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A Superstring Theory Underview*

LANCE J. DIXON†

Stanford Linear Accelerator Center
Stanford University, Stanford, California 94309

ABSTRACT

I give a brief survey of the current status of superstring phenomenology, with an emphasis on the (currently unrealized) possibility of obtaining model-independent results.

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† On leave from Joseph Henry Laboratory, Princeton University, Princeton, NJ 08544 until Sept. 1, 1989.

1. INTRODUCTION

This talk is meant to briefly and qualitatively summarize the current status of superstring phenomenology, and to mention what I think are the outstanding problems that must be solved before we can hope to get experimental predictions from string theory. The topic of this SLAC Summer Institute is "Physics at the 100 GeV Mass Scale", the region accessible to current accelerators. On the other hand, the characteristic mass scale for superstrings is the Planck mass $M_{\rm Pl} \simeq 10^{19}$ GeV/c^2. The Planck scale appears because strings provide a description of quantum gravity, and gravity is characterized by Newton's constant, which has mass dimension -2: $G_N = c\hbar/M_{\rm Pl}^2$. It will be useful to keep in mind throughout this talk the gap of 17 orders of magnitude in energy, because any predictions from superstrings will depend on analysis of effects over the entire range.

Recent work on superstrings has concentrated on answering questions in the following three areas:

(1) Fundamental understanding: Is string theory based on some new physical principle, in the same way that Yang-Mills theory, for example, is based on local gauge invariance?

(2) Formalism: Can we extend the calculational tools that have been developed for string theory to include non-perturbative effects?

(3) Phenomenology: Can we hope to make definite predictions of new (and old) phenomena at experimentally accessible energies?

I will focus on the third area in this talk (though I will digress to cover a little of the relevant formalism as well). Far below the Planck scale, strings appear to be pointlike, and string theory reduces to some *effective field theory.* At energies around 100 GeV this field theory should in turn reduce to approximately the standard model, if string theory is to describe experimental reality. At these energies, any predictions of string theory are therefore indirect, and can be summarized as predictions of the 18 or so input parameters of the standard model (or more accurately, *explanations* of those parameters that have already been measured), plus predictions of possible field theory extensions of the standard model and the parameters therein. Currently, string predictions also involve the rather arbitrary selection of a candidate string *vacuum*, or ground state, from a large set of degenerate vacua. Each such choice leads to a different effective field theory, and hence to a different model for physics at 100 GeV. Here I will describe the set of string vacua, and the prospects for obtaining "model-independent" predictions from string theory, and also (very briefly) summarize a few recent attempts to construct specific quasi-realistic‡ models.

The three general areas of research are all related. In particular, a better fundamental understanding of string theory should lead to advances in the formalism, which should in turn lead to more concrete phenomenological predictions than at present.

Superstrings[1] generated much theoretical excitement in 1984, when it was realized that they had the potential for unifying all of the known gauge forces together with gravity, and for doing so in a practically unique way. Uniqueness is of particular importance due to the large discrepancy in energy scales between strings and current experiments: Since the signatures of string theory are so indirect, it would be nice if they were quite definite. At present, we still have the same limited number of string theories as were known in 1984: *bosonic strings*,[2] *superstrings*,[3] and *heterotic strings*.[4] The bosonic and superstring theories can contain either closed strings only, or else both closed and open strings; the strings can be either orientable (have an arrow along them specifying a direction) or non-orientable (no such arrow). The heterotic string theory is a hybrid of the bosonic and superstring theories, and as a consequence it contains only closed, orientable strings. Of these theories, only the heterotic string has any phenomenological promise, at least at our present level of understanding of string theory. A possible exception to this statement is the closed+open superstring with gauge group $SO(32)$;[5] this string

‡ A very fashionable adjective in string phenomenology.

theory has received much less attention than the heterotic string (due partly to a lack of any *obvious* phenomenological promise, and partly to a relative lack of open-string formalism), and I too will slight it here.

2. Vacuum Proliferation, Particle Spectrum, and Effective Lagrangians

So we still have in 1989 only one phenomenologically viable string theory, the heterotic string (with the possible exception of the $SO(32)$ superstring). However, there has been a recent proliferation of string *vacuum states*, which appear to be degenerate in energy to all orders in (string) perturbation theory. This non-uniqueness of the vacuum can play havoc with attempts to extract definite predictions from string theory, so I will spend some time describing it.

The vacuum proliferation has come through the realization that for any two-dimensional (2d) *conformal field theory*[6] (CFT) one can construct a vacuum state for classical string theory. I will describe the relation between CFT and string vacua in more detail later. For now, we just need to know that, in any particular string vacuum state, *the full spectrum of elementary particle masses and couplings are completely specified by the CFT for that vacuum.*

Elementary particles in string theory are identified as the different rotational and vibrational eigenmodes of a string as it oscillates around a vacuum state. The $(\text{mass})^2$ of the particle is proportional to the frequency of vibration, and in the first approximation it is an integer multiple of some basic mass parameter M^2 of order M_{Pl}^2. A typical spectrum is shown in Figure 1(a). Each of the mass levels shown in the figure is highly degenerate. The lowest level is of the most interest to us because it contains all the massless particles, where "massless" means having mass much less than M_{Pl}. In particular, the massless level should contain all the particles that have been observed to date: the leptons, quarks and gauge bosons of the standard model. There must also be a "hyper-hyperfine" splitting of the mass degeneracy in Figure 1(a), as shown in Figure 1(b), in order to generate the observed particle masses of order M_W. The Higgs mechanism is usually invoked to explain this splitting, but a major task of string phenomenology is to explain why M_W is so incredibly small relative to $M_{\rm Pl}$; the latter question is essentially the familiar hierarchy problem. The remaining levels in the string spectrum contain particles with Planck-scale masses, which won't be produced in accelerators in the near future, and so they are of less direct interest. (If any of the massive particles are stable, however, they might have survived as relics of the Big Bang.)

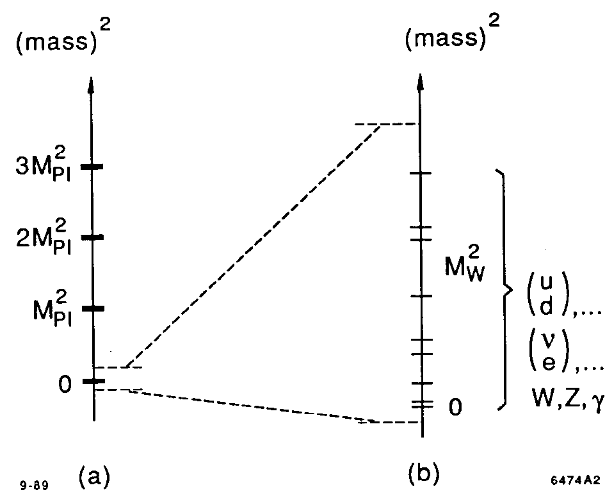

Figure 1. (a) A typical mass spectrum for a particular string vacuum state. (b) An enlarged view of the massless sector, showing the "hyper-hyperfine" splitting.

The couplings between particles in a given vacuum are found by studying how strings scatter. A four-string scattering process at tree level is shown in Figure 2(a), and at one loop (the next order in string perturbation theory) in Figure 3(a). For this process to describe the scattering of four massless particles, the four external strings should be prepared in the corresponding rotational/vibrational states. Remarkably, there is only one string scattering diagram for a particular process, at a given order in perturbation theory. For instance, the tree-level diagram in Figure 2(a) can be viewed as representing either s-, t-, or u-channel scattering, by stretching it either horizontally, vertically, or out of the page. (This property of string scattering amplitudes is called *duality*, and is the reason why string theories were termed "dual models" early in their history.) The one diagram summarizes the contributions of an infinite number of particles — both massless and massive — that can appear as intermediate states in the scattering, and in any of the channels. The individual particle contributions are represented by ordinary field theory Feynman diagrams, shown schematically in Figures 2(b) and 3(b). If the external states in a string scattering process are massless particles, and if the energy of the collision is much less than M_{Pl} (the usual case experimentally!), then the amplitude for the process can be reproduced using an *effective Lagrangian* $\mathcal{L}_{\mathrm{eff}}$ which only involves the massless fields.[7] For example, to reproduce the amplitudes in Figure 2, one needs three-particle couplings of the type shown in Figure 4(a). In addition, an infinite set of *non-renormalizable terms* (terms in $\mathcal{L}_{\mathrm{eff}}$ with dimension larger than four, whose coefficients contain inverse powers of M_{Pl}) results from exchanges of massive particles (Figure 4(b)). The latter terms are completely analogous to the four-Fermi interaction terms that reproduce the low-energy effects of W-exchange in the standard model. Similar considerations apply to loop- as well as tree-level string scattering amplitudes, as represented in Figure 3.

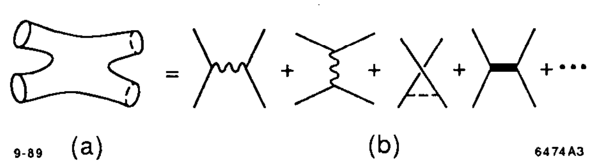

Figure 2. (a) A four-string scattering process at tree level. (b) The field theory Feynman diagrams that represent the contributions of individual particles to the amplitude. Thick lines denote massive particles; all other lines denote massless particles.

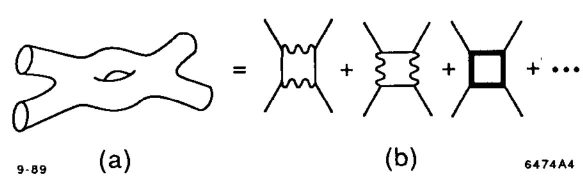

Figure 3. (a) A four-string scattering process at the one-loop level. (b) The Feynman diagrams contributing to it.

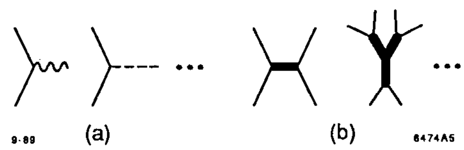

Figure 4. (a) Three-particle couplings needed to reproduce the four-string scattering amplitude at tree level. (b) Additional non-renormalizable interactions that are needed, due to exchanges of massive particles in the string amplitude.

To recap, if one chooses a particular string vacuum state (a CFT), then the string dynamics are fixed, and they can be used to obtain an effective Lagrangian $\mathcal{L}_{\rm eff}$ for the massless particles in that vacuum. The effective Lagrangian describes physics at energies just below the Planck scale, where strings start to appear point-like, in terms of a conventional field theory (albeit including non-renormalizable terms). To describe physics at the electroweak scale, one then "just" renormalizes $\mathcal{L}_{\rm eff}$ down to 100 GeV or so. All the non-renormalizable terms in $\mathcal{L}_{\rm eff}$ can (in principle) be "renormalized" using the loop-corrected string amplitudes, because the latter are actually finite in the ultra-violet. The renormalization task is complicated by the presence of many light fields and many possible intermediate stages of symmetry breaking. Finally one compares $\mathcal{L}_{\rm eff}|_{100\ {\rm GeV}}$ with the Lagrangian of the standard model, $\mathcal{L}_{\rm s.m.}$. Usually they won't agree in sufficient detail for $\mathcal{L}_{\rm eff}|_{100\ {\rm GeV}}$ to be considered realistic; in this case one can go back and pick another vacuum state, and go through the whole procedure again...!

Of course, in this *constructive* approach to string phenomenology, the results obtained may be highly model-dependent. Thus one should insist that a model give correct "postdictions" of old phenomena (namely, the host of standard model parameters that have already been measured) before taking seriously its predictions of new phenomena. I think it is fair to say that no model constructed to date satisfies this criterion. It is important to supplement the constructive approach with a *model-independent* approach, one that tries to determine the general low-energy properties common to *all* string vacua. In this way one may be able to test string theory rather than just testing specific string vacua. Unfortunately, there has been relatively little progress to date in the model-independent approach. Interplay between both approaches seems necessary in order to get the most complete understanding of what low-energy physics can be expected from strings.

3. The String Vacuum Landscape and Vacuum Cleaning

It would be nice to have a picture of what the space of string vacua looks like, in particular how different vacua are related to each other, in order to help understand how one of the vacua might be selected over the others (presumably by some non-perturbative string dynamics). Unfortunately our current picture of the string vacuum landscape is quite crude, and is only a local picture. That is, we understand the neighborhood of any particular vacuum reasonably well, through the effective Lagrangian described above, but we really have no idea of where the other vacua are in relation to it. The landscape near a vacuum state can be described roughly by plotting the effective potential $V_{\rm eff}(\phi_i)$ for the massless scalar fields (particles) ϕ_i in that vacuum. (Note that $V_{\rm eff}(\phi_i)$ is just a piece of the effective Lagrangian $\mathcal{L}_{\rm eff}(\phi_i, \partial_\mu\phi_i, A_\mu, \ldots)$.) As an example, suppose there are only two massless scalars (usually there are many more), ϕ_1 and ϕ_2, with

$$V_{\rm eff}(\phi_1, \phi_2) = \lambda_1 \phi_2^4 + \lambda_2 M_{\rm Pl}^{-2} \phi_1^2 \phi_2^4 + \text{ higher order terms},$$

such that $V_{\rm eff}$ vanishes identically whenever $\phi_2 = 0$. (See Figure 5.) Then any vacuum expectation value of ϕ_1 minimizes $V_{\rm eff}$ and provides a string vacuum state, so $\langle\phi_1\rangle$ parametrizes a line of vacua. Fields like ϕ_1 are called *moduli*, and occur frequently in string vacua.[8,9,10]

In Figure 6 I have embedded this example along with some others into what is supposed to be a more complete picture of the vacuum landscape. There is a great variety among the various vacua: they may feature different gauge groups, different numbers of moduli and/or massless fields, even different numbers of space-time dimensions. (We'll focus on those with four space-time dimensions!) The question marks in the figure reflect our almost total lack of knowledge about regions in the

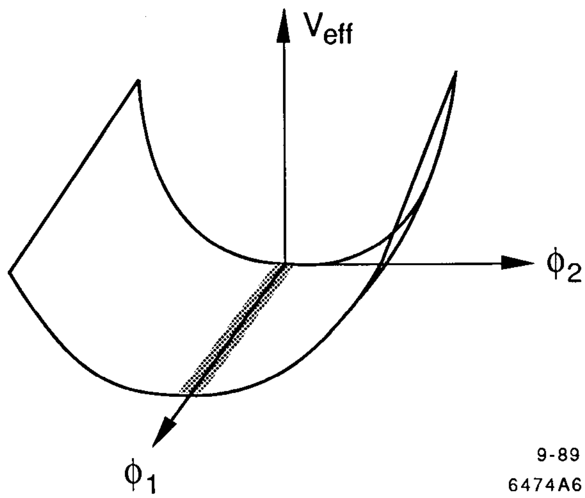

Figure 5. Plot of a simple effective potential, $V_{\rm eff}(\phi_1, \phi_2) = (\lambda_1 + \lambda_2(\phi_1/M_{\rm Pl})^2)\phi_2^4 + \ldots$. The line of vacua, parametrized by $\langle\phi_1\rangle$ and having $\langle\phi_2\rangle = 0$, is shaded.

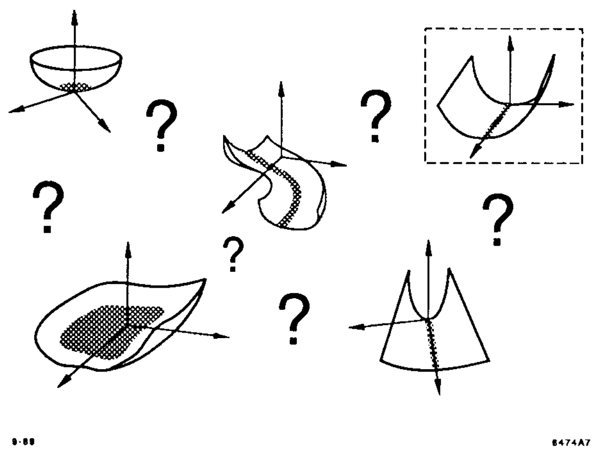

Figure 6. A rough sketch of the string vacuum landscape. The effective potential $V_{\rm eff}$ is plotted on the vertical axes, as a function of the massless fields in a given vacuum. The example of Figure 5 is in the dashed box. Again the vacua are shaded in. The question marks denote regions far from any vacua, which are not understood at all. (Here there be dragons!)

space that are not very close to any vacua. We don't even have a set of global coordinates with which to describe the space. In the neighborhood of the first example we may use the fields ϕ_1 and ϕ_2, but another vacuum (if not connected to the first one) will have another set of massless fields, say $\widetilde{\phi}_i$, which bear no obvious relation to the first set ϕ_i. Similarly the effective Lagrangian describing the low-energy physics of the first vacuum, $\mathcal{L}_{\rm eff}(\phi_i)$, has no obvious relation to that for the other vacuum, $\widetilde{\mathcal{L}}_{\rm eff}(\widetilde{\phi}_i)$. The massive fields have been omitted from the picture, because including them would make the space infinite-dimensional; nevertheless, they will certainly play a role in our future understanding of the relation between the disconnected vacua, *i.e.* in filling in the question marks in the figure.

At present, there are no dynamical criteria for preferring any one vacuum over another. All the vacua (to be more precise, all the vacua with unbroken space-time supersymmetry) remain stable vacua to all orders in string perturbation theory.[11] One generally hopes that non-perturbative effects will lift this vacuum degeneracy, but without a formalism for calculating such effects, the best one can do at present is to apply phenomenological criteria to do the "vacuum cleaning". Here is a rather minimal set of criteria, which are also relatively easy to implement (or at least check) in a given string vacuum:

- Four-dimensional space-time.
- N=1 space-time supersymmetry at the Planck scale.
- A gauge group containing $SU(3) \times SU(2) \times U(1)$.
- Massless particles with the gauge quantum numbers of the standard model (quarks, leptons, *etc.*).

It should be noted that even these criteria are not completely free of theoretical prejudice. Four-dimensional space-time is on a pretty safe footing, but unbroken space-time supersymmetry at the Planck scale is put in so that it may play a role in explaining the naturalness/hierarchy problem of why $M_W/M_{\rm Pl}$ is so small.[12] (Vacua without space-time supersymmetry are also difficult to analyze because, unlike the supersymmetric ones, they *are* generally de-stabilized by radiative corrections.) Extended ($N > 1$) supersymmetry is excluded because it forces a non-chiral theory, *i.e.* one with no parity-violating gauge interactions[13]; it is assumed that parity is not spontaneously broken at some lower energy scale. The last two criteria assume that the observed gauge bosons and fermions are fundamental string excitations, rather than (say) composites of such excitations. There are of course many more criteria that could be applied, but they are generally much more difficult to implement. For example:

- No fast proton decay.
- The correct set of quark and lepton masses.

And so on.

4. Conformal Field Theory and String Vacua

I would now like to give a brief description of how two-dimensional (2d) conformal field theory (CFT) enters into string theory, and of how one actually implements the phenomenological criteria just discussed in terms of CFT's.

A string is a one-dimensional object, so as it moves through space-time it sweeps out a 2d surface, called the *world-sheet* — analogous to the one-dimensional world-line swept out by a point particle. (See Figure 7.) The equations of motion

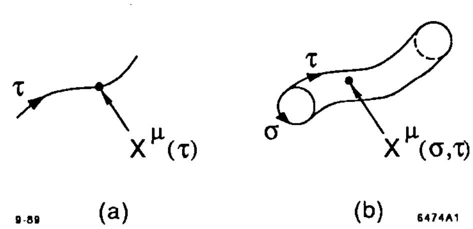

Figure 7. (a) The world-line swept out by a point particle moving through space-time. The position of the particle in space-time is given by $X^\mu(\tau)$. (b) The cylindrical world-sheet swept out by a closed string. The coordinates of the world-sheet are τ and σ; the position in space-time of a point on the world-sheet is given by $X^\mu(\sigma,\tau)$.

for a point particle can be obtained from an action, $S_{1d} = \int d\tau (ds/d\tau)$, equal to the proper length of the world-line (parametrized by τ). Similarly, the equations of motion of the world-sheet are determined by a 2d field theory, $S_{2d} = \int d\tau d\sigma \mathcal{L}_{2d}(\tau, \sigma)$, where τ and σ parametrize the world-sheet. In the simplest case S_{2d} is the area of the world-sheet, but there are many other possibilities. The Lagrangian $\mathcal{L}_{2d}$ should not be confused with the space-time effective Lagrangian $\mathcal{L}_{\text{eff}}$; also, the fields in the 2d field theory are called *world-sheet fields* in order to distinguish them from the space-time fields occurring in $\mathcal{L}_{\text{eff}}$.

When the motion of the string is quantized, the world-sheet fields become operators in the 2d quantum field theory and create eigenstates of the string's oscillations. These states can be identified as particles moving in space-time. Thus the full particle spectrum — in particular the massless spectrum — can be determined by enumerating the world-sheet fields. (Note that the space-time fields, acting as operators in the quantized effective field theory $\mathcal{L}_{\text{eff}}$, also create particle states, but only the massless particles, and only in accordance with the spectrum found by studying $\mathcal{L}_{2d}$.) It turns out that string *interactions* are also fixed uniquely by the choice of $\mathcal{L}_{2d}$; this means that $\mathcal{L}_{2d}$ determines not only the particle mass spectrum, but also all couplings between particles!

The simplest examples of world-sheet fields are those that represent the position of the world-sheet in space-time: $X^\mu(\sigma, \tau)$, where $\mu = 0, 1, 2, 3$ labels the four space-time coordinates and σ, τ are the two world-sheet coordinates. (See Figure 7.)

In an arbitrary 2d field theory, most of the remaining world-sheet fields do not have such a simple geometrical interpretation. However, many string vacua can be described as *compactifications*, in which some space-time dimensions are taken to have sizes of order the Planck length. In the prototypical example of a compactification, one dimension X^i lives on a circle with a radius R of order the

Planck length, and the rest X^μ parametrize space-time. If the X^μ are represented

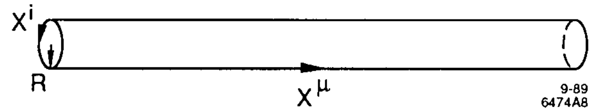

Figure 8. Compactification of a single extra coordinate X^i on a circle with radius R of order the Planck length. The four coordinates X^μ of Minkowski space-time are represented by the long direction of the "drinking straw".

by a single line, the result is the "drinking-straw" picture of Figure 8. When this compactification is used as a string vacuum, $X^i(\sigma,\tau)$ becomes a world-sheet field, just like $X^\mu(\sigma,\tau)$; it describes the position of the string in the compactified dimension. The moduli that parametrize string vacua also have a geometric interpretation in the case of a compactification: They represent the lengths of the internal dimensions, such as the radius R in the above example.

Not all 2d field theories give rise to *string vacua.* In a vacuum state, strings must not be created spontaneously — that is, all tadpole graphs must vanish. This non-trivial condition* on the 2d field theory will be satisfied if it is *conformally invariant*,[14,15] and has a few other properties to be described shortly.

I won't explain exactly why conformal invariance is required, but I should at least say what it is. The Lagrangian $\mathcal{L}_{2d}$ for the 2d field theory depends not only on the world-sheet fields X^μ, *etc.*, but also on the 2d metric $g_{\alpha\beta}$. Conformal invariance means that $\mathcal{L}_{2d}$ is invariant under a local change of scale on the world-sheet which preserves angles,

$$g_{\alpha\beta}(\sigma,\tau) \to e^{\phi(\sigma,\tau)} g_{\alpha\beta}(\sigma,\tau); \qquad (1)$$

for example the transformation of Figure 9.

⋆ The condition is somewhat trivial in the point-particle case, because one can adjust the particle interactions independently of S_{1d} so that particles are not created spontaneously. In the string case, however, the interactions are already fixed once S_{2d} is specified.

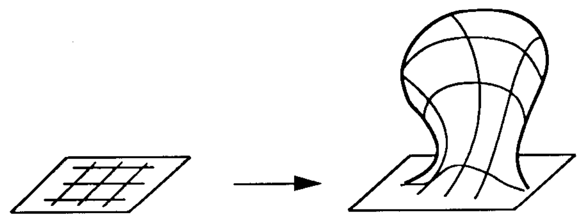

Figure 9. Example of a conformal transformation: a local change of scale on the surface that preserves angles.

In fact, 2d field theories with (1) as a classical symmetry generally develop a so-called *conformal anomaly* under the transformation at the quantum level. The details of this anomaly are not particularly important here. We just need to know that:

(a) The anomaly is characterized by a single real number c, which is additive in the sense that if a CFT consists of two non-interacting pieces, say $\mathcal{L}_{2\text{d}} = \mathcal{L}_{2\text{d}}^{(1)} + \mathcal{L}_{2\text{d}}^{(2)}$, then $c = c^{(1)} + c^{(2)}$.

(b) There is a contribution of $c = -26$ from the metric $g_{\alpha\beta}$ for the bosonic string, which must be cancelled by a contribution of $c = +26$ from world-sheet fields other than the metric $(X^{\mu}, \ldots)$; for the superstring $c = -15$ must be cancelled by $c = +15$.

(c) A single world-sheet field of the type $X^{\mu}(\sigma, \tau)$ contributes an anomaly $c = 1$, so if space-time is D-dimensional ($\mu = 0, 1, 2, \ldots, D-1$) the D fields will contribute $c = D$.

The "critical dimension" $D_c = 26$ for the bosonic string is obtained by assuming that there are no world-sheet fields other than X^{μ} (not counting the metric). For the superstring the critical dimension is $D_c = 10$. (It is 10 rather than 15 because each X^{μ} field has to be accompanied by a world-sheet superpartner ψ^{μ} which contributes an additional $c = \frac{1}{2}$). Four-dimensional (super)string vacua are constructed using only four fields X^{μ} that represent space-time coordinates (plus

ψ^μ in the super case), but also using extra "internal" fields — like X^i — in such a way that the total anomaly from fields other than the metric continues to have the correct value, either $c = 26$ or $c = 15$.

The most important world-sheet fields in a CFT, which also have the simplest behavior, are those that move in only one direction on the string, either to the left or to the right; thus they depend only on $\sigma_L = \tau - \sigma$, or only on $\sigma_R = \tau + \sigma$, and are called *left-moving* fields or *right-moving* fields, respectively. The 2d Lorentz properties of any world-sheet field can be summarized by its *scaling* or *conformal dimension*,[6] which generally will get anomalous contributions at the quantum level, and which can be split into left- and right-moving parts, denoted h and $\overline{h}$ $(h, \overline{h} \geq 0)$. The left-moving fields all have $h > 0$, $\overline{h} = 0$, while the right-moving fields have $h = 0$, $\overline{h} > 0$.

In any conformal field theory, the energy-momentum tensor $T_{\alpha\beta}$ provides an important pair of left- and right-moving fields. While $T_{\alpha\beta}$ is present in any field theory, in a CFT it is traceless, $T^\alpha_\alpha = 0$, which allows it to be split into the left- and right-moving components

$$\begin{aligned} &T(\sigma_L) \qquad \text{with } (h, \overline{h}) = (2, 0), \\ &\overline{T}(\sigma_R) \qquad \text{with } (h, \overline{h}) = (0, 2). \end{aligned}$$

The short-distance behavior of $T_{\alpha\beta}$ also determines the conformal anomaly c, because $T_{\alpha\beta}$ generates conformal transformations, in addition to its usual role as generator of rigid translations of the world-sheet. *Any* CFT with an energy-momentum tensor $T_{\alpha\beta}$ giving rise to $c = 26$ provides a vacuum for the bosonic string (at the classical level).

A superstring vacuum has a few more restrictions on it — the 2d CFT must be supersymmetric as well. In this case the energy-momentum tensor $T_{\alpha\beta}$ will have a fermionic superpartner $(T_F)_{\alpha\theta}$, which can also be split into left- and right-moving pieces:

$$T_F(\sigma_L) \qquad \text{with } (h,\overline{h}) = (\tfrac{3}{2},0),$$
$$\overline{T}_F(\sigma_R) \qquad \text{with } (h,\overline{h}) = (0,\tfrac{3}{2}).$$

Any such *super*conformal field theory with conformal anomaly $c = 15$ provides a classical superstring vacuum.

The heterotic string is constructed by combining the left-moving world-sheet fields of the bosonic string with the right-moving fields of the superstring; hence it requires the presence of T, $\overline{T}$ and $\overline{T}_F$ but not necessarily T_F. The left- and right-moving conformal anomalies are now different from each other — $c = 26$ and $\overline{c} = 15$ respectively. In fact, a very large number of such superconformal field theories are now known to exist. This embarrassment of riches is precisely *the (heterotic) string vacuum degeneracy problem.*

5. Phenomenological Constraints

So far I have discussed the constraints on the CFT that come just from consistency of (super)string propagation. Now I would like to focus on the heterotic string, and impose some of the additional "phenomenological" constraints discussed above, in the hopes of reducing the vacuum degeneracy problem somewhat. Many of these constraints require the presence (or absence) of particular world-sheet fields, often the relatively simple purely left-moving (or purely right-moving) fields.

• **Four-dimensional space-time** is implemented by requiring the two-dimensional Lagrangian to have the form

$$\mathcal{L}_{\text{2d}} = \mathcal{L}_0(X^\mu, \psi^\mu) + \mathcal{L}_{\text{internal}}(\varphi_i), \tag{2}$$

where

$$\mathcal{L}_0(X^\mu, \psi^\mu) = \frac{1}{2\pi} \int d\sigma d\tau \{\partial_\alpha X^\mu \partial^\alpha X_\mu + \overline{\psi}^\mu (1-\gamma^3)\gamma^\alpha \partial_\alpha \psi_\mu\}. \tag{3}$$

The fields X^μ describing the string's position in space-time ($\mu = 0,1,2,3$), and their world-sheet superpartners ψ^μ, are free (non-interacting) fields due to Eq. (3);

ψ^μ is a right-moving field, and X^μ can be split into left- and right-moving pieces, $\partial_{\sigma_L} X^\mu$ and $\partial_{\sigma_R} X^\mu$. The remaining "internal" fields φ_i do not interact with X^μ and ψ^μ. In the case where $\mathcal{L}_{2d}$ represents a compactification, they are the degrees of freedom associated with the (six) compactified dimensions: $X^i, \psi^i, \ldots$ The conformal anomalies associated with $\mathcal{L}_0(X^\mu, \psi^\mu)$ are $c = 4$ and $\bar{c} = 6$; hence $\mathcal{L}_{\text{internal}}$ must have

$$c_{\text{internal}} = 26 - 4 = 22, \qquad \bar{c}_{\text{internal}} = 15 - 6 = 9.$$

Additional constraints will now be imposed on $\mathcal{L}_{\text{internal}}$. The easiest constraints to implement have to do with the connection between massless particles (specifically, gravitinos, gauge bosons, quarks and leptons) and certain world-sheet fields having $(h, \bar{h}) = (1, \frac{1}{2})$ under the energy-momentum tensor for $\mathcal{L}_{2d}$. The contributions to $(h, \bar{h})$ coming from $\mathcal{L}_{\text{internal}}$ depend on the space-time Lorentz properties of the particles in question; they lead to the following restrictions on $\mathcal{L}_{\text{internal}}$:

- **Space-time supersymmetry** requires several additional right-moving fields; in particular, the right-moving component of the energy-momentum tensor $\overline{T}$ has a second superpartner $\overline{T}'_F$ in addition to $\overline{T}_F$:

$$\overline{T}'_F(\sigma_R) \qquad \text{with } (h, \bar{h}) = (0, \tfrac{3}{2}).$$

  This means the internal CFT possesses an *extended* ($N = 2$) world-sheet supersymmetry.[16,17]

- **Not too much space-time supersymmetry** (which would destroy chirality) requires the *absence* of all right-moving fields with $(h, \bar{h}) = (0, \frac{1}{2})$.[18]

- **A gauge group containing** $SU(3) \times SU(2) \times U(1)$ requires left-moving fields, one for each gauge boson:

$$J^a(\sigma_L) \qquad \text{with } (h, \bar{h}) = (1, 0),$$

  where a labels the generators of the gauge group. (See *e.g.* Ref. 19.) The properties of the J^a are completely specified by the choice of gauge group, *i.e.* the structure constants f^{abc}, and of certain positive integers k_i (one for each non-abelian factor in the gauge group) which show up in the short-distance behavior of the J^a. The k_i are important in determining which representations of the gauge group can appear in the spectrum, and in the relation between different gauge couplings (see below).

- **Massless quarks and leptons** require world-sheet fields $\Phi_i(\sigma_L, \sigma_R)$ that are *not* purely left- or right-moving (they have $(h, \overline{h}) = (1, \frac{1}{2})$). (See *e.g.* Ref. 10.) They have specific charges under the fields $J^a(\sigma_L)$, corresponding to their gauge quantum numbers, but otherwise they are not terribly well specified.

And then there are the criteria that are much harder to implement, such as **no fast proton decay** and **reasonable quark and lepton masses**. Both these quantities are related to certain CFT correlation functions $\langle \Phi_i \Phi_j \Phi_k \rangle$, but they are also very sensitive to other possible fields, to radiative corrections, and to various stages of symmetry breaking below $M_{\rm Pl}$, so in practice it is difficult to evaluate them for a specific model, let alone to implement them as general conditions.

Now that we have imposed some constraints on the CFT, we would like to see if the general properties of the allowed CFT's can in turn impose restrictions on the possible space-time effective Lagrangians. Unfortunately, not much progress has been made in this direction, at least for the heterotic string.⋆ A primitive example of a restriction on $\mathcal{L}_{\rm eff}$ is the tree-level relation[20,19]

$$\frac{g_{SU(3)}}{g_{SU(2)}} = \sqrt{\frac{k_{SU(2)}}{k_{SU(3)}}} \, . \tag{4}$$

Here $g_{SU(3)}$, $g_{SU(2)}$ are the strong and weak gauge coupling constants, and $k_{SU(3)}$, $k_{SU(2)}$ are the integers associated with the world-sheet fields J^a that were mentioned above. In a given string vacuum, each k_i takes on one specific value, but

⋆ For the closed superstring, on the other hand, one can show that there are *no* CFT's satisfying the combined constraints discussed above (though CFT's exist satisfying the individual constraints). This rules out *all* classical vacua of the closed superstring on phenomenological grounds.[18]

that value can (in principle) be any positive integer. It is instructive to compare this prediction with that for the grand unified theory based on $SU(5)$[21] (or any gauge group containing $SU(5)$ and with the same embedding of $SU(3) \times SU(2)$):

$$\frac{g_{SU(3)}}{g_{SU(2)}} = 1. \tag{5}$$

The string theory prediction (4) is clearly much less definite than the GUT prediction (5). On the other hand, in the effective Lagrangians for many string vacua the standard model gauge group is not embedded into a simple Lie group (such as $SU(5)$, $SO(10)$ or E_6); under these circumstances a generic field theory would give no gauge coupling prediction at all! In fact, almost all of the attempts to date to build phenomenological string models have involved choosing $k_{SU(3)} = k_{SU(2)} = 1$, in which case one recovers the usual GUT prediction. There are, however, no arguments excluding $k > 1$ on phenomenological or other grounds (provided k is not too large); indeed, explicit models featuring $k = 2$ have recently been constructed.[22]

Both Eqs. (4) and (5) are tree-level relations that will get loop corrections. In particular they will change significantly under renormalization from $M_{\rm Pl}$ ($M_{\rm GUT}$) down to M_W, and the change will depend on the masses of particles carrying $SU(3)$ and $SU(2)$ quantum numbers below $M_{\rm Pl}$ ($M_{\rm GUT}$).[23] This change introduces further model dependence into the string prediction, beyond the choice of $k_{SU(3)}$ and $k_{SU(2)}$, but this kind of dependence is common to GUT's as well.

The model dependence of even the primitive string prediction (4) raises the more general question of the sensitivity of models to both *discrete* and *continuous* modifications. In Figure 6, discrete modifications jump the vacuum from one "known" patch to another, whereas continuous modifications involve changing the vacuum expectation values of moduli and slide the vacuum along the troughs in a given patch. In general, models are very sensitive to discrete modifications: The gauge group and matter content — even the number of generations — can change drastically. On the other hand, models are actually rather insensitive to continuous modifications. For example, tree-level relations like (4) between gauge couplings do not change[20,17]; at least some of the matter content (usually including the quarks and leptons) remains the same[9,10]; and some of the Yukawa couplings among the matter fields even stay constant.[24] This insensitivity makes it difficult to tune low-energy predictions by continuous adjustments of the moduli, which I view as an encouraging result.

6. Specific Constructions

I would now like to make a few remarks about specific constructions of models. There are several different kinds of constructions, which overlap somewhat (different constructions can give rise to the same vacua), and which I will make no attempt to describe here.[25] They go under the names: Calabi-Yau compactifications,[8] orbifolds (symmetric[26,27] and asymmetric[28]), free fermions,[29] free bosons (or covariant lattices),[30] and tensor products of $N = 2$ superconformal field theories.[31,32] There are certainly at least thousands of models contained in these categories. Only a relative few have been analyzed in much detail. The more promising ones have certain features in common beyond the minimal criteria I mentioned previously:

- Three generations of light fermions. Four generations have been considered; one has then to prevent extra colored members of each generation from lurking at the 100 GeV scale (otherwise the $SU(3)_c$ coupling constant will fail to be asymptotically free and will blow up before the unification scale[33]). However, the very recent⋆ measurements at SLC[34] and LEP[35] of the width of the Z^0 now seem to rule out four generations with light neutrinos.
- Space-time supersymmetry is broken by non-perturbative effects in a "hidden sector". (It is generally believed that supersymmetry will not be broken perturbatively if it is present at tree-level.) The hidden sector generally consists of a strongly-interacting gauge theory with some gauginos λ but no other charged matter fields, plus some gauge singlet matter fields. The supersymmetry breaking is supposed to be triggered by a condensate of the gauginos,[36]

⋆ slightly postdating this talk, in fact!

$\langle\lambda\lambda\rangle^{1/3} \sim 10^{14}$ GeV. The particles in the hidden sector interact with observable particles only through gravitational-strength interactions, and so the scale of supersymmetry breaking in the observable sector is reduced by a factor of $M_{\rm Pl}^{-2}$, to $\sim M_W$. The masses of the superpartners of the quarks and leptons then fall in the usual range $\lesssim 1$ TeV or so.

- There are no matter fields transforming under higher-dimensional representations of $SU(3)$ (that is, only singlets and triplets are present) or of $SU(2)$ (only singlets and doublets are present). This is a quite generic feature; it is true for any vacuum in which $k_{SU(3)} = k_{SU(2)} = 1$.
- There are also usually no light "exotic" ($SU(2)$ singlet) quarks, but for a different reason: They can cause fast proton decay if they are lighter than $\sim 10^{15}$ GeV, unless certain Yukawa couplings happen to vanish.
- There are often a few additional $SU(2)$ doublets and singlets around at TeV energies.
- The gauge group that acts on the observable particles is not necessarily unified anywhere below $M_{\rm Pl}$; typical gauge groups are $SU(3)^3$ (found in a particular Calabi-Yau/$N=2$ tensor product construction[37,38,39]), $SU(3)\times SU(2)\times U(1)^n$ (found in various orbifold constructions[27]), or (flipped) $SU(5)\times U(1)$ (found in certain fermionic constructions[40]). Usually the gauge group is supposed to break spontaneously at a high "intermediate" mass scale $M_I \sim 10^{13\pm2}$ GeV, leaving at TeV energies only the standard model gauge group $SU(3)\times SU(2)\times U(1)$, plus perhaps one extra $U(1)$ factor (a Z' gauge boson).

7. Prospects

I will conclude by commenting on two of the major obstacles to extracting predictions from string theory. The first occurs in the analysis of any specific model — it is the question of how non-perturbative effects could lead to the breaking of space-time supersymmetry. As I mentioned earlier, there are often many fields (the moduli) whose potentials in V_{eff} are flat, *i.e.* their vacuum expectation values are undetermined to all orders in string perturbation theory. Some of the moduli can receive corrections to their potential from non-perturbative effects, in particular from interactions with the hidden sector gaugino condensate mentioned above, but perhaps also from other non-perturbative effects that have yet to be identified. In many cases the gaugino-induced corrections to the potentials cause vacuum expectation values for the moduli to run away to infinity! Thus we need to know: What is the full corrected potential? Does it have a stable minimum? If a minimum exists, is supersymmetry broken there? Finally, exactly how is the supersymmetry breaking manifested in the observable sector? There is a general understanding of the last question in supergravity (in terms of so-called "soft-breaking terms"), but the details can be sensitive to model-dependent parameters in $\mathcal{L}_{\text{eff}}$ (such as non-renormalizable kinetic-energy terms for observable fields, *etc.*).[12] The details can in turn be very important in determining many low-energy quantities, including the masses of superpartners but also the observed quark and lepton masses; the latter masses are sensitive to the pattern of supersymmetry breaking in models where there are additional scales of gauge symmetry breaking between M_W and M_{Pl}.[39]

A more fundamental obstacle to obtaining predictions from string theory is the issue of how one vacuum state is dynamically preferred over another. This issue cannot even be addressed until one has a formalism that can simultaneously describe two disconnected vacua, like two of the "islands" depicted in Figure 6. Until such a formalism is developed, the two ways one might hope to make phenomenological progress are: (1) to "get lucky" in finding "the right" vacuum, which would predict all the standard model parameters correctly; or (2) by trying to determine the general properties common to all the vacua, in the hopes of deciding whether *any* of them can lead to realistic physics at the 100 GeV scale. Clearly there is still a lot of theoretical work to be done before we know whether superstrings are a theory of everything or a theory of nothing.

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