a finite number of parameters. This is known as “renormalizability.” This property does not depend on parameter values and discrete choices and remains just as valid if we make the electron mass twice as large. As soon as evidence for a new term with dimension larger than 4 is found this will define a limiting mass scale M_{new} (where “new” stands for new physics). All computations would be off by unknown contributions of order Q/M_{new} , where Q is the mass scale of the process of interest. Since such new terms can be expected to exist on many grounds, including ultimately quantum gravity (with a scale $M_{\text{new}} = M_{\text{Planck}}$), the standard model is just an effective field theory valid up to some energy scale.

I. Running couplings

As a direct consequence of the renormalization procedure, the values of the constants in the Lagrangian depend on the energy scale at which they are measured. In the simplest case, the loop corrections to a gauge coupling constant have the form

$$g(Q) = g + \beta_0 g^3 \log(Q/\Lambda) + \text{higher order}, \quad (2.6)$$

where g is the coupling constant appearing in the Lagrangian, and Λ is a manually introduced ultraviolet cutoff of a momentum integral. We use $g(Q)$ as the physical coupling constant to be compared to experimental results at a scale Q . This then removes the dependence on Λ in all physical quantities to this order. But if we had used instead a different scale Q' , we would have measured a different value for the coupling constant $g(Q')$. The value of $g(Q')$ can be expressed in terms of $g(Q)$ using Eq. (2.6) and involves a term $\beta_0 \log(Q/Q')$. One can do better than this and sum up the

leading contributions (“leading logs”) of Feynman diagrams of any order in the loop expansion. This leads to the renormalization group equations, with a generic form

$$\frac{dg_i(t)}{dt} = \beta(g_i(t)), \quad (2.7)$$

where β is a polynomial in all parameters in the Lagrangian. Here $t = \log(Q/Q_0)$, where Q_0 is some reference scale.

J. Range of validity

Now that we finally know all standard model couplings including the Higgs self-coupling λ , we can see what happens to them if we assume that there is nothing but the standard model. It turns out that until we reach the Planck scale they all remain finite; all Landau poles (points where the coupling constants diverge) are beyond the Planck scale.

Note that not only the dimensionless parameters change logarithmically with Q , but also the parameter μ^2 in the Higgs potential, even though Eq. (2.6) looks different in this case: there are additional divergent contributions proportional to Λ^2 . This implies that μ^2 may get quantum contributions that are many orders of magnitude larger than its observed value. But this by itself does not invalidate the standard model, nor its extrapolation: the parameter μ^2 is a renormalized input parameter, just as all others.

K. The stability bound

The only potential problem in the extrapolation of the standard model couplings is that the Higgs self-coupling λ may become negative before the Planck scale, which may signal an instability. More precise determinations of the top quark mass and the QCD coupling are needed to be certain if λ does indeed go negative, and even if it does, it implies only a metastability of our vacuum with a lifetime that exceeds the current age of the Universe. Perhaps this is problematic for the evolution of the early Universe, but certainly not for its current state. Furthermore the problem can easily be avoided by adding a weakly coupled singlet scalar (Lebedev, 2012), and hence it does not offer a clear hint at elaborate new structures beyond the standard model.

L. Neutrino masses

The observation of neutrino oscillations implies that the classic standard model needs to be modified, because at least two neutrinos must have masses. Only squares of mass differences can be determined from these experiments. They are

$$\begin{aligned} \Delta m_{21}^2 &= (7.5 \pm 0.2) \times 10^{-5} \text{ eV}^2, \\ |\Delta m_{23}^2| &= (2.3 \pm 0.1) \times 10^{-3} \text{ eV}^2. \end{aligned}$$

In principle, neutrinos could be nearly degenerate in mass with minute differences, but from various cosmological observations we know that the sum of their masses must be less than about half an eV [see de Putter *et al.* (2012) for a recent update]. The masses can have a normal hierarchy $m_1 < m_2 < m_3$ or an inverted hierarchy $m_3 \ll m_1 < m_2$. They

are labeled 1, 2, and 3 according to their ν_e fraction in descending order.

The simplest way of accommodating neutrino masses is to add N fermions ψ_S that are standard model singlets.¹ The number N is not limited by anomaly constraints, and, in particular, does not have to be three. To explain the data one needs $N \geq 2$, but $N = 2$ looks contrived. Better motivated options are $N = 3$, for right-handed neutrinos as part of families, as in SO(10)-related grand unified theories (GUTs), or $N \gg 3$ in string models with an abundance of singlets.

As soon as singlets are introduced, not only Dirac, but also Majorana masses are allowed (and hence perhaps obligatory). The most general expression for couplings and masses is then (omitting spinor matrices)

$$\mathcal{L}_\nu = \sum_{i=1}^3 \sum_{a=1}^N \bar{\psi}_{\nu_L}^i Y_{ia} \psi_S^a + \sum_{ab}^N \mathcal{M}_{ab} \psi_S^a \psi_S^b. \quad (2.8)$$

The first term combines the three left-handed neutrino components with three (or two) linear combinations of singlets into a Dirac mass m , and the second term provides a Majorana mass matrix M for the singlets. This gives rise to a 6×6 neutrino mass matrix with 3×3 blocks, of the form

$$M_\nu = \begin{pmatrix} 0 & m \\ m & M \end{pmatrix}. \quad (2.9)$$

The mass scale of $\mathcal{M}$ is not related to any other standard model scale and is usually assumed to be large. In the approximation $m \ll M$ one gets three light neutrinos with masses of order m^2/M and N heavy ones. This is called the seesaw mechanism. It gives a very natural explanation for the smallness of neutrino masses (which are more than 8 orders of magnitude smaller than the muon mass) without unpalatable side effects. The optimal value of the Majorana mass scale is debatable and can range from 10^{11} to 10^{16} GeV depending on what one assumes about “typical” lepton Dirac masses.

If we assume $N \geq 3$ and discard the parameters of the heavy sector, which cannot be seen in low-energy neutrino physics, this adds nine parameters to the standard model: three light neutrino masses, four CKM-like mixing angles, and two additional phases that cannot be rotated away because of the Majorana nature of the fermions. This brings the total number of parameters to 28. However, as long as the only information about masses is from oscillations, the two extra phases and the absolute mass cannot be measured.

The current values for the mixing angles are

$$\begin{aligned} \sin^2(2\theta_{12}) &= 0.857 \pm 0.024, & \sin^2(2\theta_{23}) &> 0.95, \\ \sin^2(2\theta_{13}) &= 0.09 \pm 0.01. \end{aligned}$$

Note that the lepton mixing angles are not all small, unlike the CKM angles for quarks. The fact that $\theta_{13} \neq 0$ is known only since 2012 and implies that the CKM-like phase of the neutrino mixing matrix is measurable, in principle. This also rules out the once popular idea of tri-bi-maximal mixing

¹One may give Majorana masses to the left-handed neutrinos without introducing extra degrees of freedom, but this requires adding nonrenormalizable operators or additional Higgses.

(Harrison, Perkins, and Scott, 2002), removing a possible hint at an underlying symmetry.

III. ANTHROPIC LANDSCAPES

The idea that our own existence might bias our observations has never been popular in modern science, but especially during the last 40 years a number of intriguing facts have led scientists from several areas of particle physics, astrophysics, and cosmology in that direction, often with palpable reluctance. Examples are Dirac's large number hypothesis in astrophysics (Carter, 1974; Carr and Rees, 1979), chaotic inflation (Linde, 1986b), quantum cosmology (Vilenkin, 1986), the cosmological constant (Davies and Unwin, 1981; Barrow and Tipler, 1986; Weinberg, 1987), the weak scale in the standard model (Agrawal *et al.*, 1998b), quark and lepton masses in the standard model (Hogan, 2000), the standard model in string theory (Schellekens, 1998), and the cosmological constant in string theory (Bousso and Polchinski, 2000; Susskind, 2003).

This sort of reasoning goes by the generic name "anthropic principle" (AP) (Carter, 1974). In the remainder of this review, the term AP is used in the following sense. We assume a multiverse, with some physical mechanism for producing new universes. In this process, a (presumably large) number of options for the laws of physics is sampled. The possibilities for these laws are described by some fundamental theory; they are solutions to some equations. Furthermore we assume that we are able to conclude that some other sets of mathematically allowed laws of physics do not allow the existence of observers, by any reasonable definition of the latter [and one can indeed argue about that; see, for example, Gleiser (2010)].

This would be a rather abstract discussion if we had no clue what such a fundamental theory might look like. But fortunately there exists a rather concrete idea that, at the very least, can be used as a guiding principle: the string theory landscape described in the Introduction. The remainder of this section does not depend on the details of the string landscape, except that at one point we assume discreteness. However, the existence of some kind of landscape in some fundamental theory is a prerequisite. Without that, all anthropic arguments lose their scientific credibility.

A. What can be varied?

In the anthropic literature many variations of our laws of physics are considered. Often it is realized years later that a variation is invalid, because the parameter value is fixed for some previously unknown fundamental reason. One also encounters statements like we vary parameter X , but we assume parameter Y is kept fixed. But perhaps this is not allowed in a fundamental theory. So what can we vary, and what should be kept fixed?

In one case we can give a clear answer to these questions: we can vary the standard model within the domain of quantum field theory, provided we keep a range of validity up to an energy scale well above the scale of nuclear physics. Furthermore, we can vary anything, and keep anything we want fixed. For any such variation we have a quantum field

theory that is equally good, theoretically, as the standard model. For any such variation we can try to investigate the conditions for life. We cannot be equally confident about variations in the parameters of cosmology (see Sec. III.E.2).

Even though it is just an effective field theory, it goes too far to say that the standard model is just the next nuclear physics. In nuclear physics the limiting, new physics scale M_{new} is within an order of magnitude of the scale of nuclear physics. Computations in nuclear physics depend on many parameters, such as coupling constants, form factors, and nucleon-nucleon potentials. These parameters are determined by fitting to data, as are the standard model parameters. But unlike the standard model parameters, they cannot be varied outside their observed values in any way that makes sense. There is no theory of nuclear physics with twice the observed pion-nucleon coupling and everything else unchanged.

This difference is important in many cases of anthropic reasoning. Some anthropic arguments start with unjustified variations of parameters of nuclear physics. If life ceases to exist when we mutilate the laws of physics, nothing scientific can be concluded. The only admissible variations in nuclear physics are those that can be derived from variations in the relevant standard model parameters: the QCD scale Λ_{QCD} and the quark masses.

This raises an obvious question. If the standard model is just an effective field theory, made obsolete one day by some more fundamental theory, then why can we consider variations in its parameters? What if the fundamental theory fixes or constrains its parameters, just as QCD does with nuclear physics? The answer is that the relevant scale Q for anthropic arguments is that of chemistry or nuclear physics. This is far below the limiting scale M_{new} , which is more than a TeV or so. New physics at that scale is irrelevant for chemistry or nuclear physics.

If we ever find a fundamental theory that fixes the quark and lepton masses, the anthropic argument will still be valid, but starts playing a totally different role in the discussion. It changes from an argument for expectations about fundamental physics to a profound and disturbing puzzle. In the words of Ellis (2006a) "in this case the anthropic issue returns with a vengeance: (...). Uniqueness of fundamental physics resolves the parameter freedom only at the expense of creating an even deeper mystery, with no way of resolution apparent."

B. The anthropocentric trap

There is another serious fallacy one has to avoid: incorrectly assuming that something is essential for life, whereas it is only essential for our life. Any intelligent civilization (either within our own Universe or in an entirely different one with different laws of physics) might be puzzled about properties in their environment that seem essential for their existence. But that does not imply that life cannot exist under different circumstances.

Arguments based on water or DNA should be viewed with suspicion. Perhaps we do not even need fusion-fueled stars (Adams, 2008); degenerate stars (white dwarfs or neutron stars) may provide sufficient energy.

Arguments based on abundances are equally suspect. Fred Hoyle famously predicted the existence of a resonance in the

carbon nucleus that would enhance carbon production, and indeed this resonance was found. This is often referred to as a successful anthropic prediction, because carbon is essential for our kind of life. But it is in fact just a prediction based on the observed abundance of some element. Indeed, Hoyle himself did not make the link between the abundance of carbon and life until much later (Kragh, 2010).

The current status of the Hoyle state and its implications are summarized in Sec. V.B.1.f. Based on what we know we cannot claim that life is impossible without this resonance. We do not know which element abundances are required for life, nor do we know how they vary over the standard model parameter space. Perhaps there even exists a parameter region where ^8Be is stable, and the beryllium bottleneck is absent (Higa, Hammer, and van Kolck, 2008). This would turn the entire anthropic argument on its head.

If we discover that we live near an optimum in parameter space, this would be a strong indication of multiverse scanning (a unique theory is not likely to land there), but as long as the maximum is broad or other regions exist there is no need to overdramatize. Most observers observe conditions that are most favorable to their existence.

In view of the difficulties in defining anthropic constraints some have proposed other criteria that are under better control and still are a good “proxy” for life. In particular, it seems plausible that the formation of complex structures will always be accompanied by entropy production in its environment, a criterion that would certainly work in our own Universe. This “entropic principle” has led to some success for cosmological parameters (Bousso and Harnik, 2010), but seems less useful for the subtle details of the standard model parameter space.

1. Other habitable universes

Going to extremes, one can imagine habitable universes with only electromagnetic and gravitational interactions, with fundamental nuclei and electrons created by some kind of generalized baryogenesis and with only dim stars stabilized by degeneracy pressure of fermions, radiating gravitational energy built up during their collapse. These universes would still have solid matter, chemistry, and biology like ours.

A less extreme possibility is a universe without weak interactions. Harnik, Kribs, and Perez (2006) made some clever changes in the theory to mimic physics in our Universe as closely as possible, so that one can rely on our experience with conventional physics. Quarks and leptons have small masses (in Planck units), not because of a light Higgs boson, but by having extremely small Yukawa couplings. Type-II supernovae are not available, but type-Ia supernovae, whose explosions are driven by the strong interactions, can take over their role in spreading heavy elements. However, there are some serious worries: there is no known mechanism for baryogenesis,² stars are less bright, there may be no plate tectonics and volcanism (which are fueled to a large extent by weak decays), type-I supernovae may not produce enough oxygen (Clavelli and White, 2006), and there

is a potentially harmful (Cahn, 1996; Hogan, 2006) stable neutron background.

Instead of changing the quantum field theory parameters underlying our own Universe, one can also try to change cosmological parameters, such as the baryon-to-photon ratio, the primordial density perturbations, the cosmological constant, and the curvature density parameter Ω . This was done by Aguirre (2001), and also in this case regions in parameter space could be identified where certain parameters differ by many orders of magnitude, and yet some basic requirements of life are unaffected.

Alternative universes that must probably be ruled out anthropically are the exact supersymmetric ones, because supersymmetric theories are the hardest to dismiss on fundamental grounds. Fortunately, ruling them out is easy. In supersymmetric theories electrons are degenerate with scalars called selectrons. These scalars are not constrained by the Pauli principle and would all fill up the s wave of any atom (Cahn, 1996). Chemistry and stability of matter (Dyson, 1967; Lieb, 1990) would be lost. Although this may look sufficiently devastating, it has not stopped speculation about the possibility of life under these conditions; see, e.g., Clavelli (2006) and Banks (2012).

C. Is life generic in quantum field theory?

It may seem that we are heading toward the conclusion that any quantum field theory (QFT) allows the existence of life and intelligence. Perhaps any complex system will eventually develop self-awareness (Banks, 2012). Even if that is true, it still requires sufficient complexity in the underlying physics. But that is still not enough to argue that all imaginable universes are on equal footing. We can easily imagine a universe with just electromagnetic interactions and only particles of charge 0, ± 1 , and ± 2 . Even if the clouds of hydrogen and helium in such a universe somehow develop self-awareness and even intelligence, they will have little to be puzzled about in their QFT environment. Their universe remains unchanged over vast ranges of its parameters. There are no anthropic tunings to be amazed about. Perhaps, as argued by Bradford (2011), fine-tuning is an inevitable consequence of complexity and hence any complexity-based life will observe a fine-tuned environment. But this just strengthens the argument that we live in a special place in the space of all quantum field theories, unless one drops the link between complexity and life. But if life can exist without complexity, that just begs the question why the problem was solved in such a complicated way in our Universe.

If we put everything we know and everything we do not know together, the picture that emerges is one of many domains where life might exist, and many more where it definitely does not. Presumably the habitable regions are narrow in certain directions and very elongated in others. A cartoon version of such regions in part of QFT space is shown in Fig. 1, with the gray circle showing our own location and the experimental uncertainties.

This diagram represents two unrelated gedanken computations (Schellekens, 2008). The contours are the result of the anthropic gedanken computation explained previously. The

²In our own Universe we are not certain about the mechanism either, but at least we are sure that one exists.

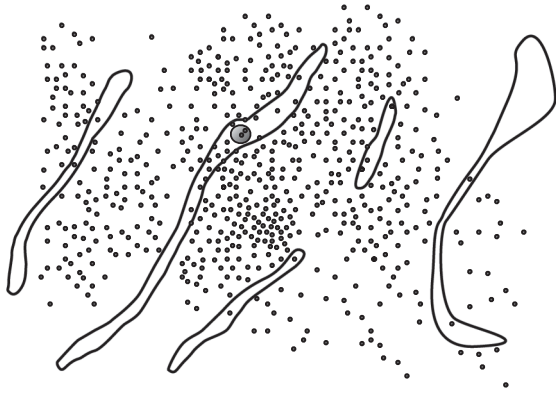


FIG. 1. Habitable regions in QFT space. The gray circle represents the experimental bounds on the standard model. The dots show the distribution of QFT points in a hypothetical landscape.

dots show the results of a very different one. They represent points in QFT space obtained from some fundamental theory, such as string theory. Here the implicit assumption is made that such a theory will lead to a discrete set of points. In this concrete setting, it is clear that the two gedanken computations are completely unrelated. The first one involves low-energy physics: nuclear and atomic physics and chemistry. The second one involves geometry and topology of manifolds with membranes and fluxes wrapped around them, and determining minima of potentials generated by all this structure. We can actually do both kinds of computations only in simple cases, but we know enough to conclude that it would take a miracle for them to match each other, if the second computation were to produce a unique answer. The obvious way out is precisely what string theory suggests: that there is not a single point, but a cloud of points, covering a substantial part of the QFT parameter space. Note that no such cloud is required for a point to land precisely in the gray, experimental circle, because unlike the anthropic contours this circle cannot be determined by a computation.

These contours are sharp lines in the case of particle physics thresholds, such as reactions that stop being exothermic or stability of essential building blocks (although there is usually a small transition region where a particle is just stable enough). In other cases they are more like contour lines of distributions. Most papers make different assumptions about the definitions of these lines (i.e., the necessary conditions for life) and consider different slices through the parameter space.

Moving out of our own location, the first line we encounter is the end of our region. There our kind of life ends, and we have to rely on speculation to know if other kinds of life are possible. This happens, for example, if one of the crucial processes in the functioning of stars is shut off. Other processes may take over, but stellar lifetimes and/or heavy element abundances may differ by orders of magnitude, and we cannot rely on experimental data to be certain that such a universe will “work.” Beyond this *terra incognita* (perhaps more appropriately called “no man’s land”) there is usually another boundary where the conditions become so adverse that any kind of complexity can be ruled out. For discussion along similar lines, see Hall and Nomura (2008). In the remainder of this review we shall not make this distinction

over and over again and use the adjective anthropic rather loosely for any parameter change that is likely to affect life, whether it is our life or life in general.

Real plots of this kind can be found in many papers, e.g., Agrawal *et al.* (1998b), Tegmark (1998), Hogan (2000), Hellerman and Walcher (2005), Tegmark *et al.* (2006), Barr and Khan (2007), Graesser and Salem (2007), Hall and Nomura (2008), Jaffe, Jenkins, and Kimchi (2009), Elor *et al.* (2010), and Barnes (2012).

Even without drawing further conclusions, it is exciting to see where we are located on the parameter space map, and to see the lines of minor and major catastrophes surrounding us. It is a bit like seeing our fragile planet in the vastness of space on the first Apollo 8 pictures. It is also a great way of appreciating how our Universe really works. If we do indeed understand that, we should be able to change something and work out the consequences.

Figure 1 was deliberately drawn in this way to illustrate a few fallacies that are perhaps blatantly obvious, but that are nevertheless repeated incessantly in the literature.

- Anthropic reasoning will never completely determine the standard model. It is quite clear that even in our own environment there are variations that have no conceivable impact on life, such as the τ mass.
- Anthropic reasoning combined with a fundamental theory is not likely to determine the standard model either. This would require the density of the cloud to match the size of the anthropic region, in such a way that precisely one point lands inside it. That would be another miracle.
- There is no reason to expect the maximum of the density distribution, even when folded with sampling probabilities, to select our vacuum. Computing these maxima is another gedanken computation that cannot be sensitive to the location of the domains,³ the other gedanken computation.
- Bounds on parameters may disappear as others are allowed to vary. Obviously the projections of the regions on the axes cover essentially everything, but if we intersect them with horizontal or vertical lines, we get narrow bounds.

If one can show that a parameter is anthropically constrained, keeping all others fixed, that is tremendous success. If one can do it while allowing others to vary, that is an even bigger success. Only in cases where strong claims are made about the actual value of a parameter (especially that it must

³Unless life in a universe somehow affects the sampling probability of its offspring. This includes science fiction ideas like scientists making copies of their own universe in experiments. A related idea was proposed by Smolin (1994), who argued that collapsing black holes creates new universe with slightly changed parameters. This would make the maximum of black hole production a point of attraction in a multiverse. However, black holes are hardly the optimal environment for life, nor a suitable device for transferring information. For further discussion, see Rothman and Ellis (1993), Barrow (2001), Smolin (2006), and Vilenkin (2006a). Note that the *existence* of a landscape is in any case a prerequisite for such a proposal.

be small), it becomes really important to ask if the smallness is a consequence of fixing other parameters.

D. Levels of anthropic reasoning

Even in the interpretation used in this review, one may distinguish several versions of the AP:

- (1) AP0: A mere tautology.
- (2) AP1: An explanation for certain fine-tunings.
- (3) AP2: A predictive method.

AP0: If the fundamental theory allows many universes that do not allow observers, we should not be puzzled to find ourselves in one that does. This is true, but not very useful.

AP1: Suppose we conclude that some variable x , *a priori* defined on an interval $[0, 1]$ has to lie in an extremely narrow band of size ϵ for observers to exist. If the fundamental theory contains N values of x evenly scattered over the interval, the chance that none of them is in the observer range is $(1 - \epsilon)^N$. For $N = M/\epsilon$ and small ϵ this goes like e^{-M} . For sufficiently large M , everyone would agree that there is nothing surprising about the existence of a point in the observer band. For concreteness, one may think of numbers like 10^{-120} for ϵ and 10^{500} for N , so that $M = 10^{380}$. The chance that a flat distribution contains no points in the observer range would then be the absurdly small number $\exp(-10^{380})$. Obviously, the fine-tuning is then explained. Note that we are talking about landscape density distributions here, not about sampling probabilities in eternal inflation (see Sec. VI for various approaches toward defining the latter).

AP2: It may be possible to go one step further and determine the most probable point where we should expect to find ourselves within the anthropic window. This requires additional information compared to AP1. We should be able to assign a probability to each point, work out the probability distribution, and determine its maximum. This brings some very serious measure problems into the discussion. What counts as an observer, and what counts as an observation? Should we sum over the entire history of the Universe, and how do we include parts of the Universe that are currently behind the horizon? How do we even define probabilities in the context of eternal inflation, where anything that can happen happens an infinite number of times? Furthermore there is the issue of “typicality” (Vilenkin, 1995a). If we can define and compute a probability distribution, should we expect to find ourselves at its maximum? Are we typical? Does statistics even make sense if we can observe just a single event?

Many criticisms of anthropic reasoning are aimed at the measure and typicality problems in AP2, and especially its use for predicting the cosmological constant. See, for example, Muller (2001), Smolin (2004), Neal (2006), Starkman and Trota (2006), Bostrom (2007), Maor, Krauss, and Starkman (2008), and Armstrong (2011) for a variety of thoughts on this issue. We return to the measure problem in Sec. VI.

Perhaps AP1 is as far as we can ever get. We may determine the boundaries of our domain and find out how a fundamental theory spreads its vacua over that domain. There is a lot of interesting physics and mathematics associated with all of these questions. In the end we may just be satisfied that we roughly understand where we are, just as we

are not especially obsessed with deriving the orbit and size of our planet in the landscape of astrophysical objects. Establishing the fundamental theory will have to be done by other means, perhaps purely theoretically, and by ruling out alternatives.

E. First signs of a landscape?

The current situation in particle physics invites an appeal to Occam’s razor. We cannot avoid asking the obvious question: Could it be that the standard model, including a minor extension to accommodate neutrino oscillations, is really all there is? Indeed, suggestions in that direction were made some time ago by Shaposhnikov and Tkachev (2006), albeit not in the context of a landscape.

It is undeniable that this state of affairs has contributed to the interest in anthropic and landscape thinking in particle physics. Could it be true that the standard model is like a dart that was thrown repeatedly at the space of all quantum field theories, until one of them landed in one of the anthropic domains of Fig. 1? This is the central question of this review.

But even in the most extreme landscape scenario, there are plenty of problems left that require a solution. It is just that the nature of the remaining problems has shifted in a remarkable way in a certain direction: most problems are now environmental, and many have anthropic implications.

One can roughly order the open problems according to their urgency, in the following way:

- No consistent theory.
- Disagreement between theory and experiment.
- Environmental, but not anthropic problems.
- Potentially anthropic problems.

In the following we make an admittedly rather artificial separation between particle physics and cosmology.

1. Particle physics

The main item in the first category is quantum gravity. The standard model does not contain gravity and adding it using standard QFT methods leads to inconsistencies.

In the second category there is a long list of deviations of low statistical significance that may one day develop into real problems, astrophysical phenomena for which there is no good theoretical model, but which may point to new particle physics, a hint of a gamma-ray line in cosmic rays at 130 GeV (Weniger, 2012) and a 4σ indication for spatial variations of the fine-structure constant (Webb *et al.*, 2011).

In the third category are all standard model parameters that have peculiar values, without any reason to hope that anthropic arguments are going to be of any help. The most important one is the CP -violating angle $\bar{\theta}$ of the strong interactions, arguably the most important standard model problem in the context of a landscape (Banks, Dine, and Gorbatov, 2004; Donoghue, 2004). Other examples of non-anthropocentric parameters with small values are the CKM angles and some of the quark mass ratios.

The last category consists of all problems related to parameters whose values do potentially have an impact on the existence of life. This includes the group structure and

representations of the standard model, the scales of the strong and the weak interactions (the “gauge hierarchy problem,” see Sec. V.C.2), the light quark masses and the electron mass (assuming the heavier fermions stay heavy), neutrino masses, and perhaps even the mass of the top quark. The environmental impact of the fermion masses is discussed in Sec. V.B.

2. Cosmology

The main cosmological parameters are the cosmological constant Λ , the density parameter Ω , the matter density fluctuations $Q = \delta\rho/\rho$, the dark-to-baryonic matter ratio ζ , the baryon-to-photon ratio η , and the parameters of inflation [see Tegmark *et al.* (2006) for a systematic survey of all parameters]. The theoretical foundations of cosmology belong to the first category defined above. There is no effective theory of cosmology where all of these parameters can manifestly be varied independently and without worrying about the impact of changes in our understanding of gravity. For example, the cosmological constant has only an observable meaning in a theory of gravity. The notion of decoupling it from gravity, as one can do for standard model parameters, does not even make sense.

Anthropic issues in cosmology will not be discussed in detail in this review, except for the cosmological constant, the focal point of much attention. Here we just briefly mention some interesting observations.

The main item in the second category is “dark matter,” or more precisely the complete set of problems that is elegantly solved if we postulate the existence of dark matter: galaxy rotation curves, the bullet cluster, structure formation, the features of the cosmic microwave background, the amount of deuterium produced in big bang nucleosynthesis, and the matter density of the Universe. There is a minority point of view that holds that these problems belong in the first category and require a modification of gravity. But should we really be so surprised if dark matter exists? Is it not a typical example of anthropocentric hubris to assume that anything that exists in the Universe must be observable by us, or made out of the same stuff that we are made of? Postulating dark matter moves this problem largely to category four, although there are still serious problems in computer simulations of galaxy formation which may point to a more fundamental problem [see Famaey and McGaugh (2013) for a list of open problems].

The dark-to-baryonic matter ratio ζ , which is ≈ 5 in our Universe, may have anthropic implications, since dark matter plays an important role in structure formation. This was first discussed for axion dark matter (Linde, 1988), because the most popular solution to the strong CP problem, the Peccei-Quinn (PQ) mechanism, predicts an additional particle, the axion, that contributes to dark matter. In contrast to the more popular WIMP dark matter,⁴ whose abundance is predicted by its interactions, axionic dark matter must satisfy constraints which are in part anthropic in nature (for more on axions see Sec. V.D). The constraints were made more precise by Hellerman and Walcher (2005), who found $\zeta < 10^5$ and

Tegmark *et al.* (2006) who concluded that $2.5 < \zeta < 10^2$, using some additional anthropic requirements. These papers also discuss the effect of other parameter variations (in particular, Q and Λ) on these bounds. Using assumptions about a multiverse measure and the number of observers per baryon, Freivogel (2010) gave an anthropic statistical prediction for ζ roughly in agreement with the observed value. Although the emphasis on all these papers is on axionic dark matter, some of the conclusions on ζ do not really depend on that.

Most other cosmological parameters are also in the fourth category. Changing any of these substantially has an impact on some feature in the history and/or current status of the Universe that would appear to be catastrophic at least for our kind of life, and hence it is at least possible that this is part of the reason we observe the values we do.

But we should not jump to conclusions. An extreme example is the smoothness and isotropy of the cosmic microwave background. This fact may be regarded as environmental, and if it were a wildly fluctuating distribution this could have a very negative impact on the prospects for life (Tegmark and Rees, 1998). But surely one cannot assume that the entire density perturbation function is tuned this way just for life to exist in one galaxy. The most popular solution to this “horizon problem” is inflation, which solves another problem with anthropic relevance, the flatness problem, but also introduces some new fine-tunings.

Inflationary cosmology offers interesting opportunities for predictions based on landscape and/or anthropic ideas, especially for observations of the CMB; see, e.g., Tegmark (2005), Holman, Mersini-Houghton, and Takahashi (2008), Ashoorioon (2010), Frazer and Liddle (2011), and Yamauchi *et al.* (2011). Furthermore, the CMB may even give direct hints at the existence of a multiverse. There is a chance of observing collisions with other bubbles in the multiverse; see, for example, Aguirre, Johnson, and Shomer (2007) and the Wilkinson Microwave Anisotropy Probe (WMAP) results presented by Feeney *et al.* (2011). Gonzalez-Díaz and Alonso-Serrano (2011) considered an even more exotic possibility involving nonorientable tunneling. In principle there might be information about other universes in the detailed structure of the cosmic microwave background, but at best only in the extreme future (Ellis, 2006b).

Anthropic predictions for the density parameter Ω were already made a long time ago by Garriga, Tanaka, and Vilenkin (1999). This work, as well as Freivogel *et al.* (2006), points out the plausibility of observing negative spatial curvature (i.e., $\Omega_k > 0$, where $\Omega_k \equiv 1 - \Omega$) in a multiverse picture. They argue that 60 e-folds of inflation are anthropically needed, and having a larger number of e-folds is statistically challenged. The current observational constraint is $|\Omega_k| < 10^{-2}$. Furthermore, Guth and Nomura (2012) and Kleban and Schillo (2012) point out that observation of even a small positive curvature ($\Omega_k < -10^{-4}$) would falsify most ideas of eternal inflation, because tunneling in a landscape gives rise to open Friedmann-Robertson-Walker universes.

That the baryon-to-photon ratio $\eta \approx 6 \times 10^{-10}$ may have anthropic implications was already observed a long time ago [see Carr and Rees (1979), Nanopoulos (1980), Linde (1985),

⁴WIMPs are “weakly interacting massive particles,” which are present, for example, in certain supersymmetric extensions of the standard model.

but also [Aguirre \(2001\)](#) for critical comments], but it is not simply a tunable free parameter. Inflation would dilute any such initial condition, as would any baryon number violating process that gets into equilibrium in the early stages of the Universe. See [Shaposhnikov \(2009\)](#) for a list of 44 proposed solutions to the baryogenesis problem. Most of these solutions generate new anthropic issues themselves.

This brief summary does not do justice to the vast body of work on string and landscape cosmology. Further references can be found in reviews of string cosmology; see, e.g., [Burgess and McAllister \(2011\)](#).

3. The cosmological constant

The cosmological constant Λ is a parameter of classical general relativity that is allowed by general coordinate invariance. It has dimension $[\text{length}]^{-2}$ and appears in the Einstein equations as [the metric signs are $(-, +, +, +)$]

$$R_{\mu\nu} - \frac{1}{2}g_{\mu\nu}R + \Lambda g_{\mu\nu} = 8\pi G_N T_{\mu\nu}. \quad (3.1)$$

Without a good argument for its absence one should therefore consider it as a free parameter that must be fitted to the data. It contributes to the equations of motion with an equation of state $P = w\rho$, where P is pressure and ρ is density, with $w = -1$ (matter has $w = 0$ and radiation $w = \frac{1}{3}$). As the Universe expands, densities are diluted as (the initial values are hatted)

$$\rho_w = \hat{\rho}_w \left(\frac{a}{\hat{a}}\right)^{-3(1+w)}. \quad (3.2)$$

As a result, if $\Lambda \neq 0$ it will eventually dominate if the Universe lasts long enough. The natural length scale associated with Λ is the size of the Universe.

The parameter Λ contributes to the equations of motion in the same way as vacuum energy density ρ_{vac} , which has an energy momentum tensor $T_{\mu\nu} = -\rho_{\text{vac}}g_{\mu\nu}$. Vacuum energy is a constant contribution to any (quantum) field theory Lagrangian. It receives contributions from classical effects, for example, different minima of a scalar potential and quantum corrections (e.g., zero-point energies of oscillators). However, it plays no role in field theory as long as gravity is ignored. It can simply be set to zero. Since vacuum energy and the parameter Λ are indistinguishable, it is customary to identify ρ_{vac} and Λ . The precise relation is

$$\frac{\Lambda}{8\pi} = \frac{G_N \rho_{\text{vac}}}{c^2} := \rho_\Lambda. \quad (3.3)$$

This immediately relates the value of Λ with all other length scales of physics, entering in ρ_Λ , which of course are very much smaller than the size of the Universe. The extreme version of this comparison is to express ρ_Λ in Planck mass per (Planck length)³, which gives a value smaller than 10^{-120} . This was clear long before ρ_Λ was actually measured.

More recently, observations of redshifts of distant type-Ia supernovae gave evidence for accelerated expansion ([Riess et al., 1998](#); [Perlmutter et al., 1999](#)), which can be fitted with the Λ parameter. Combined with more recent data on the cosmic microwave background, this indicates that the contribution of Λ to the density of the Universe is about 70% of the critical density $\rho_c \approx 9.9 \times 10^{-27} \text{ kg/m}^3$, assuming the

standard Λ CDM model of cosmology. This then leads to an “observed” value

$$\rho_\Lambda \approx +1.3 \times 10^{-123}. \quad (3.4)$$

a. Anthropic arguments

The foregoing discussion already implies that there will be an anthropic range for Λ , assuming everything else is kept fixed. Although this may have been clear to some much earlier, it appears that the first paper stating this is [Davies and Unwin \(1981\)](#). They did not make it quantitative, though. In subsequent years [Linde \(1984\)](#), [Sakharov \(1984\)](#), and [Banks \(1985\)](#) also discussed anthropic implications of $\Lambda \neq 0$. Sakharov’s paper contains the remarkable statement: “If the small value of the cosmological constant is determined by ‘anthropic selection,’ then it is due to the discrete parameters. This obviously requires a large value of the number of dimensions of the compactified space or (and) the presence in some topological factors of a complicated topological structure.”

Crude bounds on ρ_Λ in any habitable universe can already be obtained by requiring that complex objects with a large number of constituents (for example, brains) can form and fit inside the horizon in dS [see the last section of [Harnik, Kribs, and Perez \(2006\)](#)], or that nongravitational interaction time scales are much smaller than the collapse time in anti-de Sitter (AdS). This implies that if ρ_Λ can vary on Planckian scales, its observed value is in any case at least partly anthropic.

Much tighter bounds can be obtained if we fix the other parameters at their observed value. [Barrow and Tipler \(1986\)](#) pointed out that if Λ is too large and negative, the Universe would collapse before life has evolved. They used the average lifetime of a main-sequence star to get a limit. This quantity can be entirely expressed in terms of standard model parameters and the Planck mass and leads to a limit

$$|\rho_\Lambda| \lesssim \alpha^{-4} \left(\frac{m_e}{m_p}\right)^4 \left(\frac{m_p}{M_{\text{Planck}}}\right)^6 = 6.4 \times 10^{-120}. \quad (3.5)$$

Rather than theoretical lifetimes of stars, one can consider observational extremes: the minimal stellar lifetime of about 3×10^6 years and the current age of the Universe. The fastest time in which intelligent life can form must lie between these extremes. Requiring that this is less than the time of collapse $\pi\sqrt{3/\Lambda}$ gives $\rho_\Lambda > -\rho_{\text{min}}$, with

$$1.8 \times 10^{-122} < \rho_{\text{min}} < 3.8 \times 10^{-115}. \quad (3.6)$$

The limit (3.5) was argued to be valid for positive Λ as well. However, [Weinberg \(1987\)](#) pointed out that structure that has already formed will not be ripped apart by an expanding universe. Once galaxies have formed, it makes no difference how much time is needed to make stars or evolve life, because the expansion will not inhibit that from happening. He then derived a limit based on the assumption that life would not form if the Universe expands too fast to inhibit galaxy formation. The exact form of Weinberg’s bound is

$$\rho_\Lambda < \frac{500}{729} \Delta^3 \rho_0, \quad (3.7)$$

and was derived by studying the collapse of a spherical overdensity Δ using a Robertson-Walker metric. The overdensity starts expanding at $t = 0$ when the Universe has a matter

density ρ_0 . For $\rho_\Lambda = 0$ it recollapses and forms structure, but as ρ_Λ is increased a point is reached beyond which the recollapse does not occur anymore. This gives the maximum value of ρ_Λ for the overdensity Δ . The absolute upper limit in a given universe is given by determining the maximal overdensity that can occur. Since density fluctuations are distributions, there will not be a strict upper limit, but the number of galaxies that can be formed will drop off rapidly beyond a certain ρ_Λ .

In 1987 precision cosmology did not exist yet, and no theoretical estimate of the upper limit was possible. Hence an empirical estimate was made. If protogalaxies can be observed at high redshift z , when the matter density was larger by a factor $(1+z)^3$, a cosmological constant density of the same size would not obstruct galaxy formation either. In 1987 this led to an upper limit $\rho_\Lambda < 550\rho_{\text{matter}}$ from quasars at $z = 4.4$. However, meanwhile dwarf galaxies have been observed at $z = 10$, increasing the bound by almost an order of magnitude (Loeb, 2006).

b. Estimates of the value of Λ

Nowadays we can determine the density fluctuations using COBE and WMAP (and recently PLANCK) results. It is instructive to make a rough estimate using the time of matter-radiation equality as the starting point of structure formation. An order of magnitude estimate for the matter density at equality is (Hellerman and Walcher, 2005) $\rho_{\text{eq}} \approx T_{\text{eq}}^4$, $T_{\text{eq}} \approx m_p \eta(\zeta + 1)$, where $\eta = 6.3 \times 10^{-10}$ is the baryon-to-photon ratio and ζ is the cold-dark-matter-to-baryon ratio. Using for Δ the average for the fluctuations, $Q \approx 2 \times 10^{-5}$ yields $\rho_\Lambda < 7.3 \times 10^{-125}$ [with parameter values from Tegmark *et al.* (2006)]. Putting in the correct factors of the order of 1 and taking into account the contribution of neutrinos to matter-radiation equality lowers this number substantially. Clearly a more careful treatment of galactic-size density perturbations (which contribute with a third power) is needed.

Furthermore, the “bound” is not a step function. One expects a mean density of galaxies that falls off with increasing ρ_Λ . Such a function was computed by Efstathiou (1995) based on the results of COBE (but prior to the observation of accelerated expansion). Although the observation of a positive Λ in 1998 came as a shock to many, there were already several indications in that direction because of the density contribution needed for spatial flatness (as predicted by inflation) and the age of the Universe. This had already been pointed out by Weinberg (1987). The results of Efstathiou (1995) predicted a value for ρ_Λ in agreement with that expectation, although with large uncertainties, and subject to some criticisms (Vilenkin, 1995b; Weinberg, 1996). This computation was improved and done analytically rather than numerically by Martel, Shapiro, and Weinberg (1998) with similar results. Distributions for ρ_Λ based on more recent cosmological data can be found in Tegmark *et al.* (2006) and Pogosian and Vilenkin (2007).

Computations of this kind rely on several assumptions. The distribution of theoretically allowed values of ρ_Λ must be essentially flat near $\Lambda = 0$. Since $\Lambda = 0$ is not a special point from the perspective of quantum gravity, and since the relevant range is extremely small in Planck units, this seems plausible. Furthermore, the vacuum selection mechanism (for example, eternal inflation) must not prefer special values

either. This is less obvious; see Sec. VII. It is assumed that observers are correlated with galaxies, and sometimes with stars, planets, and baryons, and that we are typical observers [the “principle of mediocrity” of Vilenkin (1995a)].

The computations mentioned above assumed that only ρ_Λ varies. The possibility that Q also varies was considered by Tegmark and Rees (1998), who computed the anthropic bounds $10^{-6} < Q < 10^{-4}$ assuming $\Lambda = 0$. They also pointed out that without anthropic bounds on Q , the bound on Λ is invalid. A potentially serious problem was raised by Banks, Dine, and Gorbatov (2004), Graesser *et al.* (2004), Feldstein, Hall, and Watari (2005), and Garriga and Vilenkin (2006). Depending on models of inflation, the probability distribution can vary so steeply as a function of Q that extreme values are strongly preferred, so that the observed value $Q \approx 10^{-5}$, roughly in the middle of the anthropic range, has a very low probability of being observed (the “ Q catastrophe”). But even when both ρ_Λ and Q vary, there is a robust bound on ρ_Λ/Q^3 (Garriga and Vilenkin, 2006). See Vilenkin (2004) for a brief review of anthropic predictions for the cosmological constant.

We return briefly to the cosmological constant problem in Sec. VII, after the string theory landscape and the measure problem have been explained.

F. Possible landscapes

1. Fundamental theories

The anthropic principle discussed here is not a principle of nature and not our ultimate goal. That goal is a fundamental theory in which different quantum field theories are realized and can be sampled. The fundamental theory provides the input distributions for anthropic arguments and may in principle be falsified with the help of such arguments. But it is the fundamental theory we should try to falsify and not the anthropic principle, which is only a tool that may help us to find the theory. Once that has been achieved, the anthropic principle will only be a footnote.

We can try to decide which properties such a fundamental theory should have and which current ideas qualify. There are a few reasons to believe quantum gravity should play an essential role. In particular, one cannot discuss parameter values without dealing with the problem that they are fundamentally undetermined in a renormalizable quantum field theory. Furthermore, there are infinitely many of them in a nonrenormalizable theory like naively quantized gravity. One cannot consider changing parameters without discussing changes in vacuum energy, which can be done only in the context of gravity. So we need a fundamental theory of quantum gravity with dynamics and connectivity in the space of couplings.

2. Other landscapes?

The string theory landscape seems to fit the bill, although there is much work still to be done, and much that can go wrong. There are many ideas that are presented as competitors, and here we list a few of them to see if they qualify. We will not enter here in a discussion about the relative merits of some of these ideas as theories of quantum gravity.

Some alternative approaches to quantum gravity, for example, loop quantum gravity (Ashtekar, 1986) or dynamical triangulations (Ambjorn, Jurkiewicz, and Loll, 2004) have nothing to say about matter. Asymptotically safe gravity (Weinberg, 1976; Reuter, 1998) strongly restricts matter if quantum field theory is also required to be asymptotically safe, but cannot fix the couplings of asymptotically free gauge theories. There is no known way of physically connecting different theories. The same is true for noncommutative geometry (Chamseddine and Connes, 2008). In contrast to earlier claims it does not yield the standard model uniquely; for example, one can also obtain supersymmetric QCD (van den Broek and van Suijlekom, 2011). But it is still far from providing a useful landscape. Finite unified theories (Heinemeyer, Mondragon, and Zoupanos, 2008) also limit the possible quantum field theories, but do not yield a connected landscape. Spontaneously broken local conformal invariance was argued ('t Hooft, 2011) to be a physically motivated condition that fixes all parameters, leaving only a (perhaps denumerably infinite) number of discrete choices of gauge groups and representations.

Since all these authors agree that they do not propose an anthropic landscape, it is fair to say that in this respect string theory really is the only game in town.

3. Predictive landscapes

The existence of a landscape does not necessarily imply that all predictive power is lost. We just list some options here to counter some common philosophical objections.

Universal predictions: A large ensemble of possibilities may still have one or more universal predictions. In the case of the string landscape, what comes closest to that is a *negative* prediction, namely, the *absence* of variations in standard model parameters (see Sec. V.E). There may be other opportunities for universal predictions because of the universal existence of moduli and axions in string theory.

Sparse landscapes: If a landscape is small enough, current data may already be sufficient to find the solution that corresponds to our Universe. Having determined that, all parameters would be known exactly. The standard model data have been estimated to provide about 80 digits worth of information (Douglas and Kachru, 2007) so that a landscape of, say, 10^{30} points would realize this possibility, with many predictions left. But this is not likely to be true in the string theory landscape, if current ideas about the cosmological constant are correct. This already requires more than 10^{120} solutions, and a computation of the cosmological constant with 120 digit precision in each of them, if we want to pin down the solution exactly. See de Alwis (2007) and Deneff and Douglas (2007) for an exposition of some of the problems involved.

Friendly landscapes: It might happen that some parameters vary over a wide range, while others are sharply peaked at definite values. Toy examples of such landscapes have been constructed using scalar field potentials (Arkani-Hamed, Dimopoulos, and Kachru, 2005; Distler and Varadarajan, 2005). For a large number N of scalars, some parameters may be distributed widely, whereas others vary by a fraction $1/\sqrt{N}$. The widely distributed ones were argued to be the dimensionful ones, i.e., the weak scale and the cosmological constant. This would allow anthropic arguments for the

dimensionful parameters to be valid without eliminating the possibility for fairly sharp predictions for Yukawa couplings and hence quark and lepton masses. There might be enough variability left to allow even the anthropic constraints on those masses to be met. They might not be at the peak of their distribution, but anthropically pushed toward the tail.

Overwhelming statistics: The following example shows that the dream of an *ab initio* determination of the standard model and all its parameter values is not even necessarily inconsistent with anthropic arguments. It requires a large hierarchy of sampling probabilities, the probability for a vacuum to be selected during eternal inflation. We assume that the treacherous problem of defining these probabilities (see Sec. VI) has been solved and order the vacua according to this probability. Suppose that the m th vacuum has probability ϵ^m , where ϵ is a small number. Furthermore, assume that, on average, only one out of M vacua lands in the anthropic domain. For definiteness, we take $\epsilon = 0.1$ and $M = 1000$. The first anthropic vacuum is not likely to be the one with $m = 0$, and hence it will have a very small sampling probability, but that does not matter. The point is that the second anthropic vacuum would typically have a probability of 10^{-1000} with respect to the first. Such a scenario might be realized if one “master” vacuum dominates the population of vacua by a large statistical factor, and all other vacua are obtained from it by a sequence of tunneling events (see Sec. VI). To actually compute the dominant anthropic vacuum would require determining the master vacuum, the tunneling rates, and the anthropic domains, all of which are in principle computable without experimental input. In practice this seems utterly implausible, but in this example all observed anthropic miracles would be explained, provided the complete set of vacua is large enough and distributed in the right way, and still there would be a nearly unquestionable prediction of all parameters.

4. Catastrophic landscapes

The last scenario implicitly assumes that anthropic regions in QFT space are described by step functions, so that a given QFT either allows or does not allow life. In reality there will be smooth distributions at the boundaries, and depending on how fast they fall off there is an important potential problem: outliers in distributions may be strongly selected. To illustrate that, consider an extreme version of overwhelming statistics, suggested by Linde and Vanchurin (2010). They consider the possibility that landscape probabilities depend on the cosmological constant Λ as $\exp(24\pi^2/\Lambda)$, and that Λ can take only a discrete set of positive values $\Lambda = n/N$, $n = 1, \dots, N$. Here Λ is expressed in Planck units, and N is a large integer. In this situation, $n = 1$ is strongly favored statistically. If we define $P(n)$ as the probability for vacuum n , then we find

$$\frac{P(n)}{P(1)} = e^{-24\pi^2 N[(n-1)/n]}. \quad (3.8)$$

If the most probable vacuum $n = 1$ is ours, then $N \approx 10^{120}$, and anything else is suppressed by behemoth factors. They conclude “This means that by finding the vacuum with the smallest Λ we fix all other parameters; no additional anthropic reasoning is required.”

But this is not likely to be true. If one can define strict anthropic boundaries in field theory space, as in Fig. 1, the vacuum with smallest Λ has only a small chance of ending up within the anthropic contours. If any boundary line is in reality a contour of a Gaussian distribution, with a tail stretching over the entire parameter space, then the $n = 1$ vacuum is vastly more likely to lie somewhere in the tail. Suppose, for example, a variable x has an anthropic distribution $\propto \exp[-(x - x_0)^2/(2\sigma^2)]$, and suppose vacuum 2 happens, against all odds, to lie near the peak. Then vacuum 1 can lie $\approx \sqrt{N}$ or about 10^{60} standard deviations away from the peak and still beat vacuum 2 in overall probability.

This would be the worst possible outcome. It resembles uniqueness, but is catastrophically inferior. There would be a large landscape that does not solve any problem. It would not explain any fine-tunings, not even those of the cosmological constant itself. It is very unlikely that we would ever be able to compute the lowest Λ vacuum, because Λ would depend on all intricacies of particle physics, cosmology, and a fundamental theory, which would have to be computed with 120 digits of precision.

IV. STRING THEORY

Just as “standard model” and “anthropic principle,” “string theory” is poorly named. It owes its name to its original formulation: strings propagating through space-time and interacting by splitting and joining. But nowadays this is merely a general name for an interconnected web of theories, including some that do not have a string interpretation at all.

We introduce only a few basic concepts of string theory here. There are many excellent books on this subject, such as the classic by [Green, Schwarz, and Witten \(1987\)](#), the introductory course by [Zwiebach \(2004\)](#), the books by [Polchinski \(1998\)](#) and [Kiritis \(2007\)](#), and the recent one by [Blumenhagen, Lüst, and Theisen \(2013\)](#). These books also provide extensive references to classic string theory papers, which we omit here unless they have direct relevance to the landscape.

A. Generalities

In its most basic form, a string amplitude is derived from the following two-dimensional action:

$$S[X, \gamma] = -\frac{1}{4\pi\alpha'} \int d\sigma d\tau \sqrt{-\det \gamma} \times \sum_{\alpha\beta} \gamma^{\alpha\beta} \partial_\alpha X^\mu \partial_\beta X^\mu g_{\mu\nu}. \quad (4.1)$$

Here $X^\mu(\sigma, \tau)$ is a map from the two-dimensional surface swept out by the string (the world sheet, with coordinates σ and τ) into space-time, $\gamma_{\alpha\beta}$ is the metric on that surface, and $g_{\mu\nu}$ is the space-time metric. The parameter α' has the dimension $[\text{length}]^2$ and is related to the tension of the string as $T = 1/2\pi\alpha'$. The two-dimensional metric γ can be integrated out, so that the action takes the form of a surface area. Amplitudes are computed by performing a path integral over surfaces weighted by a factor $\exp(-iS/\hbar)$.

The modes of vibration of the propagating string are observed as particles. The particle spectrum consists of a

tachyon, a massless symmetric tensor $G_{\mu\nu}$, an antisymmetric tensor $B_{\mu\nu}$, and a scalar ϕ , the dilaton, plus an infinite tower of excitations. The interpretation of $G_{\mu\nu}$ as the graviton field implies a relation between Newton’s constant and α' :

$$G_N \propto g_s^2 (\alpha')^{(1/2)(D-2)}, \quad (4.2)$$

where g_s is the string coupling constant defined below. The parameter α' also sets the mass scale for the string excitations. Consequently, their spacing is in multiples of the Planck scale. The space-time metric $g_{\mu\nu}$ in Eq. (4.1) should be viewed as a space-time background in which the string propagates. The background can be curved, but it is subject to consistency conditions that follow from the quantization. They imply Einstein’s equations plus higher order corrections, but also restrict the number of space-time dimensions. For a flat metric, this yields the requirement $D = 26$. The other two massless fields $B_{\mu\nu}$ and a scalar ϕ can be included in a generalization of Eq. (4.1) as background fields. The dilaton couples as

$$S(X, \gamma, \phi) \propto \int d\sigma d\tau \sqrt{\gamma} R(\gamma) \phi. \quad (4.3)$$

This introduces a dependence of amplitudes on the Euler index χ of the surface as $e^{-\chi\phi}$. Hence the constant mode ϕ_0 of ϕ provides a weight factor for surfaces of different topology. This defines a loop expansion parameter: the string coupling constant $g_s = e^{\phi_0}$. It is not a constant set by hand in the action, but it is the vacuum expectation value of a scalar field. Therefore its value can be set dynamically. The only genuine parameter is α' , but this is a dimensionful quantity that sets the scale for everything else.

The bosonic string action can be generalized by adding two-dimensional fermions ψ^μ to the two-dimensional bosons X^μ , both with $\mu = 0, \dots, D-1$. Quantization consistency then requires the existence of a two-dimensional supersymmetry called *world-sheet supersymmetry* relating the bosons and the fermions. These are called fermionic strings. In flat space, they can only be consistently quantized if $D = 10$.

Another generalization is to consider two-dimensional surfaces that are not oriented, such as the Klein bottle, and surfaces with boundaries, such as the annulus. This leads to theories of open and closed strings that can exist in 26 and 10 dimensions for bosonic and fermionic strings, respectively.

Furthermore, one can make use of the fact that in free two-dimensional theories left- and right-moving modes can be treated independently. In closed string theories one can even use bosonic string modes for the left movers and fermionic ones for the right movers. These are called heterotic strings, and their flat space-time dimension is limited by the smaller of the two, namely, $D = 10$.

B. Modular invariance

Although the string theory spectrum consists of an infinite set of particles, string theory is not simply a quantum field theory with an infinite number of fields. The difference becomes manifest in the simplest closed string one-loop graph, the torus. At lowest order, the relevant integral takes the form

$$\int \frac{d^2\tau}{(\text{Im}\tau)^2} (\text{Im}\tau)^{(2-D)/2} \text{Tr} e^{2i\pi\tau(L_0 - c/24)} e^{-2i\pi\bar{\tau}(\bar{L}_0 - c/24)}.$$

The operators $L_0 - c/24$ and $\bar{L}_0 - c/24$ are the two-dimensional Hamiltonians of the left- and right-moving modes, and the trace is over the tensor product of the two Hilbert spaces. The integral in QFT would be over the entire complex upper half plane and is clearly divergent near $\tau = 0$. But in string theory the contributions to this integral consist of infinitely many identical copies of each other, and they would be overcounted if we were to integrate over the entire upper half plane. These identical copies are related by the following transformation:

$$\tau \rightarrow \frac{a\tau + b}{c\tau + d}, \quad a, b, c, d \in \mathbb{Z}, \quad ad - bc = 1. \quad (4.4)$$

The restriction to a single copy is allowed provided that the integrand is invariant under this transformation, which implies strong constraints on the spectrum of eigenvalues of L_0 and $\bar{L}_0$. These are known as modular invariance constraints.

1. Finiteness and space-time supersymmetry

Modular invariance is the real reason why closed string theory is UV finite. This holds for any closed string theory, including the bosonic string. There is a widespread belief that in order to deal with UV divergences in quantum gravity and/or quantum field theory nature must be supersymmetric at its deepest level. However, the UV finiteness of closed strings has nothing to do with space-time supersymmetry.

The τ integral may still diverge for another reason: the presence of tachyons in the spectrum. Furthermore, if the one-loop integral is nonzero, there is a dilaton tadpole, which leads to divergences at two loops and beyond because the dilaton propagator is infinite at zero momentum. But both of these problems are related to an inappropriate choice of the background and are IR rather than UV. The tachyon signals an instability, an expansion around a saddle point of the action. They are absent in certain fermionic string theories. Their absence requires fermions in the spectrum, but does not require supersymmetry.

Space-time supersymmetry automatically implies the absence of tachyons and the dilaton tadpole, but it is not an exact symmetry of nature and therefore cannot be used to argue for their absence.

2. Ten-dimensional strings

The condition of modular invariance is automatically satisfied for the bosonic string, but imposes relations among the boundary conditions of the world-sheet fermions. These conditions have several solutions: supersymmetric ones and non-supersymmetric ones, with and without tachyons.

The best-known solutions are the supersymmetric ones. There are two closed fermionic superstrings, called type IIA and type IIB, and two heterotic superstrings, distinguished by having a gauge algebra $E_8 \times E_8$ or $SO(32)$. Open string theories have to satisfy an additional constraint: cancellation of tadpoles for the $\chi = 1$ surfaces, the disk, and the cross cap. This leads to just one theory, called type I, with gauge group $SO(32)$. Apart from the type-IIA theory, all of these theories have chiral fermions in their spectrum.

C. D-branes, p forms, and fluxes

Open strings can have two kinds of boundary conditions: the Neumann boundary condition that respects space-time Poincaré invariance, and the Dirichlet boundary condition that explicitly violates it by fixing the end point of the open string to a definite space-time point. However, they can have a perfectly consistent interpretation by assuming that the open strings end on a physical object, localized in space-time and spanning a subspace of it, called a D-brane (Polchinski, 1995). In d space-time dimensions, the end points of open strings with $d - k$ Neumann boundary conditions and k Dirichlet boundary conditions sweep out a m -dimensional surface called a D_m -brane (where the “D” stands for Dirichlet and $m = d - k - 1$).

These D-branes are part of string theory as nonperturbative solutions, like solitons in field theory [see Duff, Khuri, and Lu (1995) for a review]. Since they are nonperturbative, they cannot be read off directly from the low-energy effective action of string theory, but they do betray their existence because they are sources of massless fields which do appear in the spectrum. These fields are antisymmetric tensors of rank p , called p forms. The sources for such p -form fields are membranes with $(p - 1)$ -dimensional spacelike surfaces (M_{p-1} branes) that sweep out a p -dimensional world volume V_p as they propagate. A p -form field A_p has a field strength tensor F_{p+1} , which is an antisymmetric tensor with $p + 1$ indices. All of these statements are fairly straightforward generalizations of Maxwell’s theory of electrodynamics in four dimensions, which correspond to the case $p = 1$. In this case the sources are M_0 branes (particles) that sweep out a one-dimensional world line. The relation between fields, field strengths, source branes, and their world volumes can be summarized as follows:

$$A_p \rightarrow F_{p+1} \rightarrow M_{p-1} \rightarrow V_p. \quad (4.5)$$

One can define a magnetic dual of these fields, again in analogy with electric-magnetic duality in electromagnetism. In general, this relates the field strength F_n to a field strength F_{d-n} in the following way:

$$F_{\mu_1 \dots \mu_n} = \epsilon_{\mu_1 \dots \mu_d} F^{\mu_{n+1} \dots \mu_d}. \quad (4.6)$$

In this way the field A_p is related to a field A_{d-p-2} , and the source M_{p-1} branes are dual to M_{d-p-3} branes. For electromagnetism in $d = 4$ dimensions ($p = 1$) this yields pointlike electric charges, dual to pointlike magnetic charges.

The analogy with electrodynamics extends to a quantization condition for the dual brane charges, analogous to the Dirac quantization condition for electric and magnetic charges, $eg = 2\pi k$, $k \in \mathbb{Z}$. This will play an important role in the following. On compact manifolds, these p -form fields can wrap around suitable topological cycles of the correct dimension to support them. These wrapped fields are called fluxes. An instructive toy model, using the monopole analogy, can be found in Denef, Douglas, and Kachru (2007).

In the closed string spectrum of type-II strings, p -form fields originate from the left-right combination of space-time spinors, which in their turn originate from world-sheet fermions with periodic boundary conditions along the closed string, called Ramond fermions. For this reason the part of the

spectrum containing these fermions is referred to as the “ RR sector.” In type-IIA string theories, the RR tensor fields have odd rank p , and they are sources of D_{p-1} -branes, starting with the D_0 -branes that correspond to particles. In type-IIB strings the p -form tensor fields have even rank, and the branes odd rank.

In string theory one always has two-forms $B_{\mu\nu}$ which are sourced by one-dimensional objects, the strings themselves. In ten dimensions, these are dual to five-branes. In type-II strings this gives rise to “NS5-branes,” called this because the $B_{\mu\nu}$ field originates from the combination of left- and right-moving Neveu-Schwarz fermions with antiperiodic boundary conditions along the closed string. In heterotic strings they are called heterotic five-branes.

D. Dualities, M theory, and F theory

The discovery of branes led to a plethora of proven and conjectured relations between *a priori* different string constructions. The ten-dimensional $E_8 \times E_8$ and $SO(32)$ heterotic strings can be related to each other after compactifying each of them on a circle, inverting its radius ($R \rightarrow \alpha'/R$; this is called target space duality or T duality) and giving VEVs to suitable background fields (Ginsparg, 1987). The same is true for type-IIA and type-IIB strings (Dai, Leigh, and Polchinski, 1989; Dine, Huet, and Seiberg, 1989). The $SO(32)$ heterotic string was shown to be related to the type-I $SO(32)$ string under inversion of the string coupling constant $g \rightarrow 1/g$ (strong coupling duality or S duality; Polchinski and Witten, 1996).

S duality, foreseen several years earlier by Font *et al.* (1990), produces a remarkable result for the remaining ten-dimensional theories. Type IIA is mapped to an 11-dimensional theory compactified on a circle (Townsend, 1995; Witten, 1995). The radius of the circle is proportional to the string coupling constant and is inverted as in T duality. For an infinitely large radius one obtains an uncompactified 11-dimensional theory; in the limit of small radius this compactification describes the weakly coupled type-IIA theory. The 11-dimensional theory is not a string theory. It is called “M theory.” Its field theory limit turned out to be the crown jewel of supergravity: $D = 11$ supergravity, which until then had escaped the new developments in string theory. Because of the existence of a three-form field in its spectrum it is believed that it is described by interacting two-dimensional and/or five-dimensional membranes.

A similar relation holds for the $E_8 \times E_8$ heterotic string. Its strong coupling limit can be formulated in terms of 11-dimensional M theory compactified on a line segment (Horava and Witten, 1996), the circle with two halves identified. This is sometimes called “heterotic M theory.”

Strong coupling duality maps type-IIB strings to themselves (Hull and Townsend, 1995). Furthermore the self-duality can be extended from an action just on the string coupling, and hence the dilaton, to an action on the entire dilaton-axion multiplet. This action is mathematically identical to the action of modular transformations on the two moduli of the torus, Eq. (4.4), and corresponds to the group $SL(2, \mathbb{Z})$. This isomorphism suggests a geometric understanding of the self-duality in terms of a compactification torus T_2 , whose degrees of freedom correspond to the dilaton and axion

field. An obvious guess would be that the type-IIB string may be viewed as a torus compactification of some 12-dimensional theory (Vafa, 1996). But there is no such theory. The first attempts to develop this idea led instead to a new piece of the landscape called “F theory,” consisting only of compactifications and related to $E_8 \times E_8$ heterotic strings and M theory by chains of dualities.

E. The Bousso-Polchinski mechanism

It was realized decades ago (Linde, 1984) that rank-4 field strengths of rank-3 antisymmetric tensors might play an important role in solving the cosmological constant problem. Such four-index field strengths can get constant values without breaking Lorentz invariance, namely, $F_{\mu\nu\rho\sigma} = c\epsilon_{\mu\nu\rho\sigma}$, where $\epsilon_{\mu\nu\rho\sigma}$ is the Lorentz-invariant completely antisymmetric four-index tensor. The presence of such a classical field strength in our Universe is unobservable unless we couple the theory to gravity. If we do, it gives a contribution similar to the cosmological constant Λ , in such a way that the latter is replaced by

$$\Lambda_{\text{phys}} = \Lambda - \frac{1}{48} F_{\mu\nu\rho\sigma} F^{\mu\nu\rho\sigma} = \Lambda + \frac{1}{2} c^2. \quad (4.7)$$

In string theory c is not an arbitrary real number: it is quantized (Bousso and Polchinski, 2000). This is due to a combination of the well-known Dirac quantization argument for electric charges in theories with magnetic monopoles and string theory dualities. The formula for the cosmological constant now looks something like

$$\Lambda_{\text{phys}} = \Lambda + \frac{1}{2} n^2 f^2, \quad (4.8)$$

where f is some number derived from the string theory under consideration. If instead of $F_{\mu\nu\rho\sigma}$ we were to consider an electromagnetic field, f would be something like the strength of the electromagnetic coupling e : some number of the order of 1. For generic negative values of Λ we would be able to tune Λ_{phys} to an extremely small value only if f is extremely small.

However, it turns out that string theory typically contains hundreds of fields $F_{\mu\nu\rho\sigma}$. Taking N such fields into account, the result now becomes

$$\Lambda_{\text{phys}} = \Lambda + \frac{1}{2} \sum_{i=1}^N n_i^2 f_i^2. \quad (4.9)$$

If indeed the values of f_i are distinct and incommensurate, then Eq. (4.9) defines a dense discrete set of values. Bousso and Polchinski called it a discretuum. It is an easy exercise to show that with N equal to a few hundred, and values for f_i of the order of electromagnetic couplings and small integers n_i , one can indeed obtain the required small value of Λ_{phys} , given some negative Λ .

This realizes a dynamical neutralization of Λ first proposed by Brown and Teitelboim (1987, 1988) [see Feng *et al.* (2001) for a related string realization]. This makes any field strength $F_{\mu\nu\rho\sigma}$ (and hence Λ) decay in discrete steps by bubble nucleation. This process stops as Λ approaches zero. This is analogous to the decay of an electric field between capacitor plates by pair creation of electron-positron pairs.

However, Brown and Teitelboim [as well as [Abbott \(1985\)](#) in an analogous model] already pointed out an important problem in the single field strength case they considered. First, as noted above, one has to assume an extremely small value for f . But even if one does, the last transition from an expanding dS universe to ours would take so long to complete that all matter would have been diluted (the “empty universe problem”). With multiple four-form field strengths, both problems are avoided; see [Bousso \(2008\)](#) for details.

All the ingredients used in the foregoing discussion are already present in string theory; nothing was added by hand. In particular large numbers of fields $F_{\mu\nu\rho\sigma}$ are present, and the quantization of the field strengths follows using standard arguments.

F. Four-dimensional strings and compactifications

There are essentially two ways of building string theories in four dimensions. One is to choose another background space-time geometry, and the other is to change the world-sheet theory. The geometry can be chosen as a flat four-dimensional space combined with a compact six-dimensional space. This is called “compactification.” This is not simply a matter of hand picking a manifold: it must satisfy the equations of motion of string theory and must be stable. Indeed an obvious danger is that a given manifold simply “decompactifies” to six flat dimensions. The world-sheet theory can be modified by choosing a different two-dimensional conformal field theory. In the action (4.1) and its supersymmetric analog only free bosons X or free fermions ψ are used. One can choose another two-dimensional field theory that satisfies the conditions of conformal invariance. This is called a conformal field theory (CFT). In particular, one may use interacting two-dimensional theories. Only X^μ and ψ^μ , $\mu = 0, \dots, 3$, must remain free fields.

As in ten dimensions, all four-dimensional string theories are related to others by strong-weak dualities, target space dualities, and combinations thereof. This suggests a connected landscape of four-dimensional strings.

We present here a brief sketch of the string compactification landscape. For further details see [Ibañez and Uranga \(2012\)](#) and references therein.

1. Landscape studies versus model building

The amount of work on string compactifications or four-dimensional string constructions is too vast to review here. Most of this work is focused on finding examples that match the standard model as closely as possible. This is important, at the very least as an existence proof, but it is not what we focus on in this review. Our main interest is not in finding a “model” where property X is realized, but the question if we can understand why we observe property X in our Universe, given anthropic and landscape constraints. The relative importance of these two points of view depends on how optimistic one is about the chances of finding the exact standard model as a point in the landscape.

2. General features

For phenomenological, but more importantly practical reasons most efforts have not focused on getting the standard

model, but the minimal supersymmetric standard model (MSSM). But it turns out that “minimal” is not exactly what one typically finds. Usually there are many additional fields that have not (yet) been observed. In addition to the superpartners of all the standard model particles and the additional Higgs field of the MSSM, they include moduli, axions, additional vector bosons, and additional exotic matter.

Moduli are massless scalar singlets whose presence can be understood in terms of continuous deformations of the compactification manifold or other features of the classical background fields. The vacuum expectation values of these fields generate the deformations. Typically, there are tens or hundreds of them. In the more general setting of M theory, the dilaton is part of this set as well.

Axions may be thought of as the imaginary part of the moduli, which are complex scalars in supersymmetric theories. It is useful to make the distinction, because mechanisms that give masses to moduli, as required for phenomenological reasons, sometimes leave the imaginary part untouched. Axions may provide essential clues about the landscape; see Sec. V.D.

Essentially all “raw” string spectra contain, in addition to the chiral standard model particles, large numbers of scalars and vectorlike (i.e., nonchiral) fermions. Unlike chiral fermions, they can acquire a mass if the string spectrum is perturbed, for example, by giving VEVs to moduli. If this is not generically what happens, string theory makes an incorrect prediction.

Furthermore, one often finds particles that do not match any of the observed matter representations. Notorious examples are particles with fractional electric charge or higher rank tensors. These particles may be acceptable if they are vectorlike, because one can hope that they become massive under generic perturbations.

Although superfluous particles may appear to be a curse, some of them may turn out to be a blessing. All quantum field theory parameters depend on the moduli, and hence the existence of moduli is the first step toward a landscape of possibilities. Axions can play a role in solving the strong CP problem and may also provide a significant part of dark matter. Additional gauge groups are often needed as “hidden sectors” in model building, especially for supersymmetry breaking. Extra $U(1)$ ’s may be observable through kinetic mixing ([Goodsell and Ringwald, 2010](#)). Vectorlike particles and exotics might be observed and provide evidence for string theory, although this is wishful thinking.

3. Calabi-Yau compactifications

The first examples of compactifications with chiral spectra and $N = 1$ supersymmetry were found for the $E_8 \times E_8$ heterotic string by [Candelas *et al.* \(1985\)](#). They used six-dimensional, Ricci-flat, Kähler manifolds with $SU(3)$ holonomy, called Calabi-Yau (CY) manifolds. They assumed that the $B_{\mu\nu}$ field strength $H_{\mu\nu\rho}$ vanishes, which leads to the consistency condition

$$dH = \text{Tr} R \wedge R - \frac{1}{30} \text{Tr} F \wedge F = 0. \quad (4.10)$$

This implies, in particular, a relation between the gravitational and gauge field backgrounds. This condition can be solved by using a background gauge field that is equal to the

spin connection of the manifold, embedded in an $SU(3)$ subgroup of one of the E_8 factors. In compactifications of this kind one obtains a spectrum with a gauge group $E_6 \times E_8$. The group E_6 contains the standard model gauge group $SU(3) \times SU(2) \times U(1)$ plus two additional $U(1)$'s. The group E_8 is superfluous but hidden (standard model particles do not couple to it), and may play a role in supersymmetry breaking. In these compactifications one obtains h_{11} chiral fermions in the representation (27) and h_{12} in the representation (27) of E_6 , where h_{11} and h_{12} are the topological Hodge numbers of the Calabi-Yau manifold.

The number of Calabi-Yau manifolds is large. Kreuzer and Skarke (2002) enumerated a subset associated with four-dimensional reflexive polyhedra. This list contains more than 470×10^6 topological classes with 31 108 distinct Hodge number pairs. The total number of topological classes of Calabi-Yau manifolds has been conjectured to be finite.

Strominger (1986) considered more general geometric background geometries with torsion, leading to so many possibilities that he concluded “all predictive power seems to have been lost.”

4. Orbifold compactifications

One can also compactify on a six-dimensional torus, but this does not yield chiral fermions; the same is true for the more general asymmetric torus compactifications found by Narain (1986). But string theory can also be compactified on tori with discrete identifications. The simplest example is the circle with the upper half identified with the lower half, resulting in a line segment. These are called orbifold compactifications (Dixon *et al.*, 1985) and do yield chiral fermions. These methods opened many new directions, such as orbifolds with gauge background fields (“Wilson lines”) (Ibáñez, Nilles, and Quevedo, 1987), and were soon generalized to *asymmetric orbifolds* (Narain, Sarmadi, and Vafa, 1987), where “asymmetric” refers to the way left- and right-moving modes were treated.

5. Free field theory constructions

World-sheet methods started being explored in 1986. The first idea was to exploit boson-fermion equivalence in two dimensions. In this way the artificial distinction between the two can be removed, and one can describe the heterotic string entirely in terms of free fermions (Kawai, Lewellen, and Tye, 1986; Antoniadis, Bachas, and Kounnas, 1987) or free bosons (Lerche, Lust, and Schellekens, 1987). These constructions are closely related. The free boson constructions have an elegant description in terms of even self-dual lattices for which remarkable counting formulas exist. Using such formulas and assuming a definite structure for the (bosonized) fermionic string sector, Lerche, Lust, and Schellekens (1987) arrived at a rigorous (but far from saturated) upper limit of the total number of string theories in this class: 10^{1500} .

6. Gepner models

In 1987 world-sheet constructions were extended further by the use of interacting rather than free two-dimensional conformal field theories (Gepner, 1988). The “building blocks” of this construction are two-dimensional conformal field theories with $N = 2$ world-sheet supersymmetry. These building

blocks are combined (“tensored”) in such a way that they contribute in the same way to the energy momentum tensor as six free bosons and fermions. This is measured in terms of the central charge of the Virasoro algebra, which must have a value $c = 9$. In principle the number of such building blocks is large, but in practice only a very limited set is available, namely, an infinite series of “minimal models” with central charge $c = 3k/(k+2)$, for $k = 1, \dots, \infty$. There are 168 distinct ways of adding these numbers to 9. For each of the 168 tensor combinations a number of distinct modular invariant partition functions can be constructed for a grand total of about 5000 (Fuchs *et al.*, 1990; Schellekens and Yankielowicz, 1990).

There is a close relationship between these “Gepner models” and geometric compactifications on Calabi-Yau manifolds. Exact correspondence between their spectra was found, including the number of singlets. This led to the conjecture that Gepner models are Calabi-Yau compactifications in special points of moduli space. Evidence was provided by a conjectured relation between $N = 2$ minimal models and critical points of Landau-Ginzburg models (Lerche, Vafa, and Warner, 1989; Vafa and Warner, 1989).

Modular invariance requires the left- and right-moving sectors of Gepner algebras to be the same. There is no such limitation in free CFT constructions, but these are limited by being noninteracting in two dimensions. But asymmetric and interacting CFT constructions also exist. Examples in this class were obtained using a method called “heterotic weight lifting” (Gato-Rivera and Schellekens, 2011a). In the left-moving sector one of the superconformal building blocks (combined with one of the E_8 factors) is replaced by another CFT that has no superconformal symmetry, but is isomorphic to the original building block as a modular group representation. But this is just a small step into a part of the landscape that is hard to access.

7. New directions in heterotic strings

The discovery of heterotic M theory opened many new directions. Instead of the canonical embedding of the $SU(3)$ valued spin connection of a Calabi-Yau manifold, some of these manifolds admit other bundles that can be embedded in the gauge group. In general, Eq. (4.10) is then not automatically satisfied, but in heterotic M theory one may get extra contributions from heterotic five-branes (Lalak, Pokorski, and Thomas, 1999; Lukas, Ovrut, and Waldram, 1999).

In this way one can avoid getting the standard model via the complicated route of E_6 grand unification. Some examples that have been studied are $SU(4)$ bundles (Braun *et al.*, 2006), $U(1)^4$ bundles (Anderson *et al.*, 2012), and $SU(N) \times U(1)$ bundles (Blumenhagen, Møster, and Weigand, 2006) which break E_8 to the more appealing $SO(10)$ GUTs, to $SU(5)$ GUTs, or even directly to the standard model. Extensive and systematic searches are underway that have resulted in hundreds of distinct examples (Anderson *et al.*, 2011) with the exact supersymmetric standard model spectrum, without even any vectorlike matter (but with extra gauge groups and the usual large numbers of singlets).

A more traditional orbifold approach is the “heterotic minilandscape.” This is based on a class of orbifold compactifications on a torus $T^6/\mathbb{Z}_6$ constructed so that the heterotic gauge group $E_8 \times E_8$ is broken down to different subgroups at

different fixed points, such as $SO(10)$, $SU(4)^2$, and $SU(6) \times SU(2)$. This leads to the notion of *local unification* (Forste *et al.*, 2004; Buchmuller *et al.*, 2005, 2006). The standard model gauge group is the intersection of the various “local” gauge groups realized at the fixed points. The number of three-family models in this part of the landscape is of the order of a few hundred, and there is an extensive body of work on their phenomenological successes and problems; see Lebedev *et al.* (2007) and Nilles *et al.* (2009) and references therein. But despite the name, work in this area is not really aimed at landscape distributions, but at getting the standard model.

8. Orientifolds and intersecting branes

Another way to get gauge groups in string theory is from stacks of membranes. If open strings end on a D-brane that does not fill all of space-time, a distinction must be made between their fluctuations away from the branes and the fluctuations of their end points on the branes. The former are standard string vibrations leading to gravity (as well as a dilaton and other vibrational modes of closed strings), whereas fluctuations of the end points are observable only on the brane and give rise to fermions and gauge interactions.

a. Chan-Paton groups

To get toward the standard model, one starts with type-II string theory and compactifies six dimensions on a manifold. In these theories one finds suitable D-branes coinciding with four-dimensional Minkowski space and intersecting each other in the compactified directions. These can be D5, D7, or D9 branes in type IIB and D6 branes in type IIA [some other options can be considered, but require more discussion; see, for example, Ibáñez and Uranga (2012)]. Each brane can give rise to a gauge group, called a Chan-Paton gauge group, which can be $U(N)$, $Sp(N)$ or $O(N)$ (Marcus and Sagnotti, 1987). By having several different branes one can obtain a gauge group consisting of several factors, like the one of the standard model. The brane intersections can give rise to massless string excitations of open strings with their ends on the two intersecting branes. These excitations can be fermions, and they can be chiral. Each open string end endows the fermion with a fundamental representation of one of the two Chan-Paton groups, so that the matter is in a bifundamental representation of those gauge groups.

Remarkably, a standard model family has precisely the right structure to be realized in this manner. The first example was constructed by Ibáñez, Marchesano, and Rabadan (2001) and is called the “Madrid model.” It consists of four stacks of branes, a $U(3)$ stack giving the strong interactions, a $U(2)$ or $Sp(2)$ stack for the weak interactions, plus two $U(1)$ stacks. The standard model Y charge is a linear combination of the unitary phase factors of the first, third, and fourth stacks (the stacks are labeled **a** $\cdots$ **d**)

$$Y = \frac{1}{6}Q_a + \frac{1}{2}Q_c - \frac{1}{2}Q_d.$$

This configuration is depicted in Fig. 2(a).

To build a complete model requires another topological feature, an orientifold plane, needed to cancel the tadpoles of the disk diagram. This also cancels the leading contributions to chiral anomalies. Anomalous $U(1)$ gauge bosons acquire a

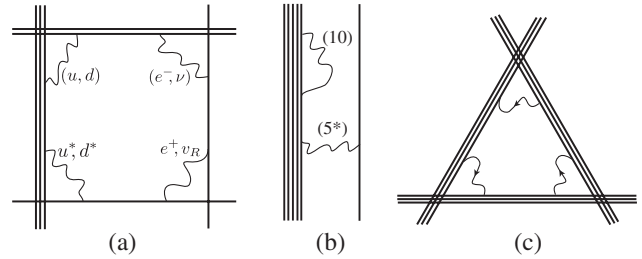


FIG. 2. Brane configurations: (a) the Madrid model, (b) $SU(5)$ GUTs, and (c) trinification.

mass by absorbing an axion field participating in a generalized Green-Schwarz mechanism. But this can also give a mass to anomaly-free $U(1)$ gauge bosons, and care must be taken that this does not happen to the standard model $U(1)$, Y . There are hundreds of papers where these conditions are solved, resulting in standard model spectra. These are called orientifold models. An extensive review of the first 5 years of this subject can be found in Blumenhagen, Cvetic *et al.* (2005).

b. The three main classes

There are other ways of getting the standard model. If there are at most four brane stacks involved, they fall into three broad classes, labeled by a real number x . The standard model generator is in general some linear combination of all four brane charges [assuming stack **b** is $U(2)$ and not $Sp(2)$], and takes the form (Anastasopoulos *et al.*, 2006)

$$Y = (x - \frac{1}{3})Q_a + (x - \frac{1}{2})Q_b + xQ_c + (x - 1)Q_d. \quad (4.11)$$

Two values of x are special. The case $x = \frac{1}{2}$ leads to a large class containing among others the Madrid model, Pati-Salam models (Pati and Salam, 1974), and flipped $SU(5)$ (Barr, 1982) models. The value $x = 0$ gives rise to classic $SU(5)$ GUTs (Georgi and Glashow, 1974). To get standard model families in this case one needs chiral antisymmetric rank-2 tensors, which originate from open strings with both their end points on the same brane. The simplest example is shown in Fig. 2(b). It has one $U(5)$ stack giving rise to the GUT gauge group, but needs at least one other brane in order to get matter in the (5^*) representation of $SU(5)$.

Other values of x can occur only for oriented strings, which means that there is a definite orientation distinguishing one end of the string from the other end. An interesting possibility in this class is the trinification model, depicted in Fig. 2(c).

c. Boundary rational CFT constructions

Just as in the heterotic string, one can construct spectra using purely geometric methods, orbifold methods, or world-sheet constructions.

World-sheet approaches use boundary CFT: conformal field theory on surfaces with boundaries and cross caps. This requires an extension of the closed string Hilbert space with “states” that describe closed strings near a boundary or in the presence of orientation reversal. An extensive formalism for computing boundary and cross cap states in (rational) CFT was developed in the last decade of the last century, starting with work by Cardy (1989), developed further by several groups, including Bianchi and Sagnotti (1990), Pradisi, Sagnotti, and Stanev (1996), Fuchs and Schweigert (1998), Huiszoon,

Schellekens, and Sousa (1999), and Behrend *et al.* (2000), culminating in a simple and general formula (Fuchs *et al.*, 2000). For an extensive review of this field, see Angelantonj and Sagnotti (2002). This was applied by Dijkstra, Huiszoon, and Schellekens (2005) to orientifolds of Gepner models and led to a large (of the order of 200.000) number of distinct string spectra that match the chiral standard model. This set provides an extensive scan over the orientifold landscape.

9. Decoupling limits

Brane model building led to an interesting change in strategy. Whereas string theory constructions were originally “top down” (one constructs a string theory and then compares with the standard model), using branes one can to some extent work in the opposite direction, “bottom up.” The idea is to start with the standard model and construct a brane configuration to match it, using branes localized at (orbifold) singularities. Then this brane configuration may be embedded in string theory at a later stage. This point of view was pioneered by Aldazabal *et al.* (2000). This is a useful approach in open string models because the gauge fields are localized on D-branes. This makes it possible to decouple gravity by sending the compactification radius to infinity. By contrast in heterotic string models both gravity and gauge interactions originate from closed string exchange, and such a decoupling limit would not make sense. Examples with $\mathbb{Z}_3$ singularities were given by Aldazabal *et al.* (2000). Berenstein, Jejjala, and Leigh (2002) considered the discrete group Δ_{27} , and Verlinde and Wijnholt (2007) used D3-branes on a del Pezzo 8 singularity.

Decoupling of gravity is an important element in recent work on F-theory GUTs (Beasley, Heckman, and Vafa, 2009a, 2009b; Donagi and Wijnholt, 2011b) obtained by compactifying F theory on elliptically fibered Calabi-Yau fourfolds. This allows the construction of models that may be thought of as nonperturbative realizations of the orientifold SU(5) GUT models depicted in Fig. 2(b), solving some of their problems, especially the absence of the top-Yukawa coupling, which is perturbatively forbidden. This has led to a revival of grand unified theories, invigorated with features of higher-dimensional theories. We return to this topic in Secs. V.A.3.d and V.B.5; see reviews by Heckman (2010), Weigand (2010), Leontaris (2011), and Maharana and Palti (2013) for further details.

The other extreme is to take the details of the standard model for granted and focus on issues like moduli, supersymmetry breaking, and hierarchies. In this case one has to assume that once the latter are solved, the standard model can be added. This is what is done in recent work on M-theory compactifications (Acharya, Kane, and Kumar, 2012). Getting chiral $N = 1$ supersymmetric spectra in M theory requires compactification on a seven-dimensional manifold with G_2 holonomy (Acharya and Witten, 2001), also known as a Joyce manifold. Much less is known about M theory than about string theory, and much less is known about Joyce manifolds than about Calabi-Yau manifolds, since the powerful tool of complex geometry is not available. For this reason the standard model is treated as input rather than output in the spirit of QFT.

Another kind of compactification that allows splitting the problem into decoupled parts is the large volume scenario

(LVS) (Balasubramanian *et al.*, 2005), originally invented for the purpose of moduli stabilization (see Sec. IV.H.1). Here both kinds of decoupling limits have been discussed, and there have also been steps toward putting both parts together (Conlon, Maharana, and Quevedo, 2009). This illustrates that focusing on decoupling limits does not mean that the original goal of a complete theory is forgotten. Indeed, there also exist *global* F-theory constructions (Blumenhagen *et al.*, 2010; Marsano *et al.*, 2013).

G. Nonsupersymmetric strings

Although the vast majority of the literature on string construction concerns space-time supersymmetric spectra, in world-sheet based methods (free bosons and fermions, Gepner models, and certain orbifolds) it is as easy to construct nonsupersymmetric ones. These spectra are generally plagued by tachyons, but by systematic searches one can find examples where no tachyons occur. This was first done in ten dimensions by Alvarez-Gaumé *et al.* (1986) and Dixon and Harvey (1986). They found a heterotic string theory with a $SO(16) \times SO(16)$ gauge group, the only tachyon-free nonsupersymmetric theory in ten dimensions, out of a total of seven. Four-dimensional nonsupersymmetric strings were already constructed shortly thereafter (Kawai, Lewellen, and Tye, 1987; Lerche, Lust, and Schellekens, 1987).

Nonsupersymmetric strings can also be constructed using orientifold methods; see, for example, Sagnotti (1995), Angelantonj (1998), Sugimoto (1999), and Gato-Rivera and Schellekens (2009). This includes the interesting possibility of having broken supersymmetry only in the open sector [“brane supersymmetry breaking” (Antoniadis, Dudas, and Sagnotti, 1999)].

Nonsupersymmetric strings can have a vacuum energy Λ of either sign. See, for example, Dienes (2006) for a distribution of values of the vacuum energy for a class of heterotic strings. Examples also exist where Λ vanishes exactly to all orders in perturbation theory (Kachru, Kumar, and Silverstein, 1999), but probably this feature does not hold beyond perturbation theory (Harvey, 1998).

Because of the lack of evidence for low-energy supersymmetry one might think that nonsupersymmetric strings are to be preferred. Unfortunately they tend to have instabilities. They all have massless scalars (at least a dilaton) that can run off toward tachyonic regions and have tadpoles that cause divergences in two-loop diagrams. There is always a dilaton tadpole. This signals that the flat background space-time that was used is not a solution to the equations of motion; instead one must use dS or AdS space with precisely the value Λ as its cosmological constant (Fischler and Susskind, 1986a, 1986b). Unfortunately this argument provides only an explanation for the presence of the tadpole, but does not provide an exact (A)dS solution.

H. The string theory landscape

A crucial test for the string landscape is the existence of (meta)stable dS vacua. They are needed for three reasons: there is evidence that our own Universe approaches such a space at late times, eternal inflation requires the existence of at least one dS vacuum, and cosmic inflation in our own

Universe may need, at least approximately, a dS space as well. Furthermore, for explanations of apparent anthropic tunings we need a large number of such spaces, and they have to be distributed in the right way.

1. Existence of de Sitter vacua

The art of constructing dS vacua is based on assembling the many ingredients of the string toolbox in a controlled way: branes, fluxes, orientifold planes, nonperturbative effects (usually in the concrete forms of “brane instantons” or gaugino condensation), world-sheet perturbative corrections, and string perturbative corrections. Fortunately, several fairly recent review articles are available; see, e.g., [Graña \(2006\)](#), [Blumenhagen, Kors *et al.* \(2007\)](#), [Douglas and Kachru \(2007\)](#), and [Denef \(2008\)](#) and the slightly more accessible one by [Denef, Douglas, and Kachru \(2007\)](#). Here we just give a brief summary and mention some recent developments.

The most explicit results have been obtained in type-IIB (and related F-theory) compactifications. One starts with a Calabi-Yau compactification. The continuous deformations of such manifolds are described by moduli of two different kinds: $h_{2,1}$ complex structure (“shape”) moduli and $h_{1,1}$ Kähler (“size”) moduli, where $h_{2,1}$ and $h_{1,1}$ are the Hodge numbers of the CY manifold. One can add three-form RR and Navier-Stokes (NS) fluxes, five-form fluxes, denoted F_3 , H_3 , and F_5 , respectively, and D3- and D7-branes.

In type-IIB theories the three-form fluxes can stabilize all complex structure moduli. This stabilization is due to a tree-level term in the superpotential that takes the form ([Gukov, Vafa, and Witten, 2000](#))

$$W_{\text{flux}} = \int (F_3 - \tau H_3) \wedge \Omega, \quad (4.12)$$

where $\tau = a + ie^{-\phi}$, with a the axion and ϕ the dilaton. The dependence on the complex structure moduli is through Ω , the holomorphic three-form of the Calabi-Yau manifold. This term also fixes the dilaton and axion. However, W_{flux} does not depend on the Kähler moduli and hence cannot fix them. This leaves therefore at least one modulus unfixed, since every CY manifold has at least one Kähler modulus.

The next step is to try and fix the size moduli with non-perturbative terms in the superpotential. These take the form $W \propto \exp(i\lambda s)$, where s is the size modulus and λ is a parameter. Such terms can be generated by instantons associated with Euclidean D3-branes ([Witten, 1996](#)) or from gaugino condensation in gauge groups on wrapped D7-branes. Assuming at least one of these effects to be present, [Kachru *et al.* \(2003\)](#) [usually referred to as Kachru, Kallosh, Linde, and Trivedi (KKLT)] obtained string vacua with all moduli stabilized. This work builds on several earlier results, such as [Dasgupta, Rajesh, and Sethi \(1999\)](#), [Klebanov and Strassler \(2000\)](#), and [Giddings, Kachru, and Polchinski \(2002\)](#), and other references therein. KKLT considered the special case $h_{1,1} = 1$, so that only one size modulus needs to be stabilized. They argued that by suitable choices of fluxes one can obtain solutions where supersymmetry is unbroken, and all world-sheet and string perturbative corrections (i.e., the α' and g_s expansion) are small. The solution obtained in this way has a negative vacuum energy and is a fully stabilized supersymmetric AdS vacuum. This is achieved by choosing fluxes so that W_{flux} is small, the

volume is large, and the dilaton (which determines the string coupling) is stabilized at a point where the coupling is small. Here “small” and “large” refer to tunings by just a few orders of magnitude.

This is, however, just a scenario, since the existence of the nonperturbative effects still needs to be demonstrated. Many would-be instantons do not contribute because of superfluous zero modes. It turns out that models with just one Kähler modulus do not work, and that instanton contributions are “not generic” ([Denef, Douglas, and Florea, 2004](#); [Robbins and Sethi, 2005](#)) but still occur sufficiently often to allow a large number of solutions.

The next step is more problematic and more controversial. One must break supersymmetry and obtain a dS vacuum (this is called “uplifting”). In KKLT this is done by adding an anti-D3-brane in a suitable location on the Calabi-Yau manifold, such that the validity of the approximations is not affected. Anti-D3-branes explicitly violate supersymmetry, and hence after introducing them one loses the control offered by supergravity. Of course, supersymmetry must be broken anyway, but it would be preferable to break it spontaneously rather than explicitly. Attempts to realize the KKLT uplifting in supergravity or string theory have failed so far ([Bena *et al.*, 2012, 2013](#)), but opinions differ on the implications of that result. There exist several alternatives to D3-brane uplifting [see, e.g., [Burgess, Kallosh, and Quevedo \(2003\)](#), [Saltman and Silverstein \(2004\)](#), and [Lebedev, Nilles, and Ratz \(2006\)](#); and also [Covi *et al.* \(2008\)](#) and [Westphal \(2008\)](#) for further references].

The result of a fully realized KKLT construction is a string vacuum that is free of tachyons, but one still has to worry about nonperturbative instability. The uplift contribution vanishes in the limit of large moduli, so there is always a supersymmetric vacuum in that limit, separated from the dS vacuum by the uplifted barrier that stabilized the AdS vacuum. One can work out the tunneling amplitude, and KKLT showed that it is generically much larger than the observed lifetime of our Universe, yet well below the theoretical upper limit in dS space, the Poincaré recurrence time. See also [Westphal \(2008\)](#) for a systematic analysis of several kinds of minima.

An alternative scenario is the LVS, already mentioned in Sec. IV.F.9. The starting point is the same: type-IIB fluxes stabilizing the complex structure moduli and the dilaton and axion. But they use α' corrections to their advantage rather than tuning parameters to minimize them. By means of suitable $(\alpha')^3$ corrections they were able to find minima where all moduli are stabilized at exponentially large volumes in nonsupersymmetric AdS vacua. The fact that α' corrections can be important at large volumes may be counterintuitive, but can be understood in terms of the no-scale structure of the underlying supergravity. For other work discussing the importance of perturbative corrections see [Becker *et al.* \(2002\)](#), [von Gersdorff and Hebecker \(2005\)](#), [Berg, Haack, and Kors \(2006\)](#), and [Bobkov \(2005\)](#). Additional mechanisms are then needed to lift the vacuum to dS. An explicit example was presented recently by [Louis *et al.* \(2012\)](#). This scenario requires special Calabi-Yau manifolds with $h_{2,1} > h_{1,1} > 1$ and a structure consisting of one large topological cycle and one or more small ones. This has

been given the suggestive name “Swiss cheese manifold.” Not every Calabi-Yau manifold has this property, but several hundreds are known (Cicoli, Kreuzer, and Mayrhofer, 2012; Gray *et al.*, 2012). A natural hierarchy can be obtained by associating standard model branes with the small cycles.

Although type-IIA and type-IIB string theories in ten dimensions differ only by a single sign flip, the discussion of moduli stabilization for the compactified theories is vastly different. This is because in type-IIA theories the available RR fluxes are even forms, and the available D-branes are D-even branes. Since there still are three-form NS fluxes one now gets flux potentials that depend on the complex structure moduli and others that depend on the Kähler moduli. As a result, all moduli can now be stabilized classically by flux potentials (DeWolfe *et al.*, 2005) [see, however, McOrist and Sethi (2012)]. Unfortunately, it can also be shown (Hertzberg *et al.*, 2007) that none of the aforementioned ingredients can be used to lift these theories to dS. There are more ingredients available, but so far no explicit examples are known [see Danielsson *et al.* (2011) for a recent attempt].

Moduli stabilization for heterotic M theory was discussed by Braun and Ovrut (2006). Supersymmetry is broken and a lift to dS achieved using heterotic five-branes and anti-five-branes. For the perturbative heterotic strings in the “mini-landscape” a scenario for moduli stabilization was presented by Dundee, Raby, and Westphal (2010). Acharya *et al.* (2006) discussed this for M-theory compactifications on manifolds with G_2 holonomy. They do not use fluxes, because in this class of models they would destroy the hierarchy. Instead, all moduli are stabilized by nonperturbative contributions generated by strong gauge dynamics. To this end they introduce two hidden sector gauge groups. A similar mechanism was applied to type-IIB theories by Bobkov *et al.* (2010). These arguments often rely on plausible but unproven assumptions about terms in potentials and nonperturbative effects. In explicit models the required terms may be absent, even though generically allowed.

2. Counting and distributions

Fluxes are characterized by integers specifying how often they wrap the topological cycles on the manifold. However, the total number of possibilities is limited by conditions for cancellation of tadpoles. For a large class of F-theory construction this condition takes the form

$$N_{D3} - N_{\overline{D3}} + \frac{1}{2\pi^4\alpha'^2} \int H_3 \wedge F_3 = \frac{\chi(X)}{24}, \quad (4.13)$$

where the first two terms denote the net contribution from D3-branes, the third one the contribution due to fluxes, and the right-hand side is a contribution (“tadpole charge”) from orientifold planes (Sethi, Vafa, and Witten, 1996); $\chi(X)$ is the Euler number of a Calabi-Yau fourfold defining the F theory under consideration. Since the flux contribution is always positive, this makes the number of possibilities finite.

This has been the starting point for estimates of the total number of flux vacua. Douglas (2004a) gave the following estimate [based on Ashok and Douglas (2004) and Denef and Douglas (2004)]:

$$N_{\text{vac}} \approx \frac{(2\pi L)^{K/2}}{(K/2)!}, \quad (4.14)$$

where L is the tadpole charge and K is the number of distinct fluxes. For typical manifolds this gives numbers of the order of 10^N , where N is of the order of a few hundred. This is the origin of the (in)famous estimate 10^{500} . Note that Eq. (4.14) should still be summed over distinct manifolds, that it only counts fluxes and no other gadgets from the string theory toolbox, and that none of these 10^{500} vacua includes the standard model, because no structure (such as intersecting D-branes or singularities) is taken into account to produce chiral matter. Indeed, the presence of chiral matter may influence moduli stabilization in a negative way (Blumenhagen, Moster, and Plauschinn, 2008).

It is worth noting that this formula turns a nuisance (a large number of moduli) into a virtue: the large number of moduli gives rise to the exponent of Eq. (4.14), and it is this large exponent that makes neutralization of the cosmological constant possible. This is not automatically true for all string compactifications and moduli stabilization mechanisms; the existence of a sufficiently large set of vacua has to be demonstrated in each case. Bobkov (2009) has shown that fluxless G_2 compactifications of M theory also yield a large discretuum of vacua.

In type-IIA constructions there are also tadpole conditions to satisfy, but in this case they do not reduce the vacuum count to a finite number. Instead it was found that supersymmetric AdS vacua exist at arbitrarily large volume, in combination with an arbitrarily small cosmological constant. This implies that the total number of vacua is infinite, but it can be made finite by making a phenomenologically inspired cut on the volume of the compactification. Acharya and Douglas (2006) presented general arguments suggesting that the number of string vacua must be finite, if one puts upper bounds on the cosmological constant and the compactification volume.

The most important contribution not taken into account in Eq. (4.14) is the effect of supersymmetry breaking. Douglas (2004a) mentioned the possibility that most of the AdS vacua might become tachyonic if such a lift is applied. Recent work indicates that this is indeed what happens. Chen *et al.* (2012) investigated this for type-IIA vacua and Marsh, McAllister, and Wrase (2012) for supergravity. They analyzed general scalar potentials using random matrices to determine the likelihood that the full mass matrix is positive definite. They found that this is exponentially suppressed by a factor $\approx \exp(-cN^p)$, where N is the number of complex scalar fields and p is estimated to lie in the range of 1.3 to 2. This suppression can be reduced if a large subset of the scalars is decoupled by giving them large supersymmetric masses. Then only the number of light scalars contributes to the suppression. Even more worrisome results were reported recently by Greene *et al.* (2013). In a study of landscapes modeled with scalar fields, they found a doubly exponential decrease of the number of metastable vacua as a function of the number of moduli, due to dramatic increases in tunneling rates.

3. Is there a string theory landscape?

It is generally accepted that a large landscape of fully stabilized supersymmetric AdS solutions exists. But these

do not describe our Universe. Not because of the observation of accelerated expansion of the Universe, but because of the much more established fact that our vacuum is not supersymmetric. Supersymmetric vacua have a vacuum energy that is bounded from above at zero. Supersymmetry breaking makes positive contributions to vacuum energy. Hence if stable nonsupersymmetric vacua exist, it would be highly surprising if their vacuum energy could not surpass the value zero. Most arguments for or against the existence of dS vacua do not really depend on the sign of the cosmological constant; $+10^{-120}$ is nearly indistinguishable from -10^{-120} . Hence one would expect distributions to behave smoothly near zero, although they may drop off rapidly.

By now there are many constructions of dS vacua, although there are always some assumptions, and it is often not possible to check the effect of higher order world-sheet or string loop corrections. But given the large number of possibilities, it requires a miracle for all of them to fail. If that is the case, there should exist some general no-go theorem that was overlooked so far.

But the mere existence of vacua with positive Λ is not enough. To make use of the Bousso-Polchinski neutralization of Λ a sufficiently dense discretuum of such vacua is needed. This mechanism relies on the fact that whatever the contribution of particle physics, cosmology, and fundamental theory, it can always be canceled to 120 significant digits by flux contributions, without making actual computations with that precision. If in reality these distributions are severely depleted in part of the range, or have a highly complicated nonflat structure, this argument fails. There might still exist a large landscape, but it would be useless.

The mighty landscape of a decade ago has been eroding at an alarming rate. The actual number of vacua is the product of large numbers divided by large suppression factors. Perhaps this will reignite dreams of a unique theory. Could it be that the product is exactly one, with the standard model and the observed cosmological constant as the only survivor? That would be an absurd example of the second gedanken computation of Sec. III.C. Any hopes that landscape erosion will reduce the number of de Sitter vacua to *one* are unfounded, but there is a risk that it will be reduced to *zero*.

More fundamental objections against the use of effective potentials in quantum gravity or the formulation of QFT and string theory in de Sitter space have been raised by Banks (2012). If these objections are valid, we may not have any theoretical methods at our disposal to deal with the apparent accelerated expansion of the Universe.

V. THE STANDARD MODEL IN THE LANDSCAPE

In this section we discuss how the main features of the standard model fit in the string theory landscape, taking into account anthropic restrictions and analytical and numerical work on landscape distributions.

A. The gauge sector

It is by now clear that string theory can reproduce the discrete structure of the standard model: the gauge group and chiral fermion representations. We cannot even begin to enumerate all the papers that succeeded in doing this.

1. Gauge group and family structure

From the landscape perspective, one might hope that the gauge group can be understood using string theory plus anthropic constraints. The anthropic constraints are hard to determine, but all three factors of the gauge group are needed for our kind of life. Electromagnetism is so essential that it is impossible to imagine life without it. One can imagine life without $SU(3)_{\text{color}}$ and only electromagnetism, but it is by no means obvious that such universes will really come to life. The weak interactions also play a crucial role in our Universe, but perhaps not in every habitable one (see Sec. III.B.1).

The choice of fermion representation is also essential, but it is even harder to determine what happens if we change it. It is possible that it is chiral in order to keep the fermions light [a plausible reason why $SU(2)_{\text{weak}}$ might be needed]. Chiral fermions have chiral anomalies that must be canceled. This fixes to some extent the particle content of a single quark and lepton family, if one insists on simplicity.

If life requires electromagnetism, a non-Abelian strong interaction group, and a chiral spectrum that becomes non-chiral after symmetry breaking at energies far below the Planck scale, perhaps the one-family standard model is the simplest option one can write down. More complicated possibilities are easy to find. For example, changing the number of colors from 3 to some odd integer N and the quark charges to p/N for suitable p , one can find an infinite series of cousins of the standard model that, for all we know, are anthropically equally valid. It is likely that in the landscape small groups are statistically favored: then $N = 3$ would be the first acceptable value. Furthermore, if small numbers of gauge group factors are favored, our standard model might be the statistically dominant anthropic choice.

There have been several studies of distributions of groups and representations in sublandscapes, but because of a lack of a sufficiently well-defined question there is no good answer either. See, e.g., Dienes (2006), Dienes *et al.* (2007), and Renner *et al.* (2012) for free fermion heterotic strings and Blumenhagen, Gmeiner *et al.* (2005), Kumar and Wells (2005), Anastasopoulos *et al.* (2006), Kumar (2006), and Balasubramanian, de Boer, and Naqvi (2010) for orientifold models. Note that all these studies, as well as others mentioned below, are for unstabilized points in supersymmetric moduli spaces. Furthermore, drawing conclusions about correlations is made difficult because of limited sampling (Dienes and Lennek, 2007, 2009).

2. The number of families

We are made out of just one family of fermions. There are no good arguments why three families should be anthropically required, although some unconvincing arguments can be pondered, based on the role of the s quark in QCD, the muon in biological mutations, the top quark in weak symmetry breaking, or the CP -violating CKM angle in baryogenesis. See also Schellekens (2008) and Gould (2010) for arguments and counterarguments.

Perhaps one day we will discover a good anthropic reason for three families. If not, the number of families was just picked out of a distribution. Multiple families are a generic feature in string theory, due to topological quantities like

Hodge numbers of compactification manifolds or intersection numbers of branes (although often this notion is muddled by attempts to distinguish families in order to explain mass hierarchies).

Landscape studies of the number of families tend to suffer from lampost artifacts: initial studies of simple models favor multiples of four or six families and disfavor three, but as more general models are studied the number three becomes less and less challenged. See, for example, [Fuchs *et al.* \(1990\)](#), [Schellekens and Yankielowicz \(1990\)](#), and [Gato-Rivera and Schellekens \(2010\)](#) versus [Gato-Rivera and Schellekens \(2011a, 2011b\)](#) for heterotic Gepner models; and [Gmeiner *et al.* \(2006\)](#) versus [Rosenhaus and Taylor \(2009\)](#) for $\mathbb{Z}_2 \times \mathbb{Z}_2$ orientifold models; see [Douglas and Taylor \(2007\)](#) for an analytical study of this case.

In a systematic scan of a class of free fermion heterotic models ([Faraggi, Kounnas, and Rizos, 2007](#)) three families occurred in about 15% of all cases. However, in a study of Gepner orientifolds with standard model gauge groups ([Dijkstra, Huiszoon, and Schellekens, 2005](#)) the number of three family spectra was about 2 orders of magnitude less than those with two families. There are many other constructions giving three families, but usually no scanning is done for other values.

Taking all these results together one may conclude that getting three families may be slightly more difficult than getting one or two, but it is at worst a landscape naturalness problem at the level of a few percent, and even this suppression may be due to the examples being too special. Therefore it is legitimate at this point to view the number of families simply as a number that came out of a distribution, which requires no further explanation.

3. Grand unification in string theory

a. Fractional charges

A remarkable feature of the quark and lepton families is the absence of fractional electric charges for color singlets. There is no evidence that free fractionally charged particles exist in nature, with a limit of less than 10^{-20} in matter ([Perl, Lee, and Loomba, 2009](#)), under certain assumptions about their charges. If indeed there are none, the global standard model gauge group is not $SU(3) \times SU(2) \times U(1)$, but $S(U(3) \times U(2))$. The reason is that the former allows representations with any real values for the $U(1)$ charge, whereas in the latter case the charges are restricted by the rule

$$\frac{t_3}{3} + \frac{t_2}{2} + \frac{1}{6} = 0 \bmod 1, \quad (5.1)$$

where t_3 is the triality of the $SU(3)$ representation and t_2 is the duality of $SU(2)$, twice the spin modulo integers. This relation implies integral charges for color-singlet states. But this is just an empirical rule. Nothing we know at present imposes such a relation. Anomaly cancellation restricts the allowed charges, but arbitrary charges, even irrational ones, can be added in nonchiral pairs or as scalar fields. In fundamental theories one may expect charges to come out quantized (due to Dirac quantization for magnetic monopoles), but that still does not imply that they are quantized in the correct way.

For almost four decades we know an excellent explanation for the empirical fact (5.1): grand unification, which embeds

the standard model in a single, simple gauge group $SU(5)$ ([Georgi and Glashow, 1974](#)). So far this idea remains just a theory. In its simplest form it made a falsifiable prediction, the decay of the proton, and this was indeed falsified.

If grand unification is a fundamental law of physics, one might hope to find a theory that unequivocally predicts it. String theory is not that theory. It seemed like that for a while in 1984, when GUTs came out “naturally” from Calabi-Yau compactifications of the $E_8 \times E_8$ heterotic string, but within a few years it became clear that GUTs are by no means the only possible outcome, and that furthermore the GUTs obtained from Calabi-Yau related compactifications do not generically break in the correct way to the standard model gauge group.

b. Heterotic strings

There are two equivalent ways of understanding why grand unification emerges so easily in $E_8 \times E_8$ heterotic strings. In Calabi-Yau compactification this comes from the embedding of the $SU(3)$ holonomy group of the manifold in one of the E_8 factors, breaking it to E_6 , an acceptable but not ideal GUT group. In world-sheet constructions this is a consequence of the “bosonic string map” ([Lerche, Lust, and Schellekens, 1987](#)) used to map the fermionic (right-moving) sector of the theory into a bosonic one, in order to be able to combine it in a modular invariant way with the left-moving sector. This automatically gives rise to a four-dimensional theory with an $SO(10) \times E_8$ gauge group and chiral fermions in the spinor representation of the first factor. This $SO(10)$ group is seen by many as the ideal GUT group. The somewhat less ideal E_6 appearing in typical Calabi-Yau compactifications is an artifact of these constructions.

But this is as good as it gets. Nothing in the structure of the standard model comes out more convincingly than this. A mechanism to break $SO(10)$ to $SU(3) \times SU(2) \times U(1)$ can be found, but it does not come out automatically. Furthermore, it works less nicely than in field theory GUTs. The heterotic string spectrum does not contain the Higgs representation used in field theory. The breaking can instead be achieved by adding background fields (Wilson lines).

But in that case the full spectrum of these heterotic strings will never satisfy Eq. (5.1), and it is precisely the deep underlying structure of string theory that is the culprit. In a string spectrum every state is relevant as is fairly obvious from the modular invariance condition. Removing one state destroys modular invariance. In this case, what one wants to remove are the extra gauge bosons in $SU(5) \subset SO(10)$ in comparison to $SU(3) \times SU(2) \times U(1)$. To do this one has to add something else to the spectrum, and it turns out that the only possibility is to add something that violates Eq. (5.1) and hence is fractionally charged ([Schellekens, 1990](#)). The possible presence of fractional charges in string spectra was first pointed out by [Wen and Witten \(1985\)](#) and the implications were discussed further by [Athanasios *et al.* \(1988\)](#).

A possible way out is that the fractional charges may all have Planck masses. They may also be vectorlike, which means that they may become massive under perturbations of the spectrum. But how often does this happen? [Assel *et al.* \(2011\)](#) made a survey of a large class of free fermionic theories with Pati-Salam spectra. They found examples with three families where all fractionally charged particles are at

the Planck mass, but only in a fraction of 10^{-5} of the chiral spectra. Gato-Rivera and Schellekens (2010, 2011a, 2011b) and Maio and Schellekens (2011) saw a similar small fraction, but examples were found only for even numbers of families. They also compared the total number of spectra with chiral and vectorlike fractional charges and found that in about 5% to 20% of the chiral non-GUT spectra the fractional charges are massless, but vectorlike. They also found some examples of fractional charges confined by an additional gauge group (i.e., not QCD).

If one assumes that in genuine string vacua vectorlike particles will always be very massive, this is a mild landscape naturalness problem. But avoiding fractional charges by chance is an unattractive solution. There may be a better way out. In orbifold models $SO(10)$ is broken using background gauge fields on Wilson lines. In this process fractional charges must appear, and therefore they must be in the twisted sector of the orbifold model. If the Wilson lines correspond to freely acting discrete symmetries of the manifold [see Witten (1985)], the twisted sector fields are massive, and hence all fractionally charged particles are heavy. This method is commonly used in Calabi-Yau based constructions [see, e.g., Anderson *et al.* (2010)], but is chosen for phenomenological reasons, and hence this does not answer the question why nature would have chosen this option. Also in the heterotic minilandscape an example was found (Błaszczuk *et al.*, 2010), but only after numerous examples with massless, vectorlike fractional charges. But they suggested another rationale for using freely acting symmetries, namely, that otherwise the standard model Y charge breaks if the orbifold singularities are “blown up.” It is not clear how that would impact models at the exact orbifold point without blowup, but at least it may point toward a solution.

In heterotic strings, the problem of fractional charges can also be avoided by considering realizations of the gauge groups in terms of higher level affine Lie algebras (Lewellen, 1990). One can even get GUT gauge groups (Kakushadze and Tye, 1997) with adjoint Higgses. But this comes out only by choice, and the same is true for the fermion representations. Generically, these will have massless higher rank tensor matter representations, which cannot occur for level 1 affine algebras.

c. Intersecting brane models

In all three classes of intersecting branes depicted in Fig. 2, fractional charges are automatically avoided for open strings with both ends on a standard model stack. But this is partly by design: these brane configurations are constructed to give at least all the particles in a standard model family, and then it turns out that there is no room anymore for additional matter. But if additional branes are added that do not contribute to the standard model gauge group (as hidden or dark matter sectors), they carry a fractional charge $\pm x \bmod 1$ [with x defined in Eq. (4.11)], so that only in the $SU(5)$ class all charges are integer.

But even in this case, one cannot speak of true unification: intersecting brane models in this class include cases (presumably the vast majority) where the $U(5)$ stack is pulled apart into a $U(3)$ and a $U(2)$ stack. This works equally well for getting the standard model representations, but without any

$SU(5)$ GUT group. This is essentially a realization of the $S(U(3) \times U(2))$ group that is sufficient to explain electric charge integrality for color singlets. This substantially weakens any claim that understanding the structure of a standard model family requires a full GUT group. Furthermore, intersecting brane GUTs allow massless symmetric rank-2 tensors (Cvetič, Papadimitriou, and Shiu, 2003), which can be avoided only by carefully hand picking spectra that do not contain them (Anastasopoulos *et al.*, 2006).

In F theory, GUT spectra were found only about 12 years after the invention of F theory, and it is therefore hard to argue that GUTs appear naturally. F-theory GUTs can be thought of as nonperturbative generalizations of the intersection brane GUTs mentioned above, and similar remarks apply. In particular, they are an option and not a prediction of string theory. However, after making this choice and putting in some information about quark masses and mixings, a remarkable group-theoretic structure emerges, which we discuss in Sec. V.B.5.

d. Coupling constant unification

It has been known for decades that the three running gauge coupling constants converge to roughly the same value at an energy scale a few orders of magnitude below the Planck scale. This requires a GUT-motivated normalization of the $U(1)$ coupling and the assumption of low-energy supersymmetry.

Just as group-theoretic unification, gauge coupling unification is not an automatic consequence of string theory, but a phenomenological input. This is illustrated in Fig. 3. Here a distribution of α_s/α_w is plotted versus $\sin^2\theta_w$ for about 200,000 intersecting brane models obtained in Dijkstra, Huiszoon, and Schellekens (2005). These spectra are of the Madrid model type depicted in Fig. 2(a). Since the gauge couplings are not related, one would not expect them to respect gauge coupling unification, and indeed they do not. One gets a broad cloud of points around the GUT point, indicated by the black circle. In this corner of the landscape, coupling unification is a mere coincidence.

In corners of the landscape with group-theoretic GUT unification, coupling unification is often problematic. This can perhaps be attributed to the fact that string theory is

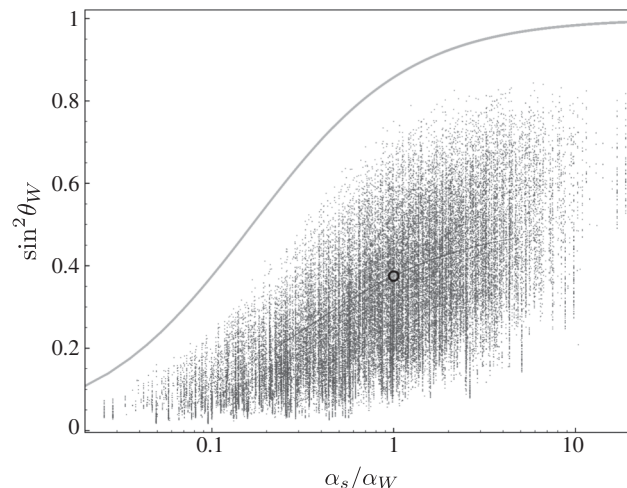


FIG. 3. Distribution of standard model couplings in a class of intersecting brane models.

simply more constraining than field theory, but it is still an indication that the perfect string GUT has not yet been found.

Heterotic GUTs predict a value for the unification scale that is substantially too large. In F theory the breaking of the SU(5) GUT group is not usually achieved by Higgses in the (24) (as in field theory) nor by Wilson lines (as in heterotic strings) but by U(1) flux in the hypercharge direction [see, however, [Marsano et al. \(2013\)](#) for an F-theory example with Wilson line breaking]. This may help to solve the notorious doublet-triplet splitting problem, but also spoils coupling unification [see [Blumenhagen \(2009\)](#) and also [Donagi and Wijnholt \(2011a\)](#) for a discussion of various contributions to thresholds]. Since there are often exotics that can contribute to the running, it may still be possible to match the observed low-energy couplings, but this turns the apparent convergence into a strange accident.

Coupling constant unification could lead to a clash between anthropic tuning and fundamental symmetries. To optimize the standard model for life, it would be better to not be constrained by a coupling constant relation, unless this is an inevitable feature of a fundamental theory. In the string landscape, it is not.

Of the three constants, g_3 is indeed anthropically constrained. It determines Λ_{QCD} and the proton mass. We discuss this in Sec. V.C. The weak coupling g_2 is much less constrained: thresholds of weak decays are much more important than the decay rates themselves. The constraints on g_1 , or almost equivalently on α , are discussed below. It does not appear to be tightly constrained, except perhaps in fine-tunings of certain nuclear levels. Unless these are much more severe than we currently know, coupling unification would not get in the way of anthropic constraints. It has two free parameters, a mass scale and the value of the unified coupling at that scale, which allow sufficient freedom to tune both Λ_{QCD} and α . Alternatively, one could argue that the value of Λ_{QCD} is tuned to its anthropic value by means of tuning of α , assuming grand unification ([Carr and Rees, 1979](#); [Hogan, 2000](#)).

e. Just a coincidence?

Standard model families have an undeniable GUT structure. One might have hoped that a bit more of that structure would emerge from a fundamental theory in a natural way, even taking into account the fact that part of this structure has anthropic relevance. GUTs can be found in several areas of string theory; see [Raby \(2011\)](#) for a review. But a compelling top-down argument in favor of GUTs is missing. Both group-theoretical and coupling unification are options in string theory, not predictions. Nevertheless, one could still speculate that grand unification is chosen in the string landscape because either GUTs are statistically favored [despite suggestions that symmetry is *not* favored ([Douglas, 2012](#))] or that it offers anthropic advantages. For example, it might turn out to play a role in inflation or baryogenesis after all, although the originally proposed GUT-based baryogenesis mechanism does not work.

But is it just a coincidence that the three running coupling constants seem to converge to a single point, close to, but just below the Planck scale? It would not be the only one. The little-known mass formula for leptons pointed out by [Koide \(1983\)](#) $m_e + m_\tau + m_\mu = \frac{2}{3}(\sqrt{m_e} + \sqrt{m_\mu} + \sqrt{m_\tau})^2$ is seen

by most as a coincidence, because it relates pole masses at different mass scales. But it predicts the τ mass correctly with 0.01% accuracy, much better than the few percent accuracy of GUT coupling unification. Another potential coincidence, allowed by the current data within 2 standard deviations, is that the self-coupling of the Higgs boson might run toward zero with vanishing β function, exactly at the Planck mass ([Bezrukov et al., 2012](#)), a behavior predicted in the context of asymptotically safe gravity [see, however, [Hebecker, Knochel, and Weigand \(2012\)](#) for an alternative idea in string theory]. Note that this coincidence is incompatible with GUT coupling unification: the latter requires low-energy supersymmetry, but the former requires a pure standard model. So at least one of these two coincidences must be just that.

4. The fine-structure constant

The fine-structure constant enters in nearly all anthropically relevant formulas, but it is often not very sharply constrained. Rather than tight constraints, one gets a large number of hierarchies of scales, such as sizes of nuclei, atoms, living beings, planets, solar systems, and galaxies, as well as time scales and typical energies of relevant processes. See [Carr and Rees \(1979\)](#), [Press et al. \(1983\)](#), [Barrow and Tipler \(1986\)](#), and [Bousso, Hall, and Nomura \(2009\)](#) for attempts to express these scales in terms of fundamental parameters, usually including α .

An example of a hierarchical condition is the requirement that the Bohr radius should be substantially larger than nuclear radii, i.e., $\alpha(m_e/m_p) \ll 1$, presumably anthropically required, but not a strong restriction on α . A stronger condition follows from the upper and lower limits of stellar masses ([Barrow and Tipler, 1986](#)):

$$\left(\frac{\alpha^2 m_p}{m_e}\right)^{3/4} N m_p \leq M_\star \leq 50 N m_p, \quad (5.2)$$

where N is the typical number of baryons in a star $N = (M_{\text{Planck}}/m_p)^3$. Requiring that the upper limit be larger than the lower one yields $\alpha^2 \leq 200(m_e/m_p)$ or $\alpha \leq 0.3$. See [Barnes \(2012\)](#) and [Tegmark \(1998\)](#) for plots of many other limits.

The value of α is constrained from above by the competition between strong and electromagnetic interactions. The electromagnetic contribution to the neutron-proton mass difference is about 0.5 MeV and proportional to α . Changing α by a factor of 3 destabilizes the proton, but this is far from determining α . In nuclei, total strong interaction binding energies scale with the number of nucleons N , electromagnetic repulsion energy scales as $\alpha N^2/R$, and R scales as $N^{1/3}$. Hence the maximum number of nucleons in a nucleus scales as $\alpha^{-3/2}$ ([Hogan, 2000](#)). Increasing α by a factor of 3 implies drastic changes, but also here a tight bound is hard to obtain. The precise location of nuclear levels is much more sensitive to α and might give tight lower and upper bounds, for example, via the beryllium bottleneck. But to draw any conclusions one has to recompute all potentially relevant nuclear levels and all types of nucleosynthesis. As a function of α , levels may not just move out of convenient locations, but also into convenient locations.

A lower bound on α can be derived from limits on the CMB fluctuations Q (Tegmark and Rees, 1998). In our Universe, $Q \approx 10^{-5}$. If Q is too large, galaxies would be too dense and planetary orbits would be disrupted too frequently; if Q is too small, the galaxies could be unable to form stars or retain heavy elements after a supernova explosion. Clearly these are not strict limits, but taking them at face value one finds that the anthropic upper limit on Q is $\approx 10^{-4}$, and scales with $\alpha^{16/7}$, whereas the lower limit is $Q \approx 10^{-6}$, scaling with $\alpha^{-1}[\ln(-\alpha)]^{-16/9}$. For smaller α the upper limit decreases and the lower limit increases. The window closes if α is about a factor of 5 smaller than $1/137.04$. This assumes everything else is kept fixed. Although the origin of the α dependence is a complicated matter, the fact that a lower bound is obtained is ultimately traceable to the need for electromagnetic cooling of matter in galaxy formation and the role of electromagnetic radiation in the functioning of the Sun. Obviously, switching off electromagnetism is bad for our health.

The competition between gravity and electromagnetism in stars is another place to look for anthropic relations. An interesting one concerns the surface temperature of typical stars compared to the ionization temperature of molecules $T_{\text{ion}} \approx \alpha^2 m_e$. These two temperatures are remarkably close. Since the former temperature depends on the relative strength of gravity and the latter does not, the coincidence implies a relation between the strength of the two interactions. Equating these temperatures gives

$$\alpha^6 \left(\frac{m_e}{m_p} \right)^2 \approx \left(\frac{m_p}{M_{\text{Planck}}} \right). \quad (5.3)$$

Numerically, both sides of this relation are 4.5×10^{-20} and 7.7×10^{-20} . Although this is close, the actual temperatures are proportional to the fourth root of these numbers so that the sensitivity is less than the formula suggests (often the square of this relation is presented, making it look even more spectacular). But does the closeness of those two temperatures have any anthropic significance? Carter (1974) conjectured that it might. Because of the temperature coincidence, typical stars are on the dividing line between radiative and convective, and he argued that this might be linked to their ability to form planetary systems [see Carr and Rees (1979) and Barrow and Tipler (1986) for a discussion]. Perhaps a more credible relation was suggested by Press *et al.* (1983), who argued that solar radiation would either be too damaging or not useful for photosynthesis if these temperatures were very different.

B. Masses and mixings

1. Anthropic limits on light quark masses

In the standard model quark masses are eigenvalues of Yukawa coupling matrices λ multiplied by the Higgs VEV v . Therefore anthropic constraints on these masses take the form of long elongated regions in the standard model (λ, v) parameter space, with rescalings in λ compensating those of v . All constraints come from the effect of changes in the quark masses on QCD and do not depend on the origin of these masses. An early discussion of the environmental impact of fermion masses can be found in Cahn (1996).

The only admissible variations in hadronic and nuclear physics are those that can be derived from variations in the relevant standard model parameters: the QCD scale Λ_{QCD} and the dimensionless ratios

$$\frac{m_u}{\Lambda_{\text{QCD}}}, \quad \frac{m_d}{\Lambda_{\text{QCD}}}, \quad \frac{m_s}{\Lambda_{\text{QCD}}}, \quad (5.4)$$

although we often just write m_u , m_d , and m_s . The strange quark is light enough to make a sizable contribution to nucleon masses by virtual processes [see Kaplan and Klebanov (1990)] and some take its variation into account (Jaffe, Jenkins, and Kimchi, 2009), even allowing it to become as light as the u and d quarks. In the limit $m_u = m_d = 0$, the chiral limit, the theory has an exact $\text{SU}(2)_L \times \text{SU}(2)_R$ symmetry, which is spontaneously broken. In this limit the pion, the Goldstone boson of the broken symmetry, is exactly massless. In the real world it has a mass proportional to $\sqrt{\Lambda_{\text{QCD}}(m_u + m_d)}$, and the pions are the only hadrons whose mass vanishes in the chiral limit. All other hadron masses are proportional to Λ_{QCD} .

In the parameter plane (5.4) one wants to know the location of several interesting anthropic boundary lines: the stability line of (^1H), the combined stability line of all forms of hydrogen, including deuterium and tritium, the stability lines of dinucleons, and the stability lines of all elements thought to be anthropically essential, as well as contour plots of all abundances. We are still very far from all that, and one can also argue about anthropic necessities. For example, deuterium and tritium can take over the role of (^1H) in biochemistry. If deuterium and all other dinucleons are unstable, synthesis of all elements from nucleons would have to start with three-body processes, but hydrogen stars could simply get hotter and denser until this happens. Keeping all these caveats in mind, we now study the location of these lines.

a. The proton-neutron mass difference

The most obvious feature of the quark masses is the extremely small up-quark mass. This is important, because the Coulomb interaction tends to make the neutron lighter than the proton, and the $m_d - m_u$ quark mass difference overcomes that. The proton-neutron mass difference can be parametrized as follows (Damour and Donoghue, 2008):

$$m_n - m_p = Z(m_d - m_u) - \epsilon_{\text{EM}}. \quad (5.5)$$

Here Z is an empirical scale factor, relating quark masses defined at some high scale to the observed mass difference. This parametrizes renormalization group running, which cannot be reliably calculated at low energy. The electromagnetic mass difference $\epsilon_{\text{EM}} \approx 0.5$ MeV is to first approximation proportional to $\alpha \Lambda_{\text{QCD}}$ [see Quigg and Shrock (2009) for more details]. For the quark masses at 2 GeV quoted by the Particle Data Group (Beringer *et al.*, 2012) one gets $Z = 0.7$.

If $m_d - m_u$ is *increased*, the neutron becomes less stable, so that it starts decaying within nuclei. Since neutrons are required for nuclear stability, this eventually implies instability of all nuclei. If $m_d - m_u$ is *decreased*, the proton becomes unstable. First the hydrogen atom becomes unstable against electron capture, for a slightly higher value the free proton can decay, and eventually all nuclei become unstable.

It is convenient to express all limits in terms of the available energy $\Delta = m_n - m_p - m_e$ in neutron decay. We assume that neutrino masses remain negligible. From electron capture and β decay of nuclei one gets, respectively, the following limits:

$$M(A, Z) - M(A, Z - 1) < \delta(\Delta) < M(A, Z + 1) - M(A, Z).$$

The masses $M(A, Z)$ used here are *atomic* masses, and hence include electron masses. The maximum variation in Δ is about ± 25 MeV (which translates to ± 35 MeV for the quark mass differences). Beyond that point no stable nuclei exist. This is a very conservative bound, which does not depend much on details of nuclear binding. Long before reaching this bound catastrophic changes occur, and there is no guarantee that the few stable nuclei can actually be synthesized.

b. Nuclear binding

While it is obvious that increasing or decreasing $m_u - m_d$ by a few tens of MeV in both directions will lead to instability of all nuclei, this is far less obvious for variations in $m_u + m_d$. An intuitive argument is suggested by the lightness of the pion. The pion mass increases with $\sqrt{m_u + m_d}$, which decrease the range of the one-pion exchange potential, and this could make nuclei less stable. But one-pion exchange is not a correct description of nuclear physics. In the literature, estimates have been given of the effect of quark mass changes on binding of heavy nuclei based on effective field theory and models for nuclear matter. Damour and Donoghue (2008) studied the binding energy per nucleon for heavy nuclei as a function of scalar and vector contact interactions. According to Damour and Donoghue, a conservative estimate for the maximum allowed increase in $m_u + m_d$ is about 64%.

c. Bounds on the Higgs VEV

The limits discussed above are often expressed in terms of allowed variations of the Higgs vacuum expectation value, under the assumption that the Yukawa couplings are kept fixed. The upper bound of Δ of 25 MeV translates into an upper bound on v/v_0 (where v_0 is the observed value) of about 20. The negative lower bound has no effect, because v cannot be negative. But if one just requires the stability of hydrogen ^1H under electron capture, the bound is $\Delta > 0$, which implies (but note that the error in $m_d - m_u$ is large)

$$\frac{v}{v_0} > \frac{\epsilon_{\text{EM}}}{Z(m_d - m_u) - m_e} \approx 0.4. \quad (5.6)$$

Here we used the method of Damour and Donoghue (2008); Hogan (2006) estimated the lower bound as 0.6 ± 0.2 using lattice results on isospin violation (Beane, Orginos, and Savage, 2007). If we also use the more model-dependent nuclear binding bounds, the window for v/v_0 is quite small, $0.4 < v/v_0 < 1.64$.

Limits on v/v_0 were first presented by Agrawal *et al.* (1998b), who estimated an upper limit $v/v_0 < 5$, from a combination of the two arguments on stability of nuclei discussed above. In this work the Higgs mass parameter μ^2 is varied over its entire range, from $-M_{\text{Planck}}^2$ to $+M_{\text{Planck}}^2$, while keeping all other parameters in the Lagrangian fixed. Then if μ^2 is negative, $v = \sqrt{-\mu^2/\lambda}$, and v/v_0 can lie anywhere between 0 and 10^{17} GeV. The anthropic range is

in any case extremely small in comparison to the full allowed range. Note that for $v/v_0 > 10^3$ a qualitative change occurs, because the stable particle will be the Δ^{++} instead of the proton; however, this is not expected to improve the odds for complex life.

The interesting and important case $\mu^2 > 0$ (no Higgs mechanism, but quarks and leptons getting a mass from the pion VEV) is also discussed in these papers; see also Quigg and Shrock (2009). The arguments against this case rest on the electron mass becoming too small, so that all matter increases in size and decreases in average density and typical biochemical temperatures are reduced.

An updated discussion of bounds on quark masses can be found in Barr and Khan (2007). They also considered the possibility of having separate up- and down-quark Higgs bosons, each with variable scales, while the Yukawa couplings are kept fixed.

d. Big bang nucleosynthesis

In our kind of Universe big bang nucleosynthesis (BBN) leads mainly to production of ^4He , ^1H , and small amounts of deuterium, tritium, and lithium. The main potential impact of BBN is therefore a destructive one: there might be too little hydrogen left. A hydrogenless universe is anthropically challenged, but there are no obvious arguments against the other extreme, a heliumless universe (Carr and Rees, 1979). Helium is needed as a stepping stone to heavier elements, but can also be made in stars.

In which extreme we end up is to a large extent determined by the electroweak freeze-out temperature (the temperature where the rate of electroweak $n \leftrightarrow p$ conversions drops below the expansion rate)

$$T_f \approx \left(\frac{G_N}{G_F^4} \right)^{1/6} = (v/M_{\text{Planck}})^{1/3} v \approx 0.66 \text{ MeV}, \quad (5.7)$$

where v is the Higgs VEV. At temperatures above T_f protons and neutrons are in thermodynamic equilibrium, and their ratio is given by a Boltzmann factor $n/p = \exp[-(m_n - m_p)/T_f]$. At T_f the ratio n/p is “frozen” and decreases only slightly because of neutron decay. After freeze-out, the outcome of BBN is determined only by strong interactions, which conserve flavor. They burn essentially all remaining baryons into helium, removing equal amounts of p and n . Hence one ends up with a fraction of hydrogen equal to $(p - n)/(p + n)$ at freeze-out. This fraction approaches the danger zone (no ^1H) if

$$\left(\frac{m_n - m_p}{v} \right) \left(\frac{M_{\text{Planck}}}{v} \right)^{1/3} \rightarrow 0. \quad (5.8)$$

This remarkable quantity involves all four interactions, since $m_n - m_p$ receives contributions from quark mass differences and electromagnetic effects. The latter are proportional to Λ_{QCD} , and in this way BBN is sensitive to changes in that scale (Kneller and McLaughlin, 2003).

There are two remarkable order of magnitude coincidences here: $T_f \approx m_n - m_p$, and the neutron lifetime τ_n is of the order of the duration of nucleosynthesis. It is not clear if these have any anthropic relevance. Increasing $m_n - m_p$ and decreasing τ_n to more natural values leads to a larger fraction of

¹H. It is almost as if these quantities are *antianthropically* tuned. The hydrogen fraction is only moderately sensitive to increases of ν , since for large ν the dependence cancels out in the first factor, and the neutron lifetime decreases. Even if we ignore the latter, an increase of ν by a factor of 1000 decreases the mass fraction of hydrogen from 75% to 6%. It is hard to argue that this would not be enough.

e. Few-nucleon systems

The stability properties of two- and three-nucleon systems certainly look fine-tuned in our Universe: Deuterium is just bound by 1.1 MeV per nucleon, diprotons and dineutrons are just not bound by about 60–70 keV. Tritium is much more strongly bound than deuterium but β decays to ³He. But a decrease of the neutron-proton mass difference by a mere 20 keV would make it stable. Once β decay is forbidden, tritium may be stable even after the deuterium stability line has been crossed, because of its higher binding energy.

Possible consequences of tritium stability on stars, apart from its potential role in chemistry, were discussed by Gould (2012). Gould speculates that changes in fusion processes in stars could affect the formation of planets.

Claims about the important impact of diproton stability on BBN, in much of the literature on anthropic tuning, are probably exaggerated, as they incorrectly assume that the diproton production cross section would be comparable to that of deuterium (Bradford, 2009; MacDonald and Mullan, 2009).

Stability of dinuclei has a big impact on stars. If the diproton were stable, the deuteron production rate could be 10 orders of magnitude larger than in our Universe, with unknown consequences (Bradford, 2009). So the diproton stability line, if it exists at all, marks the end of our region and the beginning of *terra incognita*.

The tritium stability line can undoubtedly be crossed by changing the quark masses, but for the other stability lines this cannot be decided without a more detailed look at nuclear binding. The dependence of binding on quark masses is still uncertain. For instance, it is not clear if the deuteron is bound in the chiral limit; see Beane and Savage (2003a, 2003b), and Epelbaum, Meissner, and Gloeckle (2003). For recent results and references on the impact of variations of quark masses on nuclear forces and BBN, see Berengut *et al.* (2013).⁵

Properties of few-nucleon systems are potentially anthropically relevant and appear to be fine-tuned, but too little is known about either to draw firm conclusions.

f. The triple- α process

BBN ends with a universe consisting mainly of protons, electrons, and α particles. Fusion to heavier elements is

inhibited because there are no stable nuclei with $A = 5$ or 8. Hence there are no paths with only two-particle reactions leading to heavier nuclei. The most obvious path to ¹²C is $\alpha + \alpha \rightarrow$ ⁸Be, followed by ⁸Be + $\alpha \rightarrow$ ¹²C. But ⁸Be is unstable with a lifetime of about 10^{-16} s, so this does not look promising.

There are at least three facts that improve the situation. First, the ⁸Be ground state is a very narrow resonance in the $\alpha\alpha$ channel, enhancing the first process. The narrowness of this resonance is due to a remarkable tuning of strong versus electromagnetic interactions (Higa, Hammer, and van Kolck, 2008). Second, there is a resonance of ¹²C (the second excitation level) that enhances the second process. Finally, a logical third step in this chain, ¹²C + $\alpha \rightarrow$ ¹⁶O, is not enhanced by a resonance. If that were the case, all ¹²C would be burned to ¹⁶O. Indeed, there is a resonance in ¹⁶O (at 7.10 MeV) that lies close to, but just below the ¹²C + α threshold at 7.16 MeV.

The reaction rate of the triple- α process is given by (Burbidge *et al.*, 1957)

$$r_{3\alpha} \propto \Gamma_{\gamma} \left(\frac{N_{\alpha}}{k_B T} \right)^3 e^{-\epsilon/k_B T}, \quad (5.9)$$

where $\epsilon \approx 397$ keV is the energy of the ¹²C resonance above the 3α threshold, Γ_{γ} is the width of its radiative decay into ¹²C, and N_{α} is the α -particle number density. This formula enters into the calculation of element abundances, which can be compared with observations. Assuming ¹²C synthesis takes place in the late stage of red giants at temperatures of order 10^8 K one can then fit ϵ to the observed abundances. This was done by Hoyle (1954) and led to a prediction for ϵ , which in its turn led to a prediction of an excited level of ¹²C at 7.65 MeV above the ground state. This resonance (now known as the “Hoyle state”) was indeed found. For an account of the physics and the history, see Kragh (2010).

Since the abundance of carbon is at stake, it is tempting to draw anthropic conclusions. But there are several caveats. Carbon production is obviously not maximized for the observed value of ϵ : for smaller ϵ the rate is even larger. One cannot assume that if ϵ is changed, T remains fixed. Since the triple- α process must provide energy to counterbalance gravitational pressure, it is inevitable that the star compresses to higher densities and temperatures if ϵ is increased. Furthermore, one should also take oxygen production into account. At higher temperatures ¹⁶O production starts becoming more important. The net effect is that if ϵ is increased, a larger fraction of helium is burned to ¹⁶O and a smaller fraction to ¹²C. To compute an optimum, one would have to know the optimal carbon-to-oxygen ratio for life, and without a theory, and only our own kind of life as data, this is impossible. An additional complication is that for smaller ϵ red giant-type stars produce very little ¹⁶O, but more massive, hotter stars can take over. Even if no ¹²C is formed or all of it is destroyed, there would still be heavier elements, and perhaps there can be complexity and life without carbon.

Without the Hoyle state the third excited state of ¹²C at 9.64 could take over its role, but then stars would burn at such high temperatures that even primordial ¹²C would be destroyed (Livio *et al.*, 1989). Hence the existence of the Hoyle state is indeed important for our kind of life.

⁵Many papers studying the impact of variations on BBN or the triple- α process consider *observational* constraints, for the purpose of detecting variations in constants of nature. This should not be confused with *anthropic* constraints. Another source of confusion is that some authors convert variations in the strong force to variations in α via an assumed GUT relation as explained by Calmet and Fritzsch (2002) and Langacker, Segre, and Strassler (2002). This greatly enhances the sensitivity to variations in α ; see, e.g., Coc *et al.* (2010).

However, according to Weinberg (2005) the existence of the Hoyle state in ^{12}C can be understood on the basis of collective dynamics of α particles and hence is not a major surprise.

The quantitative effect of changes of the resonance energy was studied by Livio *et al.* (1989). They varied the excitation level in large steps in numerical stellar nucleosynthesis models and found that for an upward change of 277 keV or more very little ^{12}C is produced. For an increase of 60 keV there was no significant change, whereas a decrease of 60 keV led to a fourfold increase in ^{12}C . Schlattl *et al.* (2004), using more advanced stellar evolution codes that follow the entire evolution of massive stars, found that in a band of ± 100 keV around the resonance energy the changes in abundances are small.

To decide how fine-tuned this is one wants to see the effect of standard model parameter changes. The first step in that direction was made by Oberhummer, Csoto, and Schlattl (2000), who studied the effect on the resonance energy of rescalings of the nucleon-nucleon and Coulomb potentials. They concluded that changes of 0.5% and 4%, respectively, led to changes in C or O abundances by more than an order of magnitude. However, Schlattl *et al.* (2004) weakened these conclusions. Using nuclear lattice simulations Epelbaum *et al.* (2013) concluded that ^{12}C and ^{16}O production would survive a 2% change in the light quark masses or the fine-structure constant. This band corresponds to a change of around 100 keV in the Hoyle state energy. Exactly how far one can venture outside that band is a complicated issue, since a proper treatment requires keeping track of all changes in nuclear levels, the rates of all processes, and the effect on models for stellar evolution. Processes that are irrelevant in our Universe may become dominant in others.

One can try to convert these survivability bands in terms of variations of the Higgs VEV, the common scale of the quark masses. The naive expectation is that enlarging the Higgs VEV increases the pion mass, which weakens the nuclear potential, which, according to Oberhummer, Csoto, and Schlattl (2000), increases the resonance energy and hence lowers the C/O ratio. If one focuses only on ^{12}C (assuming oxygen can be made elsewhere), this would put an upper limit on the Higgs VEV v . Indeed, Hogan (2006), using Weinberg's model of collective α -particle excitations to determine the v dependence, found an *upper* bound on v about 5% above its observed value. But Jeltema and Sher (1999), using the results of Oberhummer, Csoto, and Schlattl (2000), found a *lower* limit on v about 1% below its observed value. The discrepancy may be due to a different treatment of nuclear forces or a different slice through the parameter space: in the first work Λ_{QCD} is kept fixed, whereas in the second the strong coupling is kept fixed at the GUT scale. Then changes in v affect Λ_{QCD} because of changes in quark mass thresholds.

Expressed in terms of changes of v , the results of Epelbaum *et al.* (2013) indicated that the Hoyle state energy goes up when v is increased, but there are contributing terms with different signs and large errors. Therefore the opposite dependence is not entirely ruled out.

Even the most conservative interpretation of all this work still implies that a minute change of v with respect to Λ_{QCD} in either direction has drastic consequences. Note that the full scale of v/v_0 goes up to 10^{17} , and the variations discussed above are by just a few percent.

2. The top quark mass

The top quark may not seem a target for anthropic arguments, but it may well be important because of its large coupling to the Higgs boson, which plays a dominant role in the renormalization group running of parameters. In supersymmetric theories, this large coupling may drive the Higgs μ^2 parameter to negative values, triggering electroweak symmetry breaking [see Ibáñez and Ross (1982); since this work preceded the top quark discovery, they could only speculate about its mass].

The large top quark mass may also play an important role in the standard model, although the mechanism is less clear-cut; see Feldstein, Hall, and Watari (2006). They argued that in a landscape the top quark mass is pushed to large values to enhance vacuum stability. This issue was reanalyzed by Giudice, Perez, and Soreq (2012) using the recent data on the Higgs mass and under somewhat different assumptions. They concluded that the quark masses may be understood in terms of a broad distribution centered around 1 GeV, with the light quark masses and the top quark mass as outliers, pushed to the limits by anthropic (atomic or stability) pressures.

3. Charged lepton masses

The electron mass is bounded from above by the limits from nuclear stability already discussed in Sec. V.B. If the electron is a factor of 2.5 heavier, hydrogen ^1H is unstable against electron capture; if one can live with tritium the bound goes up to about 10 MeV. Beyond that bound most heavy nuclei are unstable as well. See Jenkins (2009) for other, less restrictive bounds, for example, the fact that a much heavier electron (by a factor of ≥ 100) would give rise to electron-catalyzed fusion in matter.

There are several arguments for smallness of the electron mass in comparison to the proton mass. The bound $(m_e/m_p)^{1/4} \ll 1$ is important for having matter with localized nuclei (Barrow and Tipler, 1986), but there is no clear limit. Limits on hierarchies of scales (e.g., Bohr radius versus nuclear radius, see Sec. V.A.4) are not very tight because the electron mass is multiplied with powers of α .

There are also lower bounds on the electron mass, but mostly qualitative ones. Lowering the electron mass enhances the Thomson scattering cross section that determines the opacity of stars. It affects the temperature of recombination and all chemical and biological temperatures. The stellar mass window (5.2) gives a bound on m_e because the lower limit must be smaller than the upper one: $m_e > 0.005\alpha^2 m_p \approx 250$ eV.

If muon radiation plays an important role in DNA mutations, then the location of the muon mass just below the pion mass would be important [see footnote 17 in Banks, Dine, and Gorbatov (2004)]. But the danger of anthropocentrism is large here.

4. Masses and mixings in the landscape

In theoretical ideas about quark masses one can clearly distinguish two antipodes: anarchy versus symmetry. In the former case one assumes that masses and mixings result from Yukawa couplings that are randomly selected from some distribution, whereas in the latter case one tries to identify

flavor symmetries or other structures that give the desired result.

The quark mass hierarchies are unlikely to come out of a flat distribution of Yukawa couplings. However, one can get roughly the right answer from scale-invariant distributions (Donoghue, 1998)

$$f(\lambda) = \rho(\lambda)d\lambda, \quad \rho(\lambda) \propto \frac{1}{\lambda}, \quad (5.10)$$

where $f(\lambda)$ is the fraction of values between λ and $\lambda + d\lambda$. A flat distribution is obtained for $\rho = \text{const.}$ Scale-invariant distributions are generated by exponentials of random numbers. In string theory, this can come out easily if the exponent is an action. A canonical example is a “world-sheet instanton,” where the action is the area of the surface spanned between three curves in a compact space. In intersecting brane models of the Madrid type shown in Fig. 2(a) this is indeed how Yukawa couplings are generated from the branes whose intersections produce the left-handed quarks, the right-handed quarks, and the Higgs boson. Note that both types of distributions require small and large λ cutoffs in order to be normalizable. In the intersecting brane picture this comes out automatically since on a compact surface there is a minimal and a maximal surface area.

The smallness of the CKM angles makes a very convincing case against flat distributions. This is illustrated in Fig. 4(a). Here 2×2 random complex matrices M are considered, with entries chosen from two different distributions. What is plotted is the distribution of the values of the rotation angle required to diagonalize the matrix (this requires separate left and right matrices, and the angle is extracted from one of them). The gray line is for a flat distribution of matrix elements $M_{ij} = r_1 + ir_2$, where r_1 and r_2 are random numbers in the interval $[-1, 1]$. The black line is for a scale-invariant distribution $M_{ij} = e^{-sr_1} e^{2\pi ir_2}$, where r_1 and r_2 are random numbers between 0 and 1, and s is a real parameter. In the figure $s = 5$ was used. As s is increased, the angle distribution starts developing a peak at small angles, but also near 90° . Clearly, small angles are unlikely for flat distributions, but not for scale-invariant ones.

This is easy to understand. If a random matrix is generated with a scale-invariant distribution, typically one matrix element will be much larger than all others and will select the required rotation. If it is on the diagonal, no rotation is needed, and if it is off diagonal, one of the two matrices will have to make a 90° rotation.

This becomes a bit more murky for 3×3 matrices, but the main trait persists in the full CKM matrix. In Fig. 4(b) we

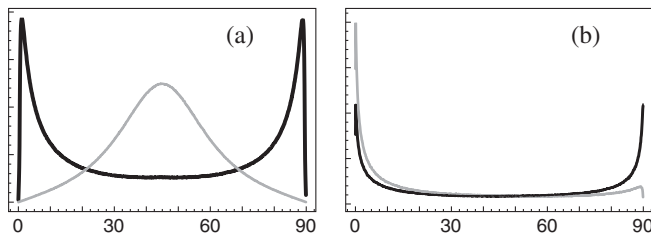


FIG. 4. Distribution of CKM angles at small and large angles for a scale-invariant distribution. The black line is for θ_{12} and θ_{23} , the gray line is for θ_{13} .

show the distribution for the three angles in the CKM matrix, with M_u and M_d distributed as above, but with $s = 12$. Only one phenomenological constraint was put in, namely, that the top quark mass must be at least 10 times the bottom quark mass; all other combinations of M_u and M_d are rejected. The largest mass was scaled to m_t by means of a common factor (the Higgs VEV). The distributions for θ_{12} and θ_{23} are indistinguishable and symmetric on the interval $[0^\circ, 90^\circ]$ and are peaked at both ends, while the distribution for θ_{13} is more strongly peaked and only near $\theta_{13} = 0$. There is a large plateau in the middle, and for θ_{12} and θ_{23} the peak is 40 times above the value at 45° . For larger values of s the peaks become more pronounced and move toward the asymptotes at 0° and 90° .

The eigenvalue distribution is even more interesting and is shown in Fig. 5. No special effort was made to fit the single parameter s to the observed quark masses and mixings; the value $s = 12$ was chosen just to get roughly in the right ballpark, for illustrative purposes only. Note that the difference between the two plots is entirely due to the requirement $m_t > 10m_b$. Renormalization group running was not taken into account. This might favor large top quark masses because of the infrared fixed point of the Yukawa couplings (Donoghue, 1998).

The angular distributions easily accommodate the observed values $\theta_{12} = 13^\circ$, $\theta_{23} = 2.38^\circ$, and $\theta_{13} = 0.2^\circ$, and the mass distributions have no difficulties with the observed mass hierarchies. Furthermore, the lowest eigenvalues have very broad distributions, so that they can easily accommodate the anthropic requirements for m_u , m_d , and the electron mass. Note that the angular distributions predict that two of the three angles are just as likely to be large ($\approx 90^\circ$) as small. Hence the observation that all three are small comes out in about one-quarter of all cases. Furthermore, there are large central plateaus.

A much more complete analysis, including renormalization group running, was done by Donoghue, Dutta, and Ross (2006). They considered more general distributions $\rho(\lambda) = \lambda^{-\delta}$, determined the optimal distribution from the quark masses, and computed the median values of the CKM matrix elements. They did indeed obtain the correct hierarchies in the angles. They also worked out the distribution of the Jarlskog invariant and found that it peaks at roughly the right value. The latter invariant was also considered by Gibbons *et al.* (2009), who introduced a natural measure on the four-dimensional coset space that is defined by the CKM matrix, $U(1)^2 \backslash SU(3)/U(1)^2$. Taking the observed quark masses into account, they obtained a likely value for J close to the observed one.

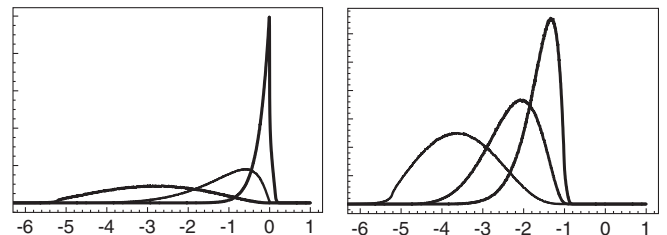


FIG. 5. Distribution of up-type (u, c, t) and down-type (d, s, b) masses. On the horizontal axis powers of ten are indicated.

An analysis that is similar in spirit was done by [Hall, Salem, and Watari \(2007, 2008\)](#). Instead of scale-invariant distributions, they assumed that Yukawa couplings derived from overlap integrals of Gaussian wave functions in extra dimensions, using a mechanism due to [Arkani-Hamed and Schmaltz \(2000\)](#) to generate hierarchies and small mixing from strongly localized wave functions in extra dimensions. An advantage of this mechanism is that wrong pairings (large mixing angles between up-type and down-type quarks of different families) are strongly suppressed. This method also accommodates all observed features of quark masses and mixings rather easily.

5. Landscape versus symmetries

The landscape ideas discussed above suggest that elaborate symmetries are not needed to understand the observed masses and mixings.

But there might be structure in the Yukawa matrices. An interesting suggestion is gauge-top unification, which is found to occur in a subset of minilandscape models. This singles out the top quark and relates its Yukawa couplings directly to the gauge couplings at the unification scale. In addition, there is a D_4 discrete symmetry relating the first two families. See [Mayorga Peña, Nilles, and Oehlmann \(2012\)](#) for further discussion and references.

In the simplest possible orientifold models, for example, the ones depicted in [Fig. 2](#), all families are on equal footing. But this is not always the case, and there are many examples where different families have their end points on different branes. This gives rise to Yukawa coupling matrices where some entries are perturbatively forbidden, but can be generated by D-brane instantons, giving rise to a hierarchy of scales. Several possibilities were investigated by [Anastasopoulos, Kiritsis, and Lionetto \(2009\)](#).

Almost the exact opposite of landscape anarchy has emerged in the context of F theory. The most striking phenomenon is a stepwise enhancement of symmetries toward E_8 . Gauge fields live on D7-branes, which have an eight-dimensional world volume. Four of these dimensions coincide with Minkowski space, and the other four wrap a four-dimensional volume in the eight-dimensional Calabi-Yau fourfold that defines F theory. Two-dimensional intersection curves of the four-dimensional curves correspond to matter, and pointlike triple intersections of matter curves correspond to Yukawa couplings. This leads to enrichment of previous GUT ideas into higher dimensions: gravity sees all dimensions, gauge groups live on eight-dimensional surfaces, matter on six-dimensional surfaces, and three-point couplings are localized in four dimensions or just a point in the compactified space.

The properties of gauge groups and matter are determined by ADE-type singularities defined by the embedding of these surfaces in the elliptically fibered Calabi-Yau fourfold. To get the required GUT group one starts with seven-branes with an $SU(5)$ singularity. The matter curves have an enhanced singularity; to get a $(\bar{5})$ of $SU(5)$ the singularity must enhance $SU(5)$ to $SU(6)$, and to get a $(\mathbf{10})$ it must enhance it to $SO(10)$. Further enhancements occur for the pointlike singularities that correspond to Yukawa couplings: to get the $\mathbf{10.5.5}$ down-quark couplings one needs an $SO(12)$ singularity, and to get the $\mathbf{10.10.5}$ up-quark couplings one needs E_6 .

The Yukawa couplings are, to first approximation, rank-1 matrices, which implies that each has one nonvanishing eigenvalue (t , b , and τ) and two zero eigenvalues. But two arbitrary rank-1 matrices will have their eigenvectors pointing in unrelated directions, and since the CKM matrix is defined by the relative orientation, it will in general not be close to 1, as it should be. This can be solved by assuming that the top and down Yukawa points lie very close to each other. If they coincide, the singularity is enhanced to E_7 [which contains both E_6 and $SO(12)$]. Finally there are arguments based on neutrino physics that suggest that the singularity must be further enhanced to E_8 ([Heckman, Tavanfar, and Vafa, 2010](#)). Although this group-theoretic structure gained attention in recent F-theory GUT constructions ([Heckman and Vafa, 2010](#)), it was described prior to that by [Tatar and Watari \(2006\)](#) in a more general setting, applied to heterotic strings, M theory, and F theory. They derived the E_7 structure requiring the absence of baryon number violation dimension-4 operators.

To get nonzero values for the other masses, a mechanism like the one of [Froggatt and Nielsen \(1979\)](#) was proposed. This works by postulating one or more additional $U(1)$'s and assigning different charges to the different families. [Heckman and Vafa \(2010\)](#) showed that similar $U(1)$ symmetries automatically exist in certain F-theory compactifications, and that they could lead to the required hierarchies and small mixing angles. These are parametrized in terms of a small parameter $\epsilon \approx \sqrt{\alpha_{\text{GUT}}} \approx 0.2$. But to actually obtain deviations from rank-1 matrices has been a fairly long struggle, since some expected contributions turned out to respect the exact rank-1 structure. For recent work and further references, see [Font et al. \(2013\)](#).

But important questions remain: Why would we find ourselves at or close to an E_8 point in the landscape? A CKM matrix close to 1 is *phenomenologically*, but not *anthropically* required. It is not clear how the exact values are distributed. One should also ask the question if, in any of the methods discussed, the quark mass hierarchies and mixings would have been even roughly predicted, if we had not known them already.

6. Neutrinos

There is a lot to say about neutrino masses in string theory and other theories, but here we focus on landscape and anthropic issues. For a summary of what is known about neutrinos see [Sec. II](#), and for a recent review of various new ideas, see [Langacker \(2012\)](#).

a. The seesaw mechanism

Neutrinos offer an interesting confrontation between “new physics” and anthropic arguments. On the one hand, small neutrino masses are explained convincingly by the seesaw mechanism, which requires nothing more than a number of singlet fermions, Yukawa couplings between these singlets, and the lepton doublets and Majorana masses for the singlets. In the string landscape the singlets are generically present because most standard model realizations are $SO(10)$ related and because singlets are abundant in nearly all string compactifications. Unlike $SO(10)$ -related singlets, generic singlets

usually do not have Yukawa couplings with charged leptons, but those couplings may be generated by scalar VEVs; see [Buchmüller *et al.* \(2007\)](#) for an explicit heterotic string example.

Majorana masses tend to be a bigger obstacle. It is not obvious that string theory satisfies the QFT lore that “anything that is allowed is obligatory,” which implies that all allowed masses are nonzero, and, in particular, that all singlets must have Majorana masses. In an extensive study of the superpotential of a class of heterotic strings, [Giedt *et al.* \(2005\)](#) found no examples of such mass terms. Even if such examples were found in other cases [see, e.g., [Buchmüller *et al.* \(2007\)](#) and [Lebedev *et al.* \(2008\)](#)], this still casts doubt on the generic presence of Majorana masses. But perhaps the examples are too special, and perhaps all singlet fermions have large masses in generic, nonsupersymmetric, fully stabilized vacua. If not, string theory is facing the serious problem of predicting, generically, a plethora of massless or light singlet fermions. Even if they do not have Dirac couplings and hence do not participate in a neutrino seesaw, this is a problem in its own right.

As with Yukawa couplings, Majorana masses can be generated by scalar VEVs, but one can also obtain Majorana masses in exact string theory. In the context of orientifold models of the Madrid type this can in principle be achieved as follows. In these models there is always a $B - L$ symmetry. Usually this symmetry is exact and leads to a massless gauge boson ([Dijkstra, Huiszoon, and Schellekens, 2005](#)). This is in disagreement with experiment, and since massless $B - L$ gauge bosons are ubiquitous in string theory, it is reasonable to ask why we do not see one in our Universe. The answer may be anthropic: $B - L$ gauge bosons lead to a repulsive force between protons and neutrons and may destabilize nuclei. There would also be drastic changes in atoms and chemistry. But let us take this for granted and consider the small set of cases where the $B - L$ symmetry is broken. In those cases a Majorana mass may be generated by nonperturbative effects due to D-brane instantons ([Argurio *et al.*, 2007](#); [Blumenhagen, Cvetič, and Weigand, 2007](#); [Cvetič, Richter, and Weigand, 2007](#); [Florea *et al.*, 2007](#); [Ibañez and Uranga, 2007](#)). This does indeed work, but in practice the relevant instanton contributions are nearly always destroyed by a surplus of zero modes ([Ibañez, Schellekens, and Uranga, 2007](#)). Even if one assumes that this is an artifact of special models, there is still another problem: instanton generated terms have logarithmically distributed scales. Since D-brane instantons have mass scales that are unrelated to those of the standard model gauge group, their scale is not linked to the standard model scale. But there is also no particular reason why it would be the large scale needed for small neutrino masses.

If a large number of singlet neutrinos is involved in the seesaw mechanism, as string theory suggests, this may have important benefits. It raises the upper limit for leptogenesis ([Eisele, 2008](#)) and also raises the seesaw scale ([Ellis and Lebedev, 2007](#)).

b. Anthropical arguments

Neutrinos are not constituents of matter, so that they do not have to obey “atomic” anthropic bounds. Nevertheless, they have a number of potential anthropic implications. In our

Universe, neutrinos play a role in big bang nucleosynthesis, structure formation, supernova explosions, stellar processes, the decay of the neutron, pions, and other particles, the mass density of the Universe, and possibly leptogenesis.

Many of these processes would change drastically if neutrino masses were in the typical range of charged leptons, but one should not jump to anthropic arguments too quickly. The fact that universes may exist where weak interactions, including neutrinos, are not even necessary ([Harnik, Kribs, and Perez, 2006](#)) underscores that point. But there are a few interesting limits nonetheless.

If the sum of all neutrino masses exceeds 40 eV, they would overclose the Universe. But there is no need to argue if this is an observational or an anthropic constraint, because for much larger masses (larger than the pion mass) they would all be unstable, invalidating any such argument. An interesting limit follows from leptogenesis ([Fukugita and Yanagida, 1986](#)), which sets an upper bound to neutrino masses of 0.1 eV ([Buchmüller, Di Bari, and Plumacher, 2003](#)). If this is the only available mechanism for generating a net baryon density, this would imply an anthropic upper bound on neutrino masses.

[Tegmark, Vilenkin, and Pogosian \(2005\)](#) gave a rationale for small neutrino masses based on galaxy formation. They argued that fewer galaxies are formed in universes with larger neutrino masses. If the distribution of neutrino masses does not favor very small values, this leads to an optimum at a finite value, which is about 1 eV (for $\sum m_\nu$). This is barely consistent with the aforementioned leptogenesis limit. Note that this mechanism favors Dirac masses. The seesaw mechanism with GUT-scale Majorana masses gives distributions that are too strongly peaked at zero.

c. Landscape distributions

In the neutrino sector one can still make predictions. Until recently this included the angle θ_{13} , which until 2012 was consistent with zero, an implausible value from the landscape perspective.

The other opportunities for prediction are the masses or at least their hierarchy. Generically, any model that gives the required large quark and lepton mass hierarchies will tend to produce hierarchies in the neutrino sector as well. Therefore it is not surprising that all work listed below prefers a normal hierarchy (the inverted hierarchy requires two relatively large, nearly degenerate masses).

The two large neutrino mixing angles are an obvious challenge for distributions that produce small quark mixing angles. But there are several ways in which neutrino masses could be different from quark and charged lepton masses. First of all, right-handed neutrinos might not belong to families the way quarks and leptons do. Second, there may be hundreds of them, not just three, and, third, the origin of their Majorana mass matrix is not likely to be related to that of the Higgs coupling.

[Donoghue, Dutta, and Ross \(2006\)](#) studied neutrino mixing angle distributions using Dirac couplings distributed like those of quarks and with three right-handed neutrinos. These were assumed to have a Majorana matrix with random matrix elements, with various distributions. They found that with these minimally biased assumptions the likelihood of

getting the observed mixing angles is only about 5% to 18%, with the latter value occurring for a small Majorana scale of about 10^7 GeV. They strongly predicted a normal hierarchy, a wide distribution of θ_{13} disfavoring the value zero, and a Majorana neutrino mass (as would be observed in neutrino-less double- β decay) of the order of 0.001 eV.

The approach studied by [Hall, Salem, and Watari \(2007, 2009\)](#), mentioned above for quarks, can accommodate neutrino mixing by assuming that wave functions of lepton doublets are less localized than those of quarks. The Majorana mass matrices are generated using overlap integrals of randomized Gaussian wave functions. This works, but is more biased toward the observed result.

Neutrino masses and mixings have also been studied in F theory ([Bouchard et al., 2010](#)). An interesting prediction is that the hierarchy is not just normal, but more concretely $m_1 : m_2 : m_3 \approx \alpha_{\text{GUT}} : \sqrt{\alpha_{\text{GUT}}} : 1$ with $\alpha_{\text{GUT}} \approx 0.04$. Using the two mass splittings this gives neutrino masses of approximately 2, 9, and 50 meV. The predicted value for θ_{13} is equal to $\sqrt{\alpha_{\text{GUT}}}$ and is compatible with the recently observed value.

C. The scales of the standard model

The classic standard model has two scales, the strong and the weak scale. To first approximation the strong scale Λ_{QCD} determines the proton mass, and the weak scale determines the masses of the quarks and leptons. The proton mass owes less than 1% of its mass to the up and down quarks. Indeed, the proton mass is nonvanishing in the limit of vanishing quark masses and is only a little bit smaller in that limit.

The weak scale and the strong scale have a rather different origin in the standard model. The former is directly related to the only dimensionful parameter in the Lagrangian, the parameter μ^2 , whereas the latter comes out as a pole in the running of the QCD coupling constant toward the IR region. This produces a dimensionful parameter Λ_{QCD} from a dimensionless one $\alpha_s = g_s^2/4\pi$. This is known as “dimensional transmutation.” At one-loop order, the logarithmic running of α_s determines Λ_{QCD} in the following way:

$$\alpha_s(Q^2) = \frac{1}{\beta_0 \ln(Q^2/\Lambda_{\text{QCD}}^2)}, \quad (5.11)$$

with $\beta_0 = (33 - 2N_f)/12\pi$, where N_f is the number of quark flavors $N_f = 6$. Here Q is the relevant energy scale. If we measure the function at one scale, it is determined at any other scale. One can invert this relation to obtain

$$\Lambda_{\text{QCD}} = Q e^{-1/[2\beta_0\alpha(Q^2)]}. \quad (5.12)$$

Note that Λ_{QCD} is a free parameter, which can be traded for $\alpha_s(Q^2)$ at some fixed scale, if desired.

Two things are remarkable about the weak and strong scales. Both are very much smaller than the Planck scale

$$M_{\text{Planck}} = \sqrt{\frac{\hbar c^5}{G_N}} = 1.2209 \times 10^{19} \text{ GeV}, \quad (5.13)$$

and they are within about 2 or 3 orders of magnitude from each other. The smallness of both scales is responsible for the extreme weakness of gravity in comparison to the other forces. This fact has important anthropic implications.

There are many ways of varying these scales while keeping other parameters fixed. Many papers on anthropic arguments in astrophysics, such as [Carr and Rees \(1979\)](#), study the effect of varying m_p/M_{Planck} . However, m_p is not a standard model parameter. It is mainly determined by Λ_{QCD} , but it is ultimately also affected by the weak scale. If we move up that scale by a few orders of magnitude while keeping the Yukawa couplings fixed, the quark masses rather than Λ_{QCD} dominate the proton mass. Many other things change as well, making it hard to arrive at a clean conclusion. If we enlarge the proton mass by enlarging Λ_{QCD} , it is not just the proton mass that changes, but also the strength of the strong coupling.

1. Changing the overall scale

The cleanest way of studying the effect of varying the QCD scale is to vary all standard model scales by the same factor L with respect to M_{Planck} . This keeps all nuclear physics and chemistry unchanged, except for the overall scale. No thresholds are crossed, and every allowed process remains allowed in rescaled universes. Hence the chemistry of life is unaffected.

It is not hard to establish the existence of an anthropic bound. Basic kinematics implies a maximum for the number of nucleons in objects with gravitation balanced by internal pressure. This maximum is $\approx (M_{\text{Planck}}/m_p)^3$ and determines the maximum number of nucleons in stars to within a factor of the order of 10 ([Carr and Rees, 1979](#)). If we increase m_p , we reach a point where the maximum is smaller than the number of nucleons in a human brain, which means that brain-sized objects collapse into black holes. If we set the necessary number of nucleons in a brain conservatively at about 10^{24} , we find a limit of $m_p \ll 10^{-8} M_{\text{Planck}}$.

These objects are just clusters of nucleons, not necessarily hot enough to have nuclear fusion. It is probably not too anthropocentric to assume that stars should ignite, not just to have stars as sources of energy but even more importantly as processing plants of elements heavier than lithium. Conditions for existence of stars in other universes were investigated by [Adams \(2008\)](#). The result is that the combined standard model scale cannot be enlarged by more than about a factor of 10 without losing nuclear fusion in stars.⁶

Variation of all standard model mass scales with respect to the Planck mass was studied by [Graesser and Salem \(2007\)](#). They considered the effect of changing the Planck mass on several cosmological processes, such as inflation, baryogenesis, big bang nucleosynthesis, structure formation, and stellar dynamics, and found that the anthropic window on the scale is narrow (less than an order of magnitude in either direction), if other cosmological parameters are kept fixed.

Therefore the smallness of the ratio m_p/M_{Planck} (in the sense of a variation of the overall scale of the standard model) is undoubtedly needed anthropically. The true distribution of the scale depends ultimately on the landscape distributions at the string scale. The fact that the strong scale seems

⁶Note that [Adams \(2008\)](#) allows variations of nuclear reaction rates beyond QCD and hence finds a larger allowed variation. Tracing the scale dependence in the computation leads to a much smaller effect.

distributed logarithmically because of dimensional transmutation [i.e., Eq. (5.12)] is not in dissonance with anthropic reasoning, which requires logarithmic tuning only to the right order of magnitude. It is harder to establish a lower bound on the overall scale, but big changes do occur if it is lowered, since astrophysical sizes, times, and temperatures scale differently than biological ones; see, for example, the discussion of the Carter conjecture in Sec. V.A.4.

2. The weak scale

The smallness of the weak scale, also known as the gauge hierarchy problem, is not just a matter of very small ratios, but now there is also a fine-tuning problem. The small parameter μ^2 gets contributions from quantum corrections or rearrangements of scalar potentials that are proportional to M^2 , where M is the relevant large scale. Hence it looks like these terms must be tuned to 30 significant digits so that they add up to the very small μ^2 we observe.

a. Anthropic bounds on the weak scale

The idea that the weak scale might be anthropically determined was suggested for the first time by [Agrawal et al. \(1998a\)](#). They considered anthropic bounds on the weak scale following from changes in quark masses, keeping the Yukawa couplings fixed, as discussed in Sec. V.B. But what happens if we allow the Yukawa couplings to vary as well?

[Donoghue et al. \(2010\)](#) computed a likelihood function for the Higgs VEV using a scale-invariant distribution function of the Yukawa couplings, determined from the observed distribution of quark masses. Using this distribution, and a flat distribution in v , both the Higgs VEV and the Yukawa couplings are allowed to vary, under the assumption that the Yukawa distribution does not depend on v . The conclusion is that values close to the observed VEV are favored.

However, [Gedalia, Jenkins, and Perez \(2011\)](#) made different assumptions. They also considered, among others, scale-invariant distributions. But scale-invariant distributions require a cutoff to be normalizable. If one assumes that values as small as $\lambda_y = 10^{-21}$ have a similar likelihood as values of the order of 1, then it is statistically easier to get three small masses (for the u and d quarks and for the electron) using small Yukawa couplings and a large Higgs VEV than the way it is done in our Universe. If, furthermore, one assumes a weakless universe as discussed by [Harnik, Kribs, and Perez \(2006\)](#), the conclusion would be that in the multiverse there are far more universes without than with weak interactions, given atomic and nuclear physics as observed. See, however, [Giudice, Perez, and Soreq \(2012\)](#) for a way of avoiding the runaway to small Yukawas and large Higgs VEVs.

If indeed in the string landscape extremely small values of Yukawa couplings are not strongly suppressed, and if weakless universes are as habitable as ours [which is not as obvious as [Gedalia, Jenkins, and Perez \(2011\)](#) claim], this provides one of the most convincing arguments in favor of a solution to the hierarchy problem: a mechanism that tilts the distribution of μ^2 toward smaller values.

b. Low-energy supersymmetry

The fact that a logarithmic behavior works for the strong scale has led to speculation that a similar phenomenon should

be expected for the weak scale. At first sight the most straightforward solution is to postulate an additional interaction that mimics QCD and generates a scale by dimensional transmutation. The earliest idea along these lines is known as “technicolor.” Another possibility is that large extra dimensions exist, lowering the higher-dimensional Planck scale to the TeV region. But the most popular idea is low-energy SUSY. The spectacular results from the Large Hadron Collider (LHC) experiments have put all these ideas under severe stress, but low-energy SUSY remains a viable possibility. For this reason this is the only option that we consider more closely here.

Low-energy SUSY does not directly explain the smallness of the Higgs parameter μ^2 , but rather the “technical naturalness” problem. In the standard model, the quantum corrections to μ^2 are quadratically sensitive to high scales. In the supersymmetric standard model, every loop contribution is canceled by a loop of a hypothetical particle with the same gauge quantum numbers, but with spin differing by half a unit, and hence opposite statistics: squarks, sleptons, and gauginos. None of these additional particles has been seen so far. Supersymmetry is at best an exact symmetry at high energies.

Rather than a single dimensionful parameter μ^2 the supersymmetrized standard model has at least two, a parameter which, somewhat confusingly, is traditionally called μ , and a scale M_S corresponding to SUSY breaking. The latter scale may be generated by dimensional transmutation, and this is the basis for SUSY as a solution to the hierarchy problem. But the additional scale μ , which can be thought of as a supersymmetric Higgs mass prior to weak symmetry breaking, requires a bit more discussion. To prevent confusion we equip the supersymmetric μ parameter with a hat.

Since μ^2 , just as $\hat{\mu}$, is merely a parameter that can take any value, it may seem that nothing has been gained. The difference lies in the quantum corrections these parameters get. For the μ^2 parameter these quantum corrections take the (simplified) form

$$\mu_{\text{phys}}^2 = \mu_{\text{bare}}^2 + \sum \alpha_i \Lambda^2 + \text{logarithms}, \quad (5.14)$$

whereas for $\hat{\mu}$ one finds

$$\hat{\mu}_{\text{phys}} = \hat{\mu}_{\text{bare}} \left(1 + \sum \beta_i \log(\Lambda/Q) + \dots \right). \quad (5.15)$$

Here “bare” denotes the parameter appearing in the Lagrangian and “phys” the observable, physical parameter, defined and measured at some energy scale Q ; Λ denotes some large scale at which the momentum integrals are cut off.

The difference between these two kinds of quantum corrections is most easily understood if one thinks of them in terms of distributions, i.e., a landscape. Indeed, the concept of naturalness, especially in the technical sense, implicitly assumes a landscape, a point also emphasized by [Hall and Nomura \(2008\)](#). If one adopts the landscape paradigm, the rationale for a natural solution of the hierarchy problem would be that the unnatural solution comes at a high statistical price $\mu^2/M_{\text{Planck}}^2 \approx 10^{-35}$. This holds for the standard model with a flat distribution of values of μ^2 between 0 and M_{Planck}^2 as suggested by the renormalization of μ^2 . On the

other hand, the renormalization of $\hat{\mu}$, proportional to $\hat{\mu}$ itself, gives no information about its distribution.

c. The supersymmetry breaking scale

Low-energy SUSY lowers the statistical price by replacing M_{Planck} by M_{SUSY} , the SUSY breaking scale. Here we define it as the typical scale of super multiplet mass splittings.⁷ This suggests that the statistical price for a small weak scale can be minimized by setting $M_{\text{SUSY}} \approx \mu$. This is the basis for two decades of predictions of light squarks, sleptons, and gauginos, which, despite being much more sophisticated than this, have led to two decades of wrong expectations. But in a landscape, the likelihood $P(\mu)$ for a weak scale μ is given by

$$P(\mu) = P_{\text{nat}}(\mu, M_{\text{SUSY}}) P_{\text{landscape}}(M_{\text{SUSY}}). \quad (5.16)$$

The first factor is the naive naturalness contribution $P_{\text{nat}}(\mu, M_{\text{SUSY}}) \propto \mu^2/M_{\text{SUSY}}^2$, and the second one is the fraction of vacua with a SUSY breaking scale M_{SUSY} .

During the last decade there have been several attempts to determine $P_{\text{landscape}}(M_{\text{SUSY}})$. One such argument, suggested by Douglas (2004b) and Susskind (2004), suggested that it increases with a power given by the number of SUSY breaking parameters (F and D terms). If true, that would rather easily overcome the $(M_{\text{SUSY}})^{-2}$ dependence of the first factor. However, this assumes that all these sources of SUSY breaking are independent, which is not necessarily correct (Denef and Douglas, 2005). Other arguments depend on the way SUSY is broken [called “branches” of the landscape by Dine, O’Neil, and Sun (2005)]. The arguments are presented in detail in Douglas and Kachru (2007). An important contributing factor that was underestimated in earlier work is the fact that vacua with broken SUSY are less likely to be stable. This can lead to a substantial suppression (Chen *et al.*, 2012; Marsh, McAllister, and Wrase, 2012). There are large factors going in both directions, but the net result is uncertain at present.

One might expect intuitively that there should be another suppression factor $\Lambda^4/M_{\text{SUSY}}^4$ in Eq. (5.16) due to the fact that unbroken SUSY can help fine-tuning the cosmological constant Λ just as it can help fine-tuning μ (Banks, Dine, and Gorbatov, 2004; Susskind, 2004). But this is wrong, basically because it is not true that $\Lambda = 0$ in supergravity. In general one gets $\Lambda \leq 0$, which must be canceled to 120 digit precision just as in the nonsupersymmetric theories. There is a branch with $\Lambda = 0$ before SUSY breaking, but this requires a large (R) symmetry, which is statistically unlikely (Dine and Sun, 2006).

Despite the inconclusive outcome there is an important lesson in all this. Conventional bottom-up naturalness arguments that make no mention of a landscape are blind to all these subtleties. If these arguments fail in the only landscape that can be discussed at present, they should be viewed with suspicion. Even if in the final analysis all uncertain factors

conspire to favor low-energy SUSY in the string theory landscape, the naive naturalness arguments would have been correct only by pure luck.

d. Moduli

There is another potentially crucial feature of string theory that conventional low-energy SUSY arguments are missing: moduli (including axions). This point was made especially forcefully by Acharya, Kane, and Kumar (2012) and earlier work cited therein.

It has been known for a long time that moduli can lead to cosmological problems (Coughlan *et al.*, 1983; de Carlos *et al.*, 1993; Banks, Kaplan, and Nelson, 1994). If they are stable or long lived, they can overclose the Universe; if they decay during or after BBN, they will produce additional baryonic matter and destroy the successful BBN predictions. For fermionic components of moduli multiplets these problems may sometimes be solved by dilution due to inflation. But bosonic moduli have potentials and will in general be displaced from their minima. Their time evolution is governed by

$$\ddot{\phi} + 3H\dot{\phi} + \frac{\partial V}{\partial \phi} = 0, \quad (5.17)$$

where H is the Hubble constant. If $V = \frac{1}{2}m^2\phi^2 + \text{higher order}$ and $H \gg m$ then the second term dominates over the third, and ϕ gets frozen at some constant value (“Hubble friction”). This lasts until H drops below m . Then the field starts oscillating in its potential and releases its energy. The requirement that this does not alter BBN predictions leads to a lower bound on the scalar moduli mass of a few tens of TeV (30 TeV, for definiteness).

Furthermore, one can argue (Acharya, Kane, and Kuflik, 2010) that the mass of the lightest modulus is of the same order of magnitude as the gravitino mass $m_{3/2}$. The latter mass is generically of the same order as the soft SUSY breaking scalar masses: the squarks and sleptons searched for at the LHC. This chain of arguments leads to the prediction that the sparticle masses will be a few tens of TeV, out of reach for the LHC, probably even after its upgrade. But there was also a successful (although fairly late and rather broad) prediction of the Higgs mass⁸ (Kane *et al.*, 2012).

However, there are loopholes in each step of the chain. Light moduli can be diluted by “thermal inflation” (Lyth and Stewart, 1996), and the mass relation between gravitinos and sparticles can be evaded in certain string theories. The actual result of Acharya, Kane, and Kuflik (2010) is that the lightest modulus has a mass smaller than $m_{3/2}$ times a factor of the order of 1, which can be large in certain cases. Hence this scenario may be generic, but is certainly not general.

The relation between $m_{3/2}$ and fermionic superparticles (Higgsinos and gauginos) is less strict and more model

⁷At least two distinct definitions of the SUSY breaking scale are used in the literature. Furthermore, several mechanisms exist for “mediation” of SUSY breaking, such as gauge and gravity mediation. The discussion here is only qualitative and does not depend on this. See Douglas and Kachru (2007) for further details.

⁸The Higgs mass ≈ 126 GeV was also correctly predicted in finite unified theories; see Heinemeyer, Mondragon, and Zoupanos (2008) and on the basis of asymptotically safe gravity, see Shaposhnikov and Wetterich (2010). Bottom-up supersymmetric models, ignoring moduli, suggested an upper limit of at most 120 GeV.

dependent. They might be lighter than $m_{3/2}$ by 1 to 2 orders of magnitude and accessible at the LHC. Gaugino mass suppression in fluxless M-theory compactifications has been discussed by [Acharya *et al.* \(2007\)](#). This was also seen in type-IIB compactifications, with typical suppression factors of the order of $\log(M_{\text{Planck}}/m_{3/2})$ ([Choi *et al.*, 2004](#); [Conlon and Quevedo, 2006](#); [Choi and Nilles, 2007](#)).

A SUSY scale of 30 TeV introduces an unnatural fine-tuning of 5 orders of magnitude,⁹ the “little hierarchy.” This tuning requires an explanation beyond the mere phenomenological necessity. The explanation could be anthropic, which would be much better than observational. A universe that seems fine-tuned for our existence makes a lot more sense than a universe that seems fine-tuned just to misguide us.

Could this explain the 30 TeV scale? Statements such as “the results of BBN are altered” or “the Universe is over-closed” if moduli are lighter do indeed sound potentially anthropic. But it is not that simple. Constraints from BBN are mostly just observational, unless one can argue that all hydrogen would burn to helium. Otherwise, what BBN can do, stars can do better. Overclosure just means disagreement with current cosmological data. Observers in universes just like ours in all other respects might observe that they live in a closed universe with $\Omega \gg 1$, implying recollapse in the future. But the future is not anthropically constrained. The correct way to compare universes with light moduli anthropically to ours is to adjust the Hubble scale so that after inflation $\Omega \approx 1$. This gives a universe with different ratios of matter densities, but it is not at all obvious that those ratios would be catastrophic for life. Without such an argument, the claim that moduli require a 30 TeV SUSY scale is much less convincing. See also [Giudice and Rattazzi \(2006\)](#) for a different view on a possible anthropic origin of the little hierarchy.

e. The cost of SUSY

Another anthropically relevant implication of low-energy SUSY is stability of baryons. Supersymmetry allows “dimension-4” operators that violate baryon number and lepton number that do not exist in the standard model: they are group theoretically allowed, but contain an odd number of fermions. If all these operators are present with $\mathcal{O}(1)$ coefficients, they give rise to anthropically disastrous proton decay. This can be solved by postulating a discrete symmetry that forbids the dangerous couplings [most commonly R parity, but there are other options; see [Berasaluce-Gonzalez *et al.* \(2011\)](#) for a systematic summary]. In the landscape global symmetries are disfavored, but R parity may be an exception ([Dine and Sun, 2006](#)). Landscape studies of intersection brane models indicate that they rarely occur ([Ibañez, Schellekens, and Uranga, 2012](#); [Anastasopoulos *et al.*, 2013](#)), but since they are anthropically required one can tolerate a large statistical price.

But apart from anthropically required tunings, SUSY is also observationally fine-tuned. There are dimension-5 operators that can give rise to observable but not catastrophic

proton decay. A generic supersymmetric extension of the standard model gives rise to large violations of flavor symmetry: for a general soft mass term, the diagonalization of squark matrices requires unitary rotations that are not related to those of the quarks. There are also substantial contributions to CP -violating processes. All of these problems can be solved, but at a statistical price that is hard to estimate and hard to justify. Moving the SUSY breaking scale to 30 TeV ameliorates some of these problems, but does not remove them.

Since SUSY has failed to fully solve the hierarchy problem, we must critically examine the other arguments supporting it. The so-called “WIMP miracle,” the claim that stable superpartners precisely give the required amount of dark matter, has been substantially watered down in recent years. On closer inspection, it is off by a few orders of magnitude ([Arkani-Hamed, Delgado, and Giudice, 2006](#)), and a “non-thermal” WIMP miracle was suggested ([Acharya *et al.*, 2009](#)) in its place. Although this is based on WIMPs produced in out of equilibrium decays of moduli, and fits nicely with string theory, two miracles are one too many. Axions are a credible dark matter candidate, and several have suggested scenarios where both kinds of dark matter are present ([Tegmark *et al.*, 2006](#); [Acharya, Kane, and Kumar, 2012](#)). But then we could also do without WIMPs altogether. Furthermore, dark matter is constrained anthropically. Although crude arguments based on the structure formation of [Hellerman and Walcher \(2005\)](#) still allow a rather large window of 5 orders of magnitude, this is not much larger than the uncertainty of the WIMP miracle. It is far from obvious that life would flourish equally well in dense dark matter environments so that the true anthropic bound might be much tighter. The other main argument, gauge coupling unification, has already been discussed in Sec. V.A.3.d. It is more seriously affected by problems at the string scale than by the upward motion of the SUSY scale, on which it depends only logarithmically.

Ideas such as *split* supersymmetry (a higher mass scale just for the superpartners of fermions) and *high scale* supersymmetry (a larger SUSY scale) are becoming more and more *salonfähig* in recent years. Perhaps counterintuitively, their scales are constrained from above by the Higgs mass measurement ([Giudice and Strumia, 2012](#)): in supersymmetric theories the Higgs self-coupling cannot become negative, as it appears to be doing. It is hard to avoid the idea that the most natural scenario is *no* supersymmetry. But that would also imply that everything we think we know about the landscape is built on quicksand. This is a big dilemma that we will hear a lot more about in the future.

D. Axions

Unlike the large gauge hierarchy, the extreme smallness of the strong CP -violating angle $\bar{\theta}$ has few anthropic implications. Apart from producing as yet unobserved nuclear dipole moments, $\bar{\theta}$ can have substantial effects on nuclear physics, including anthropically relevant features like deuteron binding energies and the triple- α process. [Ubbaldi \(2010\)](#) found the reaction rate of the triple- α process to be 10 times larger if $\bar{\theta} = 0.035$. But at best this would explain 2 to 3 of the observed 10 orders of magnitude of fine-tuning.

⁹In comparison with a weak scale of ≈ 100 GeV and expressed in terms of the square of the scale, in accordance with the scale dependence of quantum corrections.

There are several possible solutions, but one stands out because of its simplicity: the mechanism discovered by [Peccei and Quinn \(1977\)](#). It requires nothing more than adding a scalar a and a nonrenormalizable coupling:

$$\Delta\mathcal{L} = \frac{1}{2} \partial_\mu a \partial^\mu a + \frac{a}{32\pi^2 f_a} \sum_a F_{\mu\nu}^a F_{\rho\sigma}^a \epsilon^{\mu\nu\rho\sigma}, \quad (5.18)$$

where f_a is the “axion decay constant.” Since $F\tilde{F}$ (where $\tilde{F}_{\mu\nu} = \frac{1}{2} \epsilon_{\mu\nu\rho\sigma} F^{\rho\sigma}$) is a total derivative, after integration by parts the second term is proportional to $\partial_\mu a$. Hence there is a shift symmetry $a \rightarrow a + \epsilon$. This allows us to shift a by a constant $-\bar{\theta}f_a$ so that the $F\tilde{F}$ term (2.5) is removed from the action. However, the shift symmetry is anomalous with respect to QCD because the $F\tilde{F}$ term is a derivative of a gauge noninvariant operator. Through nonperturbative effects the anomaly generates a potential with a minimum at $a = 0$ of the form

$$V(a) \propto \Lambda_{\text{QCD}}^4 [1 - \cos(a/f_a)]. \quad (5.19)$$

Note that $\bar{\theta}$ is periodic with period 2π , so that the shift symmetry is globally a $U(1)$ symmetry. It was pointed out by [Weinberg \(1978\)](#) and [Wilczek \(1978\)](#) that this breaking of the $U(1)$ symmetry leads to a pseudoscalar pseudo-Goldstone boson, called an “axion.” The mass of this particle is roughly $\Lambda_{\text{QCD}}^2/f_a$, but if we take into account the proportionality factors in Eq. (5.19) the correct answer is

$$m_a = \frac{m_\pi f_\pi}{f_a} F(m_q), \quad (5.20)$$

where f_π is the pion decay constant and $F(m_q)$ is a function of the (light) quark masses that is proportional to their product. The scale f_a was originally assumed to be that of the weak interactions, leading to a mass prediction of the order of 100 keV that is now ruled out. But soon it was realized that f_a could be chosen freely, and, in particular, much higher, making the axion “harmless” or “invisible” [see [Kim \(1987\)](#) and references therein]. This works if the coupling f_a is within a narrow window. For small f_a the constraint is due to the fact that supernovae or white dwarfs would cool too fast by axion emission. This gives a lower limit $f_a > 10^9$ GeV.

The upper limit is cosmological. In the early Universe the axion field would be in a random point θ_0 in the range $[0, 2\pi]$ (“vacuum misalignment”). The potential (5.19) is irrelevant at these energy scales. During the expansion and cooling of the Universe, the field remains at that value until the Hubble scale drops below the axion mass. Then the field starts oscillating in its potential, releasing the stored energy, and contributing to dark matter densities. The oscillating axion field can be described as a Bose-Einstein condensate of axions. Despite the small axion mass, this is cold dark matter: the axions were not thermally produced. Axions may in fact be the ideal dark matter candidate ([Sikivie, 2012](#)).

The axion contribution to dark matter density is given by

$$\Omega_a \propto (f_a)^{1.18} \sin^2(\frac{1}{2}\theta_0) \quad (5.21)$$

[see [Bae, Huh, and Kim \(2008\)](#) for a recent update and earlier references]. The requirement that this does not exceed the observed dark matter density leads to a limit $f_a < 10^{12}$ GeV, unless $\theta_0 \approx 0$. This results in a small allowed window for the

axion mass $6 \mu\text{eV} < m_a < 6 \text{ meV}$. Observing such a particle is hard, but one can use the fact that axions couple (in a model-dependent way) to two photons. Several attempts are underway, but so far without positive results. The location of the axion window is fascinating. It is well below the GUT and Planck scales, but roughly in the range of heavy Majorana masses in seesaw models for neutrinos. It is also close to the point where the extrapolated Higgs self-coupling changes sign, although there are large uncertainties.

There are many string-theoretic, landscape, and anthropic issues related to axions. Candidate axions occur abundantly in string theory [see [Svrcek and Witten \(2006\)](#) for details and earlier references].

But exact global symmetries, like axion shift symmetries, are not supposed to exist in theories of quantum gravity, and hence they are not expected to exist in string theory. Therefore one expects all the candidate axions to acquire a mass. The PQ mechanism can work only if a light axion survives with couplings to QCD, and with a mass contribution from other sources that is much smaller than the QCD-generated mass.

Axions are imaginary parts of moduli, which must be stabilized, and they must somehow escape getting a mass from the stabilization. They must also survive orientifold projections and not be eaten by vector bosons in a Stueckelberg mechanism. However, in most string theories candidate axions exist that are exactly massless to all orders in perturbation theory, and which must therefore get their masses from nonperturbative effects. These effects can be expected to give rise to axion masses proportional to e^{-S} , where S is an instanton action.

It is not likely that a light axion exists just for QCD. From the string theory perspective, it seems strange that out of the large number of candidate axions just one survives. From the gauge theory perspective, many different gauge groups with many different non-Abelian factors are possible. Either they generically come with axions or QCD is a special case for no apparent reason.

This has led to the notion of an “axiverse” ([Arvanitaki et al., 2010](#)), a plethora of axions, with masses spread logarithmically over all scales; only the mass of the QCD axion is determined by Eq. (5.20). Realizations of an axiverse have been discussed in fluxless M-theory compactifications ([Acharya, Bobkov, and Kumar, 2010](#)) and in type-IIB models in the large volume scenario ([Cicoli et al., 2012](#)). Both papers consider compactifications with many Kähler moduli that are stabilized by a single nonperturbative contribution rather than a separate contribution for each modulus. Then all Kähler moduli can be stabilized, but just one “common phase” axion acquires a large mass. All remaining ones get small masses from other instantons. For supersymmetric moduli stabilization (such as the KKLT scenario, but unlike LVS) a no-go theorem was proved by [Conlon \(2006\)](#), pointing out that for each massless axion there would be a tachyonic saxion after uplifting. But [Choi and Jeong \(2007\)](#) considered a generalization of the KKLT scenario where this problem is avoided. Axions in the heterotic minilandscape were discussed by [Choi et al. \(2009\)](#). They considered discrete symmetries that restrict the superpotential, so that the lowest order terms have accidental $U(1)$ symmetries that may include a PQ symmetry.

The upper limit $f_a < 10^{12}$ GeV is problematic for axions in string theory, which generically prefers a higher scale (Svrcek and Witten, 2006). A way out of this dilemma is to assume that the misalignment angle in Eq. (5.21) is small. This is an option if the PQ phase transition occurred before inflation, so that we just observe a single domain of a multi-domain configuration with a distribution of values of θ_0 . If the phase transition occurred after inflation, we would instead observe an average of $\sin^2 \theta_0$, equal to $1/2$. To allow an increase of f_a to the GUT or string scale of about 10^{16} GeV a value of $\theta_0 \approx 10^{-3}$ is sufficient. One could even assume that this value came out “by accident,” which is still a much smaller accident than required for the strong CP problem. However, the fact that the upper limit on f_a is due to the axion’s contribution to dark matter has led to the suggestion that we live in an inflated domain with small θ_0 not by accident, but for anthropic reasons (Linde, 1991). Furthermore, the fact that this parameter is an angle and that axions are not strongly coupled to the rest of the landscape makes it an ideal arena for anthropic reasoning (Wilczek, 2004). This was explored in detail by Tegmark *et al.* (2006) and Freivogel (2010). The upper bound on the axion decay constant can be raised if there is a nonthermal cosmological history, for example, caused by decay of ≈ 30 TeV moduli (Acharya, Kane, and Kumar, 2012).

Whatever solution is proposed for the strong CP problem, it should not introduce a fine-tuning problem that is worse. Therefore models specifically constructed and tuned to have a QCD axion in the allowed window, but which are rare within their general class, are suspect. This appears to be the case in all models suggested so far. The “rigid ample divisors” needed in the M theory and type-II constructions mentioned above are not generic, and the discrete symmetries invoked in heterotic constructions may be a consequence of the underlying mathematical simplicity of the orbifold construction. But it is difficult to estimate the amount of fine-tuning that really goes into these models.

The anthropic tuning required to avoid the upper bound on f_a was discussed by Mack (2011). Mack concluded that avoiding constraints from isocurvature fluctuations in the CMB, which are observational and not anthropic, requires tuning of both θ_0 and the inflationary Hubble scale to small values. The amount of tuning is more than the 10 orders of magnitude needed to solve the strong CP problem. This problem increases exponentially if there are many axions (Mack and Steinhardt, 2011).

There are numerous possibilities for experiments and observations that may shed light on the role of axions in our Universe and thereby provide information on the string theory landscape. The observation of tensor modes in the CMB might falsify the axiverse (Fox, Pierce, and Thomas, 2004; Acharya, Bobkov, and Kumar, 2010). See Arvanitaki *et al.* (2010), Marsh *et al.* (2012), and Ringwald (2012) for a variety of possible signatures, ongoing experiments, and references.

E. Variations in constants of nature

If we assume that constants of nature can take different values in different universes, it is natural to ask if they might also take different values within our own Universe. In the

standard model the parameters are fixed (with a computable energy scale dependence) and cannot take different values at different locations or times without violating the postulate of translation invariance.

There is a lot of theoretical and observational interest in variations of constants of nature and for good reason. The observation of such a variation would have a big impact on current ideas in particle physics and cosmology. See Langacker, Segre, and Strassler (2002) for a concise review and Uzan (2003) for a more extensive one, and Chiba (2011) for an update on recent bounds and observations. The results are most often presented in terms of variations in α or the electron-to-proton mass ratio $\mu = m_e/m_p$. The best current limits on $\Delta\alpha/\alpha$ are about 10^{-17} per year, from atomic clocks and from the Oklo natural nuclear reactor. Recently a limit $\Delta\mu/\mu < 10^{-7}$ was found by comparing transitions in methanol in the early universe (about 7×10^9 years ago) with those on Earth at present (Bagdonaite *et al.*, 2013).

But in addition to limits there have also been positive observations. Using the Keck observatory in Hawaii and the Very Large Telescope (VLT) in Chili, Webb *et al.* (2011) reported a *spatial* variation of α . Earlier observations at Keck of a smaller value of α , at that time interpreted as a temporal variation (Webb *et al.*, 2001), combined with more recent VLT observations of a larger value, fit a dipole distribution in the sky. These results have a statistical significance of $(4-5)\sigma$. Because these results imply a spatial and not a *temporal* variation, a clash with other, negative, results is avoided.

There are no good theoretical ideas for the expected size of a variation, if any. In string theory, and quite generally in theories with extra dimensions, the couplings are functions of scalar fields and are determined by the vacuum expectation value of those fields, subject to equations of motion of the form (5.17). This makes it possible to maintain full Poincaré invariance and relate the variations to changes in the vacuum. For example, the action for electrodynamics takes the form

$$\mathcal{L} = -\frac{1}{4e^2} e^{-\phi/M_{\text{Planck}}} F_{\mu\nu} F^{\mu\nu}, \quad (5.22)$$

where ϕ is the dilaton field or one of the other moduli. Variations in ϕ lead to variations in α :

$$\Delta\alpha \propto \frac{\delta\phi}{M_{\text{Planck}}}. \quad (5.23)$$

All other parameters of the standard model have a dependence on scalar fields as well. Although this formalism allows variations in α , it is clearly a challenge to explain why they would be as small as 10^{-15} per year. Note that this is about 10^{-66} in Planck units, the natural units of a fundamental theory like string theory.

The observation of a variation in any standard model parameter implies a big fine-tuning problem, with little hope of an anthropic explanation: variations of fundamental parameters might have adverse effects on the evolution of life, but there is no reason why the variation has to be as small as it is. Then the most attractive way out is that within our Universe these parameters really are constants, although they must vary in the multiverse. The string theory landscape solves this problem in an elegant way, because each of its vacua is at the bottom of a deep potential, completely

suppressing any possible variations of the moduli at sub-Planckian energies.

This can be seen by considering the effect of changes in VEVs of moduli fields on vacuum energy. Here one encounters the problem that contributions to vacuum energy in quantum field theory are quartically divergent. But this cannot be a valid reason to ignore them completely as is often done in the literature on variations of constants of nature. Banks, Dine, and Douglas (2002) pointed out that if a cutoff Λ_{cutoff} is introduced in quantum field theory, then the effect of a change in α on vacuum energy V is

$$\delta V \propto \Delta\alpha (\Lambda_{\text{cutoff}})^4. \quad (5.24)$$

With $\Lambda_{\text{cutoff}} = 100$ MeV, the QCD scale, and assuming that vacuum energy should not dominate at the earliest stages of galaxy formation (corresponding to the time when quasar light was emitted), this gives a bound of $\Delta\alpha/\alpha < 10^{-37}$. If one assumes that δV depends on $\Delta\alpha$ with a power higher than 1, this bound can be reduced, but a power of at least 8 is required to accommodate the observed variation. This can be achieved only by a correspondingly extreme tuning of the scalar potential. Spatial variations are restricted by similar arguments, although less severely.

There are also constraints from “fifth forces” violating the equivalence principle. This is a general problem associated with variations in constants of nature as observed a long time ago by Dicke (1957). For a recent discussion see Damour and Donoghue (2011).

Currently the observation of variations in constants of nature is still controversial, but there is a lot at stake. Evidence for variations would be good news for half of this review, and bad news for the other half. If the parameters of the standard model already vary within our own Universe, the idea that they are constants can be put into the dust bin of history, where it would be joined almost certainly by the string theory landscape. String theory would be set back by about two decades, to the time where it was clear that there were many “solutions,” without any interpretation as vacua with a small cosmological constant.

VI. ETERNAL INFLATION

If string theory provides a large landscape with a large number of vacua, how did we end up in one particular one? The answer is eternal inflation, a nearly inevitable implication of most theories of inflation. See Guth (2000), Linde (2002), and Freivogel (2011) for more discussion and references. If there is a possibility for transitions to other universes, then this would inevitably trigger an eternal process of creation of new universes.

For different views on eternal inflation and on populating the landscape see, respectively, Hawking and Hertog (2006) and Mersini-Houghton and Perry (2012).

A. Tunneling

Vacuum decay can take place in various ways. The best-known process were described by Coleman and De Luccia (1980) and Hawking and Moss (1982). The former describes tunneling between false vacua, and the latter tunneling of a false vacuum to the top of the potential. These processes

generate the nucleation of bubbles of other vacua which expand, and then themselves spawn bubbles of still more vacua (Lee and Weinberg, 1987). Tunneling between dS vacua can occur in both directions, up and down in vacuum energy, although up-tunneling is strongly suppressed with respect to down-tunneling [see, e.g., Schwartz-Perlov and Vilenkin, 2006]

$$\Gamma_{i \rightarrow j} = \Gamma_{j \rightarrow i} \exp\left(24\pi^2 \left[\frac{1}{\Lambda_j} - \frac{1}{\Lambda_i}\right]\right). \quad (6.1)$$

The end point of tunneling may be another dS vacuum, but it may also be a Minkowski or AdS vacuum. Whether tunneling from Minkowski to AdS is possible has been disputed by Dvali (2011) and Garriga, Shlaer, and Vilenkin (2011). Minkowski vacua do not inflate, and AdS universes collapse classically in a finite amount of time. Up-tunneling from these vacua to dS space is impossible, and therefore they are called terminal vacua. They are “sinks in the probability flow” (Ceresole *et al.*, 2006; Linde, 2007). According to Bousso (2012) and Susskind (2012) their existence in the landscape may be essential for understanding the arrow of time and for avoiding the Boltzmann brain problem (see below). Even though a large portion of an eternally expanding universe ends up in a terminal vacuum, the rest continues expanding forever. A typical observer is expected to have a long period of eternal inflation in its past (Freivogel, 2011).

B. The measure problem

The word “eternal” suggests an infinity, and this is indeed a serious point of concern. As stated in many papers: “In an eternally inflating universe, anything that can happen will happen; in fact, it will happen an infinite number of times.” This, in a nutshell, is the measure problem [see Vilenkin (2006b), Guth (2007), Freivogel (2011), and Nomura (2012)]. If we want to compute the relative probability for events A and B, one may try to define it by counting the number of occurrences of A and those of B, and taking the ratio. But both numbers are infinite. It is not that hard to think of definitions that cut off the infinities, but many of them make disastrous predictions. For example, they may predict that observers—even entire solar systems with biological evolution—created by thermal or quantum fluctuations (“Boltzmann brains”) vastly outnumber ones like ourselves, with a cosmological history that can be traced back in a sensible way. Or they may predict that universes just a second younger than ours are far more numerous (the “Youngness paradox”). If these predictions go wrong, they go wrong by double exponentials, and a formalism that gives this kind of a prediction cannot be trusted for any prediction.

1. The dominant vacuum

An ingredient that could very well be missing is a theory for the initial conditions of the multiverse. It would be unduly pessimistic to assume that this is a separate ingredient that cannot be deduced from string theory (or whatever the theory of quantum gravity turns out to be). If it cannot be deduced by logical deduction, it might be impossible to get a handle on it.

But eternal inflation may make this entire discussion unnecessary, provided all vacua are connected by physical

processes. In that case, successive tunneling events may drive all of them to the same “attractor,” the longest lived dS vacuum whose occupation numbers dominate the late time distribution. This is called the “dominant vacuum” (Garriga and Vilenkin, 1998; Garriga *et al.*, 2006; Schwartz-Perlov and Vilenkin, 2006). Since tunneling rates are exponentially suppressed, this vacuum may dominate by a large factor. Then the overwhelming majority of vacua would have this attractor vacuum in its history. This would erase all memory of the initial conditions. Furthermore, Brown and Dahlen (2011) argued that despite some potential problems—vacua not connected by instantons or connected only through sinks (Clifton, Linde, and Sivanandam, 2007)—all dS vacua are reachable with nonzero transition rates. This result holds for minima of the same potential, but arguments were given for parts of the landscape with different topologies as well. See Danielsson, Johansson, and Larfors (2007), Chialva *et al.* (2008), and Ahlqvist *et al.* (2011) for a discussion of connections between Calabi-Yau flux vacua.

The dominant vacuum may sound a bit like the old dream of a selection principle. Could this be the mathematically unique vacuum that many people have been hoping for? Since it can in principle be determined from first principles (by computing all vacuum transition amplitudes), it is not very likely that it would land exactly in an anthropic point in field theory space; see Fig. 1. If the dominant vacuum is not itself anthropic, the anthropic vacuum reached from it by the largest tunneling amplitude is now a strong candidate for describing our Universe. With extreme optimism one may view this as an opportunity to compute this vacuum from first principles (Douglas, 2012). Unfortunately, apart from the technical obstacles, there is a more fundamental problem: the dominant vacuum itself depends on the way the measure is defined.

2. Local and global measures

The earliest attempts at defining a measure tried to do so globally for all of space-time by defining a time variable and imposing a cutoff. Several measures of this kind have been proposed, which we will not review here.

But a comparison with black hole physics provides an important insight why this may not be the right thing to do. There is a well-known discrepancy between information disappearing into a black hole from the point of view of an infalling observer or a distant observer. In the former case information falls into the black hole with the observer, who does not notice anything peculiar when passing the horizon, whereas in the latter case the distant observer will never see anything crossing the horizon. A solution to this paradox is to note that the two observers can never compare each others observations. Hence there is no contradiction, as long as one does not try to insist on a global description where both pictures are simultaneously valid. This is called *black hole complementarity* [and has come under some fire recently; see Almheiri *et al.* (2013) and Braunstein, Pirandola, and Życzkowski (2013) and later papers for further discussion].

The same situation exists in eternal inflation. The expanding dS space, just like a black hole, also has a horizon. In many respects, the physics is in fact analogous (Gibbons and Hawking, 1977). If it is inconsistent to describe black hole physics simultaneously from the distant and infalling

observer perspective, the same should be true here. This suggests that one should count observations only within the horizon. This idea has been implemented in somewhat different ways. The causal patch measure (Bousso, 2006) takes into account observations only in the causal past of the future end point of a world line. Several variations on this idea exist which we will not attempt to distinguish here. Remarkably, in some cases these local measures are equivalent to global ones (local or global duality); see Bousso and Yang (2009) and Bousso, Freivogel, and Yang (2009).

Using only quantum mechanical considerations, Nomura (2011) developed a picture that includes observations only by a single observer. In the end, probabilities are then defined as in quantum mechanics, as squares of absolute values of coefficients of a quantum state. In this approach, “the multiverse lives in probability space,” and this is claimed to be tantamount to the many-world interpretation of quantum mechanics. Such a relation has been pointed out by others as well (Susskind, 2003; Tegmark, 2009; Aguirre, Tegmark, and Layzer, 2011; Bousso and Susskind, 2012), but it is too early to tell whether all these ideas are converging.

The current status can be summarized by two quotes from recent papers. Nomura (2012) states emphatically “The measure problem in eternal inflation is solved,” whereas just a year earlier Guth and Vanchurin (2011) concluded “We do not claim to know the correct answer to the measure question, and so far as we know, nobody else does either.”

VII. THE COSMOLOGICAL CONSTANT IN THE STRING LANDSCAPE

The anthropic explanation for the smallness of Λ requires a fundamental theory with a distribution of values of Λ , realizable in different universes. In string theory, this is provided by the Bousso-Polchinski discretuum (see Sec. IV.E). This yields a dense set of 10^{hundreds} discrete points over the full Planckian range¹⁰ of ρ_Λ . If this set does indeed exist, it would be fair to say that string theory combined with anthropic arguments explains the first 120 digits of ρ_Λ on a particular slice through parameter space. But of course all those digits are zero.

To go beyond this we need better control of inflation to deal with variations in Q and other parameters. We also need a solution to the measure problem and a better understanding of the issues of typicality and the definition of observers. At this moment the subject is still very much in a state of flux, without clear convergence to a definitive answer. For example, using different assumptions about the measure and different ways of parametrizing observers, Bousso *et al.* (2007), De Simone *et al.* (2008), and Larsen, Nomura, and Roberts (2011) obtained cosmological constant distributions that peak closer to the observed value than earlier work using the Weinberg bound. Bousso *et al.* used the amount of entropy produced in a causal patch as a proxy for observers. De Simone *et al.* used a global measure, and Larsen, Nomura, and Roberts used the solution to the measure problem proposed by Nomura (2011); the latter two use conventional anthropic criteria.

¹⁰The smoothness of this distribution near zero is important and requires further discussion; see Schwartz-Perlov and Vilenkin (2006) and Olum and Schwartz-Perlov (2007).

An important test for solutions to the problem is whether they can explain coincidences [see, e.g., [Garriga and Vilenkin \(2003\)](#)]. The most famous of these is the “why now” problem: why do we live fairly close (within a few billion years) to the start of vacuum energy domination? By its very definition, this is an anthropic question. Another striking coincidence is the order of magnitude of the absolute value of upper and lower bounds on Λ [cf. Eq. (3.5)]. In other words, the life span of typical stars is comparable to the age of the Universe and the starting time of vacuum energy domination. This depends on an apparent coincidence between cosmological parameters and standard model parameters $\rho_\Lambda \approx (m_p/M_{\text{Planck}})^6$.

In essentially all work determining Λ one of the coincidences is input and determines the scale for the Λ distribution. For example, in work based on galaxy formation, the quantity $Q^3 \rho_{\text{eq}}$ determines that scale, but the “why now” coincidence is not solved. On the other hand, in [Bousso *et al.* \(2007\)](#) the time of existence of observers is the input scale, so that the “why now” problem is solved if ρ_Λ peaks near 1 on that scale. This then turns the proximity of the maximum ρ_Λ for galaxy formation, i.e., the Weinberg bound, into an unexplained coincidence. If the cosmological constant can be computed as a pure number, as suggested, for example, by [Padmanabhan \(2012\)](#), all these coincidences remain unexplained. The same is true if ρ_Λ can be expressed in terms of some standard model parameters or if it is determined by the lowest possible value in the discretuum (see below). In all cases additional arguments will be needed to explain these coincidences or they will remain forever as unsolved naturalness problems.

Still more coincidences are listed in [Bousso, Hall, and Nomura \(2009\)](#). They attempt to explain them by arguing that landscape distributions may drive us toward the intersection of multiple catastrophic boundaries, beyond which life is impossible. The boundaries are computed using traditional anthropic arguments in universes with standard-model-like particle physics. They conjecture that the gauge hierarchy, via the aforementioned stellar lifetime coincidence, might be related to the cosmological constant hierarchy. The latter may then find an explanation in the discreteness of the landscape, a possibility also suggested by [Bousso *et al.* \(2011a\)](#). This requires a total number of (anthropic) string vacua of about 10^{120} . A different approach to coincidences was used by [Bousso *et al.* \(2011b\)](#), who argued that the coincidences can be understood entirely in terms of the geometry of cutoffs that define the measure in eternal inflation. They used a minimal anthropic assumption, namely, that observers are made out of matter.

Several hope to avoid the anthropic argument, even though they accept the existence of a landscape, by suggesting that the probability distribution of ρ_Λ is peaked at zero. However, strong peaking near zero for pure dS spaces is not likely to work. Only gravity can measure the cosmological constant, and in the early Universe, when the ground state is selected, its value is negligible in comparison to all other contributions. See [Polchinski \(2006\)](#) for a more extensive explanation of this point.

Despite this objection, some speculate that somehow the cosmological constant is driven to the lowest positive value Λ_{min} . The value of Λ_{min} is then roughly equal to the inverse of N , the total number of vacua. For variations on this idea, see

[Kane, Perry, and Zytow \(2005\)](#) and [Linde and Vanchurin \(2010\)](#). A different proposal was made by [Kobakhidze and Mersini-Houghton \(2007\)](#), who suggested $\Lambda_{\text{min}} = 1/N^2$. [Tye \(2006\)](#) and [Sarangi, Shiu, and Shlaer \(2009\)](#) argued that due to “resonance tunneling” all vacua have very short lifetimes, except some with very small Λ . Ideas of this kind would leave all apparent anthropic tunings unexplained.

In the full set of string vacua, not just pure dS but including matter, there may well exist a unique vacuum, defined by having the smallest positive Λ . But this is not likely to be our Universe, since a unique vacuum will not satisfy the other anthropic requirements. Even if for some reason it is strongly selected, this will generate runaway behavior in other variables or leads to the kind of catastrophic predictions explained in Sec. III.F.4.

Some use an analogy with solid state physics to argue that because of tunneling the true ground state wave function is a Bloch wave. But there is an important difference. In solid state physics observation times are much larger than tunneling times, whereas in the landscape it is just the other way around. If observations are made at times much shorter than the tunneling time, this leads to collapse of the wave function and decoherence. Furthermore, in the landscape there must exist tunneling processes that change gauge groups, representations, and parameters. Therefore, these cannot be treated as superselection sectors. The best one could hope to get is a linear combination of amplitudes with different values of all standard model and cosmological parameters, which does not solve the problem of determining them.

Should we expect to understand why $\Lambda > 0$ in our Universe or is the sign just selected at random? On the one hand, from the perspective of vacuum energy in quantum field theory the point $\Lambda = 0$ is not special. Nor is it special from the anthropic perspective: life with $\Lambda < 0$ seems perfectly possible. On the other hand, classical physics and cosmology at late times are extremely sensitive to the sign: the Universe either collapses or expands. The difference in sign implies important differences in quantum physics. The definition of the S matrix in quantum field theory (and string theory) is problematic in dS. Tunneling amplitudes between vacua are singular for $\Lambda \rightarrow 0$ (see Sec. VI). In AdS spaces any possibility of life finishes at the crunch, and it matters how closely one can approach it; in dS spaces life is not limited by a crunch, but by the burning out of stars within the Hubble horizon [see [Peacock \(2007\)](#) for a discussion]. Note that many consider only positive values for Λ , and some that do not [see, e.g., [Bousso *et al.* \(2011b\)](#)] actually predict negative Λ more strongly than positive Λ . The differences between AdS and dS are too large to blindly assume that we ended up in a dS universe purely by chance.

Many other aspects of the cosmological constant problem and possible solutions are reviewed by [Weinberg \(1989\)](#), [Polchinski \(2006\)](#), and [Bousso \(2008\)](#).

VIII. CONCLUSIONS

Barring any surprises, we are facing a choice between two roads. One of them, the traditional symmetry-based road of particle physics, may ultimately lead nowhere. A uniquely determined theory of the Universe and all of its physics leaves us with profound conundrums regarding the existence of life.

The other road, leading toward a big landscape, is much more satisfactory in this respect, but is intrinsically much harder to confirm. Low-energy supersymmetry might have helped, but is a luxury we may not have. The SUSY-GUT idea, the lamppost of the symmetry road, is losing its shine. GUTs do not fit as comfortably in the string landscape as most people believe, and SUSY does not fit well with the data; the ways out are increasingly becoming epicyclical. Confusingly, the opposite is also true: GUTs still look as attractive as ever from a low-energy perspective, and the landscape, despite many arguments going both ways, may prefer low-energy SUSY after all.

Will we ever know? Here are some possible future developments that cast serious doubts on the string theory landscape.

- The evidence for a well-distributed and connected dS landscape in string theory crumbles.
- Low-energy supersymmetry is strongly predicted, but not seen at the LHC (or vice versa).
- Solid evidence for variations of constants of nature emerges.

There is movement on all of these fronts, and in 20 years we will probably have a different view on all of them. There are plenty of other possibilities for game-changing developments.

In the string theory landscape, the key concept linking all these issues is moduli. This is where all lines meet: supersymmetry breaking and its scale, variations of constants, axions and the strong CP problem, (eternal) inflation, dark matter, the cosmological constant and/or quintessence, and ultimately the existence and features of the string landscape itself.

But suppose there is no convincing experimental falsification on any of these issues. Then will we ever know? Ultimately the convincing evidence may have to come from theory alone. Of all the open theoretical issues, the measure problem of eternal inflation is probably the biggest headache. But not everything hinges on that. In the context of string theory, the following problems can be addressed without it.

- Derive string theory from a principle of nature.
- Establish its consistency.
- Prove that it has a landscape.
- Prove that the standard model is in that landscape.
- Show that all quantities are sufficiently densely distributed to explain all anthropic fine-tunings.
- Confirm that these vacua are connected by some physical process, so that they can all be sampled.

Perhaps this is as far as we will ever be able to go. We may never be able to derive our laws of physics, but we may just feel comfortable with our place in the landscape. This requires understanding our environment, not just the point where we live, but also the region around it. This can fail dramatically and cast severe doubts on certain landscape assumptions. Therefore a large part of this review has been devoted to all the impressive work that has been done in this area during the last decade. There is great physics in anthropic reasoning.

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