STATUS OF SUPERSTRING AND M-THEORY

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Abstract

The first lecture gives a colloquium-level overview of string theory and M-theory. The second lecture surveys various attempts to construct a viable model of particle physics. A recently proposed approach, based on F-theory, is emphasized.

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1 Introduction

The plan for my two lectures is as follows: the first lecture (in section 2) will give a colloquium level overview of string theory. The presentation will mention some of the history of the subject at the same time that basic concepts (such as compactification of extra dimensions, dualities, and branes) are introduced. There is a lot of ground to cover, and it will be necessary to be somewhat sketchy. Since the theme of this school concerns discoveries that might be made at the LHC, it seems appropriate to emphasize string theory approaches to particle physics model building. Therefore this is the subject of the second lecture (in section 3). It will survey several different approaches that have been developed, describing some of the successes and problems of each of them. One approach that has seen a lot of recent progress, and which looks especially promising, will be emphasized.

For reasons that will be explained, the emphasis is on supersymmetrical string theories (called superstring theories), which naturally are associated with ten-dimensional spacetime. In certain cases, the strong coupling limit gives an eleven-dimensional theory, called M-theory, which does not contain strings. For the student who wants to learn more, there are many textbooks on string theory. Some of them are [1, 2, 3, 4]. The one by Zwiebach [3] is addressed to advanced undergraduates, but it is suitable for all physicists who are not trying to become experts. The other three attempt to bring the reader up to the state of the art at the time when they were written. The subject is fast-moving and much was learned in the intervals between their appearances, so a lot of new material is included in each subsequent book, and a lot of the older material is not repeated.

This school pays homage to Sidney Coleman, a great scholar and teacher, whose contributions to theoretical physics in general, and the Erice schools in particular, are legendary. I always enjoyed interacting with him. Although I did not know him very well, we did cross paths in Aspen and elsewhere on a number of occasions. I recall him once saying that there are three things that he does not like, all of which are becoming popular: supersymmetry, strings, and extra dimensions. Obviously, my views are quite different, but this did not lessen my regard for him, nor did it harm our personal relationship. In fact, I respected his honesty, especially as he did not try to impose his prejudices on the community.

2 Lecture 1: Overview of Superstring and M-Theory

String theory arose in the late 1960s as a radical alternative to conventional quantum field theory in which the fundamental objects are one-dimensional strings rather than zero-dimensional points. This proposal arose in the context of S-matrix theory, a subject that has been much maligned, but whose imprint is indelible. From a modern perspective, it is clear that theories of point particles and theories of strings are related by dualities. Therefore the
two classes of theories, which used to seem completely at odds, actually are deeply intertwined and do not have a sharp separation. Thus, I would claim that the subject area that is currently called string theory should be viewed as the logical completion of quantum field theory. That being the case, string theory has a certain inevitability, and there is nothing radical about it at all.

2.1 Basic Concepts of String Dynamics

Let us start by sketching what it means to construct a theory based on strings. To do this it is convenient to review first point-particle theory (i.e., quantum field theory) from a first-quantization viewpoint. This is not the way that it is usually taught. The reason for doing this is that the analogous formulation of string theory is much better understood, and much easier to describe, than the string theory analog of second-quantized quantum field theory.

For a point particle the classical motion makes the invariant length of the world-line extremal. The corresponding action, proportional to this length, is given by

$$S = -m \int ds.$$  

Here $ds$ is the invariant line element, given in a general curved spacetime by $ds^2 = g_{\mu\nu}dx^\mu dx^\nu$, and $m$ is the mass of the particle in question. (We always set $\hbar = c = 1$.) The motion of a particle is given by a world line $x^\mu(\tau)$, where $\tau$ is an arbitrary parameter for the trajectory. The action is independent of the choice of this parametrization.

To pass to the quantum theory one computes an amplitude for propagation from an initial spacetime point $x^\mu_i = x^\mu(\tau_1)$ to a final spacetime point $x^\mu_f = x^\mu(\tau_2)$. As Feynman taught us, this is given by the path integral (or sum over histories)

$$A_H = \int_1^2 Dx^\mu(\tau) e^{iS}.$$  

This is slightly oversimplified in one respect. The $\tau$ reparametrization invariance is a type of gauge invariance that need to be accounted for. One possibility, which is not always the most convenient one, is to choose a gauge in which $\tau$ is time, i.e., $x^0(\tau) = \tau$.

Interactions can be incorporated in this formalism by allowing world lines to join or bifurcate and associating a coupling constant $g$ whenever this occurs. In this way, one can reproduce the perturbation expansion of a second-quantized quantum field theory. One shortcoming of this first-quantized approach is that it is not very convenient for studying nonperturbative phenomena.

We can now “invent” string theory by doing the same thing for one-dimensional extended objects. If the string’s topology is that of a circle it is called a closed string. If the topology is that of a line segment, it is called an open string. A string sweeps out a two-dimensional surface in spacetime, called the world sheet of the string. In the case of a closed string, the
The topology of the world sheet is that of a cylinder, and in the case of an open string it is that of a strip. For a string the motion makes the invariant area of the world-sheet extremal
\[ S = -T \int dA. \]

The invariant area element \( dA \) is given by a simple formula analogous to that of \( ds \) in the point-particle case. The coefficient \( T \) is identified as the tension (or energy per unit length) of the string. The world-sheet is described by embedding functions \( x^\mu(\sigma, \tau) \). Here \( \sigma \) is interpreted as a coordinate along the string (periodic in the case of a closed string) and \( \tau \) is a timelike parameter. The choice of coordinates is arbitrary in that there is now a two-dimensional reparametrization invariance.

The basic idea of the extension to the quantum theory is the same as before. Namely, one defines amplitudes for string propagation by a path integral, properly taking account of the reparametrization invariance. However, there are some very interesting differences between the two cases: First of all, after covariant gauge fixing, the world-sheet theory has a conformal symmetry. (The conformal group in two dimensions is infinite dimensional.) Consistency of the quantum theory requires cancellation of a conformal anomaly. The most straightforward way to achieve this is to take the spacetime dimension to be 26. This defines critical string theory. Alternative approaches give noncritical string theory.

The second important difference from the point-particle case concerns interactions. In the case of strings, these are uniquely determined by the free string theory. They arise for purely a topological reason. For example, the process in which a single closed string turns into two closed strings is given by the pants diagram. This is a smooth world sheet without any singularity associated with the string junction. Such junctions need to be incorporated in string path integrals. The smoothness of the world sheet has a remarkable consequence: loop amplitudes have no ultraviolet divergences! In the case of point-particle theories these can be traced to the world line junctions, which are short-distance singularities in the one-dimensional network of world lines. Moreover, at least in the case of oriented closed strings, there is just one Feynman diagram at each order of the perturbation expansion!

The actions \( S \) described above can be generalized to objects with \( p \) dimensions, called \( p \)-branes. However, the quantum analysis breaks down for \( p > 1 \), because their world-volume theories are nonrenormalizable. Later we will argue that various \( p \)-branes do arise in string theories and M-theory as nonperturbative excitations. The significance of this nonrenormalizability does not concern the existence of \( p \)-branes. Rather, it means that they cannot be treated as the fundamental objects on which to base a perturbation expansion. The fact that this is possible for strings is the feature that distinguishes them from higher-dimensional \( p \)-branes. However, at strong coupling, when a perturbation expansion is not helpful, this distinction evaporates and all branes are more or less equal.
2.2 A Brief History of String Theory

2.2.1 Hadronic String Theory

String theory arose in the late 1960s in an attempt to understand the strong nuclear force. This is the force that holds quark and gluons together inside the strongly interacting particles (hadrons). The story has a long and complex history that has not been assembled yet in a single scholarly study. However, many of the participants in this history presented their own personal recollections in a 2007 meeting at the Galileo Institute in Florence. My contribution to the proceedings is available [5].

Veneziano’s discovery of a simple formula for a 2-particle to 2-particle scattering amplitude with many attractive features (such as Regge behavior) quickly led to the realization that a theory based on strings, rather than point-like particles, could account for various observed features of the strong nuclear force [6]. The original string theory, sketched in the previous subsection, contains only bosons, and (as discussed above) it is consistent for 26 spacetime dimensions. 25 of the dimensions are spatial and 1 is time.

Another string theory containing fermions (as well as bosons) was constructed in 1971 by Pierre Ramond, André Neveu, and me [7, 8]. It requires 10 dimensions. Its development led to the discovery of supersymmetry, a symmetry that relates bosons and fermions. Strings in theories with this symmetry are called superstrings.

The unrealistic dimension of spacetime was obviously a problem, but not the only one. Another one was that the spectrum of string excitations in both of the string theories includes massless particles, whereas all hadrons have positive mass. In the early 1970s a better theory of the strong nuclear force, called quantum chromodynamics (or QCD), was developed. As a result, string theory fell out of favor. What had been an active community with several hundred participants rapidly shrunk to a backwater of theoretical physics kept alive by only a few diehards.

2.2.2 Unification

One of the massless particles in the closed-string spectrum of the 26-dimensional bosonic string theory, as well as in the closed-string spectrum of ten-dimensional superstring theories, has precisely the right properties to be the graviton – the particle responsible for transmitting the gravitational force. In other words, it has spin two, which means that its polarization tensor belongs to a symmetric and traceless representation of the little group. Moreover, the requisite local symmetries are incorporated, namely general coordinate invariance and local Lorentz invariance. These properties ensure precise agreement with general relativity at energies low compared to the string scale.

To be completely honest, it should also be mentioned that there is also a massless scalar (called the dilaton). This field is very important, because its vacuum expectation value
determines the string coupling constant. Somehow, the dynamics should give mass to this field to avoid contradictions with experiments that show that the long-range gravitational force is pure spin two to high accuracy. Presumably, this is achieved at the same time as the extra dimensions are compactified. I will discuss this more later.

The open-string spectrum contains massless spin one gauge fields, like those in the standard model. Again, the usual gauge symmetries are incorporated, which ensures agreement with Yang–Mills gauge theory at energies low compared to the string scale.

For several years, string theorists (including myself) tried to modify string theory so as to give mass to these fields, since we knew that there are no massless hadrons. However, these attempts always led to inconsistencies. Eventually it became clear that these were essential features of the theory, and that we should let the theory guide us, rather than trying to impose our will on it. Accordingly, in 1974 Joël Scherk and I proposed to use string theory as a unified theory of all forces (including gravity), rather than just the strong nuclear force \cite{Scherk:1974yz}. Thus, we stumbled upon an approach to achieving Einstein’s dream – a unified theory of all fundamental forces. The connection to general relativity at low energies was described independently by Yoneya \cite{Yoneya:1974sj,Yoneya:1974jd}, though he did not propose using string theory as a unified theory.

The interpretation of string theories as unified theories of gravity and other forces has several advantages, which were immediately apparent. All prior attempts to describe quantum corrections to Einstein’s theory of gravity assumed point particles. They gave nonrenormalizable ultraviolet divergences, which makes the theory unacceptable, at least within the context of perturbation theory. String theory, on the other hand, was known to be UV finite to all orders in perturbation theory. The intuitive reason for this, is that the string world sheet is smooth, without any short-distance singularities, no matter how complicated the topology.

A second advantage of using string theory as a theory of gravity is that the extra dimensions can be an advantage rather than a disadvantage. Extra spatial dimensions can be compact in a gravitational theory, where the spacetime geometry is determined by the dynamics. So the geometry can appear four-dimensional at low energies. It then becomes a dynamical question whether the theory has acceptable quantum vacua of this type.

### 2.2.3 The Size of Strings

When strings were supposed to describe hadrons their typical size needed to be roughly that of a typical hadron, namely $10^{-13}$ cm. However, in a relativistic quantum theory of gravity there is a characteristic length, known as the Planck length

$$L_P \sim \sqrt{\hbar G/c^3} \sim 10^{-33} \text{cm}.$$
This corresponds to a mass scale of about $10^{19}$ GeV. This is the natural first guess for the fundamental string length scale in a string theory containing gravity. This shrinkage by 20 orders of magnitude was a very large conceptual change. Fortunately, the mathematics remained pretty much the same. More recently, it has been appreciated that this first guess might be modified for various reasons. There are possible corrections involving factors of the coupling constant, the volume of compactified extra dimensions, or even warpage in the geometry of extra dimensions. Therefore, it makes sense for experimentalists to look for signs of the string scale at the LHC, though I would be astonished if it did show up.

2.2.4 The First Superstring Revolution

String theory became a hot subject in the mid-1980s following some important discoveries: anomaly cancellation [14], heterotic strings [15], and Calabi–Yau compactification [16]. By the time the dust settled, it appeared that there are exactly five consistent superstring theories, each of which requires ten spacetime dimensions and supersymmetry. Three of them, called Type I, Type IIA, Type IIB, were introduced by Green and me [13] building on earlier work of Gliozzi, Scherk, and Olive [12]. The other two, called $SO(32)$ heterotic and $E_8 \times E_8$ heterotic, were formulated in [15]. They incorporate the gauge groups identified in [14]. The anomaly-free type I superstring also has $SO(32)$ gauge symmetry.

Each of these five theories has no adjustable dimensionless parameters. All dimensionless parameters arise either

- *dynamically* as the expectation values of scalar fields

  or

- as *integers that count something* such as topological invariants, physical objects (branes), or quantized fluxes.

It appeared that the quest for a unique theory was almost at hand. The only mystery was why there should be five theories when only one is required.

One scheme looked particularly promising. Specifically, the $E_8 \times E_8$ heterotic theory has consistent vacuum solutions in which six spatial dimensions form a compact *Calabi–Yau manifold*. Calabi–Yau manifolds are a class of six-dimensional manifolds that can be described as having three complex dimensions. The precise mathematical statement is that they are Kähler manifolds with $SU(3)$ holonomy.\(^2\)

The other four dimensions form Minkowski spacetime, $M_{3,1}$. Thus, the ten-dimensional spacetime is a product space

$$M_{10} = CY_6 \times M_{3,1}.$$\(^2\)

\(^2\)Kähler manifolds are complex manifolds that have no torsion. The restriction to the $SU(3)$ subgroup of the generic $U(3)$ holonomy group corresponds to the vanishing of the first Chern class. It also implies the existence of a covariantly constant spinor, which in turn implies that some supersymmetry is preserved.
Intuitively, one can imagine a Calabi–Yau space attached to every point in four-dimensional spacetime. If the CY space is small, then its structure cannot be probed at low energy. Experimental detection requires energies greater than the inverse size of the compact CY space. I will say more about these constructions in my second lecture.

2.2.5 The Second Superstring Revolution

In the mid-1990s there was a remarkable burst of progress in addressing some of the issues that we have raised. A few of the key initial contributions were [17, 18, 19]. The main lessons were the following:

- There is just one theory! What had been viewed as five theories are actually five different corners of a space of solutions to a unique underlying theory. The five superstring theories are related by various surprising equivalences, called dualities:

  **T-duality** \((R \rightarrow 1/R)\) relates the two type II superstring theories, and also the two heterotic string theories. Thus, for example, if the type IIA superstring theory is considered in a spacetime that is nine-dimensional Minkowski spacetime times a circle of radius \(R (M_{8,1} \times S^1)\), this is equivalent to the type IIB superstring theory in a spacetime that is nine-dimensional Minkowski spacetime times a circle of radius \(1/R\). (Here we are using units in which the fundamental string length scale \(L_s\) has been set equal to one.) Thus, the two type II theories in ten-dimensional Minkowski spacetime are two limiting cases of a continuum of consistent possibilities that are connected by letting the radius \(R\) range from zero to infinity. This radius is determined by the vacuum expectation value of a modulus field (one of the components of the ten-dimensional metric).

  **S-duality** \((g_s \rightarrow 1/g_s)\) relates the type I superstring theory and the \(SO(32)\) heterotic string theory. The string coupling constant \(g_s\) is given by the vacuum expectation value of a modulus field called the dilaton. Prior to this discovery one only knew how to analyze string theory at weak coupling using perturbation theory. S-duality implies that this pair of theories is continuously connected by varying \(g_s\) from zero to infinity. It also implies that the strong coupling expansion of one theory is given by the weak coupling expansion of the other theory. The type IIB superstring theory is related to itself in a similar manner, which determines its strong coupling expansion as well.

- Knowing the strong coupling expansion of three of the five superstring theories, it was natural to explore the strong coupling behavior of the other two superstring theories. The conclusion is quite striking. At strong coupling the type IIA and the \(E_8 \times E_8\) theories become 11-dimensional. More precisely, the eleventh dimension has size proportional to \(g_s L_s\). In the IIA case it is circle whereas in the \(E_8 \times E_8\) case it is a line.
interval with two endpoints. In the latter case the 11-dimensional spacetime has two ten-dimensional faces. It turns out that one $E_8$ gauge group is localized on one face and the other $E_8$ on the other face, so that they are spatially separated.

In either the type IIA or $E_8 \times E_8$ theory the $g_s \to \infty$ limit leads to the same 11-dimensional Minkowski spacetime theory. This 11-dimensional theory, which is still not yet very well understood, is called $M$-theory. M-theory is approximated at low-energies by 11-dimensional supergravity, a theory that was formulated many years earlier [20].

- New stable objects, called $p$-branes, arise nonperturbatively. $p$ is the number of spatial dimensions they occupy. Thus, in this notation, a point-particle is a 0-brane, a string is a 1-brane, and so forth. $p$-branes can carry generalized conserved charges that are sources for antisymmetric tensor gauge fields with $p + 1$ indices, which generalize the Maxwell field in the $p = 0$ case. The conserved charges satisfy generalized Dirac quantization conditions. In $D$ spacetime dimensions, the magnetic dual of a $p$-brane is a $(D - p - 4)$-brane.

The main categories of $p$-branes are called D-branes, M-branes, and NS-branes. D-branes are characterized by Dirichlet boundary conditions for open strings. In the type IIA theory stable D-branes exist for even values of $p$ and in the type IIB theory they exist for odd values of $p$. M-theory has a three-form gauge field, which can couple electrically to an M2-brane and magnetically to an M5-brane. The NS5-brane is the magnetic dual of the fundamental string in the heterotic and type II superstring theories.

### 2.3 AdS/CFT: Holographic Duality

In 1997 Maldacena proposed a new class of dualities relating certain string theory and M-theory solutions in anti de Sitter geometries (AdS$_{d+1}$ times a compact space) to conformally invariant quantum field theories (CFT$_d$) [21]. The first piece of evidence for such a relationship is the fact that the conformal symmetry group in $d$ dimensions, $SO(d, 2)$, is the same as the isometry group of anti de Sitter space in $d + 1$ dimensions. The precise relationship between amplitudes/correlation functions in the two pictures was spelled out [22, 23]. Such dualities are called holographic, because string theory in AdS$_{d+1}$, which has $d + 1$ dimensions, is equivalent (dual) to a conformally invariant quantum field theory in $d$ dimensions: the

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3These gauge fields can also be regarded as differential forms $A = \frac{1}{(p+1)!} A_{\mu_1 \mu_2 \cdots \mu_{p+1}} dx^{\mu_1} \wedge dx^{\mu_2} \wedge \cdots \wedge dx^{\mu_{p+1}}$. In this notation the coupling of the gauge field to the brane world volume is given by the simple expression $\mu \int A$. The parameter $\mu$ is the brane charge.

4The type I superstring does not carry a conserved charge, and therefore it can break. It does not have a magnetic dual.

5Anti de Sitter space is a maximally symmetric spacetime with constant negative curvature.
dual CFT\textsubscript{d}. This is reminiscent of recording a three-dimensional image on a two-dimensional emulsion. The reason that two constructions that look so radically different can nonetheless be equivalent is that when one of them is weakly coupled, the other is strongly coupled. If it were easy to understand both of them at the same time, there would be a paradox.

This is an enormous subject, and I will just say a little bit about it here. Three maximally symmetric examples are given in Maldacena’s original paper: Type IIB superstring theory in an $AdS_5 \times S^5$ background geometry with $N$ units of five-form flux threading the sphere is dual to $\mathcal{N} = 4$ super Yang–Mills theory with an $SU(N)$ gauge group. This by far the most studied example. The leading term in the ’t Hooft large-$N$ expansion of the gauge theory (the planar approximation) is supposed to be dual to the string theory in the tree approximation. This much of the duality is quite well understood, and I think it is likely that it will be completely proved someday. The higher-order terms in the $1/N$ expansion correspond to string loop corrections. This duality can be generalized to examples with less supersymmetry in which the five-sphere is replaced by another five-dimensional Sasaki–Einstein space. It is also possible to replace the $AdS$ space by a space that is only asymptotically anti de Sitter. In this case the conformal field theory is perturbed by relevant operators.

The other two maximally supersymmetric examples relate M-theory in an $AdS_4 \times S^7$ background geometry with flux to a three-dimensional conformal field theory and M-theory in an $AdS_7 \times S^4$ background geometry with flux to a six-dimensional conformal field theory. Within the past year there has been a great deal of progress in understanding the first case (as well as an orbifold generalization) \cite{24}. On the other hand, there has been very little progress in understanding the second M-theory duality.

The general AdS/CFT framework is being used to construct brane configurations that capture many of the essential features of QCD, even though it isn’t precisely the same theory \cite{25}. From the modern viewpoint, such constructions can be regarded as spinoffs of string theory rather than its central goal. Even so, this brings string theory back to its historical origins. It also makes the reasons for the early failure to achieve the original goal of describing hadrons much clearer. For one thing, it is now apparent that a string theory dual of QCD must involve at least one extra spatial dimension \cite{26}. There are also spinoffs to other fields: 1) There have been striking applications of AdS\textsubscript{5}/CFT\textsubscript{4} duality to studies of the quark-gluon plasma (RHIC, LHC), especially the computation of the viscosity to entropy density ratio $\eta/s$ \cite{27, 28}. Taking the field theory at finite temperature corresponds to including a black hole in the dual AdS geometry! 2) Similar constructions could have condensed matter applications. Systems that show promise are ones in which there is a phase transition at which the conformal field theory is strongly coupled. Then the dual black-hole description could be helpful. A possible application of this type is to high $T_c$ superconductors. For a recent overview see \cite{29}. 3) Nonrelativistic versions of AdS/CFT may be relevant to Bose-condensed cold-atom systems under special conditions \cite{30, 31}. 
2.4 Some Other Topics in String Theory

2.4.1 Brane Worlds

As we already mentioned, the defining property of D-branes is that strings can end on them. Away from D-branes, a string must have the topology of a circle, but the string can break provided that its ends are attached to D-branes. The lowest excitation mode of the open string is a massless gauge field. This has the crucial consequence that Yang–Mills gauge theories (like the standard model) can live on stacks of D-branes. For example, in the case of $N$ coincident type II D-branes, the world-volume theory of the branes is a supersymmetric $U(N)$ gauge theory.

This picture raises new possibilities for the physics of the extra dimensions, which is rather different from the usual Kaluza–Klein picture. Rather than having a wave function that is uniformly spread over the compact extra dimensions, a low-energy field can have its wave function concentrated on (or close to) a D-brane, which is a defect in the extra dimensions that fills the noncompact spacetime dimensions. The simplest such possibility is that the observable Universe is actually a stack of D3-branes, which are points embedded in the six extra spatial dimensions. A generalization of this approach uses intersecting stacks of higher-dimensional D-branes. For example, one could consider stacks of D6-branes wrapping 3-cycles in the six extra dimensions.

2.4.2 Flux Compactifications

As has already been mentioned, the type II superstring theories contain various massless antisymmetric-tensor (or differential-form) gauge fields

$$A_n = \frac{1}{n!} A_{\mu_1 \mu_2 \ldots \mu_n} dx^{\mu_1} \wedge dx^{\mu_2} \wedge \ldots \wedge dx^{\mu_n}.$$ 

$p$-branes with $p = n - 1$ are sources for these fields. These differential-form gauge fields have gauge-invariant field strengths of the form $F_{n+1} = dA_n$, which are invariant under gauge transformations of the form $\delta A_n = dL_{n-1}$. This is a straightforward generalization of Maxwell theory (the $n = 1$ case).

If the compact dimensions contain nontrivial $(n + 1)$-cycles, $C_{n+1}$, it is possible to have (quantized) flux threading the cycle:

$$\int_{C_{n+1}} F_{n+1} = 2\pi N, \quad N \in \mathbb{Z}.$$ 

Such flux can be present without a $p$-brane source. The possibility of constructing superstring vacua containing such fluxes greatly increases the number of possible vacua, even though there are constraints that must be satisfied. Mike Douglas analyzed one particular CY compactification of the type IIB theory and estimated the number of distinct flux vacua to be about $10^{500}$ [32]. For a review of this subject, see [33].
The moduli problem — the occurrence of massless scalar fields $\phi_i$ with continuously adjustable vacuum expectation values — can be solved in the context of flux compactification: The fluxes induce a nontrivial potential energy function for the moduli $V(\phi_i)$. This landscape function has isolated minima, which are what Douglas counted. The moduli are massive at an isolated minimum.

2.4.3 Warped Compactification

One of the important properties of flux compactifications is that they give rise to warped geometries. This means that the ten-dimensional geometry is no longer a direct product of the form

$$\mathcal{M}_{10} = M_{3,1} \times K_6.$$ 

Instead, the four-dimensional Minkowski part of the ten-dimensional metric is multiplied by a warp factor $h(y)$ that depends on the position $y$ in the internal manifold

$$ds_{10}^2 = h(y)dx \cdot dx + ds_6^2.$$ 

In 1999 Randall and Sundrum proposed that an exponential $y$ dependence could give a large ratio between the warp factors at the positions of two 3-branes, which they called the standard model brane and the Planck brane [34]. They proposed that this could solve the hierarchy problem. This kind of exponential warping is exactly what one has in the case of an AdS spacetime. The new idea was to only consider a slice of AdS spacetime between two branes. This is only an approximation to what can be achieved in honest string theory solutions. Nonetheless, the RS scenario can be made rather precise in the context of flux compactifications with branes added [35]. If this is how Nature works, quantum gravity effects might be accessible at the LHC. The experimentalists intend to look for them.

2.5 String Cosmology

An important question, which was not discussed much in the 20th century, is “What does string theory have to say about cosmology?” Nowadays, string cosmology is a very active research area. There are even entire conferences devoted to it. Clearly, progress in this area requires understanding time-dependent solutions of string theory, which is technically challenging.

There is a lot of evidence that the very early Universe underwent a period of inflation during which the scale factor grew exponentially by a factor of at least $e^{60}$ (“60 e-foldings”). In simple field theory models this is described in terms of a “slowly rolling” scalar field called an inflaton. Proposals for the string-based origins of inflation, or possible alternatives, are being explored extensively. One scenario is based on CY compactification with flux and warped throats in the geometry [36, 37]. In this scenario inflation takes place as a D3-brane
moves down a throat, attracted to an anti-D3-brane at the bottom until they collide and annihilate. A scalar mode of an open string connecting the branes is the inflaton. The annihilation releases the brane tension energy. It heats up the Universe to start the hot big bang epoch. All sorts of strings are produced, and some might survive to be observable as cosmic superstrings [38].

3 Lecture 2: Superstring Phenomenology

Since the focus of this school is the physics of the LHC, it seems appropriate to discuss attempts to relate string theory to the particle physics phenomenology in greater detail than other topics. That is the purpose of this lecture. However, before getting into the details of such efforts, let me emphasize that the unique feature of string theory is that it unifies gravitation and particle physics. Moreover, its most successful prediction (or postdiction, if you wish) so far is that Gravity Exists!

It is not yet clear how predictive string theory is for particle physics. The discovery of the string theory landscape has led to some pessimism in this regard. However, as I will argue, there are good reasons to believe that if one makes certain experimentally motivated assumptions and inputs some experimental details, string theory could be quite predictive. I don’t think it is reasonable to expect to predict all experimental facts from first principles. It would already be a great success, if one could predict a large number of facts starting from a small number of experimentally motivated assumptions. How far one can go in this direction remains to be seen.

In perturbative string theory there are two basic ways to obtain non-abelian gauge symmetry in four dimensions: 1) It is already present in 10 dimensions (type I and heterotic strings) and broken partially by compactification. 2) It arises as the symmetry of compact extra dimensions (Kaluza–Klein). The Kaluza–Klein approach does not give anything realistic. Nonperturbatively, there are other possibilities that we will discuss later.

3.1 Perturbative Heterotic String

The gauge groups of the standard model and grand unified models embed nicely in $E_8$ but not in $SO(32)$. In the period 1985–95 only one scheme looked promising for phenomenology: Calabi–Yau compactification of the $E_8 \times E_8$ heterotic string. This entails taking the ten-dimensional geometry to be a product of four-dimensional Minkowski spacetime and a six-dimensional Calabi–Yau space, $\mathcal{M}_{10} = \text{CY}_6 \times M_{3,1}$. There must also be nontrivial gauge fields, satisfying certain consistency conditions (related to anomaly cancellation), in the internal space. This gives a GUT-like effective four-dimensional theory at low energy with $\mathcal{N} = 1$ supersymmetry. These effective theories have a number of attractive features:
• They have the structure of *supersymmetric grand unified theories*. The advantages of low-energy supersymmetry and grand unification are therefore naturally incorporated.

• Each solution has a definite number of families of quarks and leptons determined by the topology of the CY space.

• The standard model gauge symmetry can be embedded in one $E_8$ factor. Then there is a *hidden sector*, associated to the second $E_8$ factor. Supersymmetry can break dynamically in the hidden sector. This breaking is communicated by singlet fields, which have gravitational strength interactions, to the visible sector. This is usually called *gravity mediation*, though more than gravity is involved.

• There are several good *dark-matter* candidates: The lightest supersymmetric particle (or LSP) is absolutely stable in schemes with unbroken R-parity. In gravity mediation the leading dark-matter candidates are the lightest neutralino (a mixture of the partners of the photon, $Z$ boson, and Higgs bosons) and hidden-sector particles.6

These developments created a lot of excitement in the mid-80s. In my view a large part of the theory community made a phase transition from being too skeptical of string theory to being too optimistic about the short-term prospects for constructing a realistic theory. The successes were qualitative, and there were many problems and puzzling questions:

• Hundreds of Calabi–Yau manifolds were known (now many thousands). Which one of them, if any, is the *right one*? Is there a theoretical principle for deciding?

• Why are there four other consistent superstring theories? Can they give interesting solutions, too?

• The CY compactification scenario described above is based on perturbation theory. What new *nonperturbative features* appear at strong coupling?

• These CY solutions typically give a large number of massless scalar fields (called *moduli*).7 The moduli fields have gravitational strength interactions, but long-range scalar forces of that strength are ruled out by standard tests of general relativity. Thus, all these CY solutions are unsatisfactory unless there is some way of giving mass to the moduli. (Recall that in the first lecture I reported that this can be achieved in flux compactifications of type II superstring theories.)

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6In other mediation schemes, such as gauge mediation, the gravitino (the supersymmetry partner of the graviton) and saxions (supersymmetry partners of axions) are candidates.

7Moduli fields have a flat potential. This means that the vacuum energy does not depend on their values.
• In a gravitational theory the value of the minimum of the potential also matters, because it determines the \textit{vacuum energy density}, which is the effective cosmological constant. For many years it was assumed (without a good theoretical argument) that this should be zero. However, about a decade ago it was discovered in studies of distant supernovas (and confirmed by other means) that the vacuum energy density (or dark energy) is very small and positive. The value is roughly $10^{-120}$ in Planck units, which is the smallest measured number in all of science. Typically, one would expect to compute a result of order unity, so it is a great challenge to try to account for this result.\footnote{If the true value were zero, and there were no dark energy, this would still be a big problem.} Supersymmetry can help, but once it is broken at the TeV scale, one still expects to obtain a value of order $10^{-60}$ in Planck units, which is 60 orders of magnitude too large.

• The ratio of the Planck scale to the GUT scale comes out an order of magnitude too small.

• Extra family-antifamily pairs typically occur.

\section*{3.2 Other Approaches}
Following the second superstring revolution in the mid-1990s several new approaches to particle phenomenology opened up. These exploit discoveries such as

• 11d M-theory
• Singular geometries (e.g. orbifolds and orientifolds)
• D-branes and other branes
• Flux compactifications and warped geometries
• Nonperturbative type IIB superstring solutions (F-theory)

\subsection*{3.2.1 Heterotic M-Theory}
Hořava and Witten discovered that the $E_8 \times E_8$ heterotic string at strong coupling is 11-dimensional \cite{horava2000}. In particular they showed that anomaly cancellation requires that $E_8$ gauge fields live on each boundary component \cite{witten2000}. In heterotic M-theory the geometry is a solid slab with two ten-dimensional boundaries. One $E_8$ lives on each boundary.

Calling the coordinate that runs across the slab $y$, the shape and size of the $CY_6$ manifold can be $y$-dependent. In this way one incorporates nonperturbative improvements of the perturbative heterotic constructions. Taking this possibility into account improves the range
of possible ratios of the Planck scale to the GUT scale [41]. Fairly realistic examples have been constructed by the Penn group. (See [42, 43] and references therein). However, despite their many realistic features, these examples still have massless moduli.

3.2.2 M-Theory on a $G_2$ Manifold

The strong coupling limit of type IIA superstring theory is eleven-dimensional M-theory, which is approximated at low energies by 11-dimensional supergravity. To obtain $\mathcal{N} = 1$ supersymmetry in four dimensions at low energies requires a compactification of the form

$$\mathcal{M}_{11} = K_7 \times M_{3,1},$$

where $K_7$ is a compact 7-manifold of $G_2$ holonomy. ($G_2$ is the smallest of the five exceptional Lie groups. It has 14 generators.) If the $G_2$ manifold is smooth, then the effective four-dimensional theory at low energy has no non-abelian gauge symmetry or parity violation. To overcome this problem it is necessary to consider $G_2$ manifolds with certain types of singularities [44]. While the necessary types of singularities are well understood, there has been little progress on constructing compact 7-manifolds of $G_2$ holonomy that contain such singularities. Such manifolds, having an odd dimension, are obviously not complex manifolds. The math is much tougher when techniques of complex analysis cannot be used. Thus no quasi-realistic models of this type have been constructed. Nonetheless, it is possible to make some general statements, as in [45], for example.

3.2.3 Intersecting D-branes in Type II

Consider Type IIA superstring theory compactified on a six-torus. Suppose that a stack of $M$ D6-branes wraps a 3-cycle of the torus and fills the four-dimensional spacetime. This gives a supersymmetric $U(M)$ gauge theory. Suppose now that a second stack of $N$ D6-branes wraps a different 3-cycle that intersects the first one at $n$ points. Then the four-dimensional theory has $U(M) \times U(N)$ gauge symmetry, and it contains $n$ bifundamental chiral matter multiplets transforming as $(M, \bar{N})$.

It is striking that constructions that are this simple can capture roughly the types of groups and representations that appear in the standard model. Moreover, here are generalizations involving orbifolds and orientifolds that are quite realistic. But, even then, there are some challenging problems that remain:

- Not easy to get the correct $U(1)$ of weak hypercharge and no other $U(1)$ groups at low energy. Some $U(1)$s are lifted by a four-dimensional version of the anomaly cancellation mechanism. However, this does not automatically guarantee that just the desired $U(1)$ group survives.
• There are many massless moduli. Localizing the branes at singularities eliminates some of them. However, it is hard to remove all of them, as is required. One approach to overcoming this problem is to have the brane intersections be localized at singularities of the geometry. As an example of this approach see [46].

• This type of scheme does not incorporate grand unification. It is not certain that we should, since the hints of grand unification could prove to be illusory. However, the more conservative guess is that it is not an illusion.

For reviews of this remarkable subject see [47, 48, 49].

3.2.4 Type II on a Calabi–Yau with Flux

Calabi–Yau compactification of Type II superstring theories

\[ M_{10} = CY_6 \times M_{3,1} \]

give \( \mathcal{N} = 2 \) supersymmetry in four dimensions. Also, it does not give non-abelian gauge symmetry or chiral matter. So, without further measures, this is completely unrealistic. By including D-branes and/or flux one can overcome these difficulties. In fact, as discussed in the first lecture, the flux even helps to stabilize moduli. It also gives warping of the geometry, as was also discussed.

This approach allows one to embed Randall–Sundrum models in a string-theoretic setting. Recall that RS proposed that the warping could used to solve the hierarchy problem and that it could make quantum gravity effects accessible at LHC energies. In this setting one can break supersymmetry by adding anti-D3-branes. As we discussed, this approach leads to interesting string-theoretic models of inflation in which the origin of the big bang is described as a brane-antibrane collision. However, it has not yet yielded quasi-realistic particle physics models as far as I am aware.

The very large number of vacua that typically arise in flux compactifications may have discouraged some people from trying to construct viable particle physics models. However, there is still a considerable effort underway exploring the possibilities for constructing realistic examples. See, for example, [50, 51]. This large number of vacua has motivated the suggestion that various parameters of Nature (such as the cosmological constant) should be studied statistically on the landscape. This approach seems to assume implicitly that Nature corresponds to a more or less random vacuum. This in turn is motivated by some notion about how Universes are spawned in the Multiverse in a process of eternal inflation. The story gets even more entangled when the anthropic principle is brought into the discussion. Some people are enthusiastic about this type of reasoning, but I find it fundamentally defeatist. Even though in some sense it may be correct, I remain optimistic that there is a more reductionist way of understanding particle physics. The next section describes an approach that is more in that spirit.
### 3.3 F-Theory Local Models

#### 3.3.1 Local Versus Global Models

In *global models* one fully specifies the compactification geometry and all background fluxes and other fields. This completely determines a vacuum and hence the physics at all scales – up to and beyond the Planck scale. The heterotic string constructions that we have discussed are examples of global models. *Local models* are less ambitious. They are supposed to describe particle physics in a limit where gravity is turned off. More precisely, the ratio of the Planck scale to the unification scale, $M_{\text{Pl}}/M_{\text{GUT}}$ can be made arbitrarily large.

In certain classes of models in which matter fields are restricted to branes, the low-energy physics only depends on the geometry in the vicinity of the branes (in the extra dimensions). If one can decompactify the extra dimensions in directions transverse to the branes, this has the effect of increasing the four-dimensional Planck mass while keeping other scales fixed. Then one only needs to focus on the geometry in the immediate vicinity of the branes losing sensitivity to features in the geometry that are far away from the relevant branes. This approach is sometimes called *geometric engineering*.

Such local models are more amenable to a “bottom-up” construction. The basic idea is that one should first try to “engineer” a phenomenologically viable local model that correctly describes all nongravitational physics up to the unification scale to good accuracy. Then one should try to find a “global completion” of this local model that fully describes the geometry of the compact dimensions and brings gravity back into the mix. It is plausible (but not demonstrated, as far as I know) that the vast majority of local models do not have a global completion. Thus the existence of a global completion would be a confirmation that the original local model was on the right track. Furthermore, the global completion is likely to require that parameters (and other features) are correlated even though they were not correlated in the original construction of the local model. Altogether, one could imagine that the construction of a viable local model would be somewhat predictive, and that the construction of its global completion would be more predictive.

#### 3.3.2 What is F-Theory?

Type IIB superstring theory [13] is the setting for F-theory. Type IIB superstring theory is also the framework where most of the recent progress in flux compactifications, moduli stabilization, and SUSY breaking has been made. F-theory, which utilizes nonperturbative type IIB vacua containing 7-branes, was introduced in 1996 by Vafa [52]. F-theory configurations that give four-dimensional models contain various 7-branes wrapping compact 4-cycles of the six-dimensional compact space $B$. They also fill the four-dimensional Minkowski spacetime.

The type IIB theory in ten-dimensional Minkowski spacetime contains two massless scalar
fields, often called $C_0$ and $\phi$, which are conveniently combined to form a complex scalar field

$$\tau = C_0 + ie^{-\phi}.$$  

The field $C_0$ is a Ramond-Ramond zero-form, which couples electrically to D-instantons and magnetically to D7-branes. It behaves, at least in some respects, like a ten-dimensional analog of an axion. The field $\phi$ is called the dilaton. It has the special property that the value of $e^\phi$ is the string coupling strength, i.e., $g_s = \langle e^\phi \rangle$.

Type IIB superstring theory has a nonperturbative discrete gauge symmetry given by $[18]$ \[ \tau \to \tau' = \frac{a\tau + b}{c\tau + d}, \quad \text{where} \quad \begin{pmatrix} a & b \\ c & d \end{pmatrix} \in \text{SL}(2, \mathbb{Z}). \]

This means that $a, b, c, d$ are integers satisfying $ad - bc = 1$. Other fields transform at the same time. A special case is the transformation $\tau \to -1/\tau$. For $C_0 = 0$ this maps $\phi \to -\phi$ and hence $g_s \to 1/g_s$, which is S-duality. The group $\text{SL}(2, \mathbb{Z})$ is also the modular group of the torus. As often happens in string theory, things that look the same actually are the same. In this context this means that $\tau$ can be interpreted as the complex structure of a torus. So the picture is that to every point $y \in B$ one associates a torus with the complex structure $\tau(y)$. As $\tau$ approaches certain special values the torus becomes singular. (Mathematicians say that the torus degenerates.) This happens generically on the 4-cycles where the 7-branes are located.

When one analytically continues $\tau(z)$ along a curve encircling a 7-brane it comes back to a transformed value given by an $\text{SL}(2, \mathbb{Z})$ transformation. The 7-branes are classified by these transformations (more precisely the conjugacy classes), so there are an infinite number of distinct types of 7-branes. They can be described by a pair of integers $(p, q)$ without any common divisors. Type IIB superstring theory also has an infinite number of different types of strings with a similar classification scheme. In fact, a $(p, q)$ 7-brane is one on which a $(p, q)$ string can end. A special case of this statement is that a fundamental string can end on a D7-brane. The associated $\text{SL}(2, \mathbb{Z})$ transformation in this case is $\tau \to \tau + 1$. In this case the dilaton is constant and only $C_0$ varies, so it is possible to have weak coupling and use perturbative string techniques.

For a consistent F-theory solution, the varying field $\tau(z)$ and the compact six-dimensional space $B$ are required to define an elliptically fibred Calabi-Yau four-fold (with section). Then, the low-energy effective four-dimensional theory has $\mathcal{N} = 1$ supersymmetry. One subtlety is that this construction makes it seem as if the compact space is eight-dimensional, since that is the dimension of a Calabi-Yau four-fold. If this were the case the entire spacetime dimension would be 12. However, this is definitely not the case. The tori that describe the fibres have a complex structure, given by $\tau(z)$, but they have zero area. This can be understood as a limit of tori with finite area, but a complete explanation requires invoking a duality relating F-theory to M-theory, which I will not present here. There had been a
considerable amount of progress in understanding F-theory before this year [53, 54, 55], but it was not clear how to go about making realistic models.

### 3.3.3 Construction of Local F-Theory Models

The F-theory approach to phenomenology has been pursued this year by various groups starting with Donagi and Wijnholt and Beasley, Heckman, and Vafa [56] – [65]. The new proposal, which has given the subject a new lease on life, is to focus on models in which one can define a limit in which gravity is decoupled. The criterion is that it should be possible to make the dimensions transverse to the 4-cycles wrapped by the 7-branes arbitrarily large. Equivalently, it should be possible to contract the 4-cycles to points while holding the six-dimensional volume fixed. Such contractible 4-cycles must be positive curvature Kähler manifolds. These are fully classified and are given by manifolds called del Pezzo manifolds (or del Pezzo surfaces), which are denoted dPₙ. The integer n takes the values 0 ≤ n ≤ 8.

The del Pezzos have a close relationship with the exceptional Lie algebras Eₙ. The basic idea is that they contain 2-cycles whose intersections are characterized by the Eₙ Dynkin diagram. By this type of F-theory construction, one can construct an SU(5) or SO(10) SUSY-GUT model.

Constructions that involve 7-branes of various types are much more subtle – and also more interesting than ones that only involve D7-branes. D7-branes are mutually local. A stack of N of them gives U(N) gauge symmetry. Matter fields at intersections (due to stretched open strings) are bifundamental. However, different kinds of 7-branes are mutually nonlocal. As a result, there are stacks (corresponding to the ADE classification of singularities) that can give U(N), SO(2N) or even Eₙ gauge symmetry. Exceptional gauge groups, spinor representations of SO(10), and nonzero top-quark Yukawa couplings in SU(5) are possible for F-theory vacua. In constructions that only involve D7-branes these phenomena can only arise nonperturbatively. Since the top quark Yukawa coupling must be quite large to account for the large top quark mass, the F-theory viewpoint seems preferable.

The crucial assumption of the local F-theory approach – the existence of a decoupling limit in which four-dimensional gravity can be turned off – may or may not be physically correct. The fact that $M_{\text{Pl}}/M_{\text{GUT}}$ is in the range 100–1000, depending on the precise definition of each of these mass scales, makes it plausible. There are an infinite variety of types of 4-cycles that 7-branes could wrap if one did not make this assumption. Therefore this principle picks out a tiny corner of the string theory landscape. In this way, the doom and gloom that the discovery of this landscape has generated is swept aside, at least if the basic assumption is correct.

In particular, it allows one to construct an SU(5) or SO(10) SUSY-GUT with three families. The SU(5) case is more straightforward. There are three basic rules:

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9There is also one other possibility, which is a product of two two-spheres.
• Each stack of 7-branes on a 4-cycle $S$ determines a gauge group. One stack should give the GUT gauge group. Other “matter 7-branes” are also needed.

• Pairs of 4-cycles intersect on 2-cycles $S \cap S'$. Each such 2-cycle is a Riemann surface that has associated chiral matter fields given by zero-modes of the Dirac operator. The number of such zero modes determines the number of families.

• Three 2-cycles can intersect at points. Products of wave functions at these intersections determine the Yukawa couplings.

These three rules can be summarized by saying that the gauge fields live in eight dimensions, the matter fields live in six dimensions, and the interactions take place in four dimensions. Gravity, which is ignored in a local construction, lives in ten dimensions.

### 3.3.4 Gauge Symmetry Breaking

One of the strengths of the local F-theory approach to model building is the way that symmetry breaking can be incorporated. In conventional SUSY-GUTS based on four-dimensional quantum field theory, the breaking of the $SU(5)$ GUT symmetry to the $SU(3) \times SU(2) \times U(1)$ standard model gauge group requires the introduction of Higgs fields belonging to large representations of $SU(5)$. This is definitely ugly, and does not work very well. In string theory such representations are not present in the low mass spectrum, but there are other available mechanisms. In the case of the heterotic string constructions the standard method introduced in [16] is to associate nontrivial Wilson lines to noncontractible cycles of the Calabi–Yau geometry.

In the local F-theory models there are no noncontractible cycles and a different mechanism for breaking the $SU(5)$ gauge symmetry to the standard model gauge group is required. Fortunately, one is available. The proposal of Vafa and collaborators is introduce nonzero flux. In doing this one has to choose the $U(1)$ subgroup of $SU(5)$ that this flux corresponds to. By choosing it to correspond to the weak hypercharge $U(1)$ subgroup one gets exactly the desired symmetry breaking pattern. This flux is called hyperflux in the recent literature. One also has to choose the 2-cycle inside the del Pezzo surface that carries this flux. This is an elaborate and beautiful story that I don’t intend to explain in detail here. One crucial point is that by making appropriate choices it is possible to arrange to obtain complete $SU(5)$ multiplets when one wants them and incomplete multiplets when that is desired. In the case of the quarks and leptons one wants complete $\bar{5}$ and $10$ multiplets. In the case of the Higgs bosons one wants only the doublet pieces of the $5$ and $\bar{5}$ multiplets, eliminating the triplets, which would give rapid proton decay. All this can be elegantly arranged.

One also needs to incorporate electroweak symmetry breaking. The standard Higgs mechanism is triggered by the breaking of supersymmetry. Supersymmetry requires that are two Higgs doublets, of course.
3.3.5 Supersymmetry Breaking, etc.

Local models are incompatible with gravity-mediated SUSY breaking. However, for gauge-mediated SUSY breaking both the visible and the hidden sectors, as well as their mediation, can be described in a single effective field theory decoupled from gravity. According to most experts, such models are favored for suppressing flavor-changing neutral currents. When the model is extended to incorporate a global completion and bring gravity back in the story, there will be a small admixture of gravity mediation. A proposal in the recent literature is that gauge-mediated supersymmetry breaking arises due to stringy instantons (Euclidean D3-branes that wrap the 4-cycle $S$) on an auxiliary 7-brane that has been dubbed a Peccei-Quinn 7-brane.

Beasley, Heckman, and Vafa [57, 59] have also proposed interesting mechanisms by which local F-theory models can address each of the following:

- R-parity conservation
- Peccei–Quinn symmetry and axion properties
- the $\mu$ and $\mu/B_\mu$ problems
- the size of neutrino masses
- proton decay
- hierarchies of Yukawa couplings

There are many more details of this approach that I cannot go into here, including singularity enhancement, classification of fluxes, nonexistence of heterotic and M-theory duals, determination of superpotentials, threshold corrections, dark matter, sparticle spectrum, Higgsing the Peccei–Quinn symmetry, messenger fields, etc. This subject is still in its infancy, but it shows great promise. It will be very interesting to see what predictions can be made before the experimental results pour in and whether they turn out to be correct.

4 Conclusion

The understanding of string theory continues to progress on many fronts at an impressive rate. Many challenges remain, and it will undoubtedly require many decades to answer some of the deepest questions. However, there is no reason to be pessimistic. The subject is already proving useful for various purposes ranging from fundamental mathematics to particle physics phenomenology and cosmology. New experimental and observational results will be very helpful in stimulating further progress.
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References

[1] M. B. Green, J. H. Schwarz, and E. Witten, *Superstring Theory* in 2 Volumes, Cambridge Univ. Press, 1987.

[2] J. Polchinski, *String Theory* in 2 Volumes, Cambridge Univ. Press, 1998.

[3] B. Zwiebach, *A First Course in String Theory*, Cambridge Univ. Press, 2004.

[4] K. Becker, M. Becker, and J. H. Schwarz, *String Theory and M-Theory: A Modern Introduction*, Cambridge Univ. Press, 2007.

[5] J. H. Schwarz, “The Early Years of String Theory: A Personal Perspective,” arXiv:0708.1917 [hep-th].

[6] G. Veneziano, “Construction of a Crossing-Symmetric Regge-Behaved Amplitude for Linearly Rising Regge Trajectories,” *Nuovo Cim.* **57A**, 190 (1968).

[7] P. Ramond, “Dual Theory for Free Fermions,” *Phys. Rev.* **3**, 2415 (1971).

[8] A. Neveu and J. H. Schwarz, “Factorizable Dual Model of Pions,” *Nucl. Phys.* B **31**, 86 (1971).

[9] J. Scherk and J. H. Schwarz, “Dual Models for Non-Hadrons,” *Nucl. Phys.* B **81**, 118 (1974).

[10] T. Yoneya, “Quantum Gravity and the Zero Slope Limit of the Generalized Virasoro Model,” *Nuovo Cim. Lett.* **8**, 951 (1973).

[11] T. Yoneya, “Connection of Dual Models to Electrodynamics and Gravidynamics,” *Prog. Theor. Phys.* **51**, 1907 (1974).

[12] F. Gliozzi, J. Scherk, and D. Olive, “Supersymmetry, Supergravity Theories and the Dual Spinor Model,” *Nucl. Phys.* B **122**, 253 (1977).

[13] M. B. Green and J. H. Schwarz, “Supersymmetrical String Theories,” *Phys. Lett.* B **109**, 444 (1982).
[14] M. B. Green and J. H. Schwarz, “Anomaly Cancellation in Supersymmetric $D = 10$
Gauge Theory and Superstring Theory,” Phys. Lett. B 149, 117 (1984).

[15] D. J. Gross, J. A. Harvey, E. J. Martinec and R. Rohm, “The Heterotic String,” Phys. 
Rev. Lett. 54, 502 (1985).

[16] P. Candelas, G. T. Horowitz, A. Strominger and E. Witten, “Vacuum Configurations 
for Superstrings,” Nucl. Phys. B 258, 46 (1985).

[17] A. Sen, “Strong - Weak Coupling Duality in Four-Dimensional String Theory,” Int. J. 
Mod. Phys. A 9, 3707 (1994) [arXiv:hep-th/9402002].

[18] C. M. Hull and P. K. Townsend, “Unity of Superstring Dualities,” Nucl. Phys. B 438, 
109 (1995) [arXiv:hep-th/9410167].

[19] E. Witten, “String Theory Dynamics in Various Dimensions,” Nucl. Phys. B 443, 85 
(1995) [arXiv:hep-th/9503124].

[20] E. Cremmer, B. Julia and J. Scherk, “Supergravity Theory in 11 Dimensions,” Phys. 
Lett. B 76, 409 (1978).

[21] J. M. Maldacena, “The Large $N$ Limit of Superconformal Field Theories and Super-
gravity,” Adv. Theor. Math. Phys. 2, 231 (1998) [arXiv:hep-th/9711200].

[22] S. S. Gubser, I. R. Klebanov and A. M. Polyakov, “Gauge theory correlators from 
non-critical string theory,” Phys. Lett. B 428, 105 (1998) [arXiv:hep-th/9802109].

[23] E. Witten, “Anti-de Sitter space and holography,” Adv. Theor. Math. Phys. 2, 253 
(1998) [arXiv:hep-th/9802150].

[24] O. Aharony, O. Bergman, D. L. Jafferis and J. Maldacena, “N=6 Superconformal Chern-
Simons-Matter Theories, M2-branes and their Gravity Duals,” JHEP 0810, 091 (2008) 
arXiv:0806.1218 [hep-th].

[25] T. Sakai and S. Sugimoto, “Low Energy Hadron Physics in Holographic QCD,” Prog. 
Theor. Phys. 113, 843 (2005) [arXiv:hep-th/0412141].

[26] A. M. Polyakov, “The Wall of the Cave,” Int. J. Mod. Phys. A 14, 645 (1999) 
arXiv:hep-th/9809057.

[27] G. Policastro, D. T. Son and A. O. Starinets, “The Shear Viscosity of Strongly Cou-
pled $N = 4$ Supersymmetric Yang-Mills Plasma,” Phys. Rev. Lett. 87, 081601 (2001) 
arXiv:hep-th/0104066.
[28] P. Kovtun, D. T. Son and A. O. Starinets, “Viscosity in Strongly Interacting Quantum Field Theories from Black Hole Physics,” Phys. Rev. Lett. **94**, 111601 (2005) [arXiv:hep-th/0405231].

[29] S. Sachdev and M. Mueller, “Quantum Criticality and Black Holes,” [arXiv:0810.3005 [cond-mat.str-el]].

[30] K. Balasubramanian and J. McGreevy, “Gravity Duals for Non-relativistic CFTs,” Phys. Rev. Lett. **101**, 061601 (2008) [arXiv:0804.4053 [hep-th]].

[31] D. T. Son, “Toward an AdS/Cold Atoms Correspondence: a Geometric Realization of the Schroedinger Symmetry,” Phys. Rev. D **78**, 046003 (2008) [arXiv:0804.3972 [hep-th]].

[32] M. R. Douglas, “The Statistics of String/M Theory Vacua,” JHEP **0305**, 046 (2003) [arXiv:hep-th/0303194].

[33] M. R. Douglas and S. Kachru, “Flux Compactification,” Rev. Mod. Phys. **79**, 733 (2007) [arXiv:hep-th/0610102].

[34] L. Randall and R. Sundrum, “A Large Mass Hierarchy from a Small Extra Dimension,” Phys. Rev. Lett. **83**, 3370 (1999) [arXiv:hep-ph/9905221].

[35] S. B. Giddings, S. Kachru and J. Polchinski, “Hierarchies from Fluxes in String Compactifications,” Phys. Rev. D **66**, 106006 (2002) [arXiv:hep-th/0105097].

[36] S. Kachru, R. Kallosh, A. Linde and S. P. Trivedi, “De Sitter Vacua in String Theory,” Phys. Rev. D **68**, 046005 (2003) [arXiv:hep-th/0301240].

[37] S. Kachru, R. Kallosh, A. Linde, J. M. Maldacena, L. P. McAllister and S. P. Trivedi, “Towards Inflation in String Theory,” JCAP **0310**, 013 (2003) [arXiv:hep-th/0308055].

[38] J. Polchinski, “Cosmic String Loops and Gravitational Radiation,” [arXiv:0707.0888 [astro-ph]].

[39] P. Hořava and E. Witten, “Heterotic and Type I String Dynamics from Eleven Dimensions,” Nucl. Phys. B **460**, 506 (1996) [arXiv:hep-th/9510209].

[40] P. Hořava and E. Witten, “Eleven-Dimensional Supergravity on a Manifold with Boundary,” Nucl. Phys. B **475**, 94 (1996) [arXiv:hep-th/9603142].

[41] E. Witten, “Strong Coupling Expansion of Calabi-Yau Compactification,” Nucl. Phys. B **471**, 135 (1996) [arXiv:hep-th/9602070].
[42] V. Bouchard and R. Donagi, “An SU(5) Heterotic Standard Model,” Phys. Lett. B 633, 783 (2006) [arXiv:hep-th/0512149].

[43] V. Braun, Y. H. He, B. A. Ovrut and T. Pantev, “The Exact MSSM Spectrum from String Theory,” JHEP 0605, 043 (2006) [arXiv:hep-th/0512177].

[44] B. Acharya and E. Witten, “Chiral Fermions from Manifolds of G(2) Holonomy,” arXiv:hep-th/0109152.

[45] B. S. Acharya and K. Bobkov, “Kahler Independence of the G2-MSSM,” arXiv:0810.3285 [hep-th].

[46] H. Verlinde and M. Wijnholt, “Building the Standard Model on a D3-brane,” JHEP 0701, 106 (2007) [arXiv:hep-th/0508089].

[47] R. Blumenhagen, M. Cvetic, P. Langacker and G. Shiu, “Toward Realistic Intersecting D-Brane Models,” Ann. Rev. Nucl. Part. Sci. 55, 71 (2005) [arXiv:hep-th/0502005].

[48] R. Blumenhagen, B. Kors, D. Lust and S. Stieberger, “Four-Dimensional String Compactifications with D-Branes, Orientifolds and Fluxes,” Phys. Rept. 445, 1 (2007) [arXiv:hep-th/0610327].

[49] D. Malyshev and H. Verlinde, “D-branes at Singularities and String Phenomenology,” Nucl. Phys. Proc. Suppl. 171, 139 (2007) [arXiv:0711.2451 [hep-th]].

[50] D. Lust, S. Reffert and S. Stieberger, “Flux-induced Soft Supersymmetry Breaking in Chiral Type IIB Orientifolds with D3/D7-Branes,” Nucl. Phys. B 706, 3 (2005) [arXiv:hep-th/0406092].

[51] F. Marchesano and G. Shiu, “Building MSSM Flux Vacua,” JHEP 0411, 041 (2004) [arXiv:hep-th/0409132].

[52] C. Vafa, “Evidence for F-Theory,” Nucl. Phys. B 469, 403 (1996) [arXiv:hep-th/9602022].

[53] R. Friedman, J. Morgan and E. Witten, “Vector Bundles and F Theory,” Commun. Math. Phys. 187, 679 (1997) [arXiv:hep-th/9701162].

[54] M. Bershadsky, A. Johansen, T. Pantev and V. Sadov, “On Four-Dimensional Compactifications of F-Theory,” Nucl. Phys. B 505, 165 (1997) [arXiv:hep-th/9701165].

[55] B. Andreas and G. Curio, “On Discrete Twist and Four-Flux in N = 1 Heterotic/F-theory Compactifications,” Adv. Theor. Math. Phys. 3, 1325 (1999) [arXiv:hep-th/9908193].
[56] R. Donagi and M. Wijnholt, “Model Building with F-Theory,” [arXiv:0802.2969] [hep-th].
[57] C. Beasley, J. J. Heckman and C. Vafa, “GUTs and Exceptional Branes in F-theory - I,” [arXiv:0802.3391] [hep-th].
[58] H. Hayashi, R. Tatar, Y. Toda, T. Watari and M. Yamazaki, “New Aspects of Heterotic–F Theory Duality,” Nucl. Phys. B 806, 224 (2009) [arXiv:0805.1057] [hep-th].
[59] C. Beasley, J. J. Heckman and C. Vafa, “GUTs and Exceptional Branes in F-theory - II: Experimental Predictions,” [arXiv:0806.0102] [hep-th].
[60] J. J. Heckman, J. Marsano, N. Saulina, S. Schafer-Nameki and C. Vafa, “Instantons and SUSY breaking in F-theory,” [arXiv:0808.1286] [hep-th].
[61] J. Marsano, N. Saulina and S. Schafer-Nameki, “Gauge Mediation in F-Theory GUT Models,” [arXiv:0808.1571] [hep-th].
[62] R. Donagi and M. Wijnholt, “Breaking GUT Groups in F-Theory,” [arXiv:0808.2223] [hep-th].
[63] J. Marsano, N. Saulina and S. Schafer-Nameki, “An Instanton Toolbox for F-Theory Model Building,” [arXiv:0808.2450] [hep-th].
[64] J. J. Heckman and C. Vafa, “F-theory, GUTs, and the Weak Scale,” [arXiv:0809.1098] [hep-th].
[65] J. J. Heckman and C. Vafa, “Flavor Hierarchy From F-theory,” [arXiv:0811.2417] [hep-th].