String theory: an update

J. de Boer

Institute for Theoretical Physics, University of Amsterdam
Valckenierstraat 65, 1018 XE Amsterdam, The Netherlands

An overview of some of the developments in string theory over the past two years is given, focusing on four topics: realistic (standard model like) models from string theory, geometric engineering and theories with fluxes, the gauge theory-gravity correspondence, and time dependent backgrounds and string theory. Plenary talk at ICHEP’02, Amsterdam, July 24-31, 2002.

1. INTRODUCTION

In this talk I would like to discuss some of the developments that took place in string theory over the last two years, after the previous ICHEP meeting that took place in Osaka. Of course, it is impossible to give a detailed account of everything that happened in the field. In fact, it is not even clear precisely what the field is. According to some people, “anything that appears on hep-th" is a good first approximation to the term “string theory." Therefore, in order to have some degree of organization and limit the material, I have selected four themes. Within each of the four themes, there has been a significant amount of activity over the past two years. The themes are (1) getting realistic (standard model like) models from string theory, (2) geometric engineering and theories with fluxes, (3) the gauge theory-gravity correspondence, and (4) time dependent backgrounds and string theory. I will discuss each of the themes, in this order, in the four sections that follow.

Presenting such a general talk about string theory is a difficult enterprise, and therefore I would like to start with a disclaimer. Since the purpose of this talk is to give a flavor of what has been happening in string theory over the past two years, many interesting subjects will not be discussed. This by no means implies that these subjects are less interesting or important, there simply is too much material to cover and a selection has to be made. In a similar vein, the list of references will be extremely incomplete, and I apologize in advance for all omitted references, and for all places where I did not give appropriate credit.

An introduction to string theory and some general background material may be found e.g. in the books [1], review articles [2], popular articles [3] and the web page [4].

2. REALISTIC MODELS FROM STRING THEORY

The basic idea of string theory is to replace all point particles in nature by small vibrating strings. Different ways in which the string can vibrate manifest themselves as different particles. Thus all particles become different excitations of the same object. The typical size of a string is called the string length $l_s$ and typically of the order of the Planck length, i.e. $l_s \sim 10^{-35}$ cm. The main success of string theory is that it can be made into a finite theory and quite generally contains both gauge theories and gravity. It is therefore a candidate theory that unifies all fundamental forces in nature. In particular it provides us with a finite theory of quantum gravity. As a field theory, gravity is not renormalizable, with the divergences coming from the fact that it is a theory of point-particles interacting at points in space-time. In string theory, there is no longer a well-defined point where strings interact, and this renders the theory finite.
It is not possible to directly see strings, as there are no experiments that involve the large energies of order $10^{19}$ GeV needed to resolve distances of order $10^{-33}$ cm. It is therefore an important question how we might be able to experimentally test string theory. Roughly, there are three classes of possible experimental signatures.

1. The very high energies that directly probe the string length scale $l_s$ cannot be realized in an experiment done on earth, but do certainly occur at extreme situations in the universe [5]. In particular, the physics of the early universe and the physics of black holes will crucially depend on a fundamental theory of quantum gravity. Observations of the universe e.g. by means of the cosmic microwave background, by means of (the still to be observed) gravitational radiation, or by means of high energy gamma rays may therefore provide experimental signatures of string theory.

2. At low energies, the standard model is very successful in describing the interactions of fundamental particles. String theory should be able to reproduce the standard model, and ideally also constrain its free parameters. This would be quite spectacular but has not yet been achieved. A less ambitious project would be to obtain a reasonable supersymmetric extension of the standard model, postponing the problem of supersymmetry breaking. Evidence for the existence of such a supersymmetric extension would for example be the discovery of a supersymmetric partner of one of the known particles. Because supersymmetry is a generic prediction of string theory, this would certainly be enthusiastically received by the string theory community.

3. One can contemplate other string theory models, such as ones involving “brane worlds” or “large extra dimensions,” where the string length is much larger than $10^{-33}$ cm. In these scenarios, a bewildering set of possible experiments has been proposed in the literature, including measurements of violations of Newtonian gravity at length scales less than 0.2 mm, and missing energy signatures and black hole production at accelerators; see also section 2.7.

In addition, it is worth noticing that in many cases it is difficult to say anything useful about each of the three points above. That still leaves an important theoretical experiment, namely

The theory should be self-consistent, and in particular one should be able to embed it consistently in string theory.

Often, it is amazingly difficult to even perform this theoretical experiment.

In this section, we want to mainly focus on the second point in the list above, in particular we would like to discuss the prospects for getting the standard model or one of its supersymmetric extensions from string theory. In the last two years, many new models have been proposed, and to explain how these fit together, it is perhaps best to first go back in time.

Soon after its discovery, it was realized that superstring theory is only consistent in ten dimensions. This is not a problem, since as long as six of the ten dimensions are compact and very small, the theory will at low energies look like a four-dimensional theory. The details of the four-dimensional theory depend crucially on the way in which the six extra dimensions are compactified. In particular, the amount of supersymmetry that remains in four dimensions depends on the choice of 6d geometry. Among the several known ten-dimensional superstring theories, there was only one that could be compactified in such a way that only $N = 1$ supersymmetry in four dimensions remained. This string theory is called the heterotic string, and the six-dimensional geometry that is needed is a so-called Calabi-Yau manifold [6].

In 1995 the picture changed dramatically in what is known as the second superstring revolution [7]. It was realized that the different ten-dimensional string theories were not inequivalent theories, but were all different weakly coupled limits of a single theory. This is illustrated in figure 1, taken from [8].
This picture illustrates the existence of one big theory, that depends on many parameters. Depending on the choice of parameters, there is at most one weakly coupled description of the theory, as illustrated by the peaks in the picture. In the middle of the picture, all coupling constants are of order one, and there is no good weakly coupled description. Since the picture is connected, it is possible to move from one theory to the other by changing parameters. In particular, the theory one starts out with will become strongly coupled by the time the other theory becomes weakly coupled. These strong-weak coupling dualities play an important role in string theory. A simple example is the duality between M-theory and the type IIA superstring. Actually, M-theory is a somewhat peculiar element in this picture, because it is not really a string theory. It is a theory whose low-energy limit is eleven-dimensional supergravity. Above eleven-dimensions, supersymmetry necessarily involves fields of spin larger than two, and such theories (with finitely many higher spin fields) do not exist. There are indications that membranes play an important role in M-theory, but if so their role is certainly different from that of strings in string theory. In any case, the claim is that M-theory compactified on a circle of radius $R$ yields the type IIA superstring. This is certainly true at low energies, as one can explicitly check that eleven-dimensional supergravity compactified on a circle yields ten-dimensional type IIA supergravity, which is the low-energy limit of type IIA superstrings (hence the name). The string coupling constant (usually denoted by $g_s$) that measures the strength of string interactions is proportional to the radius $R$. For small $R$, the coupling constant is small, and type IIA string theory is the weakly coupled description. At large $R$, the coupling constant is large, type IIA becomes strongly coupled, and a corresponding weakly coupled description is in terms of M-theory.

The picture in figure 1 describes theories with a lot of supersymmetry. One can try to make a similar “duality web” that describes the situation of string theory compactified down to four dimensions with $N = 1$ supersymmetry remaining. A caricature of such a picture is given in figure 2. This figure is necessarily incomplete, since there are many more possible weakly coupled possibilities than are shown in figure 2. In addition, it is not even known for sure whether the picture should be connected or in reality consists of several disconnected components. Nevertheless, we will use this picture as a guiding principle to discuss some of the new ways in which one can obtain $N = 1$ supersymmetric theories in four dimensions (theories with fluxes will be discussed in section 3). Ideally, one would like to study models with no supersymmetry at all (like the standard model), but supersymmetry breaking remains an interesting and difficult problem in string theory. It is possible to explicitly break supersymmetry, but then it is often difficult to examine whether the theory is stable or not. Therefore we will restrict attention to $N = 1$ supersymmetric theories in what follows.

2.1. heterotic on CY$_3$

As we mentioned before, an $N = 1$ theory in four dimensions was first obtained by compactifying the heterotic string on a Calabi-Yau manifold. The 3 in CY$_3$ refers to the number of complex dimensions of the Calabi-Yau manifold. These models have many appealing properties. They
can naturally accommodate GUT groups such as $SU(5)$, $SO(10)$ and $E_6$, it is not too difficult to get three generations of chiral fermions in four dimensions, gauge coupling unification is quite natural, etc. The GUT scale is two or three orders of magnitude smaller than the Planck scale. There is a large number of models that contain the standard model fields, but no model that precisely gives the MSSM (minimally supersymmetric standard model). A generic problem of these models is the large number of massless scalar fields that they possess; these parameterize the different shapes and sizes of the Calabi-Yau manifold. It is possible that a potential for these scalar fields is dynamically generated once supersymmetry is broken, but unfortunately we do not know the precise mechanism for supersymmetry breaking. For more discussion of low energy phenomenology in string theory, see e.g. [9].

2.2. Horava-Witten

The Horava-Witten scenario [10] is based on a strong-weak coupling duality, as figure 2 suggests. The duality in question is that between the heterotic string on one side, and M-theory compactified on an interval on the other side. Although it may sound strange to compactify a theory on an interval, there is nothing wrong with this. It is similar to doing field theory in a finite box, though one has to be careful to choose appropriate boundary conditions for all fields at the endpoints of the interval. It turns out that consistency requires the introduction of additional boundary degrees of freedom that live only at the endpoints of the interval. These boundary degrees of freedom can freely propagate in the remaining ten dimensions, but are stuck in the eleventh dimension. In particular, they can freely propagate in the four dimensions that remain once the theory is further compactified on a Calabi-Yau space, and to the low-energy observer they look like conventional four-dimensional degrees of freedom. Sometimes one says that there are “end of the world branes” located at the endpoints of the interval that carry additional degrees of freedom.

The duality with the heterotic string is a strong-weak coupling duality in the sense that for a small interval, the heterotic string coupling is small and that is the weakly coupled description, whereas for large string coupling the interval becomes larger and the M-theory point of view is more appropriate.

The phenomenology of these models is quite rich. It includes all heterotic compactifications, but there are additional possibilities. One can tune parameters in such a way that the GUT scale and the Planck scale coincide. The study of the Horava-Witten models involves some rather heavy geometric machinery, but there are indications that models with three generations and gauge group $SU(3) \times SU(2) \times U(1)$ can be obtained. For further discussion, see [11] and references therein.

2.3. M on $G_2$

The only theory whose low-energy limit is eleven dimensional supergravity is M-theory. In order to obtain a four-dimensional theory, seven dimensions need to be compactified. In case of the heterotic string, we needed a special type of manifold, namely a Calabi-Yau manifold, in order to have unbroken $N = 1$ supersymmetry in four dimensions. Similarly, in the case of M-theory, we need a special type of seven
manifold, so-called manifolds of $G_2$ holonomy, to have unbroken $N = 1$ supersymmetry in four dimensions\(^2\). The geometry of such seven-dimensional manifolds is more complicated and less well-understood than the geometry of Calabi-Yau manifolds. Several examples of compact seven manifolds of $G_2$ holonomy are known \[12\]. One also often studies non-compact manifolds of $G_2$ holonomy. These are not useful to build a realistic theory in four dimensions, because on non-compact manifolds particles can have arbitrarily small momenta in the non-compact directions. These manifest themselves as a continuum of particles in four dimensions, which clearly is not a very realistic feature. However, non-compact manifolds provide useful toy models to understand the structure of singularities. Singularities in $G_2$ manifolds are of crucial importance because they are the only source of chiral fermions in such compactifications, a fact that further complicates matters and makes it difficult to obtain realistic models; for more discussion see e.g. \[13\].

### 2.4. $F$ on CY\(_4\)

In this title $F$ refers to F-theory, introduced in \[14\]. F-theory is a twelve-dimensional theory, but in contrast to M-theory it is not generally covariant in twelve-dimensions. One may even wonder what is really meant by the word theory in “twelve-dimensional theory.” If nothing else, F-theory is a convenient way to describe in purely geometric terms certain strongly coupled type IIB superstring compactifications. To obtain a four-dimensional theory with $N = 1$ supersymmetry, we need to compactify F-theory on an eight-dimensional (or four complex-dimensional) Calabi-Yau manifold. Though the number of possible compactifications of F-theory is very large, we are not aware of any specific phenomenologically appealing model that can only be realized in this context.

\(^2\)The $G_2$ refers to the exceptional Lie group $G_2$. It can be defined as the subgroup of $SO(7)$ that preserves a single spinor in the eight-dimensional spinor representation of $SO(7)$. This spinor is responsible for the unbroken $N = 1$ supersymmetry in four dimensions.

### 2.5. duality between branes and geometry

Before continuing with our discussion of the remaining possibilities given in figure 2, we would like to briefly discuss the notion of branes, and how theories with branes can be dual to theories without branes. More information on this subject can be found in most of the reviews listed under \[12,13\].

D-branes are certain extended objects in string theory, that were introduced by Polchinski in \[15\]. They are labelled by the number of dimensions of the object, so that a D0 brane is like a particle, a D1 brane is like a string, a D2 brane is like a membrane, etc. There are two ways to think about D-branes. On the one hand, they are solitonic solutions of the equations of motion of low-energy closed string theory. On the other hand, they are objects in open string theory with the property that open strings can end on them. Open strings have a finite tension, and their center of mass cannot be taken arbitrarily far away from the D-brane. As a consequence, the degrees of freedom of the open string can effectively only propagate in a direction parallel to the brane: one says that they are confined to the brane, or that they live on the brane. The open string spectrum can be reproduced directly from the soliton in the closed string description via a collective coordinate quantization.

As a very crude analogy, one can think about two ways to describe a monopole. On the one hand, one can think of it as the ’t Hooft-Polyakov monopole, in which case it is an extended soliton solution of the Yang-Mills-Higgs equations of motion. On the other hand, one can view a monopole as a point particle, on which magnetic field lines can end. Both descriptions have their advantages, as do the open and closed string descriptions of D-branes.

D-branes play a crucial role in string theory and in particular in string dualities. For example, type IIB string theory has a strong-weak coupling duality that inverts the the type IIB string coupling constant, $g_s \leftrightarrow 1/g_s$. Under this duality, the roles of fundamental strings and D1 branes are interchanged (both are one-dimensional objects). Therefore, D branes appear to be as fundamental as strings themselves. There is a close analogy
between this duality and the electric-magnetic duality of the Maxwell equations, with strings playing the role of electric charges, and D-branes the role of magnetic charges.

Interestingly, D branes can not only be dual to other branes and strings, but also to non-trivial geometries without branes. To explain this peculiar statement, we will use the duality between M-theory compactified on a circle, and type IIA string theory. M-theory at low energies was described by eleven-dimensional supergravity. When compactifying this on a circle, the eleven dimensional metric was part of the eleven-dimensional metric. When compactifying this on a circle, the eleven dimensional supergravity was described by eleven-dimensional super-type IIA string theory. M-theory at low energies was described by eleven-dimensional supergravity. When compactifying this on a circle, the eleven dimensional metric was part of the eleven-dimensional metric.

Now type IIA string theory has objects that are charged with respect to this gauge field. In four dimensions, charges can be measured by looking at the flux of the electric or magnetic field through a two-sphere $S^2$ surrounding the charge. In other words, we compute $\int_{S^2} F$, with $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$ the field strength. A point particle in ten dimensions on the other hand does not carry a charge under such an $F$, because we cannot surround it by a two-sphere. We can surround it by an eight-sphere, and if we would have an eight-index field strength we could integrate over the eight sphere and use that to define a charge. However, besides string excitations, type IIA string theory also has D-branes, and the list of possible D-branes includes a D6-brane. A D6 brane can be surrounded by a two sphere in ten dimensions. In general, in $d$ dimensions a $Dp$ brane can be surrounded by a $d-2-p$ sphere. It is easy to verify this directly in four and lower dimensions. Thus, a D6 brane can carry a charge with respect to the gauge field $A_\mu$ of ten-dimensional type IIA supergravity, measured by the flux through the two-sphere. One can show that it indeed does carry such a charge. In other words, a D6-brane is surrounded by a non-trivial configurations of the gauge field $A_\mu$.

What does a D6-brane look like in eleven-dimensions? In eleven dimensions, the gauge field $A_\mu$ was part of the eleven-dimensional metric. Therefore, from the eleven-dimensional perspective, a D6-brane is surrounded by a non-trivial geometry. Since the gauge field is no longer there, the D6 brane is no longer charged under any-thing. It has become a purely geometrical object. One can show that the geometry is the product of seven-dimensional Minkowski space and an Eguchi-Hanson space, which is a gravitational instanton in four dimensions.

2.6. intersecting branes

D-branes can also be used to construct many new string theories with $N = 1$ supersymmetry in four dimensions. The idea is to take any existing string theory, and to add branes to it. If we are interested in keeping Poincaré invariance in four dimensions, the branes need to have at least three space dimensions, so that they can completely fill the four noncompact dimensions where the low-energy observer lives. In the remaining compact six dimensions (seven for M-theory) the branes can have any shape and size that is compatible with supersymmetry. At low energies, besides the fields that we get from the string theory we started out with, we also get degrees of freedom from the various branes. In their presence, the theory necessarily contains open strings that start and end on the branes, and these give rise to additional degrees of freedom in four dimensions. The new degrees of freedom never contain gravity, but typically provide gauge fields and matter fields. Gravity still has its origin in the original closed string theory. The different origin of gauge fields and gravity plays an important role in the large extra dimension scenario in the next subsection.

There are many possible gauge groups and matter fields that one can get from branes. For example, nonabelian gauge symmetries can be obtained by putting several branes on top of each other. Open strings can stretch from each of the branes in the stack to any of the other branes. There are $N \times N$ different open strings and these transform in the adjoint representation of $U(N)$, where $N$ is the number of branes. Their degrees of freedom include a nonabelian $U(N)$ gauge field in four dimensions. A more complicated possibility is to have branes that intersect in the six (or seven) compact dimensions. A simplified picture of such intersections is two orthogonal planes that intersect along a line. Along the line, new degrees of freedom are localized due to the open strings
that stretch from one of the planes to the other. A similar story applies in higher dimensions, where the new fields that are confined to the intersections often give chiral matter in \( d = 4 \).

One of the simplest concrete models with intersecting branes is to start with the type IIA string theory compactified on a six-torus, i.e., each of the six compact dimensions is taken to be a circle. One can take D6 branes and add these to the theory. Since three of the six dimensions of the brane need to be along the three noncompact dimensions of our world, the remaining three have to be put on the six-torus; one can for example select three of the six compact directions, and demand that the brane fills those three. By selecting different sets of three of the six compact directions, we can put different sets of D6 branes on the six-torus, and this collection of intersecting D6 branes gives rise to interesting theories in four dimensions.

There are several advantages to such an approach. It is possible to obtain \( N = 1 \) theories in four dimensions without having to rely on the complicated geometry of Calabi-Yau manifolds, but one can work with things as simple as a six-torus. The matter content is relatively easy to control, simply by analyzing the open string that can stretch between the different branes. In this way, the chiral spectrum of the standard model and its minimal supersymmetric extension have been obtained. It is even possible to choose brane configurations that explicitly break supersymmetry, but the stability of these configurations is not clear. Another feature of intersecting brane models is the stability of the proton, because baryon number becomes a gauge symmetry.

In the past two years extensive work on these models has been done. We refer the reader to [16] and references therein for further information and details.

2.7. large extra dimensions

The large extra dimension scenario [17] is strictly speaking part of the intersecting brane story, but it is worth to point out how this idea fits in the previous discussion.

As we explained, D-branes give rise to extra degrees of freedom, typically gauge fields and matter. Gravity still comes from the closed string sector. Now if we imagine that all gauge fields and matter fields of the standard model originate from a certain set of D-branes, then in some sense we are living on these branes. The only way to probe the compact directions transversal to the branes is by means of gravity. Since gravity has been poorly tested on length scales less than 0.2 mm [18], there is no immediate contradiction with experimental data if we take the compact directions transversal to the brane to be quite large, up to sizes of the order of \( \sim 0.1 \) mm (but see [19]). The standard model interactions are confined to the brane and do not directly see the extra dimensions. This asymmetric treatment of gravity and gauge interactions is the basic idea behind the large extra dimension scenario. The ten dimensional string length can be much larger than the \( 10^{-33} \) mm mentioned before, while the Newton constant in four dimensions remains as small as it is, and therefore these models are sometimes called models with a low string scale. For more discussion, references, and experimental implications, see e.g. [20].

2.8. summary

Though all different possibilities we described look quite different, they are all part of a huge duality web. This is an important lesson. Compactifications with branes, as well as the large extra dimension scenario are not ideas that are completely disjoint from the original heterotic string on a Calabi-Yau manifold. They merely represent different strongly coupled versions of each other. In particular, branes provide a generic ingredient in string theory compactifications.

Nevertheless, the heterotic string on a Calabi-Yau manifold (and the Horava-Witten scenario) are still the most promising. They are the only models that naturally incorporate gauge coupling unification. In the other models gauge coupling unification is more artificial.

A problem shared by virtually all models is the existence of additional low energy degrees of freedom beyond those of a reasonable supersymmetric extension of the standard model. This can perhaps be resolved once the mechanism of supersymmetry breaking is better understood.
Despite this, there are several possible experimental signatures that do not depend very much on the details of the model, and the search for such signatures, both theoretically and experimentally, is clearly an important and urgent problem.

3. GEOMETRIC ENGINEERING AND THEORIES WITH FLUXES

We postponed the discussion of theories with fluxes to this separate section, because there has been quite a lot of recent activity in this area.

First we need to explain the terms “theories with fluxes” and “geometric engineering.” The starting point is the type II string compactified on a Calabi-Yau manifold. Recall that the heterotic string, when compactified on a Calabi-Yau manifold, gave rise to an \( N = 1 \) theory in four dimensions. Type II strings have twice as many supersymmetry as the heterotic string, and therefore yield four dimensional theories with \( N = 2 \) supersymmetry instead.

Geometric engineering, a term which appears to have been introduced in [21], refers to the process of constructing a suitable Calabi-Yau geometry that produces at low energies an a priori given \( N = 2 \) gauge theory in four dimensions. Type II strings have twice as many supersymmetry as the heterotic string, and therefore yield four dimensional theories with \( N = 2 \) supersymmetry instead.

Geometric engineering, a term which appears to have been introduced in [21], refers to the process of constructing a suitable Calabi-Yau geometry that produces at low energies an a priori given \( N = 2 \) gauge theory in four dimensions. It turns out that many gauge theories can be geometrically engineered. Typical gauge theories that appear are of “quiver” or “moose” type. Such theories have gauge groups that are a product of different factors, and matter fields that transform non-trivially under one or two of the gauge groups. Matter fields transforming under three or more of the gauge groups do not appear. Interestingly, this is related to the fact that strings have only two endpoints.

It is possible to break these \( N = 2 \) theories to \( N = 1 \) by adding branes, adding fluxes, or both. The dualities between these possibilities have been explored in great detail by many authors.

We already discussed the possibility of adding branes. Adding fluxes refers to the following [22]. Type II string theory has several massless tensor fields besides the graviton. For example, type IIA string theory also contains a gauge field and a three-index antisymmetric tensor field. This gauge field was crucial in the relation between type IIA string theory and M-theory. Now suppose that the Calabi-Yau manifold looks like, say, a two-sphere. Then we can turn on a non-zero field strength for the gauge field, in such a way that there is a non-zero magnetic flux through the two-sphere. In other words, the integral of the field strength over the two-sphere is non-zero. The gauge field configuration is similar to the gauge fields that surround a magnetically charged particle sitting in the middle of a two-sphere. Of course, here the \( S^2 \) is empty, and in particular there is nothing in the interior of the two-sphere and in particular there no actual magnetically charged particle that sits there.

More generally, the non-trivial geometry of the Calabi-Yau manifold can be used to give simultaneous expectation values to several of the field strengths of the massless tensor fields that appear in type II string theory.

Such fluxes give rise to a non-zero energy density, and this energy will back-react on the geometry. Thus, after the fluxes are turned on, the geometry of the space will no longer be that of a Calabi-Yau space. The corrections to the Calabi-Yau geometry are suppressed by one over the volume, so as long as the size of the Calabi-Yau is sufficiently large this is not a problem.

Type II string theories with fluxes have many interesting properties:

1. They have a rich duality structure. The full mathematical structure of fluxes and branes remains to be uncovered, but it is clear that very advanced mathematics will play an important role, see e.g. [23].

2. The gauge kinetic terms and superpotential of the low-energy effective field theory in four dimensions can often be computed exactly using topological string theory. Topological string theory is a reduced version of ordinary string theory, where many of the degrees of freedom have disappeared. In particular, as the name topological suggest, the metric is no longer a degree of freedom, and in fact there are no local degrees of freedom at all. Perhaps the simplest example
of a topological theory is Chern-Simons theory in 2 + 1 dimensions, which does not depend on any metric in 2 + 1 dimensions, and the only observables of the theory are Wilson lines. Interestingly, certain correlation functions in topological string theory are identical to those in the full string theory, and since topological string theory is so much simpler this has allowed for the explicit computation of a whole class of correlation functions. The gauge kinetic terms and superpotential in four dimensions belong to this class. Besides this important physical application, topological string theory calculations have also yielded many new results in mathematics. For more, see [24].

3. As we stressed in section 2, many string compactifications yield additional massless degrees of freedom at low energies beyond those of a (reasonable supersymmetric extension of the) standard model. In particular, there are often additional massless scalar fields, called moduli, that parameterize the shape and size of the Calabi-Yau manifold. When we compactify the type II string, and also include some fluxes, a superpotential is generated that can explicitly be computed [22]. This superpotential gives rise to mass terms for many of the moduli fields, which can thereby be removed from the low-energy effective field theory. It is much easier to analyze the superpotential in these models than for example in the heterotic string compactified on a Calabi-Yau manifold, and this is a significant advantage.

4. Another generic feature of compactifications with fluxes is the appearance of “warped” compactifications. In usual Kaluza-Klein like compactifications, the full ten-dimensional background is a direct product of a four-dimensional Minkowski space-time and a six-dimensional Calabi-Yau manifold. Four-dimensional Poincaré invariance can be preserved if we allow for a more general “warped” setup, where the metric on Minkowski space is rescaled by an overall factor that depends non-trivially on the coordinates on the Calabi-Yau manifold. For example, suppose that our world is a circle and the total space is a cone. The size of the circle varies depending on where we are on the cone, but rotational invariance remains unbroken. In warped compactifications, there is a relation between the energy scale in four dimensions and the additional coordinates of the Calabi-Yau manifold. Moving along the Calabi-Yau manifold changes the overall factor of the metric on Minkowski space and therefore also the energy. The close relation between extra dimensions and the energy scale in four dimensions also appears in the duality between gauge theories and gravity that we discuss in the next section. See [20,22] for more discussion and literature.

5. The various dualities that have been established in the context of type II strings with branes and/or fluxes have provided topological versions of the gauge theory-gravity correspondence which we discuss in the next section, with interesting mathematical and physical applications [25].

One such recent application is the following [26]: Consider a supersymmetric $N = 1$ gauge theory in four dimensions with classical superpotential $W_0$. Assume that $W_0$ has an isolated critical point where the gauge group is broken to a product of $U(N_i)$ factors. Pure $U(N_i)$ $N = 1$ gauge theories exhibit confinement at low energies which is signaled by a condensation of the gluino condensate field $S_i = \text{tr}(\lambda\lambda)$, where $\lambda$ is the gluino, the superpartner of the $U(N_i)$ gauge field. The gluino condensate in pure $N = 1$ theories is described by an effective quantum superpotential $W_{\text{eff}}(S_i)$ known as the Veneziano-Yankielowicz superpotential [27]. Returning to the case where a classical superpotential $W_0$ breaks $U(N)$ to a product of $U(N_i)$ factors, one may wonder what the analogue of the Veneziano-Yankielowicz superpotential for such a theory is. In [26] it is conjectured that this quantum superpotential can be computed by simply summing planar diagrams in
0+0 dimensional system, namely the matrix theory with action $W_0$. In many examples this can be proven using various string dualities and properties of topological strings. A complete proof directly on the level of field theory has been given in a special case in \cite{28} and in much more generality in \cite{29,30}. There has also been independent progress in all-order field theoretic instanton calculations in $d = 4$, $N = 2$ gauge theories \cite{31}, which may be related.

4. GAUGE THEORY-GRAVITY CORRESPONDENCE

The gauge theory-gravity correspondence refers to an amazing equivalence between certain theories with gravity, and certain theories without gravity. One particular example of such a duality, as originally conjectured by Maldacena \cite{32}, is the exact equivalence between type IIB string theory compactified on $AdS_5 \times S^5$, and four-dimensional $N = 4$ supersymmetric Yang-Mills theory. The abbreviation $AdS_5$ refers to an anti-de Sitter space in five dimensions, $S^5$ refers to a five-dimensional sphere. Anti-de Sitter spaces are maximally symmetric solutions of the Einstein equations with a negative cosmological constant. The large symmetry group of 5d anti-de Sitter space matches precisely with the group of conformal symmetries of the $N = 4$ super Yang-Mills theory, which for a long time has been known to be conformally invariant. In view of this, the gauge theory-gravity correspondence is often referred to as the AdS/CFT duality, where CFT stands for conformal field theory.

Anti-de Sitter space can be roughly thought of as a product of four-dimensional Minkowski space times an extra radial coordinate. The metric on Minkowski space is however multiplied by an exponential function of the radial coordinate, and Anti-de Sitter space is therefore an example of a warped space: in a suitable local coordinate system, $ds^2 = dr^2 + e^{2r}(\eta_{\mu\nu}dx^\mu dx^\nu)$. The limit where the radial coordinate goes to infinity and the exponential factor blows up is called the boundary of Anti-de Sitter space. This boundary is the place where the dual field theory lives. One can indeed verify that string theory excitations in anti-de Sitter space extend all the way to the boundary. In this way one obtains a map from string theory states to states in the field theory living on the boundary.

Is is very hard to directly prove the equivalence between type IIB string theory on $AdS_5 \times S^5$, and four-dimensional $N = 4$ super Yang-Mills theory. For one, we do not have a good definition of non-perturbative type IIB string theory. Even at string tree level, we do not (yet) know how to solve the theory completely. From this perspective, it is perhaps better to view $N = 4$ super Yang-Mills theory as the definition of non-perturbative type IIB string theory on the $AdS_5 \times S^5$ background.

A weaker form of the AdS/CFT correspondence is obtained by restricting to low-energies on the string theory side. At low-energies, type IIB string theory on $AdS_5 \times S^5$ reduces to type IIB supergravity on $AdS_5 \times S^5$. The corresponding limit on the gauge theory side is one where both $N$ and $g^2_{YM}N$ become large, where $N$ is the rank of the $U(N)$ gauge group of the $N = 4$ supersymmetric gauge theory (not to be confused with the $N$ appearing in $N = 4$), and $g^2_{YM}$ is the gauge coupling constant. The equivalence between type IIB supergravity on $AdS_5 \times S^5$ and $N = 4$ gauge theory in the large $N$ limit has been very well tested by now.

The AdS/CFT correspondence is related to two deep ideas in physics.

The first of these is the idea that large $N$ gauge theory is equivalent to a string theory \cite{33}. The perturbative expansion of a large $N$ gauge theory in $1/N$ and $g^2_{YM}N$ has the form of a string loop expansion, with the string coupling $g_s$ equal to $1/N$. Through some peculiar and not completely understood mechanism, Feynman diagrams of the gauge theory are turned into surfaces that represent interacting strings (but see \cite{34}). Apparently, this is precisely what happens in the AdS/CFT correspondence.

The second is the idea of holography \cite{35,36}. This idea has its origin in the study of the thermodynamics of black holes. It was shown by Bekenstein and Hawking \cite{37} that black holes can be viewed as thermodynamical systems with a temperature and an entropy. The temperature is di-
rectly related to the black body radiation emitted by the black hole, whereas the entropy is given by $S = A/4G$, with $G$ the Newton constant and $A$ the area of the horizon of the black hole. With these definitions, Einstein’s equations of general relativity are consistent with the laws of thermodynamics. Since in statistical physics entropy is a measure for the number of degrees of freedom of a theory, it is rather surprising to see that the entropy of a black hole is proportional to the area of the horizon. If gravity would behave like a local field theory, one would have expected an entropy proportional to the volume. A consistent picture is reached if gravity in $d$ dimensions is somehow equivalent to a local field theory in $d - 1$ dimensions instead. Both have an entropy proportional to the area in $d$ dimensions, which is the same as the volume in $d - 1$ dimensions. The analogy of this situation to that of an hologram, which stores all information of a 3d image in a 2d picture, led to the name holography. The AdS/CFT correspondence is holographic, because it states that quantum gravity in five dimensions (forgetting the compact five sphere) is equivalent to a local field theory in four dimensions.

One of the questions that the AdS/CFT immediately raises is that of the interpretation of the extra fifth dimension in the field theory. It turns out that it is closely related to the energy scale. From the 5d gravitational point of view, low-energy processes in field theory stay close to the boundary of AdS, whereas high-energy processes penetrate deeper in the interior. One can even show that the invariance under 5d general coordinate transformations implies the Callan-Symanzik renormalization group equations in the field theory. Thus, from the 5d point of view, the renormalization group is on an equal footing with 4d Poincaré invariance.

The AdS/CFT correspondence is also an example of a weak/strong coupling duality. Depending on the choice of parameters, either AdS or the CFT is a weakly coupled description of the system, but never both at the same time. Gauge theory is a good description for small $g^2_{YM}$ and small $g^2_{YM}$, whereas string theory is good for large $g^2_{YM}$ and small $g^2_{YM}$. Therefore, the AdS/CFT correspondence can be applied in two directions.

We can use string theory to learn about gauge theory, and we can use gauge theory to learn about string theory.

One of the most difficult and unsolved problems in the AdS/CFT correspondence is to reconstruct 5d local gravitational physics directly from the dual 4d field theory point of view. In particular, we would like to know in what way the local gravitational description breaks down. Does such a breakdown occur in a local way at the Planck length, or in a non-local way at much larger length scales? The AdS/CFT correspondence seems to prefer the second answer, which is also the answer that may provide a resolution to the black hole information paradox. This paradox is based on the fact that semiclassically, everything that falls into a black hole is converted into purely thermal radiation, with no memory of the object that fell in except for its mass and perhaps a few other quantum numbers. Such a process contradicts the usual rules of quantum mechanics, and we can either give up on quantum mechanics or give up on the semiclassical approximation to quantum gravity; AdS/CFT prefers the latter.

The emergence of a concrete duality between a theory with gravity and a theory without gravity is one of the most important results of string theory. Below, we summarize some of the developments in this area over the past two years. For more, see the reviews.

### 4.1. high energy scattering/deep inelastic scattering

At first sight, the AdS/CFT correspondence, or any duality between string theory and gauge theory, seems at odds with the known fact that the scattering of glueballs at high energies is hard, whereas string scattering at high energies is soft, due to their extended nature. The resolution sits in the fact that AdS is a warped space. When an object moves away from the boundary of AdS, its size is exponentially reduced. Very high energy processes in the gauge theory are described by strings which propagate a long distance from the boundary of AdS before they interact. By that time, the size of the strings has been exponentially reduced, and this compensates for the soft-
ness of string interactions to make it into a hard process in the gauge theory \[41\]. Besides such effects, which are due to the geometric warping of AdS, other gauge theory processes crucially involve strong gravity physics like black hole formation \[42\]. It is also possible to study the physics of deep inelastic scattering and the parton model from the AdS/CFT point of view \[43\].

4.2. towards a QCD string?

A more involved version of the AdS/CFT correspondence is the one given in \[44\]. It was discovered by studying branes stuck in singularities in string theory. The gauge theory that appears is an \(N = 1\) theory in four dimensions with gauge group \(SU(N) \times SU(N + M)\). There are two chiral superfields \(A_i\) in the \((N, \overline{N} + M)\) representation of the gauge group, and two chiral superfields \(B_i\) in the \((\overline{N}, N + M)\) representation. In addition, there is a nontrivial superpotential of the form \(W \sim \epsilon^{ij} \epsilon^{kl} \text{tr}(A_i B_k A_j B_l)\).

This field theory has a remarkable property: it has running gauge couplings, but does not become free at either low or high energies. The gauge coupling becomes strong either way. Strongly coupled \(N = 1\) theories in four dimensions often admit a dual weakly coupled description, a duality known as Seiberg duality \[45\]. The same is true here: both at low energies and at high energies there exist dual descriptions. However, these dual descriptions have the same problem: they are not weakly coupled at either low or high energies. Again, they admit suitable dual descriptions. The full picture that emerges is that of an infinite "cascade" of gauge theories, that continues indefinitely at high energies, with an ever increasing rank of the gauge group, but terminates at low energies once e.g. the rank of one of the gauge groups becomes one. At that point, the gauge theory becomes confining. Strictly speaking we need an infinite amount of fine tuning of irrelevant operators to obtain this infinite cascade, but quite remarkable, the dual description of this gauge theory quite naturally sees the same cascade. The ranks \(N\) and \(M\) of the gauge group become non-trivial functions of the radial coordinate of the dual AdS-like geometry. This also confirms once more the interpretation of the extra fifth dimension in the AdS/CFT correspondence as an energy scale in the field theory.

The AdS-like geometry that is dual to this infinite cascade has several nice features. String theory on this background exhibits (i) confinement, (ii) glueballs and baryons with a mass scale that emerges through dimensional transmutation, exactly as in the gauge theory, (iii) gluino condensates that break the \(\mathbb{Z}_{2M}\) chiral symmetry to \(\mathbb{Z}_2\), and (iv) domain walls separating different vacua.

The gauge theory at low energies reduces to a pure \(N = 1\) supersymmetric Yang-Mills theory. Does the dual geometry therefore provide a dual string theory for pure supersymmetric Yang-Mills theory, the long sought for QCD string? Not really, because it has new degrees of freedom beyond those of the field theory that appear at \(\Lambda_{QCD}\). This is a generic problem in trying to find weakly coupled string theory descriptions of gauge theories. To decouple the additional degrees of freedom, we need to make the curvature of the AdS-like geometry large, while keeping the string coupling \(g_s\) small. String theory in a strongly curved background is described by a strongly coupled \(1 + 1\) dimensional field theory. The structure of the sigma models relevant for the AdS/CFT correspondence is not very well understood, but there has been progress in this direction recently (see \[46\] and references therein), and the prospect of finding a string theory dual of QCD remains an exciting possibility.

4.3. other string effects in gauge theories: large quantum numbers and pp-waves

Instead of trying to find a precise string theory dual description of pure \(N = 1\) supersymmetric Yang-Mills theory, it is also interesting to look for more qualitative stringy behavior in gauge theories.

One place to find such behavior is to look at states with a large scaling dimension proportional to \(N\), the rank of the gauge group. Many gauge theories have baryons with such scaling dimensions, and it turns out that they are not described by strings but by branes in the dual geometrical description \[47\]. Thus, it is also possible to discover branes in gauge theory.

A related example is to consider operators with
a large spin $s$, like for example $\text{tr}(\Phi D_{\mu_1} \ldots D_{\mu_s} \Phi)$, where $\Phi$ is some field that transforms in the adjoint representation of $U(N)$. Such operators correspond to folded rotating closed strings in the dual geometrical description. One can compute the scaling dimension of these operators both in the field theory and in the dual geometrical description. This confirms the equivalence between the two, as one finds in both cases that it behaves like $s + \log s$ [45].

A more complete way to recover string theory from a gauge theory has been described in [49]. The idea is to take a particular scaling limit of the AdS/CFT correspondence. This scaling limit, when applied to the AdS geometry, yields a different geometry known as a “pp-wave”. In fact, many geometries admit scaling limits in which they reduce to pp-waves, as originally shown by Penrose [50]. String theory on the pp-wave, in the absence of string interactions, can be exactly solved, and in particular the free string spectrum can be obtained.

On the field theory side, the same limit can be taken. In this limit only a subset of the operators of the full $N = 4$ super Yang-Mills theory survive, namely those for which the scaling dimension $\Delta$ and a certain global $U(1)$ quantum number $J$ have the property that $\Delta + J$ scales as $N^{1/2}$, while $\Delta - J$ is kept finite, as one takes $N \to \infty$.

Interestingly, this set of operators is in one-to-one correspondence with the set of free string states. This is the first example where a complete string spectrum has been obtained from a gauge theory, albeit in a special scaling limit.

The ground state of the string theory (in lightcone quantization) is described in the gauge theory by the operator $\text{tr}(Z^J)$, where $Z$ is a complex scalar field in the adjoint representation of the gauge group with charge $J = +1$ under the distinguished global $U(1)$ symmetry. The simplest excited states of the string are operators of the form $\sum_i a_i \text{tr}(Z^i \Phi Z^{J-i})$ where the $a_i$ are phases. The string appears from this point of view as composed of “string bits.” The string bits are the operators $Z_i$, and the string is composed of a string of $J$ bits. Exciting the string amounts to introducing impurities like $\Phi$ that are distributed with phases (i.e. a discrete momentum) along the chain of $Z_i$.

A similar discretized picture of string theory can be obtained in several other gauge theories as well, see e.g. [51]. It is also presently being investigated whether one can correctly recover string interactions, or even the full string field theory, from the gauge theory [52].

5. TIME DEPENDENT BACKGROUNDS AND STRING THEORY

The celebrated type IA supernovae measurements [53] of a few years ago that showed the existence of a small but nonzero positive cosmological constant are partly responsible for a renewed interest in time-dependent backgrounds in string theory. It is much more difficult to obtain a small nonzero cosmological constant than a cosmological constant that is strictly zero. The latter can arise for example due to an underlying symmetry. A small positive cosmological constant on the other hand introduces a new scale in the theory, and this has to be put in by hand or it should be explained by some unknown mechanism in terms of the existing length scales. Although various mechanisms have been proposed in the literature, no completely satisfactory explanation has been given. Therefore, most attention has recently been focused on trying to understand time-dependent backgrounds and cosmology at a more conceptual level. In particular, many cosmological scenario’s involving branes have been proposed (see e.g. [54]). Though popular in the media, their status is mainly phenomenological, as it is often difficult to embed them in string theory. One of the reasons is that string theory does not allow any freedom in the choice of brane tensions and/or interbrane interactions. In addition, in string theory such models often lead to singularities that are hard to study.

In the remainder, we discuss some attempts to obtain and/or understand time dependent backgrounds in string theory.

5.1. time-dependent orbifolds

One of the nice features of string theory is that it can deal with certain types of singulari-
ties. In particular, if one starts with a smooth string theory background and then makes some discrete identifications, a procedure known as orbifolding, the resulting singularities are well-understood and usually completely under control. This leads to the question whether there any orbifold constructions in string theory that provide a good toy model of cosmological singularities, and/or of time-dependent backgrounds. Many examples of such orbifolds have been studied recently [55].

One of the simplest orbifolds one can think of is made by starting with flat space. This has Poincaré invariance, so one can try to make discrete identifications with respect to some finite subgroup of the Poincaré group. For example, one can make an identification under a translation in a given direction, in which case that direction becomes compact and turns into a circle. It is more interesting to consider orbifolds that also involve the time direction in some non-trivial way, so that the orbifold theory becomes time-dependent. To some extent, the physics of such time-dependent orbifolds can be extracted from the ambient theory in flat space. However, the lack of invariance under time translations, the non-existence of a good Wick rotation (and therefore of a $+i \epsilon$ prescription in propagators), and the absence of a Hamiltonian that is bounded from below make the interpretation of these time dependent orbifolds rather confusing. In addition, in case the group of identifications is not finite, it has been argued that the resulting orbifold backgrounds are generically unstable [56]. This happens roughly because a single particle in the orbifolded background corresponds to infinitely many particles in the unorbifolded space, that are all mapped into each other under the discrete identifications. This infinite set of particles in the unorbifolded space tend to form a black hole. To evade this, one needs a large number of non-compact transversal directions, which makes the models much less realistic.

5.2. S-branes

The D-branes we discussed in section 2.5 can move in time, and are often stationary. One may wonder whether it is possible to make extended objects that are localized in time, and use these to generate interesting time-dependent backgrounds in string theory. Such objects are not the same as instantons. Though instantons are thought of as objects localized in time, they are solutions of the Euclidean theory, not of the Minkowski theory, and here we want solutions of the Minkowski equations of motion. Since all directions along such localized objects are space-like, they have been called S-branes. Some interesting recent work in this direction has been done [57], but at present, no sufficiently stable S-branes have been found.

5.3. more speculative ideas

As we review below, it is difficult to obtain a well-controlled solution of string theory with a positive cosmological constant. This has led to the suggestion that perhaps we should modify string theory in a more drastic way. For example, perhaps we should start to think about non-local string theories. Such theories have been discussed in [58] as a useful framework to discuss particle creation in string theory. Another idea is to make sense of a version of string theory where some of the fields are allowed to take imaginary values [59,60]. Such theories can much more easily accommodate solutions with a positive cosmological constant than ordinary string theories.

5.4. de Sitter space

Recall from section 4 that Anti-de Sitter space played a crucial role in the AdS/CFT duality. It was a maximally symmetric solution of the Einstein equations with a negative cosmological constant. Similarly, there exists something called de Sitter space, which is a maximally symmetric solution of the Einstein equations with a positive cosmological constant. The metric of de Sitter space is of the form $ds^2 = -dt^2 + \cosh^2 \frac{t}{a} d\Omega_3$, in other words it describes a three-sphere that has its minimum size at $t = 0$ and expands exponentially in the future and in the past.

During inflation the universe expanded exponentially, and was approximately described by a de Sitter space. At this very moment the universe again appears to be entering a de-Sitter phase, driven by the small but nonzero cosmological con-
stant that has been observed. It is therefore natural to wonder whether there is a solution of string theory that somehow involves de Sitter space, or more generally, any solution of the Einstein equations with a positive cosmological constant.

This is remarkably difficult to achieve\(^3\). Solutions with a negative cosmological constant are easily generated by turning on field strengths for some of the tensor fields of string theory, but this never leads to a positive cosmological constant. There are no-go theorems\[^{62}\] that state that there is no smooth solution of supergravity that involves de Sitter space and a compact internal space. These no-go theorems do not necessarily apply to string theory to which supergravity is only a low-energy approximation. Indeed, the no-go theorems assume a certain positive energy condition which is violated in string theory: string theory has negative tension brane-like objects (so-called orientifolds), which could be crucial in obtaining solutions with a positive cosmological constant; see\[^{63}\] for a recent attempt involving non-critical strings\(^4\). Unfortunately, a well-controlled solution of string theory with a positive cosmological constant remains out of reach for now.

Still, if one assumes that de Sitter solutions of string theory exist, one may ask whether they would admit a dual field theory description like we have in the case of the AdS/CFT correspondence. Such a dual description would be a very powerful tool in analyzing the physics of de Sitter space. Preliminary results in this direction indicate that if anything, de Sitter space should be dual to rather peculiar conformal field theories\[^{64,65}\]. This leads to an interesting picture. In the case of the AdS/CFT correspondence, the extra radial dimension of AdS had the interpretation of an energy scale in the dual field theory. In de Sitter space, the role of the radial coordinate of AdS is taken over by time, and time should therefore somehow be related to an energy scale in the dual field theory. This leads to the speculation that the transition in the universe from the inflationary de Sitter phase to the present de Sitter phase is described by some sort of renormalization group flow in the putative dual field theory\[^{66}\]. If true, this would provide a completely new way to think about the time evolution of the universe.

5.5. the inflaton as a tachyon

Tachyons are particles with negative mass squared, and as free particles they are unphysical. In interacting theories they do not have to be unphysical at all. They often simply indicate that one is not expanding around the true vacuum of the theory, which is e.g. what happens in the standard model when expanding around zero Higgs field. If we start with a theory not in its vacuum state, it will undergo some dynamical process. For instance, it can decay to its ground state, while radiating all energy in the initial state away to infinity.

In string theory, we frequently find tachyons in the spectrum. Probably the best known example is the purely bosonic string that lives in 26 dimensions. It has a tachyon in its spectrum and is therefore often discarded as being inconsistent. Of course, it is possible that the bosonic string has some runaway potential for the tachyon with no minimum at all; but it is equally well possible that the bosonic string does have a stable ground state, it is simply beyond our present capabilities to determine what such a ground state should be.

Luckily, there are many string theory setups with tachyons where the tachyon is under a reasonable amount of control. A good example is a system consisting of D-branes and anti D-branes. Anti D-branes are extended objects with the opposite quantum numbers compared to D-branes, and anti D-branes and D-branes can annihilate each other. This instability is reflected in string theory by the presence of a tachyon, and has been studied in great detail\[^{67}\].

It is an interesting question whether we can find time dependent solutions in string theory by creating an unstable initial state, and letting it evolve in time. The time dependence is then generated by the dynamics of the tachyon degrees of freedom. Some time dependent solutions describing a tachyon rolling down a potential have been

\(^3\)Supergravity solutions of this type are e.g. discussed in\[^{59,61}\].

\(^4\)Other attempts to obtain interesting time-dependent backgrounds from non-critical string theory can for example be found in\[^{68}\].
found in [68]. The endpoint of the process is a gas of excited string states. The dynamics of such processes can be captured by a simple effective field theory with action $S \sim \int e^{-aT} \sqrt{1 + (\partial_\mu T)^2}$ [69], where $T$ is the tachyon degree of freedom. It is, however, hard to get realistic models in this way, since there are no naturally small parameters in the theory and the slow-roll conditions of inflation will typically not be satisfied. In addition, the nature of the gas of string states that remains was calculated in an approximation that will break down sooner or later, because the closed string states that can also be radiated away have not been precisely taken into account [70].

Alternatively, one may try to use the tachyon to describe the reheating of the universe after inflation has ended. This is however harder to make explicit in string theory. For some further discussions, see [71,72].

5.6. string theory signatures in cosmology

Is there a possibility to see signatures of string theory in cosmology? As we go back in time, there is a moment in the history of the universe when the temperature of the universe was so high that classical gravity is no longer a good approximation, and in order to describe the universe before this time a theory of quantum gravity is needed. In particular such a theory of quantum gravity should explain the initial conditions for the various fields in the universe. These initial conditions are reflected in the spectrum of the cosmic microwave background radiation. Since we have no precise string theory description of the early universe, we can only estimate the magnitude string theory effects would have on the cosmic microwave background. Naively, string theory effects will affect the power spectrum by terms of order $(H/M)^2$, where $H$ is the Hubble constant and $M$ the scale of new physics. However, this assumes a standard choice of vacuum state for the fields. If we drop this assumption the effects can be larger, of the order of $(H/M)$. With an optimistic choice of $M$, this could be at the treshold of observational limits. It would be an absolutely tremendous achievement if an experimental signature of string theory could be obtained in this way! It will nevertheless be difficult to disentangle any such effects from the data, probably the spectrum of power fluctuations will not be enough and the spectrum of tensor fluctuations will also be needed. Besides this, there is still a lively debate going on whether non-standard choices of vacuum states are physically acceptable or not; see e.g. [73] for further discussion of this issue.

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