Chapter 1

The String Theory Landscape

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Perhaps the most important way string theory has affected the perspective of particle physics phenomenology is through the "string theory landscape". We discuss the evidence supporting its existence, describe the regions of the landscape that have been explored, and examine what the string theory landscape might imply for most Standard Model problems.

1. Changing Perspectives

This chapter consists of three parts. In the present section, we will give a brief historical overview of the birth of the string theory landscape and we will explain its main features. We also present the main arguments in favor of its existence, both from a top-down (string theory) and a bottom up (the Standard Model) point of view. Furthermore we underline the important changes this concept has on the perspectives for string phenomenology. In section 2 we present an overview of the various methods of constructing string theories in four dimensions, by direct construction and by compactification. In section 3 we discuss how far one can get towards understanding the Standard Model from the landscape perspective, in comparison to the traditional, symmetry-based approach.

Part of this chapter is based on the review article [1].

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1.1. *The Age of Symmetries*

At the time when string theory started being considered as a theory of all interactions including gravity, the theoretical work on the Standard Model had reached its final form. Several decades of experimental work were still needed to establish it, but theorists started moving ahead using the concepts that had led to so much success in understanding the three non-gravitational interactions. The most important of these concepts was symmetry. It revolutionized the understanding of fundamental physics. From the seemingly hopeless chaos of nuclear and hadronic physics a very simple description of those interactions had emerged in just about two decades: a spontaneously broken gauge theory which was called the "Standard Model".

The name does not suggest much confidence in this idea, and indeed nobody saw the Standard Model as more than an intermediate stage, at best an approximate description of nature, eventually to be supplanted by something even simpler and mathematically more elegant. Indeed, by using the same group-theoretical methods that turned out to be so successful in the description of the Standard Model, new theories were found that looked more attractive. The highest achievable goal appeared to be supersymmetric Grand Unified Theories (susy-GUTs), which got even more credibility in the early nineties, when the precision results of LEP suggested that the three gauge couplings evolved to a common value at the very interesting energy scale of about 10^{16} GeV, a few orders of magnitude below the Planck scale. Even today, despite the fact that the experimental evidence has not (yet) shown up, it is hard to believe that this could all be just a coincidence.

A few years earlier string phenomenology had entered the scene. When it did, it seemed to point at even grander symmetries. In 1984, the first results obtained by compactifying the just-discovered $E_8 \times E_8$ heterotic string suggested the ultimate unification. The Standard Model appeared to emerge (almost) uniquely from the jewel of Lie algebra theory, the Lie algebra E_8.

1.2. *The Birth of the Landscape*

But in the remainder of the eighties there was the beginning of a slow shift away from the notions of symmetry and uniqueness that were considered almost self-understood until then. History will decide if this was the beginning of a paradigm shift or just prematurely giving up on uniqueness. But the evidence that the former is true is mounting.

Perhaps we will conclude one day that these beautiful ideas have always carried the seeds of their own destruction. A Standard Model family fits beautifully in the (5)+(10) of $SU(5)$ and even more beautifully in the (16) of $SO(10)$, but the Standard Model Higgs field does not. Furthermore, even if $SU(5)$ or $SO(10)$ exist as symmetries at short distances, there is no unique path to the Standard Model at the weak scale: in $SU(5)$ models there are two minima of the GUT Higgs potential, one leading to the Standard Model, and one to a $SU(4) \times U(1)$ gauge theory. In $SO(10)$ models the number of options increases.

The idea of low energy supersymmetry was also plagued by serious problems from the very beginning. It introduces light bosons into the spectrum that lead to rapid proton decay. By contrast, a sufficiently long life-time of the proton is automatic in the Standard Model. On general grounds, one would expect low energy supersymmetry to give rise to flavor violations that should have been observed a long time ago already. These, as well as other problems can be evaded by additional assumptions, but it is disturbing that the pieces of the puzzle do not fall into place more easily. Furthermore, although susy and GUTs are well-motivated answers to important questions, they have never led to a substantial simplification of the Standard Model.

1.2.1. *String Vacua*

The uniqueness of string theory was also in doubt right from the start. The $E_8 \times E_8$ heterotic string was not unique, but part of a small set of 10- and 11-dimensional supersymmetric theories which were initially taken less seriously.

But more importantly, there was an explosion of compactifications and four-dimensional string constructions in the two years following 1984 [2–7]. Already as early as 1986 it became customary to think of the different string theories or compactifications as "vacua" or "ground states" of a fundamental theory (see for example the last line of [8] or discussion at the end of [9]; here one also finds the remark that perhaps our universe is merely a sufficiently long-lived metastable state). The proliferation of "string vacua" has not stopped since then. Here and in the following we use the word "vacuum" for the metastable state that correctly describes our Universe, and all its analogues with different gauge theories. The proper definition is itself a difficult issue, especially in de Sitter space, but if no well-defined description exists that matches our Universe, string theory would be wrong anyway.

1.2.2. *Moduli*

Soon it became clear that these "string compactifications" or "four-dimensional strings" had continuous deformations that can be described by vacuum expectation values of massless scalar singlets, called "moduli". Typically, there are tens or hundreds of them. All quantum field theory parameters depend on the moduli, and hence the existence of moduli is a first step towards a plethora of possibilities.

These singlets generate unobserved fifth forces and their presence is cosmologically unacceptable [10, 11], but so is the fact that supersymmetry is unbroken. For more than a decade, this left room for the possibility that the abundance of string vacua would be reduced to just a few, maybe just one, once these problems were solved. But in the beginning of this century considerable progress was made towards solving the problem of moduli stabilization and – to a lesser extent – that of supersymmetry breaking.

The large number of available ingredients (fluxes, D-branes, orientifold planes and various perturbative and non-perturbative effects) led to the nearly inevitable conclusion that if there was one solution, there were going to be many more. Almost two decades after 1984 the denial phase reached its end, marked by an influential and somewhat provocative paper by L. Susskind [12], who also gave the subject its current name, the "string landscape".

1.2.3. *The Cosmological Constant*

Remarkably, these developments were driven to a considerable extent by observation: the discovery of an accelerated expansion of the universe in 1998 [13, 14]. The most straightforward interpretation is that we live in a universe with positive vacuum energy density, which acts like a cosmological constant, and implies that we live in de Sitter (dS) rather than flat Minkowski space. Contrary to some statements in the literature, there was never any difficulty in getting positive vacuum energy in string theory. Some of the aforementioned papers from 1986 built non-supersymmetric strings, and some of those string theories have positive vacuum energy. At that time this feature was merely observed, but not yet considered to be of any interest. However, we have little computational control over non-supersymmetric strings, and at the moment the only viable path to string theory in dS space is to "up-lift" supersymmetric AdS vacua with negative vacuum energy.

The explanation of the observed accelerated expansion requires not only positive vacuum energy, but an extremely small amount of it. It is about 120 to 60 orders of magnitude less than its natural scale, depending on whether one compares to the Planck scale or the weak scale. In non-supersymmetric string theory vacuum energy comes out as a sum of positive and negative contributions of Planckian size. Everything we know suggests that this will give rise to numbers of order one in units of the natural value. If the contributions in the sum are random, the chance of finding a result near the observed value is about 1 in 10^{120}. This would imply that one needs an ensemble of at least 10^{120} vacua to have a chance of finding one like ours. With any smaller ensemble, the existence of the small observed value would be a bizarre coincidence.

1.2.4. *The Bousso-Polchinski Mechanism*

It was realized decades ago [15] that anti-symmetric tensor fields $A_{\mu\nu\rho}$ might play an important rôle in solving the cosmological constant problem. Such four-index field strengths can get constant values without breaking Lorentz invariance, namely $F_{\mu\nu\rho\sigma} = c\epsilon_{\mu\nu\rho\sigma}$. If we couple the theory to gravity, it gives a contribution to the cosmological constant Λ:

$$\Lambda = \Lambda_0 - \frac{1}{48} F_{\mu\nu\rho\sigma} F^{\mu\nu\rho\sigma} = \Lambda_0 + \frac{1}{2} c^2, \tag{1}$$

where Λ_0 is the cosmological constant in the absence of anti-symmetric field strength contributions. In string theory c is not an arbitrary real number: it is quantized [16]. These quantized fields are called “fluxes”. It turns out that string theory typically contains hundreds of fields $F_{\mu\nu\rho\sigma}$, which we label by $i = 1, \ldots, N$. The resulting formula for Λ is

$$\Lambda = \Lambda_0 + \frac{1}{2} \sum_{i=1}^{N} n_i^2 f_i^2. \tag{2}$$

where the f_i are numbers derived from the string theory under consideration. One would expect the values for the real numbers f_i to be different. If the values of f_i are distinct and incommensurate, then Eq. (2) defines a dense discrete set of values. Bousso and Polchinski called it a “discretuum”. This realizes a dynamical neutralization of Λ first proposed by [17, 18]. See also [19] for a related realization of this idea in string theory.

The discretuum is populated by some physical process that is able to connect the different string vacua. The mechanism proposed for this is tunneling by bubble nucleation in eternal inflation, a near inevitability in most models of inflation. See [20, 21] for reviews and references. This area is – so far – less deeply connected to string theory, and therefore we will not discuss it in detail, except to mention that it leads to the very thorny issue of the multiverse measure problem. See [22–25] for various ideas about this.

1.2.5. *Existence and Distribution of de Sitter Vacua*

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 might be, 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 non-flat structure, this argument would fail. There might still exist a huge landscape, but it would be useless.

We will consider here only type-IIB (and related F-theory) compactifications where the most explicit results have been obtained. For references to work in other areas and more details see [1, 26–30].

In type-IIB theories one starts with a Calabi-Yau compactification with h_{21} complex structure ("shape") moduli and h_{11} Kähler ("size") moduli, where h_{21} and h_{11} are the Hodge numbers of the CY manifold (see section 2.3 for more details on Calabi-Yau compactifications). One can add to this background configuration a choice of gadgets from the string theory toolbox, such as 3-form RR and NS fluxes, 5-form fluxes, denoted $F3, H3$ and $F5$ respectively, and D3 and D7 branes.

The 3-form fluxes can stabilize all complex structure moduli. This is due to a tree-level term in the superpotential that takes the form [31]

$$W_{\text{flux}} = \int (F_3 - \tau H_3) \wedge \Omega \ , \tag{3}$$

where $\tau = a + ie^{-\phi}$, and a is 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. Since every CY manifold has at least one Kähler modulus, this leaves therefore at least one modulus unfixed.

One may 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 a parameter. Such terms can be generated by instantons associated with Euclidean D3-branes [32] or from gaugino condensation in gauge groups on wrapped D7 branes. If at least one of these effects is present, string vacua with all moduli stabilized can be obtained [33]. This work, usually referred to as "KKLT", builds on several earlier results, such as [34–36] and references cited therein. The solution obtained in this way has a negative vacuum energy, and is a fully stabilized supersymmetric AdS vacuum. However, the required instanton contributions may not exist in all cases. They are not generic [37] and may even be so rare that one only gets a "barren landscape" [38].

The next step is more problematic and more controversial. One must break supersymmetry and obtain a dS vacuum (this is called "up-lifting"). 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. Attempts to realize the KKLT uplifting in supergravity or string theory have failed so far [39, 40], but opinions differ on the implications of that result. There exist several alternatives to D3-brane uplifting [41–45].

The result of a fully realized KKLT construction is a string vacuum that is free of tachyons, but one still has to worry about non-perturbative 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.

An alternative scenario (called the LARGE volume scenario or LVS) was presented in Ref. [46]. The starting point is type-IIB fluxes stabilizing the complex structure moduli and the dilaton and axion. By means of suitable $(\alpha')^3$ corrections these authors were able to find minima where all moduli are stabilized at exponentially large volumes in *non*-supersymmetric AdS vacua. Additional mechanisms are then needed to lift the vacuum to dS. An explicit example was presented recently [47].

The existence of the required dense distribution of vacua is still disputed, and some even question the existence of *any* such vacuum. Recent work seems to indicate that in the vast majority of cases the AdS vacua become tachyonic after uplifting [48, 49]. Another potential problem is a dramatic increase in tunneling rates as a function of the number of moduli [50]. These effects may dramatically reduce the number of dS vacua, rendering the Bousso-Polchinski argument inadequate. In [51] several criticism of the landscape are presented, including the use of effective potentials and of Coleman-de Luccia [52] tunneling between dS vacua in theories of gravity.

1.2.6. *Vacuum Counting*

The KKLT construction has been the starting point for estimates of the total number of flux vacua [53–55].

$$N_{\rm vac} \approx \frac{(2\pi L)^{K/2}}{(K/2)!}, \tag{4}$$

where L is a number of order 1 to 100 (the "tadpole charge") and K the number of distinct fluxes. For typical manifolds this gives numbers of order 10^N, where N is of order a few hundred. A often quoted estimate is 10^{500}.

It is noteworthy 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), and it is this large exponent that makes neutralization of the cosmological constant possible. All the ingredients used in the foregoing discussion are already present in string theory. Since all Standard Model parameters depend on the moduli, this results in a large distribution of options covering the environment of the Standard Model in QFT.

1.3. *A Paradigm Shift?*

If we were to accept that our laws of physics are picked out of a huge ensemble, and that the parameters have such special values just by coincidence, this would imply the end of science. Then the entire Standard Model could just be a random item from a huge ensemble. It is indeed remarkable that in the current state of particle physics, many of the remaining problems *could* be just environmental: the Standard Model provides an adequate description, but often with strange parameter values. Some genuine problems remain (such as dark matter and the mechanisms behind inflation and baryogenesis), but most other problems that are often discussed should really be called "worries". This means that we cannot be completely sure that there exists a solution. Perhaps these problems only exist in our minds. This includes the choice of the Standard Model gauge group, the choice of matter representations, charge quantization, the number of families, quark and lepton mass hierarchies, the smallness of neutrino masses, the gauge hierarchy problem and dark energy.

1.3.1. *Anthropic arguments*

All these worries only exist because there are minds to worry about them. It is quite plausible that this would not be the case if we allow the parameters and discrete choices of the Standard Model to vary. In an ensemble like the string landscape many such variations can occur, and it is inevitable that worrying minds will only worry about the small subset in which they can exist. This statement is one of several possible formulations of the "anthropic principle". It is a misnomer for several reasons, and in this formulation it is certainly not a principle of physics, like the equivalence principle. It is output, not input. One could even choose to ignore it, but then one would miss several potential explanations of some of the environmental problems.

Without its embedding in the string theory landscape, anthropic reasoning might also be called the end of physics, but in combination the two concepts merely are a complete change of course for traditional physics. The notion of symmetries as a fundamental concept is replaced by a combination of anthropic arguments and information about distributions of parameters in a mathematically well-defined ensemble. The lack of evidence for "new physics" may imply that we have reached the historical moment where this change is occurring. But particle physics is an experimental science, and the huge number of experimental results coming up in the next few years may revive the notion of symmetries, and postpone the emergence of a first glimpse of a landscape indefinitely.

1.3.2. *Derivability vs. Uniqueness*

One often finds criticisms of string theory like: "String theory was supposed to explain why elementary particles could only have the precise masses and forces that they do". In reality there has never been even a shred of evidence that string theory was going to lead to that. It is also nearly impossible to find quotes of this kind even in the earliest string papers. People making such statements are simply projecting their own expectations for a fundamental theory on string theory.

These expectations reflect the traditional uniqueness paradigm of particle physics: the hope that one day we will be able to derive all laws of physics, in particular the Standard Model and all its parameters. The fine structure constant α was expected to be given by some simple formula. This hope can be illustrated by famous quotes by Einstein, Feynman and others. But two concepts are often confused in this discussion, *uniqueness* and *derivability*. According to our current understanding of string theory, it is not really the uniqueness of the underlying theory that is at stake, but the uniqueness of the vacuum. If the vacuum is not unique, the Standard Model and its parameters can not be derived by purely mathematical manipulations. Additional information about the choice of vacuum, either phenomenological or anthropic, must be provided.

The fact that the Standard Model and it couplings fit nicely in an $SU(5)$ susy-GUT is often proclaimed as a strong hint in favor of uniqueness. But once again a distinction between uniqueness and derivability must be made. While GUTs may indeed point in the direction of a unique theory, SU(5) GUTs also gave the first hint *against* derivability, since they have two physically distinct vacua. The non-uniqueness of the return path towards lower energies first encountered in GUTs becomes much worse if one starts from the loftier vantage point of $E_8 \times E_8$ heterotic strings. It is like climbing a mountain: eventually one may reach a unique point, the top, but there may exist many other paths downwards, leading to other valleys.

Since string theory emerged during the height of the symmetry era, it is not surprising that it was first seen as the realization of the dream of uniqueness, interpreted as derivability. But in reality, string theory has been sending us exactly the opposite message almost from the beginning. One day, this may be recognized as its most important contribution to science.

1.3.3. *Evidence outside String Theory*

From now on we will use the term "uniqueness" rather than "derivability", because that is what is commonly used. There has never been any evidence in favor of the uniqueness paradigm, the idea that the Standard Model has to be derivable. But there are several pieces of circumstantial evidence of the contrary, even without string theory.

Theories of inflation typically lead to multiple instances of new universes, a "multiverse". Even without inflation, what argument do we have to suggest that our own universe is unique, in any sense of the word? And if it is not unique, what argument do we have to tell us that the other universes must have exactly the same laws of physics as ours? The only fact that makes our universe and laws of physics unique is that they are the only ones we can observe.

The possible existence of a plethora of scalars provides another reason to question the uniqueness of the Standard Model. We have seen particles of various spins and Standard Model couplings, but only recently we may have observed the first Lorentz scalar, the Higgs boson. We have only been able to see it because it is unusually light and because it comes from a field that is not a Standard Model singlet. But experimentally we know nothing about Lorentz scalar fields that are gauge singlets. If they exist, they would appear in the Lagrangian as polynomials, modifying all dimensionless parameters. Then all parameter values are vacuum expectation values of scalar fields. One may still hope that this vacuum expectation value is somehow uniquely determined, but in almost all examples we know (including the Higgs potential of $SU(5)$ GUTs) scalar potentials tend to have more than one local minimum. One cannot say more without a more concretely defined theory, but in the only theory where such potentials can be discussed, string theory, the number of local minima appears to be astronomical.

Finally, the fact that the Standard Model is anthropically tuned provides evidence against the idea of its derivability. This is because the Standard Model stands out as a very special region in parameter space where nuclear physics and chemistry lead to complex structures we call "life". It would require an uncanny miracle for the two unrelated computations to give compatible results. Especially the last argument suggests that the ultimate fundamental theory – assuming such a notion makes sense and that we have enough intelligence and information to determine it – must have a large ensemble of physically connected vacua. This allows a process like eternal inflation to sample all these vacua, occasionally producing a universe within the anthropic domain.

1.3.4. *Uniqueness in the String Landscape?*

One may still hope that the resulting scalar potential somehow has a mathematically unique local minimum, but that would be pure wishful thinking. One may even hope for a unique global minimum. However, it is not clear what to minimize, because vacuum energy is not bounded from below. However, if vacuum energy takes discrete values, there is a – presumably unique – vacuum with the lowest positive vacuum energy. Could that describe our universe? Could it be that in the sampling process of universes this particular one is somehow preferred? Another notion is the "dominant master vacuum", the state that dominates eternal inflation because it is most frequently sampled, often by a huge factor [56]. This is determined by its stability against tunneling to other states, as well as the likelihood of others tunneling into it. Both of these – the vacuum with lowest positive vacuum energy or the dominant master vacuum – have a sense of uniqueness, but the anthropic tuning argument makes it very unlikely that they happen to have the properties that allow life to exist. There is no need for that, once there is a landscape. Then the anthropic genie is already out of the bottle. All that is required is that there exist metastable vacua in the anthropic zone of parameter space, and that their sampling frequency is non-zero. There is no reason why they should dominate the multiverse statistics. The dominant master vacuum does have an important positive feature, because one may argue that most vacua have it in its history. Then it can serve as a kind of eraser of initial conditions. Most observers would find themselves in the anthropic universe most frequently reached by tunneling from the dominant master vacuum. See [57] for some speculations regarding this idea.

1.4. *Changing Perspectives on String Phenomenology*

The existence (or non-existence) of the string theory landscape has an important impact on string phenomenology. One may distinguish at least three different attitudes. The first is that we should simply find the exact point in the landscape that corresponds to the Standard Model, use current data to fix it completely, and then make an indefinite number of predictions for future experiments and observations. This includes all work on explicit "string model building" in many corners of the landscape, reviewed in the next section. The second is trying to extract generic predictions from classes of models rather than individual ones. An example of work in this category is the study of a class of M-theory models, reviewed in [58] (see also section 2.11). The third is to try and understand the Standard Model by considering landscape distributions in combination with anthropic constraints. This is the point of view we take in section 3. These three points of view are not mutually exclusive. Their relative importance depends on how optimistic one is about the chances of finding the exact Standard Model as a point in the landscape.

2. The Compactification Landscape

We will present here just a brief sketch of the string compactification landscape. For further details we recommend the very complete book [59] and references therein.

2.1. *World-sheet versus space-time*

In their simplest form, fermionic string theories live in ten flat dimensions. In addition there is an 11-dimensional theory that is not described by interacting strings, but closely related to string theory, known as M-theory. But in any case, to make contact with the real world we have to find theories in four dimensions.

There are essentially two ways of doing that: to choose another background space-time geometry, or to change the world-sheet theory. The geometry can be chosen as a flat four-dimensional space-time combined with a compact six-dimensional space. This is called "compactification". The world-sheet theory can be modified by choosing an appropriate two-dimensional conformal field theory. In D-dimensional flat space a string theory is described by D free two-dimensional bosons X^μ, and, if it is a fermionic string, also by D free fermions ψ^μ. Instead, one can choose another two-dimensional field theory that satisfies the same conditions of conformal invariance, called a conformal field theory (CFT). In particular one may use interacting two-dimensional theories, as long as X^μ and ψ^μ $\mu = 0, \ldots 3$ remain free fields.

The simplest compactification manifold is a six-dimensional torus. This can easily be described both from the space-time and the world-sheet point of view. The resulting theories only have non-chiral fermions in their space-time spectrum. The same is true for the more general asymmetric torus compactifications of the heterotic string with 6 left-moving and 22 right-moving "chiral" bosons found by Narain [2].

The chirality problem is easily solved by a simple generalization that yields a valid compactification manifold, namely a torus with discrete identifications. These are called orbifold compactifications [60]. These methods opened many new directions, such as orbifolds with gauge background fields ("Wilson lines") [61], and were soon generalized to *asymmetric orbifolds* [7], where "asymmetric" refers to the way left- and right-moving modes were treated. Just as torus compactifications, orbifolds can be viewed from both a space-time and a world-sheet perspective. Some orbifold compactifications can be understood as singular limits of geometric Calabi-Yau compactifications, which historically were discovered a little earlier (see section 2.3). With more complicated compactifications, the connection between the world-sheet and space-time perspectives becomes more and more difficult to make.

2.2. *General Features*

2.2.1. *Massive and massless modes*

Before introducing some of the earliest string constructions and compactifications in a little more detail, we will give a summary of the kind of spectra they generically produce. Here we will assume a supersymmetric spectrum. This already implies the prediction of a large number of particles that have not (yet) been observed. All types of particles listed below typically occur in non-supersymmetric string spectra as well, but in that case it is even less clear what their ultimate fate is, since the stability of these string theories is not understood.

Any string theory contains infinitely many additional particles: massive string excitations, Kaluza-Klein modes as in field theory compactifications, and winding modes due to strings wrapping the compact spaces.

Their masses are respectively proportional to the string scale, the inverse of the compactification radius or the compactification radius itself. In world-sheet constructions the different kinds of modes are on equal footing, and have Planckian masses. In geometric constructions one can consider large volume limits, where other mass distributions are possible. But in any case, of all the modes of the string only the massless ones are relevant for providing the Standard Model particles, which will acquire their masses from the Higgs mechanism and QCD, as usual.

Among the massless modes of string theories one may find some that match known particles, but usually there are many that do not match anything we know. This may just be an artifact of the necessarily primitive methods at our disposal. Our intuition from many years of four-dimensional string model building may well be heavily distorted by being too close to the supersymmetric limit, and by algebraically too simple constructions. Some of the additional particles may actually be a blessing, if they solve some of the remaining problems of particle physics and cosmology. The art of string phenomenology is to turn all seemingly superfluous particles into a blessing, or understand why their presence is not generic.

In addition to moduli (already introduced in subsection 1.2.2) and axions (to be discussed in section 3.6) string spectra generically include:

2.2.2. *Chiral fermions and mirrors*

All charged Standard Model fermions are chiral, and hence they can only acquire a mass after weak symmetry breaking. Therefore one can say that the weak interactions protect them from being very massive. It is very well possible that for this reason all we have seen so far at low energy is chiral fermionic matter.

In attempts at getting the Standard Model from string theory, it is therefore reasonable to require that the chiral spectra match. In general one finds additional vector-like matter, whose mass is not protected by the weak interactions. Typically, if one requires three chiral families, one gets $N+3$ families and N mirror families. If the N families "pair off" with the N mirror families to acquire a sufficiently large mass, the low energy spectrum agrees with the data.

2.2.3. *Additional vector bosons*

Most string spectra have considerably more vector bosons than the twelve we have seen so far in nature. Even if the presence of $SU(3)$, $SU(2)$ and $U(1)$ as factors in the gauge group is imposed as a condition, one rarely finds just the Standard Model gauge group. In heterotic strings one is usually left with one of the E_8 factors. Furthermore in nearly all string constructions additional $U(1)$ factors are found. A very common one is a gauged $B-L$ symmetry.

Additional gauge groups are often needed as "hidden sectors" in model building, especially for supersymmetry breaking. Extra $U(1)$'s may be observable trough kinetic mixing [62] with the Standard Model $U(1)$, via contributions to the action proportional to $F_{\mu\nu}V^{\mu\nu}$, where F is the Y field strength, and V the one of the extra $U(1)$'s.

2.2.4. *Exotics*

One often finds particles that do not match any of the observed matter representations, nor their mirrors. Notorious examples are color singlets with fractional electric charge or higher rank tensors. These are generically called "exotics". If there are exotics that are chiral with respect to $SU(3)\times SU(2)\times U(1)$, these spectra should be rejected, because any attempt to make sense of such theories is too far-fetched to be credible. These particles may be acceptable if they are vector-like, because one may hope that they become massive under generic perturbations.

2.3. ***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 in [63], even before the aforementioned mathematically simpler orbifold compactifications. The compactification manifolds are six-dimensional, Ricci-flat Kähler manifolds with $SU(3)$ holonomy, called Calabi-Yau manifolds. The $B_{\mu\nu}$ field strength $H_{\mu\nu\rho}$ was assumed to vanish, which leads to the consistency condition

$$dH = \text{Tr } R \wedge R - \frac{1}{30}\text{Tr } F \wedge F = 0. \quad (5)$$

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 rôle in supersymmetry breaking.

If some dimensions of space are compactified, ten-dimensional fermion fields are split as

$$\Psi_+(x, y) = \Psi_L(x)\Psi_+(y) + \Psi_R(x)\Psi_-(y) \quad (6)$$

where x denotes four-dimensional and y six-dimensional coordinates, $+/-$ denotes chirality in ten and six dimensions, and L, R denote chirality in four dimensions. The number of massless fermions of each chirality observed in four dimensions is determined by the number of zero-mode solutions of the six-dimensional Dirac equation in the background of interest. These numbers are equal to two topological invariants of the Calabi-Yau manifold, the Hodge numbers, h_{11} and h_{12}. As a result one obtains h_{11} chiral fermions in the representation (27) and h_{12} in the $(\overline{27})$ of E_6. The group E_6 is a known extension of the Standard Model, an example of a Grand Unified Theory, in which all three factors of the Standard Model are embedded in one simple Lie algebra. It is not the most preferred extension; a Standard Model family contains 15 or 16 (if we assume the existence of a right-handed neutrino) chiral fermions, not 27. However, since the 11 superfluous fermions are not chiral with respect to $SU(3) \times SU(2) \times U(1)$, they can acquire a mass without the help of the Higgs mechanism, in the unbroken Standard Model. Therefore these masses may be well above current experimental limits.

The number of Calabi-Yau manifolds is huge. A subset associated with four-dimensional reflexive polyhedra has been completely enumerated [64]. This list contains more than 470 million topological classes with 31,108 distinct Hodge number pairs.

In 1986 Strominger [3] considered more general geometric background geometries with torsion. This gave rise to so many possibilities that the author concluded "*all predictive power seems to have been lost*".

2.4. *Free Field Theory Constructions*

Several other methods were developed around the same time as Calabi-Yau compactifications and orbifolds. Narain's generalized torus compactifications lead to a continuous infinity of possibilities, but all without chiral fermions. Although this infinity of possibilities is not really a surprising feature for a torus compactification, Narain's paper was an eye-opener because, unlike standard six-dimensional torus compactifications, this approach allowed a complete modification of the gauge group.

More general world-sheet methods started being explored in 1986. Free field theory constructions allowed a more systematic exploration of certain classes of string theories. It became clear very quickly that also in this case there was a plethora of possibilities. But unlike Narain's constructions, these theories can have chiral fermions, and furthermore they did not seem to provide a continuum of options, but only discrete choices. With the benefit of hindsight, one can now say that all these theories *do* have continuous deformations, which can be realized by giving vacuum expectation values to certain massless scalars in the spectrum. Since these deformed theories do not have a free field theory descriptions, these deformations are not manifest in the construction. They are the world sheet construction counterparts of the geometric moduli. This does however not imply that the plethora of solutions can simply be viewed as different points in one continuous moduli space. Since many spectra are chirally distinct, it is more appropriate to view this as the discovery of a huge number of distinct moduli spaces, all leading to different physics.

An important tool in these free-field theory constructions is 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 [4, 6] or free bosons [5]. These constructions are closely related, and there is a huge area of overlap: constructions based on complex free fermions pairs can be written in terms of free bosons. However, one may also consider real fermions or free bosons on lattices that do not allow a straightforward realization in terms of free fermions.

2.4.1. *Free fermions*

Both methods have to face the problem of finding solutions to the conditions of modular invariance, a one-loop consistency condition. In the fermionic constructions this is done by allowing periodic or anti-periodic boundaries on closed cycles on the manifold for all fermions independently. Modular transformations change those boundary conditions, and hence they are constrained by the requirements of modular invariance. These constraints can be solved systematically (although in practice usually not exhaustively). Very roughly (ignoring some of the constraints), the number of modular invariant combinations is of order $2^{\frac{1}{2}n(n-1)}$ for n fermions. There are 44 right-moving and 18 left-moving fermions, so that there are potentially huge numbers of string theories. In reality there are however many degeneracies.

In-depth explorations [65] have been done of a subclass of fermionic constructions using a special set of free fermion boundary conditions that allows spectra with three families to come out. This work focuses on Pati-Salam model. Other work [66, 67] explores the variations of the "NAHE" set of free fermion boundary conditions. This is a set of fermion boundary vectors proposed in [68] that are a useful starting point for finding "realistic" spectra.

2.4.2. *Free Bosons: Covariant Lattices*

In bosonic constructions the modular invariance constraints are solved by requiring that the momenta of the bosons lie on a Lorentzian even self-dual lattice. This means that the lattice of quantized momenta is identical to the lattice defining the compactified space, and that all vectors have even norm. Both conditions are defined in terms of a metric, which is $+1$ for left-moving bosons and -1 for right-moving ones. These bosons include the ones of Narain's torus, plus eight right-moving ones representing the fermionic degrees of freedom, ψ^μ and the ghosts of superconformal invariance. These eight bosons originate from the *bosonic string map* (originally developed for ten-dimensional strings [69]) used to map the entire fermionic sector of the heterotic string to a bosonic string sector [5]. Then the Lorentzian metric has signature $((+)^{22}, (-)^{14}))$, and the even self-dual lattice is denoted $\Gamma_{22,14}$. This is called a *covariant lattice* because it incorporates space-time Lorentz invariance for the fermionic string. Since the conditions for modular invariance are invariant under $SO(22, 14)$ Lorentz transformations, and since the spectrum of L_0 and $\bar{L}_0$ is changed under such transformations, their would appear to be a continuous infinity of solutions. But the right-moving modes of the lattice are strongly constrained by the requirement of two-dimensional supersymmetry, which is imposed using a non-linear realization [70] (other realizations exist, see for example [71, 72]). This leads to the so called "triplet constraint" [4]. This makes the right-moving part of the lattice rigid. The canonical linear realization of supersymmetry, relating X^μ to ψ^μ, on the other hand leads to lattices $\Gamma_{22,6} \times E_8$ with complete Lorentz rotation freedom in the first factor, which is just a Narain lattice.

2.5. *An Early Attempt at Vacuum Counting.*

Several of the 1986 papers make attempts at getting a rough idea about the number of solutions. This is fairly straightforward for free fermions with periodic and anti-periodic boundary conditions, as explained above. However, the main problem is that not all solutions are different. Indeed, in general there are huge degeneracies among solutions that reduce the estimate by large factors. We will explain here a counting estimate used for covariant lattice constructions, because it give an interesting insight in the growth of the number of possibilities. However, this should not be confused with counting moduli stabilized points in potentials. Indeed, all these string theories have unstabilized moduli.

An interesting estimate exists for even self dual lattices, which has the advantage that it only counts distinct solutions. Unfortunately, heterotic strings are based on lorentzian lattices, for which there are no such theorems. In fact, these lattices are unique up to Lorentz transformations, but the Lorentz transformations modify the heterotic spectrum. However, covariant lattices that lead to chiral spectra have a rigid right-moving sector that forbids continuous Lorentz transformations.

The rigidity of the right-moving part of the lattice discretizes the number of solutions, which is in fact finite for a given world-sheet supersymmetry realization. A very crude attempt to estimate the number of solutions was made in [5], and works as follows. One can map the right-moving bosons to a definite set of 66 left-moving bosons, while preserving modular invariance. This brings us into the realm of even self-dual *Euclidean* lattices, for which powerful classification theorems exist.

Such lattices exist only in dimensions that are a multiple of eight, and have been enumerated for dimensions $8, 16$ and 24, with respectively 1,2 and 24 solutions (in 8 dimensions the solution is the root lattice of E_8, in 16 dimensions they are $E_8 \oplus E_8$ and the root lattice of D_{16} plus a spinor weight lattice, and in 24 dimensions the solutions were enumerated in [73]).

There exists a remarkable formula (the "Siegel mass formula") which gives information about the total number of distinct lattices Λ of dimension $8k$ in terms of :

$$\sum_{\Lambda} g(\Lambda)^{-1} = \frac{1}{8k} B_{4k} \prod_{j=1}^{4k-1} \frac{B_{2j}}{4j} \tag{7}$$

Here $g(\Lambda)$ is the order of the automorphism group of the lattice Λ and B_{2j} are the Bernoulli numbers. Since the automorphisms include the reflection symmetry, $g(\Lambda) \geq 2$. If we assume that the lattice of maximal symmetry is D_{8k} (the root lattice plus a spinor, which is a canonical way to get an even self-dual lattice)) we have a plausible guess for the upper limit of $g(\Lambda)$ as well, namely the size of the Weyl group of D_{8k}, $2^{8k-1}(8k)!$. This assumption is incorrect for $k = 1$, where the only lattice is E_8, and $k = 2$, where the lattice $E_8 \times E_8$ wins against D_{16}, but for $k = 3$ and larger the Weyl group of D_{8k} is larger than the automorphism group of the lattice $(E_8)^k$. For $k = 3$ the assumption has been checked in [74] for all 24 Niemeier lattices. Making this assumption we get

$$\frac{1}{4k} B_{4k} \prod_{j=1}^{4k-1} \frac{B_{2j}}{4j} < N_{8k} < 2^{8k-1}(8k-1)!\; B_{4k} \prod_{j=1}^{4k-1} \frac{B_{2j}}{4j} \tag{8}$$

which for $k = 11$ gives $10^{930} < N_{88} < 10^{1090}$ (in [75] this number was estimated rather inaccurately as 10^{1500}; all numbers quoted here are based on an exact computation).

From a list of all N_{88} lattices one could read off all the free bosonic CFTs with the world-sheet supersymmetry realization discussed above. In particular, this shows that the total number is finite. However, there is a very restrictive subsidiary constraint due to the fact that 66 of the 88 bosons were obtained from the right moving sector. Those bosons must have their momenta on a $D_3 \times (D_7)^9$ lattice and satisfy an additional constraint inherited from world sheet supersymmetry, the triplet constraint. Perhaps a more reasonable estimate is to view this as a lattice with 32 orthogonal building blocks, $D_3 \times (D_7)^9 \times (D_1)^{22}$, which should be combinatorically similar to $(D_1)^{32}$ then the relevant number would be N_{32}, which lies between 8×10^7 and 2.4×10^{51}. But unlike N_{88}, N_{32} is not a strict limit, and furthermore is still subject to the triplet constraint.

All of this can be done explicitly for 10 dimensional strings. Then one needs the lattices of dimension 24, and eight of the 24 lattices satisfy the subsidiary constraints for ten-dimensional strings [75], namely the presence of a D_8 factor.

2.6. *Unexplored landscapes: Meromorphic CFTs.*

The concept of chiral conformal field theories and even self-dual lattices can be generalized to interacting theories, the so-called *meromorphic* conformal field theories [76]. These can only exist if the central charge c (the generalization of the lattice dimension to CFT) is a multiple of 8. For $c = 8$ and $c = 16$ these meromorphic CFTs are just chiral bosons on even self-dual lattices, but for $c = 24$ there 71 CFT's are conjectured [77] to exist including the 24 Niemeier lattices (most of them have indeed been constructed). Gauge symmetries in the vast majority of the heterotic strings in the literature (for exceptions see for example [78]) are mathematically described in terms of affine Lie algebras, a kind of string generalization of simple Lie-algebras, whose representations are characterized by a Lie-algebra highest weight and an additional integer parameter k called the *level*. In the free boson theories the only representations one encounters have $k = 1$, and the total rank equals the number of compactified bosons in the left-moving sector, 22 for four-dimensional strings, and 24 for Niemeier lattices. All even self-dual lattices are direct sums of level 1 affine algebras plus a number of abelian factors (U(1)'s), which we will call the gauge group of the theory. In meromorphic CFT's the restriction to level one is removed. The list of 71 meromorphic CFTs contains 70 cases with a gauge group whose total central charge is 24, plus one that has no gauge group at all, the "monster module". Just one of these yields an additional ten-dimensional string theory with tachyons and an E_8 realized as an affine Lie algebra at level 2. This solution was already known [79], and was obtained using free fermions.

The importance of the meromorphic CFT approach is that it gives a complete classification of all solutions without assuming a particular construction method. In four dimensions the same method can be used. For example, from a list of meromorphic CFTs with $c = 88$ all four-dimensional string theories with a given realization of world-sheet supersymmetry (namely the same one used above) can be obtained, independent of the construction method. Unfortunately next to nothing is known about meromorphic CFTs for $c \geq 32$. It is not known if, like lattices, they are finite in number. Their gauge groups can have central charges that are not necessarily 0 or the total central charge of the meromorphic CFT. It is not known if the gauge groups are typically large or small. There is an entire landscape here that is totally unexplored, but hard to access.

So far this method of mapping a heterotic theory to a meromorphic CFT has only been applied to a world-sheet supersymmetry realization using the triplet constraint. But this can be generalized to other realizations of world-sheet supersymmetry, including perhaps the ones discussed in the next section.

The point we are trying to make here is that despite many decades of work, we are probably still only able to see the tip of a huge iceberg.

2.7. *Gepner Models.*

In 1987 world-sheet constructions were extended further by the use of interacting rather than free two-dimensional conformal field theories [80]. 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 huge, but in practice only a very limited set is available, namely the "minimal models" with central charge $c = 3k/(k+2)$, for $k = 1 \ldots \infty$. There are 168 distinct ways of adding these numbers to 9, so that only a few members of the infinite set are actually used.

With the constraints of superconformal invariance solved, one now has to deal with modular invariance. In exact CFT constructions the partition function takes the form

$$P(\tau, \bar{\tau}) = \sum_{ij} \chi_i(\tau) M_{ij} \bar{\chi}_j(\bar{\tau}) \quad (9)$$

where χ_i are characters of the Virasoro algebra, traces over the entire Hilbert space built on the ground state labeled i by the action of the Virasoro generators L_n:

$$\chi_i(\tau) = \mathrm{Tr} e^{2\pi i \tau (L_0 - c/24)} \quad (10)$$

The multiplicity matrix M indicates how often the ground states $|i\rangle|j\rangle$ occurs in the spectrum. Its entries are non-negative integers, and it is severely constrained by modular invariance. Note that in (9) we allowed for the possibility that the left- and right-moving modes have a different symmetry (a different extension of superconformal symmetry) with different sets of characters χ and $\bar{\chi}$. But then the conditions for modular invariance are very hard to solve. They can be trivially solved if the left and right algebras are the same. Then modular invariance demands that M must commute with the matrices S and T that represent the action of the modular transformations $\tau \to -1/\tau$ and $\tau \to \tau + 1$ on the characters. This has always at least one solution, $M_{ij} = \delta_{ij}$.

However, assuming identical left and right algebras is contrary to the basic idea of the heterotic string. Instead Gepner model building focuses on a subset, namely those spectra that can be obtained from a symmetric type-II spectrum by mapping one of the fermionic sectors to a bosonic sector. For this purpose we can use the same bosonic string map discussed above. This results in a very special and very limited subset of the possible bosonic sectors.

Using the discrete symmetries of the building blocks, for each of the 168 tensor combinations, a number of distinct modular invariant partition functions can be constructed, for a grand total of about five thousand [81]. Each of them gives a string spectrum with a gauge group $E_6 \times E_8$ (or occasionally an extension of E_6 to E_7 or E_8) with massless chiral matter in the representations (27) and ($\overline{27}$) of E_6, exactly like the Calabi-Yau compactifications discussed above.

Indeed, it was understood not long thereafter that there is a close relationship between these "Gepner models" and geometric compactifications on Calabi-Yau manifolds. Exact correspondences between their spectra were found, including the number of singlets. This led to the conjecture that Gepner Models are Calabi-Yau compactifications in a special point of moduli space. Evidence was provided by a conjectured relation between $N = 2$ minimal models and critical points of Landau-Ginzburg models [82, 83].

Getting the right number of families in this class of models has been challenging, since this number turns out to be quantized in units of six or four in nearly all cases that were studied initially. The only exception is a class studied in [84].

2.8. *Dualities, M-theory and F-theory*

In general, four-dimensional string theories are related to others by maps like S-duality [85] (strong-weak dualities due to inversion of coupling constants), T-duality (transformations involving inversion of compactification radii) and combinations thereof. This suggests a connected "landscape" of four-dimensional strings. However, this is still largely based on anecdotal evidence. A complete picture of the four-dimensional string landscape is still very far away.

In ten (and eleven) dimensions, the picture is better understood. Under

S-duality, Type-IIA string theory is mapped to an 11-dimensional theory compactified on a circle [86, 87]. The 11-dimensional theory is not a string theory. It is called "M-theory". Its field theory limit is $D = 11$ supergravity.

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 [88], the circle with two halves identified. This is sometimes called "heterotic M-theory".

Strong coupling duality maps type-IIB strings to themselves [89]. 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, 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 twelve-dimensional theory [90]. 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.

2.9. *New Directions in Heterotic strings*

2.9.1. *New embeddings.*

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, condition (5) is then not automatically satisfied, but in heterotic M-theory one may get extra contributions from heterotic five branes [91, 92].

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 [93], $U(1)^4$ bundles [94] and $SU(N) \times U(1)$ bundles [95], 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 [96] with the exact supersymmetric Standard Model spectrum, without even any vector-like matter (but with extra gauge groups and the usual large numbers of singlets). However, the gauge group contains extra $U(1)$'s and an E_8 factor, and large numbers of gauge singlets, including unstabilized moduli. There can be several Higgs multiplets. To break the GUT groups down to the Standard Model, background gauge fields on suitable Wilson lines are used. For this purpose one needs a manifold with a freely acting (*i.e.* no point on the manifold is fixed by the action) discrete symmetry. One then identifies points on the manifold related by this symmetry and adds a background gauge field on a closed cycle on the quotient manifold (a Wilson line).

2.9.2. *The Heterotic Mini-landscape.*

The Heterotic Mini-landscape is a class of orbifold compactifications on a torus $T^6/\mathbf{Z}_6$ cleverly constructed so that the heterotic gauge group $E_8 \times E_8$ is broken down to different subgroups in different fixed points, such as $SO(10)$, $SU(4)^2$ and $SU(6) \times SU(2)$. This leads to the notion of *local unification* [97–99]. The Standard Model gauge group is the intersection of the various "local" gauge realized at the fixed points. Fields that are localized near the fixed points must respect its symmetry, and hence be in complete multiplets of that group. Unlike field theory GUTs, these models have no limit where $SO(10)$ is an exact global symmetry. In this way one can make sure that matter families are in complete spinor representations of $SO(10)$, while Higgs bosons need not be in complete representations of $SO(10)$, avoiding the notorious doublet splitting problem of GUTs. The number of 3-family models in this part of the landscape is of order a few hundred, and there is an extensive body of work on their phenomenological successes and problems, see for example [100, 101] and references therein.

2.9.3. *Heterotic Gepner Models*

As explained above, the original Gepner models are limited in scope by the requirement that the left and right algebras should be the same. There is no such limitation in free CFT constructions, but they are limited in being non-interacting in two dimensions. What we would like to have is asymmetric, interacting CFT constructions. Examples in this class have been obtained using a method called "heterotic weight lifting" [102]. 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. This opens up an entirely new area of the heterotic string landscape. It turns out that the difficulty in getting three families now disappears.

2.10. *Orientifolds and Intersecting Branes*

The Standard Model comes out remarkably easily from the simplest heterotic strings. But that is by no means the only way. One may also get gauge groups in string theory 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 endpoints 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 endpoints are only observable on the brane, and give rise to fermions and gauge interactions.

2.10.1. *Chan-Paton groups.*

The possibility of getting gauge theories and matter from branes sparked another direction of research with the goal of getting the Standard Model from open string theories. To get towards the Standard Model, one starts with type-II string theory, and compactifies six dimensions on a manifold. This compactified manifold may have a large radius, as in the brane world scenario, but this is optional. 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. Each such brane can give rise to a gauge group, called a Chan-Paton gauge group, which can be $U(N)$, $Sp(N)$ or $O(N)$ [103]. 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 bi-fundamental representation of those gauge groups.

Remarkably, a Standard Model family has precisely the right structure to be realized in this manner. The first example is the so-called "Madrid model" [104]. 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 stack (the stacks are labeled **a** ... **d**)

$$Y = \frac{1}{6}Q_{\mathbf{a}} + \frac{1}{2}Q_{\mathbf{c}} - \frac{1}{2}Q_{\mathbf{d}}.$$

This configuration is depicted in Fig. 1(a).

2.10.2. *The three main classes.*

There are other ways of getting the Standard Model. 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 [105]

$$Y = (x - \frac{1}{3})Q_{\mathbf{a}} + (x - \frac{1}{2})Q_{\mathbf{b}} + xQ_{\mathbf{c}} + (x-1)Q_{\mathbf{d}}. \tag{11}$$

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 [106] and flipped $SU(5)$ [107] models. The value $x = 0$ gives rise to classic $SU(5)$ GUTs [108]. To get Standard Model families in this case one needs chiral anti-symmetric rank-2 tensors, which originate from open strings with both their endpoints on the same brane. The simplest example is shown in Fig. 1(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 $(\overline{5})$ representation of $SU(5)$.

Other values of x can only occur 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. 1(c).

Note that it was assumed here that there are at most four branes participating in the Standard Model. If one relaxes that condition, the number of possibilities is unlimited.

2.10.3. *Orientifolds.*

An important issue in open string model building is the cancellation of tadpoles of the disk diagram. These lead to divergences and can lead to chiral anomalies. These tadpoles can sometimes be canceled by adding another object to the theory, called an orientifold plane. In fact, the usual procedure is to start with an oriented type-II string, and consider an involution

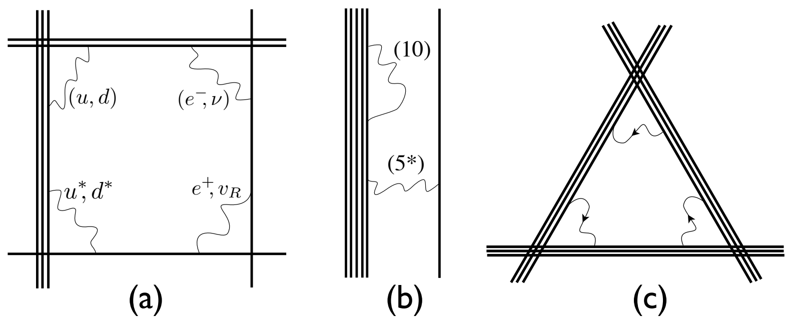

Fig. 1. Brane configurations: (a) the Madrid model, (b) SU(5) GUTs and (c) Trinification.

of the world-sheet that reverses its orientation. Then one allows strings to close up to that involution. In terms of world-sheet topology, this amounts to adding surfaces with the topology of a Klein bottle. The combination of torus and Klein-bottle diagram acts like a projection on the closed string theory, removing some of its states. In most cases, removing states from string theory comes at a price: other states must be added to compensate what was removed. In this case, this rôle is played by open strings. These ideas were pioneered in [109, 110]

Orientifold model building has been very actively pursued during the first decade of this century (see [111] for a review).

2.10.4. *Anomalies, axions and massive abelian vector bosons.*

Canceling all tadpoles between the disk and crosscap diagram removes most anomalies, but some factorized anomalies remain which can then be canceled by the Green-Schwarz mechanism [112] involving tree-level diagrams with exchange of axions. In contrast to perturbative heterotic strings the anomaly factorizes in terms of several factors. These anomalies are then canceled by a Green-Schwarz mechanism involving multiple axions, which are available in the Ramond-Ramond sector of the closed theory.

In four dimensions, a factorized anomaly always involves a $U(1)$. The corresponding $U(1)$ vector bosons acquire a mass by "eating" the axion, which provides the missing longitudinal mode. String theory will always remove anomalous symmetries in this manner, but it turns out that this can happen for non-anomalous $U(1)'s$ as well. This can be traced back to anomalies in six dimensions [113].

2.10.5. *Boundary RCFT constructions.*

Just as in the heterotic string, one can construct orientifold spectra using purely geometric methods, orbifold methods or world-sheet constructions. Most work in the literature uses the second approach.

World-sheet approaches use boundary CFT: conformal field theory on surfaces with boundaries and crosscaps. This requires an extension of the closed string Hilbert space with "states" (in fact not normalizable, and hence not in the closed string Hilbert space) that describe closed strings near a boundary, or in the presence of orientation reversal. An extensive formalism for computing boundary and crosscap states in (rational) CFT was developed in the last decade of last century, starting with [114], developed further by several groups [115–119], culminating in a simple and general formula [120]. For an extensive review of this field see [121]. This was applied to orientifolds of Gepner models [122], and led to a huge (of order 200.000) number of distinct string spectra that match the chiral Standard Model. This set provides an extensive scan over the orientifold landscape.

These spectra are exact in perturbative string theory and not only the massless but also all massive states are known explicitly. There are no chiral exotics, but in general there are large numbers of the ubiquitous vector-like states that plague almost all exact string spectra. All tadpoles are canceled, but in most cases this requires hidden sectors. However, there are a few cases where all tadpoles cancel entirely among the Standard Model branes (hence no hidden sector is present) and furthermore the superfluous $B - L$ vector bosons acquires a mass from axion mixing. These spectra have a gauge group which is exactly $SU(3) \times SU(2) \times U(1)$ (there are a few additional vector bosons from the closed sector, but the perturbative spectrum contains no matter that is charged under these bosons; this is the same as in the type IIA string, which contains a vector boson that only couples to non-perturbative states, D0-branes).

2.11. *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 in [123], who found examples with $\mathbb{Z}_3$ singularities. See *e.g.* [124, 125] for other kinds of singularities.

One extreme possibility is to decouple gravity by sending the compactification radius to infinity. In heterotic string models both gravity and gauge interactions originate from closed string exchange, and such a decoupling limit would not make sense.

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.

Both points of view are to some extent a return to the "old days" of quantum field theory. On the one hand, the techniques of branes and higher dimensions are used to enrich old ideas in GUT model building; on the other hand, string theory is treated as a "framework", analogous to quantum field theory, where gauge groups, representations and couplings are input rather than output.

Decoupling of gravity is an important element in recent work on F-theory GUTs [126–128] obtained by compactifying F-theory on elliptically fibered Calabi-Yau fourfolds. This allows the construction of models that may be thought of as non-perturbative realizations of the orientifold $SU(5)$ GUT models depicted in Fig. 1(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. See the reviews [129–132] for further details.

An example in the second category is recent work in the area of M-theory compactifications [58]. Getting chiral N=1 supersymmetric spectra in M-theory requires compactification on a seven dimensional manifold with G_2 holonomy [133], 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 [46], originally invented for the purpose of moduli stabilization (see section 1.2.5). Here both kinds of decoupling limits have been discussed, and there have also been steps towards putting both parts together [134]. 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 [135, 136].

2.12. *Non-supersymmetric strings*

Although the vast majority of the literature on string constructions 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 non-supersymmetric ones. In fact, it is easier, because space-time supersymmetry is an additional constraint. 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 in [137, 138]. These authors found a heterotic string theory with a $SO(16) \times SO(16)$ gauge group, the only tachyon-free non-supersymmetric theory in ten dimensions, out of a total of seven. Four-dimensional non-supersymmetric strings were already constructed shortly thereafter [5, 79]. Non-supersymmetric strings can also be constructed using orientifold methods, see for example [139–143].

Non-supersymmetric strings can have a vacuum energy Λ of either sign. See for example [144] for a distribution of values of the vacuum energy for a class of heterotic strings. There also exist examples where Λ vanishes exactly to all orders in perturbation theory [145] but probably this feature does not hold beyond perturbation theory [146].

One might think that in the absence of any evidence for low energy supersymmetry, and because of the evidence in favor of an accelerated expansion of the universe, non-supersymmetric strings with a positive cosmological constant are a better candidate for describing our universe than the much more frequently studied supersymmetric ones. But the absence of supersymmetry is a serious threat for the stability of these theories, even in the absence of tachyons in the perturbative spectrum.

3. The Standard Model in the Landscape

In this chapter we will 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. We focus on questions related to susy-GUTs, where the stress between symmetry and landscape anarchy has been building up in the last few years.

3.1. *The Gauge Group*

It is by now abundantly clear that string theory can reproduce the discrete structure of the Standard Model: the gauge group $SU(3) \times SU(2) \times U(1)$ with chiral fermion representations. Indeed, the gauge group is easy to get in many construction methods: Heterotic Calabi-Yau and orbifold compactifications with GUT symmetries broken by Wilson lines, orientifold models with various kinds of intersection branes, strings at singularities, free fermion and free boson constructions, heterotic Gepner models and Gepner orientifolds, higher level heterotic strings with symmetry breaking by the standard adjoint Higgs, F-theory with Y-flux, and others. See section 2 for references to all this work.

But this work also demonstrates very clearly that there is nothing special about the Standard Model from the top down perspective, except that it is rather simple. Many other gauge theories and representations are possible as well, although both are limited in size. Unlike quantum field theory, string theory only allows small representations. Furthermore, the size of the gauge group tends to be limited by the conformal anomaly in closed strings, or dilaton tadpole cancellation in open strings. This is not a theorem: there are remarkable exceptions with very large gauge groups, but it seems plausible that these are far out in the tail of landscape distributions, and hence statistically very rare. The Standard Model gauge group does have one remarkable feature, namely that it fits beautifully in a Grand Unified gauge theory. We will discuss that below in section 3.2.3.

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. Rigorous evidence of such statements is unlikely to emerge soon, as it requires to work out the full nuclear and atomic spectrum as well as astrophysics, nucleosynthesis and baryogenesis for alternatives of the Standard Model. But given what we know, the presence of electromagnetic and strong interactions are well-motivated anthropic assumptions.

The weak interactions also play a crucial rôle in our universe, but perhaps not in every habitable one. See [147] for a detailed discussion of a “weakless” universe that may yield acceptable nuclear and atomic physics even though the weak scale is pushed towards the Planck scale. Perhaps the main rôle of the weak interactions is to provide chirality to fermions, protecting them from getting a large mass. For this to be true, one has to be able to argue that in the string theory landscape it is easier to have a single light scalar than several light fermions. See [148] for a discussion along these lines.

3.2. *Family structure and Charge Quantization*

3.2.1. *Quantum Field Theory*

A Standard Model family is described by the following reducible $SU(3) \times SU(2) \times U(1)$ representation:

$$(3,2,\frac{1}{6}) + (3^*,1,\frac{1}{3}) + (3^*,1,-\frac{2}{3}) + (1,2,-\frac{1}{2}) + (1,1,1) \qquad (12)$$

where we ignore singlets. This occurs three times, and in addition to this there is a Higgs field in the representation $(1,2,-\frac{1}{2})$. At first, this looks arbitrary and unintuitive, but on closer examination some structure becomes apparent. For example, the three entries of each irreducible term multiply to an integer. This fact implies that all color singlet bound states of the broken $SU(3) \times U(1)$ spectrum have charges that are integer multiples of the electron charge. This fact is not explained in the Standard Model. An arbitrary representation has the form (R_3, R_2, q) where R_n is an $SU(n)$ representation and q a real number. But the observed representations satisfy the rule

$$\frac{t_3}{3} + \frac{t_2}{2} + \frac{1}{6} = 0 \text{ mod } 1, \qquad (13)$$

where t_3 is the triality of the $SU(3)$ representation and t_2 the duality of $SU(2)$, twice the spin modulo integers. Group-theoretically this means that all observed representations are in fact representations of the group $S(U(3) \times U(2))$, which has the same Lie algebra as $SU(3) \times SU(2) \times U(1)$.

It seems clear that the family structure is more than just an environmental fact. Some of it is explained by the consistency conditions imposed by anomaly cancellation. This implies that four cubic traces and a linear one must vanish. There is also a global $SU(2)$ anomaly [149] and perhaps one should impose a string-inspired non-abelian $SU(2)$ anomaly [150]. These anomaly cancellation conditions are sufficient to explain charge quantization if one assumes that there is just a single family with the observed $SU(3) \times SU(2)$ content. But quantum field theory offers no reason for these assumptions, and the fact that there are three families ruins the argument anyway.

3.2.2. *String Theory*

String theory makes important contributions towards understanding the family structure. First of all, it limits the choice of representations to only a handful of options. Secondly, as far as anyone knows, string theory always implies absolute charge quantization. By this we mean that the charges are rational numbers, though not necessarily the right ones. And thirdly, string theory provides a rationale for anomaly cancellation that is somewhat more deeply rooted than the ad-hoc rules of quantum field theory.

3.2.3. *Grand Unification*

Usually Grand Unification is invoked to explain Eq. (12). In the context of quantum field theory, this would offer a plausible explanation for the fact that particles fit in $S(U(3) \times U(2))$ representations. There is no good motivation in QFT to consider just $S(U(3) \times U(2))$. From the traditional symmetry-based perspective, assuming that what we see is a broken $SU(5)$ (or larger) gauge theory looks like a natural idea. But it is by no means perfect. It does not explain the family structure, but it just limits the allowed combinations of $SU(3) \times SU(2) \times U(1)$ representations. It does not explain why a breaking to $SU(3) \times SU(2) \times U(1)$ is preferred, and it has difficulties accommodating the Higgs field. In $SU(5)$ the Standard Model Higgs has an $SU(3)$ triplet partner which must remain heavy and cannot have a vacuum expectation value: the doublet triplet splitting problem. The GUT hypothesis would get a lot more credibility if a second non-trivial coincidence is established: the renormalization group convergence of the coupling constants to a single value at the GUT scale. This does not hold if one extrapolates the current low-energy couplings, but it would work if more or less standard supersymmetric partners of all Standard Model particles are discovered at a future run of LHC.

Since many Standard Model realizations in string theory look superficially like GUT theories, one might have expected that all facts mentioned in the foregoing paragraphs are naturally combined to get a satisfactory explanation of family structure and charge quantization. But this has never worked as easily as it should.

3.2.4. *Grand Unification in Heterotic Strings*

The oldest examples studied are compactifications of the heterotic string. 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, which contains $SO(10)$ (and hence $SU(5)$ as a subgroup). In world-sheet constructions this is a consequence of the "bosonic string map" [5] 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. The bosonic string map takes the fermionic sector of a heterotic or type-II string, and maps it to a bosonic sector. The world-sheet fermions ψ^μ transform under the D-dimensional Lorentz group $SO(D-1,1)$. The bosonic string map replaces this by an $SO(D+6)\times E_8$ affine Lie algebra, which manifests itself as a gauge group in space-time. In [5] this trick was used to map the problem of finding modular invariants to the already solved problem of characterizing even self-dual lattices. 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.

With only slight exaggeration one can state that this ideal GUT group, $SO(10)$, emerges uniquely from the heterotic string. All we had to do is specify the space-time dimension, $D = 4$, and apply the bosonic string map, and we get $SO(10)$ for free.

3.2.5. *Fractional Charges in Heterotic Spectra*

A mechanism to break $SO(10)$ to $SU(3)\times SU(2)\times U(1)$ can also be found, but it does not come out automatically. Furthermore, it does not work as nicely as in field theory GUTs, because 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 (13), 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 would like to remove are the extra gauge bosons in $SU(5)$ 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 (13) and hence is fractionally charged [151]. The possible presence of unconfined fractional charges in string spectra was first pointed out in [152] and the implications were discussed further in [153].

The occurrence of fractional charges in heterotic string spectra has been studied systematically for free fermion constructions and for heterotic Gepner models. All these models realize the gauge group in the canonical heterotic way, as a subgroup of $SO(10)$ (which may be further extended to E_6). There is a total of four distinct subgroups that one may encounter within $SO(10)$. These subgroups are further subdivided into several classes, distinguished by the minimal electric charge quantum that occurs in their spectra. These charge quanta are *not* determined by group theory in quantum field theory, but by affine Lie algebras in string theory. This gives a total of eight possibilities, with charge quanta given in curly brackets:

$$\begin{array}{rc} SU(3) \times SU(2) \times U(1) \times U(1) & \{\frac{1}{6}, \frac{1}{3}, \frac{1}{2}\} \\ SU(3) \times SU(2)_L \times SU(2)_R \times U(1) & \{\frac{1}{6}, \frac{1}{3}\} \\ SU(4) \times SU(2)_L \times SU(2)_R & \{\frac{1}{2}\} \end{array}$$

plus $SU(5) \times U(1)$ and $SO(10)$, which automatically yield integer charges. This classification applies to all constructions in the literature where the Standard Model is realized with level 1 affine Lie algebras, with a standard Y charge normalization, embedded via an $SO(10)$ group. The minimal electric charge *must* be realized in the spectrum, but it is in principle possible that fractionally charged particles are vector-like (so that they might become massive under deformations of the theory), have Planck-scale masses or are coupled to an additional interaction that confines them into integer charges, just as QCD does with quarks.

Fractional charges can be avoided by looking for spectra where all these particles have Planckian masses. In [65] a large class of free fermionic theories with Pati-Salam spectra was investigated. These authors did find examples with three families where all fractionally charged particles are at the Planck mass, but only in about 10^{-5} of the chiral spectra. In [102, 154–156] a similar small fraction was seen, but examples were only found for even numbers of families. These authors also compared the total number of spectra with chiral and vector-like fractional charges, and found that about in 5% to 20% of the chiral, non-GUT spectra the fractional charges are massless, but vector-like. They also found some examples of confined fractional charges.

In a substantial fraction of explicitly constructed string vacua the fractionally charged particles are vector-like. If one assumes that in genuine string vacua vector-like particles will always be very massive, this provides a way out. There is a more attractive possibility. 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 [157]), the twisted sector fields are massive, and hence all fractionally charge particles are heavy. This method is commonly used in Calabi-Yau based constructions, *e.g.* [158], 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 mini-landscape an example was found [159]. These authors suggested another rationale for using freely acting symmetries, namely that otherwise the Standard Model Y charge breaks if the orbifold singularities are "blown up".

But even though there are ways out, it is disappointing that charge quantization comes out less easily than it does in field theoretic $SU(5)$ GUTs, without string theory.

3.2.6. *GUT Unification in Brane Models.*

There is another important region in the landscape where $SU(5)$ GUTs can be obtained, namely intersecting brane models. The simplest possibility is to intersect a stack of $U(3)$ branes with a stack of $U(2)$ branes. The entire Standard Model group can be embedded in these two groups, but to get the matter representations one needs not only bi-fundamentals (from strings stretching between the two stacks) and rank-2 anti-symmetric tensors, but also $U(3)$ and $U(2)$ vectors. They would come from endpoints of an open string, but then additional neutral branes are needed for the other endpoint to end on. The resulting configuration is exactly as shown in fig. 5b, but with the $U(5)$ stack split in $U(3)$ and $U(2)$.

It has been known for a long time already that $SU(5)$ GUTs can be obtained from configurations like 5b [160]. These authors noticed however that solutions to the tadpole conditions do not generically lead to the expected anomaly free representation $(\overline{5}) + (10)$, but to more complicated solutions involving the symmetric tensor (15). One can also start with a split stack, but then more input would seem to be required. This includes not only the brane configuration, but also the exact embedding of the $U(1)$ group in $U(3) \times U(2)$ and the particle assignment. In other words, a set of allowed massless open string states is hand-picked to match the Standard Model spectrum. With all these assumptions, one can indeed find numerous solutions [105].

Interestingly, these $U(3) \times U(2)$ intersecting brane models do provide a convincing rationale for the gauge group $U(3) \times U(2)$, which is not easily found in field theory models. But how do we get the restriction to $S(U(3) \times U(2))$ to give the correct charge conjugation, and how do we justify the choice of representations in a Standard Model family? For example, if even in the $U(5)$ limit symmetric tensors are hard to avoid, with split stacks there are even more possibilities. Note that by "split stacks" we do not necessarily mean a $U(5)$ stack with branes move apart. There are more general possibilities, where the $U(3)$ and $U(2)$ stacks occupy unrelated cycles on a compactification manifold.

It turns out that there is an extremely simple answer to this question if one allows a mild anthropic condition [148]. It turns out that for all anomaly free matter configurations, and for all possible $U(1)$ embeddings the electromagnetic $U(1)$ is chiral after Higgs symmetry breaking, or there remain massless charged leptons in the spectrum, with one exception: the Standard Model, with a number of families of the form (12). The motivation for these conditions is that a chiral $U(1)$ will be broken by the color group, and massless charged leptons can be pair-produced without limit, so that the entire universe becomes an opaque plasma of lepton-antilepton pairs [161]. Although we cannot prove that life is impossible without a massless photon or in an opaque plasma, the circumstances for our kind of life – indeed, any kind of life based on atomic physics – are so adverse that one can certainly defend this as a well-motivated anthropic assumption.

One can say that the assumption of symmetry at high energies has been traded for these anthropic assumptions, and remarkably, the latter are more powerful. Even the Higgs choice does not have to be put in, but is determined. Indeed, unless we see evidence for gauge coupling unification because of new matter bending the coupling constant curves, full $SU(5)$ unification has nothing useful to offer.

The argument can be extended to more general $U(M) \times U(N)$ stacks with a few additional assumptions. The group $SU(M)$ is assumed not to be broken by the Higgs, and to be a strong interaction group that is asymptotically free, and dominates over the other gauge interactions. The Higgs is assumed to give mass to all charge fermions, not just the leptons (above we just required the quarks to become non-chiral, not necessarily massive). These simple conditions have a few solutions: the Standard Model, a series of models without leptons, a few cases with $SU(2)$ "color" and no conserved baryon number, and a few models with just electromagnetism and no strong interactions. Finally, there is an $SU(4) \times U(1)$ model, broken by a Higgs boson to $SU(3)_{\rm color} \times U(1)_{\rm em}$. This uses the alternative breaking pattern of $SU(5)$ to $SU(4)\times U(1)$ instead of the Standard Model. However, it has baryon-number violating weak interactions that are probably fatal. All other alternatives appear to have fatal problems for life as well, and the Standard Model really stands out as the optimal, and probably unique anthropic solution within this class of brane models.

In this class, $SU(5)$ symmetry is not needed to explain charge quantization, nor the structure of a family. In the heterotic case, it does not work as an explanation of charge quantization. Similar remarks may apply to F-theory GUTs, where the GUT group is present by choice, and not because it is required. There are some corners in the landscape where $SU(5)$ really works as in field theory. The only example we know are heterotic string theories with GUT group with affine level larger than 1 [162].

3.3. *The Number of Families*

The string theory landscape does not offer, according to our current understanding, an answer to the question why we observe three families. Although early constructions, for example the first Gepner models, had some difficulties getting three families (the number was predominantly a multiple of four or six [81, 163]), further work showed that the number of families in heterotic strings has a slow exponential fall-off, with the number three appearing not much less frequently than 2 (see *e.g.* [154, 164]). In orientifold models the fall off seems to be much faster [122].

There is no convincing anthropic argument for three families. We are built out of just one family. The most often mentioned feature is that three families are needed for CP violation in the CKM matrix, which in its turn might be required for baryogenesis, which is obviously anthropically relevant. But CP violation in the CKM matrix is not believed to be sufficient. The top quark plays an interesting rôle in the running of couplings, and the stability of our vacuum under tunneling depends in a remarkable way on both the top and the Higgs mass. Perhaps this points to an important rôle for the third family, but then why does the second family exist? The s-quark is not completely irrelevant in QCD, and the muon affects biological mutations, but neither of these arguments provides a convincing reason.

3.4. *Quark and Lepton Mass Hierarchies*

An area where there is an interesting rivalry between symmetry-based ideas and landscape anarchy is the understanding of quark and lepton mass hierarchies. Already long ago Grand Unified theories predicted interesting mass relations for the second and third family fermions. Certain string approaches, such as the heterotic mini landscape point to top Yukawa-gauge unification. Using F-theory compactifications, part of the structure of the observed masses and mixing angles can be nicely explained.

But on the other hand, one can also get a long way towards the correct distribution of quark and lepton masses by assuming statistical distributions of Yukawa couplings. Clearly, flat distributions will not work, because the quark and lepton masses have an unmistakable hierarchical structure, and the mixing angles are small, and seem to get smaller as the hierarchies get larger. However, scale invariant distributions (with a cut-off fit to the data) [165, 166] or Gaussian overlaps [167, 168] work rather well. They even lead generically to small mixing angles, but it is not automatic in all cases that the mass eigenvalues of the up and down quarks are ordered correctly. Then in an alternative universe the three charge $\frac{2}{3}$ quarks would predominantly couple to a different permutation of the three charge $-\frac{1}{3}$ quarks, and only in one-sixth of all universes with $SU(3) \times SU(2) \times U(1)$ the Standard Model ordering would be observed. One may view this either as a minor statistical problem, or an indication that something essential is missing. For a more detailed discussion and references see [1]

This issue is far from settled. Without prior knowledge of the answer, none of the aforementioned ideas would have given an accurate description of the quark and lepton spectrum, even if the existence of three families is provided as information, and even if the anthropic constraints on the light fermions are used. Furthermore it is plausible that the top quark mass, in combination with weak symmetry breaking and the Higgs mass, plays an important rôle that remains to be elucidated.

The smallness of neutrino masses has a well-known natural explanation, the see-saw mechanism. This is so natural and requires so few changes to the Standard Model that it is generally seen as the most plausible kind of beyond the Standard Model physics. Indeed, all that is required are additional singlets (right-handed neutrinos) having their natural mass. This mass is not proportional the the Higgs vev, and hence one would expect it to be large. How large, and how it is distributed depends on several assumptions, and there are also several anthropic issues related to neutrino masses. See [1] for more details.

3.5. *The Scales of the Standard Model*

The 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 $M_{\rm weak}$ determines the masses of the quarks and leptons. The proton mass owes less than 1% of its mass to the up and down quarks. However, the smallness of both scales with respect to the Planck scale has an important environmental impact.

In quantum field theory, the strong scale is said to be determined by "dimensional transmutation", which turns a dimensionless coupling constant into a scale. The appearance of a scale from a dimensionless theory is due to quantum loop effects. The relation is:

$$\Lambda_{QCD} = Q\ e^{-1/(2\beta_0\alpha_s(Q^2))}, \tag{14}$$

where α_s is the strong coupling constant, Q the scale where it is defined and β_0 the leading coefficient of the β-function. This relation does not determine the scale, since $\alpha_s(Q^2)$ is input, but from a landscape perspective it affects the distribution of scales in such a way that a large hierarchy is easy to obtain. How easy depends on the distribution of α_s in a fundamental theory, but what we know about the landscape suggests that this argument is valid.

This then leaves the weak scale $M_{\rm weak}$ to worry about. In contrast to Λ_{QCD}, the μ^2 parameter in the Standard Model Higgs potential receives quadratic quantum corrections from higher scales. This worry has been the focus of decades of work on natural solutions to the hierarchy problem. All of these solutions lead to predictions of new particles near the weak scale, although some can be pushed a few orders of magnitude higher. So far no such particles have been found. One can take the point of view that we simply have to be more patient. After all, the top quark and the Higgs boson also required a lot of patience. But the current situation clearly demands a reassessment of the arguments. During the year 2013, after the existence of the Higgs boson was convincingly established, there has been a lot of discussion about this. The different lines of argument are roughly as follows.

- There is no hierarchy. Some people argue that the hierarchy should be viewed as a misconception in quantum field theory. They point to the fact that unlike logarithmic corrections, quadratic corrections are not present in every regularization scheme, and argue that they may not really be physical; see *e.g.* [169–172]. However, these papers deal with just the Standard Model in quantum field theory, and ignore potential BSM physics (such as GUTs) and certain BSM physics (gravity). There is general agreement that the existence of new massive particles beyond the weak scale implies a quadratic hierarchy for μ^2 [173]. One can try to get around this by following a minimalistic approach in which new physics beyond the weak scale is completely avoided, as in [174]. But then one still has to deal with gravity. It is possible that the naive notion that since there is a Planck scale, there must be new physics at that scale is wrong. Perhaps nature is fundamentally scale invariant (as suggested for example in [175–177]; see however [178] for criticism). However, this escape route is closed off in string theory, which clearly predicts not just new physics but concrete new particles at the Planck scale. This still leaves the next option:

- The Planck scale is at the weak scale. This possibility exists if there are large extra dimensions at a length scale far larger than typical length scale of particle physics, even as large as 0.1 mm. Such a scenario was proposed in [179], but it predicts observable gravitational physics, perhaps even black holes, at the weak scale. This idea is put under severe stress by the latest LHC results.
- The weak/GUT or weak/Planck scale hierarchy is real and is generated by new dynamics (technicolor, compositeness) at the weak scale. In this approach the weak scale is generated analogously to the strong scale by dimensional transmutation. Also in this case absence of new physics at the LHC is a serious problem. Furthermore it is hard to build credible examples where the entire Standard Model (including quark and lepton masses) is reproduced.
- The weak/GUT or weak/Planck scale hierarchy is real and is "protected" by low energy supersymmetry. Since low energy supersymmetry does not determine the weak scale, this still requires an additional mechanism to generate it, and often this is related to dimensional transmutation (*e.g.* gaugino condensation). This idea is facing similar observational problems as the previous two.
- The hierarchy is mainly anthropic (see below).

From the string theory perspective, the second option is very attractive since it would give direct access to gravitational physics; from the landscape perspective the question is how often large extra dimensions occur in conjunction with a Standard Model with a weak scale as we observe. The third option is unattractive from the landscape point of view since it would imply that our vision of fundamental physics would be blurred by an additional "onion shell" which we have to peel off first (since technicolor gained some renewed interest due to an interpretation as a dual of a Randal-Sundrum model [180], some people view this as an appealing prospect).

At this moment, low energy supersymmetry is usually considered to be the most attractive of options 2, 3 and 4. It is a well-motivated idea because it not only controls the hierarchy, but also provides additional particles needed for the convergence of the running coupling constants as well as dark matter candidates. Furthermore string theory may come with supersymmetry built in. The latter statement holds for all string theories with a controlled perturbative expansion. Without supersymmetry, quantum corrections diverge beyond one loop. Supersymmetry may be required for a fundamental understanding of quantum gravity, or it may just be a calculational tool we need because our current understanding is too primitive.

Since, despite good arguments in its favor, no sign of low energy supersymmetry has been seen so far, we have to ask which mistakes may have been made in arriving at the overly optimistic expectations. We will mention three here, labeled as "anthropic", "landscape" and "string theory".

Anthropic. The idea that the hierarchy might be anthropic was not even mentioned during three of the four decades of discussion of naturalness, and during the past decade was mentioned only to ridicule it. And yet it is true. It is true in the sense that in an ensemble of theories with a range of gauge hierarchies, (intelligent) life can exist only for theories with a large hierarchy. The simplest argument is based on the weakness of gravity. The largest structures that can exist without being crushed into a black hole have N building blocks, where $N = (M_{\rm Planck}/m)^3$ and m is the mass of the building block (the proton mass in our universe). This already requires a hierarchy of nine orders of magnitude for something with the complexity of a human brain to exist. More detailed arguments with stronger (and more debatable) assumptions pin the weak scale down with a precision of less than ten per cent (see [1] for more details and references). But even if the entire hierarchy of 18 orders of magnitude is fully understood anthropically, this still does not imply that the hierarchy is explained by anthropic arguments alone, as suggested in [181] and [182].

Landscape. One may still ask the question if the underlying physics helps in getting the anthropically required hierarchy. In a landscape of variants of the Standard Model with a fixed Planck scale and a full range of values of μ^2, flatly distributed, the chance of getting an anthropically acceptable theory is about 10^{-34}. That is the fraction of theories with a μ^2 parameter that is small enough, if the parameter μ^2 can be though of as a sum of uncorrelated terms of order $(M_{\rm Planck})^2$. So with a large enough landscape, one could consider the problem solved. But this is not true if the landscape contains other vacua where the statistical penalty of 10^{-34} does not have to be paid. Technically natural theories escape this penalty. However, then the question arises what the cost is in getting technical naturalness.

These questions can only be addressed in a context where the relative abundance of theories can be compared. In particular, it makes no sense in quantum field theory. Indeed, the very concept of technical naturalness is at best a poorly defined intuitive notion without the context of an ensemble of theories, a landscape. The string theory landscape certainly contains supersymmetric theories; indeed, these are the only ones under computational control. To predict if they dominate, one would have to estimate the ratio of the number of supersymmetric and non-supersymmetric theories, given the hierarchy. Supersymmetric theories start with a huge advantage of a factor 10^{34}, but they may still lose by being far less abundant, or because they have to satisfy additional anthropic constraints to avoid fast, catastrophic proton decay. Unfortunately, answering this question is not feasible at present. What has been tried is comparing supersymmetric theories with different supersymmetry breaking scales. Very roughly (for many more details see [27]) one would like to compute

$$P\left(\frac{M_{\rm weak}}{M_{\rm susy}}\right) P\left(\frac{M_{\rm susy}}{M_{\rm Planck}}\right) , \qquad (15)$$

the probability for getting a weak scale given a supersymmetry breaking scale times the probability for getting a certain supersymmetry breaking scale given the Planck scale. The first factor is $(M_{\rm weak}/M_{\rm susy})^2$ by the usual naturalness argument. It varies between 1 and 10^{-34} if we move $M_{\rm susy}$ from the weak scale to the Planck scale, and this is the basis for the prediction $M_{\rm susy} \approx M_{\rm weak}$. But all QFT-based, bottom up naturalness arguments completely ignore the second factor, which is not even defined in QFT. The first attempts to compute it in the string landscape produced the result that it was proportional to $M_{\rm susy}/M_{\rm Planck}$ to a power larger than 2, so that large susy scales dominate over small ones by a large factor [183, 184]. Meanwhile that conclusion has been shown to be too simplistic [27, 185], and furthermore an important contributing factor was underestimated in earlier work, namely that vacua with broken susy are less likely to be stable. This can lead to a huge suppression [48, 49]. But one conclusion remains: even if the second factor in Eq. (15) is difficult to compute in a well-defined setting, this does not imply that it can be ignored in approaches where it cannot even be defined.

String Theory. An obvious weakness of the "MSSM" hypothesis is the first M, which stands for "minimal". There is no good fundamental reason to expect minimality, but dropping this restriction implies a substantial loss of predictive power. If the supersymmetric Standard Model is realized in string theory, the result is rarely minimal, see section 2.2. But while most of the additional particles can at least be avoided in special constructions, one kind is essentially inevitable: moduli.

It has been known for a long time that moduli can lead to cosmological problems [10, 11, 186]. If they decay during or after BBN they will produce additional baryonic matter and destroy the successful BBN predictions. Bosonic moduli have potentials, and will in general be displaced from their minima. Their time evolution is governed by the equation

$$\ddot{\phi} + 3H\dot{\phi} + \frac{\partial V}{\partial \phi} = 0, \tag{16}$$

where H is the Hubble constant. If $V = \frac{1}{2}m^2\phi^2$ + higher order terms, 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). For higher masses moduli decay can reheat the universe sufficiently to restart BBN from electroweak equilibrium.

Furthermore one can argue [187] 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. Gaugino masses can be one or two orders of magnitude less. 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 would still be a good chance to observe gauginos. Circumstantial evidence in favor of this scenario is that it prefers a Higgs mass near the observed value [188], whereas bottom-up supersymmetric models, ignoring moduli, suggested an upper limit of at most 120 GeV.

But there is one worrisome point. The 30 TeV bound on moduli masses is not an anthropic bound. Observers in universes with moduli masses below that bound would be deeply puzzled that their attempts at computing BBN abundances gave incorrect answers. They might see a helium abundance of only 19% while their computations predicted 24%, but nothing we know suggests that this has adverse effects on life. Now if supersymmetry prefers a lower scale because of naturalness, this would imply that universes with deeply puzzled observers should dominate universes with observers enjoying successful BBN predictions, such as ourselves.

None of these three mistakes is obviously fatal, but taking into account the "anthropic landscape of string theory" and all its implications definitely lowers the confidence level of predictions of low energy supersymmetry.

3.6. *Axions and the Strong CP problem*

The Standard Model Lagrangian contains a term

$$\theta \frac{g_3^2}{32\pi^2} \sum_{a=1}^{8} F^a_{\mu\nu} F^a_{\rho\sigma} \epsilon^{\mu\nu\rho\sigma} \ . \tag{17}$$

where the sum is over the eight generators of $SU(3)$. The parameter θ, an angle with values between 0 and 2π, is not an observable by itself. By making suitable phase rotations of the fermions its value can be changed, but then these phase rotations end up in the mass-matrices of the quarks. In the end, this leads to one new physical parameter, $\bar{\theta} = \theta - \arg \det\,(M_u M_d)$, where M_u and M_d are the quark mass matrices. A non-zero value for this parameter would produce a non-zero dipole moment for the neutron and certain nuclei, which so far has not been observed. This puts an upper limit on $\bar{\theta}$ of about 10^{-10}. Since there is no anthropic argument in favor of such a small value – the smallest one can argue for anthropically [189] is 10^{-3} – this is the one of the most serious naturalness problems in the Standard Model. No one would argue that we observe such a small value just by chance.

This problem has a rather simple solution, the Peccei-Quinn [190] mechanism. It works by postulating an additional pseudoscalar boson with a dimension 5 interaction with the QCD vector bosons

$$\Delta\mathcal{L} = \frac{1}{2}\partial_\mu a \partial^\mu a + \frac{a}{32\pi^2 f_a}\sum_a F^a_{\mu\nu}F^a_{\rho\sigma}\epsilon^{\mu\nu\rho\sigma}, \tag{18}$$

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 $a \to a + \epsilon$. This allows us to shift a by a constant $-\bar{\theta}f_a$ so that the $F\tilde{F}$ term (17) 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 non-invariant operator. Through non-perturbative effects the anomaly generates a potential with a minimum at $a = 0$ of the form

$$V(a) \propto \Lambda^4_{\mathrm{QCD}}\left(1 - \cos(a/f_a)\right). \tag{19}$$

Note that $\bar{\theta}$ is periodic with period 2π, so that the shift symmetry is globally a $U(1)$ symmetry. It was pointed out in [191, 192] that this breaking of the $U(1)$ symmetry leads to a pseudo-scalar pseudo-Goldstone boson, which was called "axion".

The mass of this particle is roughly $\Lambda^2_{\mathrm{QCD}}/f_a$, but if we take into account the proportionality factors in (19) the correct answer is

$$m_a = \frac{m_\pi f_\pi}{f_a}F(m_q), \tag{20}$$

where f_π is the pion decay constant and $F(m_q)$ 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 order 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 [193] 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. There is an upper limit of $f_a < 10^{12}$ GeV because if f_a were larger the contribution of axions to dark matter would be too large. This results in a small allowed window for the axion mass: 6 μeV $< m_a <$ 6 meV.

The upper limit of f_a is interesting since on the one hand it is large, but on the other hand it does not quite reach the string scale. The non-renormalizable interaction in the axion Lagrangian points to new physics at that scale, and from a string theory perspective the natural candidate for such new physics would be string theory itself. But then the scale seems uncomfortably low in comparison to typical string scales. Indeed, in [194] the difficulties are examined and possible ways out are discussed.

One way out is suggested by the fact that the amount of axion dark matter is proportional to $\sin^2\theta_0$, where θ_0 is the initial misalignment angle of the axion potential. In deriving the upper bound, one assumes that our universe emerges from a configuration with random alignments, resulting in an average of $\sin^2\theta_0$. This would give a value of $\frac{1}{2}$ for $\langle\sin^2\theta_0\rangle$.

The fact that the 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 [195]. It is possible that our universe comes from a single inflated region with a small value of θ_0. For a larger value of θ_0 (given an axion decay constant at the string scale) too much dark matter would be produced. One has to argue that, even though the likelihood of living in a region with small θ_0 is small, this is compensated by the fact that more observers will find themselves in such regions, because larger dark matter densities are detrimental for the existence of life. The most likely reason for that is galaxy formation and the density of matter in galaxies, both of which depend on the dark matter fraction. See [196–199] for further discussion. It has even been argued that finding a high scale axion would provide evidence for the multiverse and the string theory landscape [200]. The upper bound on the axion decay constant can also be raised if there is a non-thermal cosmological history, for example caused by decay of heavy moduli [58].

Candidate axions occur abundantly in string theory, but their survival as light particles is affected by the moduli stabilization mechanism. They may be thought of as phase factors of complex fields. The real parts of those fields must be stabilized. They would otherwise give rise to fifth forces or affect BBN predictions. But thanks to their derivative couplings, axions are far less constrained. However, not every mechanism to stabilize moduli is capable of giving mass to the real part and leave the imaginary part unaffected. For example, stabilization by fluxes or by instanton induced terms in the superpotential gives mass to both the moduli and the corresponding axions.

Axions that survive moduli stabilization may in principle play the role of PQ axions that solve the strong CP problem, provided they do not acquire masses by other mechanisms. The usual folklore that gravity does not allow exact global symmetries suggests that all axions will eventually get a mass. Otherwise their presence as massless particles would imply the existence of an exact, but spontaneously broken global symmetry, with axions as Goldstone bosons. Just as the QCD instanton get a mass from non-perturbative QCD effects, all other axions should get a mass as well. The Peccei-Quinn mechanism works as long as there is an axion coupling to QCD with a mass contribution from any other sources that is at least ten orders of magnitude smaller than the QCD contribution.

It is clear from the previous paragraph that if string theory produces a PQ axion, it is likely that it produces many other axions in addition to this. Since their masses are generated by non-perturbative effects, it is natural to expect them to be distributed in a scale invariant way, spanning many orders of magnitude. This plenitude of axions has been called the "axiverse" [200]. Since the masses of the additional axions (not involved in the PQ mechanism) are not limited to the QCD window, this provides ample opportunities for observations in many mass regions.

Realizations of an axiverse have been discussed in fluxless M-theory compactifications [201] and in type-IIB models in the LARGE Volume Scenario [202]. Both papers consider compactifications with many Kähler moduli that are stabilized by a single non-perturbative 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. For supersymmetric moduli stabilization (such as the KKLT scenario, but unlike LVS) a no-go theorem was proved in [203]. Axions in the heterotic mini-landscape were discussed in [204]. They consider discrete symmetries that restrict the superpotential, so that the lowest order terms have accidental $U(1)$ symmetries that may include a PQ symmetry.

There are numerous possibilities for experiments and observations that may shed light on the rôle 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 [201, 205]. See [200, 206, 207] for a variety of possible signatures, ongoing experiments and references.

4. Conclusions

The past four decades have been a magnificent golden age for particle physics, but also for the leading approach to understanding it in terms of a fundamental theory of quantum gravity, string theory. Now we appear to be approaching an interesting and perhaps decisive moment in history. The Standard Model seems complete, and string theory has led to a mathematical challenge of monstrous proportions: the string theory landscape.

Whether one likes it or not, successful string phenomenology requires taming this monster. This will almost certainly involve a reconsideration of the questions we wish to answer. In the future, the current time may be remembered as the transition from the era of symmetries to a new era with different ways of thinking about fundamental problems.

In experimental physics, we are faced with a situation where positive results may emerge but are not guaranteed, and where negative results tell us fairly little. Apart from the traditional "new physics", examples of positive results that would have an impact on the landscape are variations of constants of nature, the observation of axions, dark matter of any kind, neutrino Majorana masses, sterile neutrinos, massive vector bosons, and many serendipitous discoveries ranging from desirable (*e.g.* proton decay or magnetic monopoles) to totally unexpected (*e.g.* something like faster-than-light neutrinos). With the large number of experiments and astrophysical observations still underway, it seems unthinkable that the Higgs particle will turn out to be the last discovery in particle physics. But if it is, this has to be viewed as circumstantial evidence in favor of a landscape. Any of the aforementioned positive results can be good or bad for the landscape idea, but there is no gold-plated experiment that verifies or falsifies it.

Further supportive or damaging evidence must come from pure theory. The (non)-existence of a broadly spread distribution of de Sitter vacua in string theory is a decidable issue, but it has turned out to be very difficult to reach a conclusion. The fundamental principle behind string theory still eludes us. The quantization of vibrating strings is not an acceptable fundamental starting point, and does not describe everything we call string theory anyway. The best hope for acceptance of the landscape idea is that it is derived from a fundamental theory of gravity which in its turn is derived from a plausible principle of nature. Meanwhile, we can try to bridge the gap between the Standard Model and the string landscape. We must convince ourselves that our Universe is indeed contained in the string theory landscape. We can explore our environment in the landscape, to see if we can understand why we observe the Standard Model and all of its features, especially the puzzling ones. This requires determining landscape distributions in various regions, and using anthropic arguments where possible.

This may all be postponed until the indefinite future if new physics still emerges during the next few years. If that new physics is due to large extra dimensions and a low higher-dimensional Planck scale, we can explore quantum gravity directly. This has always seemed too good to be true, and probably it is. If the new physics is low energy supersymmetry, which is still a well-motivated option, perhaps the winding road to the string landscape becomes a broad avenue all the way to the Planck scale. Most other options may provide more decades of exciting particle physics, but the historic moment we might be witnessing now would have passed and may never return.

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