CERN-TH.4474/86

SUPERSTRING PHENOMENOLOGY

J. Ellis

CERN, 1211 Geneva 23, Switzerland

ABSTRACT

Attempts to link the superstring with the reality of physics in four dimensions at energies ≤ 1 TeV are discussed. Phenomenological considerations bearing on the choice of compactification scheme are presented, and possible experimental signatures are mentioned.

Invited talk given at the
Second Nobel Symposium on Elementary Particle Physics:
'Unification of Fundamental Interactions',
Marstrand, Sweden, 2-7 June 1986

CERN-TH.4474/86
June 1986

Introduction

Almost everybody at this meeting agrees that superstring theory [1] is the greatest thing since sliced bread and certainly does everything [2] except make coffee. My task is to discuss how it can be used to make coffee too, i.e. how to make contact with reality, hereafter defined to be physics in four dimensions at energies less than 1 TeV. Some might think this a foolhardy enterprise in the absence of a better fundamental understanding of the theory. However, I think that superstring phenomenology is an appropriate activity for several reasons. One is that no matter how beautiful and convincing we may find the theory, we can call it science and convince the outside world only if we can provide experimental proof that it makes correct predictions, as well as postdictions. These pieces of experimental evidence might either be in the form of a ‘smoking gun’—something so cataclysmically remarkable that the only possible culprit is the superstring, or else an accumulation of circumstantial evidence as happened with the Standard Model. Another reason for doing superstring phenomenology is that, even if the underlying theory is unique, the way in which its underlying symmetry breaks in our part of the Universe may be ambiguous not only because of our temporary ignorance, but even in principle. At the moment, many four-dimensional classical vacua seem equally plausible, and phenomenology may be able to tell theorists that some of them are not viable for practical reasons. Finally, we must still admit the possibility that the superstring is actually incorrect, either in total or in part. It would still be a useful subject to study, because it throws out many creative new ideas which generate new approaches to some classical phenomenological problems such as grand unification, flavour, and the gauge hierarchy.

Superstring phenomenology is already a vast subject, and I cannot do all of it justice in this single talk. Accordingly, after a general outline of the usual strategy for making contact with reality in Section 1, and reminding you of some relevant features of four-dimensional supergravity models, I subsequently select some specific topics for further discussion. Section 2 contains phenomenological arguments which may help us figure out how nature has chosen to compactify the superstring: $(2,0)$ manifold compactifications are disfavoured [3–5] whilst a rank 5 low-energy gauge group is preferred [6]. Section 3 discusses the possible additional neutral currents [7] which may appear at low energies, and how they are constrained by current experiments. Section 4 speaks of possible supersymmetric relics from the Big Bang, the new light cast on them by phenomenological approaches to the superstring [8], and how they might be detected in the laboratory through the neutrino flux generated by their annihilations in the Sun [9], or by the small energies they deposit when scattering elastically in a detector [10]. (These two sections are somewhat compressed in the written version of this talk, since they have been discussed in detail in another conference talk I wrote recently [11]). Section 5 discusses the new matter particles which would appear in models based on the superstring, concentrating on gauginos, higgses and Higgs bosons [12], which are a particular interest of Gösta Ekspång, the only experimentalist brave enough to attend this Symposium. Finally, Section 6 summarizes some past and possible future history.

1. Strategy for phenomenology

The strategy often used to get from the fundamental formulation of the superstring at energies of order 10^{19} GeV down to the reality of four-dimensional physics at energies $\lesssim 10^3$ GeV is shown in Fig. 1. The first step is normally [13–15] to integrate out the massive string modes, obtaining the ten-dimensional Einstein–Yang–Mills supergravity theory with an (in principle) infinite sequence of terms of higher orders in $(1/m_{\text{P}}^2)^n$. The next step is usually to compactify the six surplus dimensions of $D = 10$ supergravity, thereby arriving at a four-dimensional supergravity theory. In order to preserve supersymmetry in four dimensions, as desired to safeguard the gauge hierarchy, one must compactify on a space with $SU(3)$ holonomy [13] or its discrete equivalent. The approach first favoured [13] was to compactify on a Ricci flat Kähler manifold with $SU(3)$ holonomy—called a Calabi–Yau manifold [16]—on which the spin and gauge connections are identified. Since this first approach, alternatives have been proposed in which either the gauge connection A is no longer identified naïvely with the spin connection ω [17], and/or the assumption of a manifold is relaxed to include spaces with a finite set of conical singularities, namely orbifolds [18]. One of the jobs for us phenomenologists, discussed in Section 2, is to see which of these compactification schemes is preferred. The resulting four-dimensional supergravity theory has a gauge group which is [7, 13] some subgroup of E_6 (in the case of Calabi–Yau) selected by the Hosotani [19] mechanism of symmetry breaking by expectation values of Wilson loop integrals in a non-Abelian generalization of the Bohm–Aharonov phase factors, or an analogous subgroup of $SO(10)$ or $SU(5)$ in spaces with $\omega \neq A$ [17, 20]. Another job for us phenomenologists is to see if there is any preference among the class of four-dimensional gauge groups, as also discussed in Section 2. The four-dimensional supergravity theory has a rather special ‘no-scale’ structure [21, 22], so that the effective scalar potential has flat directions corresponding to the undetermined moduli of the compactification space [23–26]. The degeneracy of the potential along these flat directions may be broken by higher-order non-perturbative or radiative corrections, and the latter provide a possible mechanism for generating the gauge hierarchy [22, 27] as discussed later in this section.

The two steps outlined above may actually be short-circuited by going directly from the superstring to a $d = 4$ supergravity theory. Many phenomenological consequences of compactification do not [25] depend critically on the assumption of an intermediate $D = 10$ supergravity stage, which is just as well since the scales of compactification and of string excitation are probably quite close to each other and to the Planck mass [28, 29].

The next phenomenological step is the breaking of supersymmetry, presumably spontaneously with the favoured trigger being the condensation of gauginos in the hidden E_8 sector leading to a non-zero gravitino mass $m_{3/2} \neq 0$ [30]. Thereafter, higher-order perturbative radiative corrections presumably [29, 31] pass the supersymmetry breaking on to the sparticles in the observable sector, giving them masses $m_{\tilde{g}}, m_{\tilde{t}}, m_{\tilde{b}}, m_H \lesssim 10^3$ GeV which are low enough to protect the gauge hierarchy. Thus the end product is a low-energy effective gauge theory in four dimensions with softly broken supersymmetry, and at least one extra neutral gauge boson with several extra matter particles and sparticles [7, 13] if one compactifies à la Calabi–Yau [16].

To set the stage for our discussion of dynamical models and the viability of different scenarios for compactification, let us first recall some aspects of the structure of four-dimensional supergravity models [32]. They are characterized by a Kähler potential

$$K(\phi^i, \phi_j^\dagger) = K_D(\phi^i, \phi_j^\dagger) + (K_F(\phi^i) + h.c.) \quad (1a)$$

where K_F is an F-term,

$$K_F(\phi^i) = \ln P(\phi^i) \quad (1b)$$

with P corresponding to the superpotential of global supersymmetry, and K_D is a D-term which fixes the kinetic energy terms:

$$K_i^j \partial^\mu \phi^i \partial_\mu \phi_j^\dagger : K_{i_1 \dots i_n}^{j_1 \dots j_n} \equiv \frac{\partial^{n+m} K}{\partial \phi^{i_1} \dots \partial \phi^{i_n} \partial \phi_{j_1}^\dagger \dots \partial \phi_{j_n}^\dagger} \quad (2)$$

The effective scalar potential is [32]

$$V(\phi^i, \phi_j^\dagger) = e^K \left(K_i (K^{-1})^i_j K^j - 3 \right) + (\text{D-terms}) \quad (3)$$

where the gravitino mass $m_{3/2} = e^{K/2}$. To include gauge fields in the specification we need a chiral function $f_{ab}(\phi^i)$ which appears in the F-term $f_{ab} W^a W^b$. Its expectation value fixes the gauge coupling $g^2 = 1/\text{Re } f$ and the CP-violating parameter $\theta \propto \text{Im } f$, and its derivatives yield a gaugino mass [32]

$$m_{1/2} = \frac{m_{3/2}}{4\text{Re } f} K_i (K^{-1})^i_j f^j \quad (4)$$

Assuming that the effective cosmological constant in four dimensions vanishes and taking the limit $m_P \rightarrow \infty$ with $m_{3/2}$ fixed, the potential (3) yields a low-energy theory with softly broken supersymmetry [33]:

$$V \simeq \sum_i \left| \frac{\partial P}{\partial \phi^i} \right|^2 + \sum_i m_{0_i}^2 |\phi^i|^2 + \left(\sum_{i,j,k} A_{ijk} \phi^i \phi^j \phi^k + h.c. \right) \quad (5)$$

where the effective scalar masses $m_{0_i} \propto m_{3/2}$ and the model-dependent coefficients $A_{ijk} \sim O(1) \times m_{0_i}$.

How does gauge symmetry breaking arise in a model like (5), which has $m_{0_i}^2 > 0$ at least in simple models? An important idea for generating gauge symmetry breaking is dimensional transmutation [34], as expressed through the renormalization of the soft supersymmetry-breaking parameters $\tilde{m} \equiv \{m_{0_i}, A_{ijk}, m_{1/2}\}$:

$$\mu \frac{\partial \tilde{m}^2}{\partial \mu} = \left[O\left(\frac{\alpha}{\pi}\right) \text{matrix} \right] \tilde{m}^2 \quad (6)$$

The idea is illustrated in Fig. 2a: a scalar (mass)² parameter which is positive (or zero) at the Planck scale is renormalized upwards (downwards) by gauge (Yukawa) couplings as the renormalization scale μ is decreased below m_P [35]. If the Yukawa couplings of a Higgs field overcome its gauge couplings, then it is possible that m_H^2 is driven to zero at some scale μ_0 [34]:

$$\mu_0 = m_P \exp\left(-\alpha^{(1)}/\alpha_\lambda\right) : \alpha_\lambda \equiv \frac{\lambda^2}{4\pi} \quad (7)$$

where the leading role is often played by the t-quark Yukawa coupling. Then a Higgs vacuum expectation value $\langle 0|H|0 \rangle = O(\mu_0)$ forms and hence m_W is also given by Eq. (7). In the particular case of no-scale models [22], the boundary conditions provided by the theory of everything (TOE) are $m_{0i}, A_{ijk} = 0, m_{1/2} \neq 0$ at $\mu = m_P$, in which case the evolution of m_H^2 should be as shown in Fig. 2b. Note that a scalar field N with no gauge interactions would have m_N^2 driven negative at $\mu \approx m_P$ as illustrated in Fig. 2c, leading to the danger that $\langle 0|N|0 \rangle = O(m_P)$ as discussed in Section 2. In conventional supergravity models, where the Higgs has gauge interactions, one can easily get $m_W \approx \mu_0 \approx 10^2$ GeV from a formula such as Eq. (7) with $m_t \approx \lambda_t \langle 0|H|0 \rangle \approx 40$ GeV [22, 36].

No-scale supergravity models with flat directions : $V = 0$, have special forms of the Kähler potential K [Eqs. (1)]. In the case of a single matter chiral field z , there is a unique possibility [2]

$$K = -3 \ln(z+z^+) \Rightarrow V(z) = 0 \quad (8)$$

The most symmetric generalization to many fields is [22]

$$K = -3 \ln(z+z^+ - \phi^i \phi_i^+) \Rightarrow V(z, \phi^i) = 0 \quad (9)$$

and more complicated examples [24, 26] give the same physics as Eq. (9) in the limit $|\phi^i| \ll |z| = O(m_P)$ [25]. Although the potential [Eqs. (8) and (9)] is flat at the tree level, higher-order radiative corrections will in general give a non-trivial z dependence: $V(z)$ as seen in Fig. 3 [27]. In this way one can hope to fix dynamically $\langle 0|z|0 \rangle$ and hence $m_{3/2} = m_{3/2}(z)$ and eventually the other supersymmetry-breaking mass parameters

$$\tilde{m} = \tilde{m}(z) \sim \mu_0 = m_P \exp\left(-O^{(1)}/\alpha\right) : \tilde{m} \approx \mu_0 \approx \langle 0|H|0 \rangle \quad (10)$$

An exciting development was the realization [23] that no-scale models at the classical level emerged naturally from compactification of the superstring. The flat directions correspond to undetermined

moduli of the space of compactification [25]. In the simple case of an overall rescaling of the compactified dimensions: $g_{\mu\nu}^{(6)} \rightarrow e^\sigma g_{\mu\nu}^{(6)}$, at the classical level [22],

$$K_D = -\ln(S+S^*) - 3\ln(T+T^*) + (\phi\text{-dependent terms}) \quad (11a)$$

$$f_{ab} = S \delta_{ab} \quad (11b)$$

where

$$\begin{aligned} S &= e^{3\sigma} \phi^{-3/4} + 3i\sqrt{2}b : \quad \epsilon_{\mu\nu\rho\sigma} \partial^\sigma b = \phi^{-3/2} e^{6\sigma} H_{\mu\nu\rho} \\ &\quad (\mu, \nu, \rho, \sigma = 1, \dots, 4) \\ T &= e^\sigma \phi^{3/4} + \phi^i \phi_i^* - i\sqrt{2}a : \quad \epsilon_{I\bar{J}} a = B_{I\bar{J}} \\ &\quad (I, \bar{J} = 5, \dots, 10) \end{aligned} \quad (12)$$

As will be discussed later, semi-classical world-sheet instanton effects [37] can modify the Kähler potential [Eq. (11a)] [3, 4, 25], although not f_{ab} [Eq. (11b)] [34], whilst radiative corrections can also modify f_{ab} [Eq. (11b)] [29] and may be able [38] to determine dynamically the supersymmetry-breaking scale $\tilde{m}$ [Eq. (10)] as proposed previously.

2. How to compactify?

As already mentioned, in order to protect the gauge hierarchy one wishes to retain supersymmetry upon compactification, which implies that one must have SU(3) holonomy [13] or its discrete equivalent. One's options then include manifolds [13] or orbifolds [18], and identification [13] or non-identification [17] of the spin connection ω with the gauge connection A (corresponding to standard or non-standard embedding of the discrete point group in the gauge group). Here I will mainly discuss the latter choice, but afterwards will make a few comments about the manifold/orbifold choice. If $\omega = A$, the world-sheet σ -model has (2,2) supersymmetry and the gauge group is $E_6 / \{\text{Wilson loops}\}$. Matter fields have the quantum numbers of **27** representations of E_6 , plus some gauge singlets, and possibly some survivors of split $(\mathbf{27} + \overline{\mathbf{27}})$ representations. Each **27** contains the 15 fields of a conventional Standard Model or SU(5) generation: $\{Q \equiv (u,d)_L\} + \{L \equiv (\nu, \ell)_L\} + u^c + d^c + \ell^c$, as well as a field ν^c with the same SO(10) properties as a right-handed neutrino, Higgs doublets $H + \bar{H}$, charge $|1/3|$ colour triplets $D + D^c$, and a neutral SO(10) singlet field N . The possible superpotential terms of these fields, together with their possible phenomenological roles, are [7]

$$\begin{aligned} P \ni & \underbrace{Qd^c \bar{H} + Qu^c H}_{m_q} + \underbrace{L\ell^c \bar{H}}_{m_\ell} + \underbrace{L\nu^c H}_{m_\nu} \\ & + \underbrace{H\bar{H}N}_{\text{mix Higgses}} + \underbrace{DD^c N}_{m_D} \\ & + \underbrace{(DQQ + D^c u^c d^c)}_{D \text{ decays: } B(D) = -2/3} \quad \text{OR} \quad \underbrace{(D^c LQ + D\ell^c u^c + D\nu^c d^c)}_{D \text{ decays: } B(D) = 1/3, L(D) = -1} \end{aligned} \quad (13)$$

Note that the first two terms in the last row conserve baryon number if one assigns $B(D) = -\frac{2}{3}$, whilst the last two do if $B(D) = +\frac{1}{3}$. Baryon number conservation would be strongly violated if both sets of terms were present, leading to rapid baryon decay [7]. Accordingly, one set or the other must vanish. Moreover, experiment tells us that the neutrino mass terms $L\nu^c H$ must be either absent or very small. It is possible to find simple discrete symmetries which impose acceptable combinations of Yukawa couplings [39], but no deep sources for them in the topology [40] of some complaisant compactification space are known. Here I will not discuss further the origins of these mysterious symmetries, but assume them as given.

If one compactifies in such a way that $\omega \neq A$ [17], the world-sheet σ -model has (2,0) supersymmetry, and one can reduce the four-dimensional gauge group to $SO(10)$ or $SU(5)/\{\text{Wilson lines}\}$. The $SO(10)$ matter fields can occupy $\mathbf{16} + \overline{\mathbf{16}} + \mathbf{10} + \mathbf{1}$ representations, whilst matter fields in $SU(5)$ models can occupy $\mathbf{5} + \overline{\mathbf{5}} + \mathbf{10} + \overline{\mathbf{10}} + \mathbf{1}$ representations [20]. In both cases it is again necessary to postulate mysterious discrete symmetries to forbid proton decay and neutrino masses, so they are no better than E_6 models in this respect.

Such (2,0) compactifications have a series of theoretical and phenomenological problems [3–5] as we now discuss. First, world-sheet instantons destroy the flatness of the effective scalar potential following from the Kähler potential (11a) [3, 4]. Flatness of the scalar potential could be preserved in the semi-classical approximation only if

$$\Delta K^{(1)} = \frac{g(\tau) + g^*(\tau^\dagger)}{\tau + \tau^\dagger} \quad (14)$$

with $g(\tau)$ holomorphic [5]. However, instantons in general do not give this form [3, 4], and holomorphic instantons in particular give [5]

$$\Delta K^{(1)} = -\frac{1}{2} \tau^{\dagger 2} (1 + \tau) e^{-\tau} + e^{-\tau} \quad (15)$$

The fact that $V(\tau) \neq 0$ means that such manifolds are not solutions to the equations of motion for generic values of τ . It may be that they collapse to orbifolds [18] with (2,0) supersymmetry, which are known to be solutions to the equations of motion, but it is not possible to prove this on the basis of a semiclassical calculation of the effective potential due to a single world-sheet instanton. A second theoretical problem is that when gauge non-singlet fields ϕ^i are included, instantons in (2,0) models give them supersymmetry-breaking masses $m_{0i} = f_i(\tau)$ which are presumably $O(m_P)$ [3, 5]. If present, they would destroy any hope of obtaining a gauge hierarchy which requires $m_{0i} \ll m_P$: to proceed further, we must assume that these large m_{0i} can be evaded somehow. It should be noted that both the above difficulties are absent in $\omega = A$ (2,2) compactifications, which generate no $V(\tau)$ [3, 4], and leave $m_{0i} = 0$ [3]. A third theoretical problem with $\omega \neq A$ (2, 0) compactifications is the possibility [41] that the Green–Schwarz anomaly cancellation mechanism [2] may fail if $\omega \neq A$. We recall that it entails writing the three-index tensor field strength

$$H = -dB + (CS)_G^3 - (CS)_L^3 \quad (16)$$

where $(CS)_{G,L}^3$ are the gauge and Lorentz Chern-Simons three-forms. For the mechanism to work, $[(CS)_G^3 - (CS)_L^3 - H]$ must be a trivial element of the cohomology $H^3[P(M,G)]$, where $P(M,G)$ is the principal gauge bundle on the compactification space M . It is clear that $[(CS)_G^3 - (CS)_L^3 - H]$ is trivial if $\omega = A$ so that $(CS)_G^3 = (CS)_L^3$, but might not be if $\omega \neq A$. This point should [41] be checked in any given compactification proposed. A fourth, and more phenomenological, problem is that of $H\bar{H}$ mixing. The dimensional transmutation mechanism described in Section 1 generates directly a vacuum expectation value for the Higgs H coupled to charge $+2/3$ quarks in formula (13), but not the other Higgs $\bar{H}$ coupled to charge $-1/3$ quarks and leptons. To get $\langle 0|\bar{H}|0\rangle \neq 0$ also, we need to couple them together, which can only be done by an $H\bar{H}N$ coupling as in formula (13), leading to an $A\lambda H\bar{H}N$ term in the effective scalar potential. We must postulate an $SU(5)$ or $SO(10)$ singlet field and trilinear coupling to play an analogous role in $(2,0)$ models [5]. To get a suitable vacuum with $\bar{v} \equiv \langle 0|\bar{H}|0\rangle \propto v \equiv \langle 0|H|0\rangle \approx m_W$, we need $x \equiv \langle 0|N|0\rangle = O(m_W/\lambda)$. How can this come about? The effective superpotential in an $SU(5)$ or $SO(10)$ model may contain a $\lambda_N N^3$ coupling. In this case the effective potential for the N field is [5]

$$V(N) = m_N^2 |N|^2 + \lambda^2 |N^2|^2 + 2\text{Re}(A_N \lambda_N N^3) \quad (17)$$

In the absence of gauge couplings for the N field, the $H\bar{H}N$ or N^3 Yukawa couplings automatically drive $m_N^2 < 0$ at some scale $\mu_0 \approx m_P$ as in Fig. 2c, leading to $\lambda_N x = O(|m_N|) = O(m_W)$. However, there are three possible degenerate vacua: $\Delta(\arg N) = 2\pi/3$, reflecting the discrete Z_3 symmetry of the potential (17). As the Universe cools, it can be expected to form domains in each of these vacua, and the intervening domain walls would be cosmologically unacceptable [5]. One way to avoid them would be to break the Z_3 symmetry of Eq. (17) by postulating higher-order terms $\lambda_N^{(n)} N^n : n \geq 4$ in the superpotential. These would have the effect of breaking the degeneracy of the Z_3 vacua, which would suppress the unwanted domains if it were large enough [42]. This would only be if $|\lambda_N^{(4)}| = O(1)$ [5]. Another possibility you might think of is to remove the $\lambda_N N^3$ coupling. However, in this case, since $m_N^2 \leq 0$ at $\mu \approx m_P$, one gets $x = \langle 0|N|0\rangle = O(m_P)$, which is unacceptably large [5]. Although it might be possible to avoid some, if not all, of the above arguments, $(2,0)$ compactifications [17] seem to be disfavoured, and we focus hereafter on $(2,2)$ compactifications.

Which $(2,2)$ compactifications are phenomenologically acceptable? The Hosotani mechanism [19] using Wilson loops can break E_6 to a subgroup of rank 5 or 6 [7]. To get to the Standard Model, which has rank 4, we must reduce the rank by 1 in the first case or by 2 in the second case. Reducing the rank by 1 can easily be done using $x = \langle 0|N|0\rangle \neq 0$ [38]. The $H\bar{H}N$ and $DD^c N$ couplings in formula (13) can easily generate $\mu_0 : m_N^2 \leq 0$ at energies $O(1)$ TeV, leading to $x = O(1 \text{ TeV})$ as

required for the $H\bar{H}$ mixing discussed above. This means in turn that the additional neutral gauge boson Z' has a mass $\lesssim O(1)$ TeV. Scenarios realizing the dimensional transmutation idea of Section 1 in this way are described in detail elsewhere [39].

To make a phenomenological model based on rank 6, one must reduce the rank by 2 using both $x = \langle 0|N|0\rangle \neq 0$ and $y \equiv \langle 0|\nu^c|0\rangle \neq 0$. However, there are severe obstacles to generating $y \neq 0$ [6]. To get $m_{\nu^c}^2 < 0$ in a dimensional transmutation scenario, one needs some Yukawa coupling to be large, and the only candidates in formula (13) are λ_ν , $L\nu^c H$ and $\lambda_D D\nu^c d^c$. However, λ_ν must be small or zero because of the neutrino mass problem discussed earlier. As for λ_D , it induces d/D mixing leading to [6]

$$m_d = O(\bar{\nu} x/y) \quad (18)$$

Since $m_d \geq O(m_e, m_u)$, one cannot have $y \gg x$, and since Higgs mixing requires $x = O(1 \text{ TeV})$, one must also have y and $m_D = O(1 \text{ TeV})$. The resulting d/D mixing gives rise to flavour-changing neutral currents [39], which are acceptably small only if $\lambda_D \lesssim 10^{-2}$. This is too small to drive $m_{\nu^c}^2 < 0$, which requires $\lambda_D \gtrsim 0.1$ [6]. We conclude that presumably $y = 0$ and one should stick to compactifications yielding a rank 5 gauge group in four dimensions [6]. For this we need [7] a Calabi-Yau manifold with a non-Abelian discrete symmetry group, but no suitable example has yet been exhibited.

What [43] of orbifold compactifications? They are solutions to the equations of motion, but it is unknown whether they have $m_{0_i} = 0$ automatically. The $H\bar{H}$ mixing problem for manifolds with $\omega \neq A$ is likely to be similar for orbifolds with non-standard embeddings of the point group. Unfortunately, there do not seem to be any interesting orbifold compactifications with the standard embedding and three generations [18]. The above arguments favouring a rank 5 gauge group are also likely to apply to orbifold models, in which it might be easier to get rank 5 because of the additional possibility for gauge symmetry breaking through the embedding of the full space group in the gauge group [18]. Orbifold compactifications are very enticing, but the jury [43] is still pondering its verdict on their phenomenological acceptability.

I conclude this section with a few comments about axions in four-dimensional models obtained from the superstring. At the classical level, only S couples to $W^a W^b$ via $f_{ab} = \delta_{ab} S$ in Eq. (11b), leading to an $\text{Im} S F \tilde{F}$ coupling and making $\text{Im} S$ a candidate to be an axion [23]. The imaginary part of T is completely decoupled at the classical level, since K [Eq. (11a)] depends on $\text{Re} T$ alone, and $\partial f / \partial T = 0$. World-sheet instantons in general induce interactions for $\text{Im} T$ [3, 4, 37]. For example, holomorphic world-sheet instantons in (2,2) models as well as in (2,0) models have action T , and so generate interactions $\propto e^{-T}$ as seen in Eq. (15) [4, 5]. Therefore $\text{Im} T$ is not a Goldstone boson. However, world-sheet instantons do not affect f_{ab} [Eq. (11b)] [3, 4], and in particular do not introduce any T dependence in f_{ab} , so do not allow one to regard $\text{Im} T$ as an axion. It has been suggested [29, 44, 45], that radiative corrections might change the form of f_{ab} . Indeed, the renormalizations of the gauge coupling constants in four dimensions can only be expressed as a

renormalization of f_{ab} . It has been suggested [45] that this cannot occur because $f_{ab} W^a W^b$ is an F-term. However, the situation is not the same as that for superpotential terms for which the well-known no-renormalization theorems [46] have been proved. By partial integration, one can rewrite $\int d^2\theta f_{ab} W^a W^b$ in terms of

$$\int d^4\theta \text{Ref} \left[D^\alpha V W_\alpha + \bar{D}_\alpha V \bar{W}^\alpha + V \bar{D}_\alpha \bar{W}^\alpha \right] \quad (19)$$

where $W_\alpha \equiv \bar{D}\bar{D}D_\alpha V$. Thus the gauge kinetic term is equivalent to a D-term which is renormalized in the usual way. Indeed, this is how the gauge coupling is renormalized [47] in globally supersymmetric theories. This mechanism can introduce some dependence on the gauge non-singlet fields into f_{ab} , but no dependence on T [29]. It is possible that higher genus string configurations might also change the form of f_{ab} , and perhaps introduce some T dependence, but this has not been demonstrated. The Green-Schwarz [2] anomaly density (which is neither supersymmetric nor generally covariant) couples T to gauge fields, but this is irrelevant [47]. Such a term is not generated in the heterotic string theory. Even if present in other string theories, its only job is to cancel the gauge variation of the one-loop effective action in the ten-dimensional supergravity theory [2]. Therefore the overall sum of the local Green-Schwarz term and the one-loop effective action is gauge-invariant, and its functional form has no reason to resemble the gauge-variant Green-Schwarz term in isolation [47]. A full cancellation of the action induced by the genus 1 term in the topological expansion for the heterotic string might renormalize f_{ab} , but this is not yet available. In any case $\text{Im } T$ is not an axion, although some field ($\text{Im } S + \dots$) is. It is not clear how this single field can solve the strong CP problem [48].

[The next three sections are somewhat brief, since a full description is already available in Ref. 11.]

3. Additional neutral currents?

Calabi-Yau compactification yields a four-dimensional gauge group which is at least as large as $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_E$, where the extra $U(1)_E$ hypercharges Y_E of the known particles are fixed [7], as is the magnitude of the extra fine structure constant α_E [49]. Moreover, we have given dynamical arguments [6] in favour of this unique rank 5 model in the previous section. The only thing that is not fixed is the mass of the extra neutral gauge boson Z' , which is $O(x)$ and hence should be $\lesssim 1$ TeV. The mass squared matrix for the neutral bosons in this minimal rank 5 model is [44]

$$m_{Z^0}^2 \begin{pmatrix} 1 & \sin\theta_w \frac{4v^2 - \bar{v}^2}{3(v^2 + \bar{v}^2)} \\ \sin\theta_w \frac{4v^2 - \bar{v}^2}{3(v^2 + \bar{v}^2)} & \sin^2\theta_w \frac{25x^2 + 16v^2 + \bar{v}^2}{9(v^2 + \bar{v}^2)} \end{pmatrix} \quad (20)$$

where we treat the vacuum expectation values v , $\bar{v}$, and x as phenomenological parameters to be fixed or constrained by experiment. There are several phenomenological bounds on x and $m_{Z'}$.

3.1 Direct searches at the CERN $p\bar{p}$ Collider

The non-observation of $Z' \rightarrow e^+e^-$ or $\mu^+\mu^-$ by the UA1 and UA2 Collaborations tells us that if the Z' can decay into three full E_6 generations of particles and sparticles [49, 50]

$$m_{Z'} \gtrsim 110 \text{ GeV} \quad (21a)$$

whilst if it can only decay into three generations of $SU(5)$ particles [50]

$$m_{Z'} \gtrsim 140 \text{ GeV} \quad (21b)$$

3.2 Shift in the Z^0 mass

Mixing between the Z^0 and the Z' will in general reduce the Z^0 mass so that

$$\overline{\sin^2 \theta_W} \equiv 1 - \frac{m_W^2}{m_Z^2} < \left(\frac{38.65 \text{ GeV}}{m_W} \right)^2 = \sin^2 \theta_W \quad (22)$$

The observed value of the Z^0 mass agrees well with the prediction of the minimal Standard Model, so one has [38] a lower bound $x/v \gtrsim 3.5$ for typical values of $\bar{v}/v \approx 1/2$, or

$$m_{Z'} \gtrsim 230 \text{ GeV} \quad (23)$$

3.3 $\nu N \rightarrow \nu X$

Deep-inelastic νN scattering experiments also agree with the Standard Model, allowing one to deduce [49, 51, 52]

$$m_{Z'} \gtrsim 150 \text{ GeV} \quad (24)$$

3.4 e^+e^- annihilation

Measurements of the forward-backward asymmetries at PEP and PETRA do not [49] at present provide an important constraint on $m_{Z'}$, but future higher-energy measurements at the SLC or LEP could provide the best bounds on $m_{Z'}$.

3.5 $\nu e \rightarrow \nu e$

A compilation of world data on this type of reaction yields [53]

$$m_{Z'} \gtrsim 133 \text{ GeV} \quad (25)$$

and a new experiment now starting to take data should be able to improve this significantly [54].

3.6 Primordial nucleosynthesis

The renormalizable superpotential (13) contains a possible source of Dirac masses for the neutrinos, but no Majorana mass terms. Therefore the right-handed neutrinos ν^c must have the same masses as the left-handed neutrinos ν_L , which are either very small or zero for some as yet unknown reason. The ν^c have no $SU(2)_L \times U(1)_Y$ interactions and so do not appear in Z^0 decays (which currently provide the bound $N_{\nu_L} < 5.4 \pm 1.0$), and they only make a small contribution to the reaction $e^+e^- \rightarrow \gamma + \text{nothing}$ (which allows $N_{\nu_L} < 8$). However, the ν^c would have been present in the early Universe, and would have affected primordial nucleosynthesis, which is normally quoted [55] as providing the bound,

$$N_{\nu_L} < 4 \quad (26)$$

We must ask ourselves [56]: a) do we believe this bound? and if we do, b) to get g_{ν^c} sufficiently small compared with g_{ν_L} , how large must $m_{Z'}$ be compared with m_{Z^0} ?

There are three crucial ingredients [55] in the upper bound on N_{ν_L} , namely i) an upper bound on the abundance of ^4He : $Y < 0.254$; ii) a lower bound on the abundance of deuterium and ^3He : $X > 10^{-4}$; and iii) a lower limit on the neutron half-life of 10.4 minutes. However, each of these ingredients is subject to some uncertainty [56]. i) There are large systematic errors [57] in observations of the ^4He abundance Y , and we prefer to take $Y < 0.26$. ii) Some observations [58] suggest that $D + ^3\text{He}$ may be destroyed by first-generation stars, rather than created as in the Standard Model for the chemical evolution of the galaxy. Therefore, to be on the safe side we take $X < 5 \times 10^{-4}$. iii) Some experiments [59] on the neutron half-life report $\tau_{1/2} \approx 10.2$ minutes, so we take this as our lower bound. Each of these three effects increases slightly the allowable number of equivalent left-handed neutrinos, and we [56] finish with

$$N_{\nu_L} < 5.5 \quad (27)$$

Moreover, we do not think that $N_{\nu_L} = 6$ can be completely excluded, in view of the astrophysical uncertainties. However, our astrophysical friends [60] do not agree.

The density of the ν^c can be suppressed relative to the ν_L if they decouple at a higher temperature T_d : effectively [56]

$$N_{\nu_L} < \begin{Bmatrix} 5.5 \\ 4 \end{Bmatrix} \text{ if } T_d > \begin{Bmatrix} \frac{1}{2} m_\mu \\ 300 \text{ MeV} \end{Bmatrix} \quad (28)$$

The decoupling temperature T_d will be higher if the ν^c interactions are weaker, i.e. if $m_{Z'} \gg m_Z$. Numerically, we find [56]

$$m_{Z'} \gtrsim \begin{Bmatrix} 330 \text{ GeV} \\ 780 \text{ GeV} \end{Bmatrix} \text{ if } N_{\nu_L} < \begin{Bmatrix} 5.5 \\ 4 \end{Bmatrix} \quad (29)$$

However, the cosmological bound on ρ_{ν_L} could perhaps be relaxed even further, or perhaps some clever person will find a way of making $m_{\nu^c} \gg m_{\nu_L}$. Thus, although the bound (29) is stronger than the previous particle physics bounds (21) and (23) to (25), a particle physicist could be forgiven for regarding it as less certain.

4. Relics from the Big Bang?

These are expected in most supersymmetric theories [61], because there is a discrete R-parity which tells us that the highest supersymmetric particle (called LSP, denoted by χ) is stable [62]. Usually the LSP is a neutral spin- $1/2$ fermion, a mixture of $\tilde{\gamma}/\tilde{Z}^0/\tilde{H}^0/\tilde{\bar{H}}^0$, and weighs at least a few GeV, thereby becoming a candidate for cold dark matter [61]. In the sort of low-energy model inspired by the superstring which we have been discussing here, the LSP could also contain admixtures of the $\tilde{B}_E$ and the $\tilde{N}$ [8]. When we diagonalize the 6×6 neutral fermion mass matrix, we find that the LSP has $m_\chi \gtrsim 15$ GeV and contains large $\tilde{\gamma}$ and $\tilde{Z}^0$ components— if it only had large Higgsino components, the relic ρ_χ would be far above the critical density ρ_c . As it is, we find [8]

$$\Omega_\chi \equiv \rho_\chi/\rho_c = 2^{0 \pm 1} \quad (30)$$

for a wide range of parameters, if we take the present Hubble expansion rate H_0 to be $50 \text{ km s}^{-1} \text{ Mpc}^{-1}$. Alternatively, if we take $\Omega_\chi = 1$ as suggested by inflation, we deduce [8]

$$H_0 = 50 \times 2^{0 \pm \frac{1}{2}} \text{ km s}^{-1} \text{ Mpc}^{-1} \quad (31)$$

which is consistent with observation.

There are two possible ways of detecting these superstring relics in the laboratory. One involves their capture in the Sun and subsequent mutual annihilations there which could provide an observable neutrino flux at the Earth's surface, with typical energies $\langle E_\nu \rangle$ a fraction of m_χ [9]. The annihilation neutrino flux could be comparable with the cosmic-ray neutrino flux if $m_\chi > m_t$ [8]. A dedicated search using the directionality of the apparatus to look at the Sun might eventually be able to explore the possibility that $m_\chi > m_t$. A more direct way to look for the LSP would be to search for its elastic scattering off matter in the laboratory [10], which would deposit several keV of recoil

energy. Taking a local relic density $f_\chi \approx 0.3 \text{ GeV/cm}^3$ and a mean velocity $\langle v \rangle \approx 300 \text{ km/s}$ relative to the Earth, we find [8]

$$\left. \begin{aligned} \text{events/ton/day} &= 23, 8, 6 \\ \text{for } m_\chi &= 20, 30, 40 \text{ GeV} \end{aligned} \right\} \quad (32)$$

Thus detecting the superstring LSP could be possible.

5. Extra matter particles?

For definiteness, we will again work in the minimal model with a rank 5 gauge group, three generations of 27 fields, and no light ‘survivors’ from split $(27 + \overline{27})$ multiplets. The extra matter fields include the spartners of conventional particles, and their masses are in calculable ratios if $m_{1/2} \gg m_0$ at $\mu \sim m_P$ as we believe [29]:

$$m_{\tilde{q}} : m_{\tilde{g}} : m_{\tilde{l}_L} : m_{\tilde{l}_R} : m_{\tilde{\tau}} = 1.95 : 1 : 0.7 : 0.4 : \frac{1}{2} \quad (33)$$

+ D-terms

Compiling all the experimental bounds on sparticle masses, and also using phenomenological constraints such as $m_\tau^2 > 0$ (in order to avoid lepton number violation) we arrive at the lower bounds on new particle masses shown in Table I [38]. Also shown are qualitative upper bounds which follow from requiring the electroweak vacuum to be reasonably stable against radiative corrections [38]. This strongly suggests that the new particles weigh less than a few hundred GeV, but cannot be quantified reliably.

The new particle spectrum abounds [12] in mixtures of gauginos and higgsinos, whose mass matrices are simplified if one works in a field basis where only one each of the H , $\bar{H}$, and N fields has a vacuum expectation value [63]. There are four *charged fermions*, two of which are mixtures between the $\tilde{W}^\pm$ and the partners $\tilde{H}_3^\pm$ of the combinations of Higgs fields with non-zero vacuum expectation values, the lighter of which has [12]

$$\tilde{m}_3^\pm \gtrsim 436 \text{ eV} \quad (34)$$

and two $\tilde{H}_{1,2}^\pm$ which are partners of the other Higgs fields, and have masses [12]

$$\tilde{m}_{1,2}^\pm \gtrsim 236 \text{ eV} \quad (35)$$

There are a total of 12 *neutral fermions*, and the LSP mentioned in the previous section is the lightest mass eigenstate of a 6×6 submatrix mixing the gauginos with $\tilde{H}_3^0$, $\tilde{\bar{H}}_3^0$, and $\tilde{N}_3^0$. The mass eigenstates of the other 6×6 submatrix are all heavier than this LSP [12]:

$$\tilde{m}^0 \geq m_\chi \geq 15 \text{ GeV} \quad (36)$$

There are three *charged Higgs bosons*, one of which tends to be heavier than the $W^\pm$ as in the minimal supersymmetric Standard Model [12]:

$$m_3^\pm \geq 90 \text{ GeV} \quad (37)$$

whilst the other two could be as light as the present experimental limit

$$m_{1,2}^\pm \geq 20 \text{ GeV} \quad (38)$$

As for *neutral Higgs bosons*, one expects 16 even in our minimal model. There is one scalar H^0 for which we can establish a lower limit [12],

$$m_{H^0} \geq 65 \text{ GeV} \quad (39)$$

and a pseudoscalar a with [12]

$$m_a \geq 100 \text{ GeV} \quad (40)$$

Figure 2 shows values of the H_L^0 and a masses for likely values of the parameters x/v , $\bar{v}/v$, $m_{1,2}$ and the $H\tilde{H}N$ Yukawa coupling λ . Some of the other neutral scalars could be light, but in general they would have masses $O(x)$.

The light charged fermions mentioned above [formulae (34) and (35)] could be copiously produced at LEP, and the lighter charged Higgses [38] could also be detectable. The neutral Higgs H^0 would be produced via the reaction $e^+e^- \rightarrow Z^0 + H^0$ with a cross-section similar [12] to that for the minimal Higgs in the Weinberg–Salam model [64]. Also, if toponium θ is heavier than the H^0 or the a , the radiative decay $\theta \rightarrow H^0 + \gamma$ or $a + \gamma$ would occur at rates comparable [12] to that expected in the Weinberg–Salam model [64]. Thus there may be several interesting experimental possibilities at LEP!

Also included in the **27** representations are D particles, for whose scalar and fermionic components we have only found the experimental limit [65]

$$m_{D_{0, \frac{1}{2}}} \geq 20 \text{ GeV} \quad (41)$$

It is possible, in principle, for one of the D_0 scalars to be lighter than the fermion $D_{1/2}$. The fermions $D_{1/2}$ can be pair-produced in hadron-hadron collisions with cross-sections the same as those of t -quarks of the same mass. Likewise, D_0 scalars can be pair-produced at the same rate as conventional squarks. To the extent that UA1 can explore $m_t, m_{\tilde{q}} \lesssim 40$ GeV, they can also explore $m_{D_0, D_{1/2}} \leq 40$ GeV. In the case where the DQQ and $D^c u^c d^c$ couplings are present in the superpotential (13), the single production mechanism $p\bar{p} \rightarrow D_0 + X$ is possible, whilst if the DLQ etc. couplings are present, single production $ep \rightarrow D_0 + X$ is possible [65]. In the former case, $D_0 \rightarrow \bar{q} + q$ and $D_{1/2} \rightarrow \bar{q} + q + \chi$ decays dominate, whilst in the latter case $D_0 \rightarrow q + \ell$ and $D_{1/2} \rightarrow q + \ell + \chi$ decays dominate.

Thus we have the following list [65] of possible signatures for superstring particles: bumps in $(\ell + \text{jet})$ or in $(\text{jet} + \text{jet})$ combinations ($D_0?$), missing energy events due to $q + \nu$ ($D_0?$) or $\bar{q} + q + \chi$ ($D_{1/2}?$), and $\ell + \text{jet} + \text{missing energy}$ events due to $D_{1/2} \rightarrow q + \ell + \chi$. These are in addition to the possible second neutral gauge boson, the shift in the mass of the first Z^0 , and possible additional neutral currents discussed in Section 3, as well as the possibility of detecting superstring relic particles either indirectly through their neutrinos from the Sun, or directly through their elastic scattering in the laboratory, as discussed in Section 4.

6. Summary

Superstring phenomenology is a very uncertain subject, and some might characterize it as premature or even foolhardy. However, physics is a science, and our theories must be subjected to experimental test and verification. We should not be able to convince ourselves of the validity of even the most elegant theory in the absence of experimental evidence for it. It may be that we already know this evidence, and that the superstring will soon give a triumphant postdiction of some or all of the 20 parameters of the Standard Model. Or it could be that there is some dramatic new particle or interaction waiting out there to be discovered. Unfortunately, we cannot be sure what low-energy four-dimensional physics is predicted by the superstring until we know how to compactify it. It could be that there is no unique theoretically viable compactification, and that we must choose the one appropriate to our piece of the Universe by quite phenomenological arguments. Even if there is a unique theoretically viable compactification, such phenomenological considerations may help the topological theorists to find it. In Section 2 of this talk I have addressed this issue and given some reasons why (2,2) compactifications leading to a rank 5 gauge group in four dimensions may be favoured [3–6]. Unfortunately, we have no known example of an interesting compactification of this type. So far we still have considerable ambiguity in the possible low-energy manifestations of the superstring. We may well have additional neutral currents and/or matter fields, as I have described in Sections 3, 4, and 5. Looking further afield, some compactification schemes may predict fractionally charged particles and multiply charged monopoles with masses $O(m_P)$ [66]. None of these signatures is unambiguous: the superstring may not predict them, and if they are seen it is in general possible to find alternative explanations for them.

Nevertheless, we can hope that the history of gauge theories shown in Table II will repeat itself. Perhaps the lone experimentalist at this meeting or some of his colleagues will be inspired to look for

some signature of the superstring (cf. Gargamelle). Perhaps some circumstantial evidence in favour of the superstring will be found (cf. neutral currents), to be followed by another experimental evolution (cf. the November J/ψ revolution). As John Schwarz likes to say [67], the superstring teaches us to expect miracles.

REFERENCES

- [1] Green, M.B., *Surv. High Energy Phys.* **3**, 127 (1983);
Schwarz, J.H., *Phys. Rep.* **89C**, 223 (1982).
- [2] Green, M.B., and J.H. Schwarz, *Phys. Lett.* **149B**, 117 (1984) and **151B**, 21 (1985).
- [3] Ellis, J., C. Gómez, D.V. Nanopoulos and M. Quirós, *Phys. Lett.* **173B**, 59 (1986).
- [4] Dine, M., N. Seiberg, X.-G. Wen and E. Witten, Princeton University preprint (1986).
- [5] Ellis, J., K. Enqvist, D.V. Nanopoulos, K. Olive, M. Quirós and F. Zwirner, CERN preprint TH. 4449/86 (1986).
- [6] Campbell, B.A., J. Ellis, D.V. Nanopoulos and K. Olive, CERN preprint TH. 4469/86 (1986).
- [7] Witten, E., *Nucl. Phys.* **B258**, 75 (1985);
Dine, M., V. Kaplunovsky, M. Mangano, C. Nappi and N. Seiberg, *Nucl. Phys.* **B259**, 519 (1985);
Breit, J.D., B.A. Ovrut and G. Segrè, *Phys. Lett.* **158B**, 33 (1985);
Cecotti, S., J.-P. Derendinger, S. Ferrara, L. Girardello and M. Roncadelli, *Phys. Lett.* **156B**, 318 (1985).
- [8] Campbell, B. A., J. Ellis, K. Enqvist, J.S. Hagelin, D.V. Nanopoulos and K.A. Olive, *Phys. Lett.* **173B**, 270 (1986).
- [9] Silk, J., K. Olive and M. Srednicki, *Phys. Rev. Lett.* **55**, 257 (1985) and Univ. of Minnesota preprint UMN-TH-553 (1986).
- [10] Goodman, M., and E. Witten, *Phys. Rev.* **D30**, 3059 (1985).
- [11] Ellis, J., CERN preprint TH. 4439/86 (1986).
- [12] Ellis, J., D.V. Nanopoulos, S.T. Petcov and F. Zwirner, CERN preprint TH. 4454/86 (1986).
- [13] Candelas, P., G.J. Horowitz, A. Strominger and E. Witten, *Nucl. Phys.* **B258**, 46 (1985).
- [14] Fradkin, E.S. and A.A. Tseytlin, *Phys. Lett.* **160B**, 69 (1985) and *Nucl. Phys.* **B261**, 1, (1985).
Callan, C.G., E. Martinec, D. Friedan and M. Perry, *Nucl. Phys.* **B262**, 593 (1985).
- [15] Callan, C.G., talk at this Symposium.
- [16] Calabi, E., 'Algebraic Geometry and Topology: a Symposium in Honour of S. Lefschetz', p. 78, Princeton University Press, 1957.
Yau, S.T., *Proc. Nat. Acad. Sci.* **74**, 1798 (1977).
- [17] Witten, E., *Nucl. Phys.* **B268**, 79 (1986).
- [18] Dixon, L., J.A. Harvey, C. Vafa and E. Witten, *Nucl. Phys.* **B261**, 678 (1985) and Princeton University preprint (1986).
- [19] Hosotani, Y., *Phys. Lett.* **129B**, 193 (1983).
- [20] Greene, B.R., K.H. Kirklin and P.J. Miron, Oxford Univ. Dept. of Theor. Phys. preprint 6/86 (1986).
- [21] Cremmer, E., S. Ferrara, C. Kounnas and D.V. Nanopoulos, *Phys. Lett.* **133B**, 61 (1983).
- [22] Ellis, J., C. Kounnas and D.V. Nanopoulos, *Nucl. Phys.* **B247**, 373 (1984).
- [23] Witten, E., *Phys. Lett.* **155B**, 151 (1985).
- [24] Burgess, C.P., A. Font and F. Quevedo, *Nucl. Phys.* **B272**, 661 (1986).

- [25] Ellis, J., C. Gómez and D.V. Nanopoulos, Phys. Lett. **171B**, 203 (1986).
- [26] Ferrara, S., C. Kounnas and M. Porrati, U.C. Berkeley preprint UCB-PTH-86/5 (1986).
- [27] Ellis, J., A. B. Lahanas, D.V. Nanopoulos and K. Tamvakis, Phys. Lett. **134B**, 429 (1984);
Ellis, J., C. Kounnas and D.V. Nanopoulos, Nucl. Phys. **B241**, 406 (1984).
- [28] Dine, M. and N. Seiberg, Phys. Rev. Lett. **55**, 366 (1985);
Kaplunovsky, V., Phys. Rev. Lett. **55**, 1036 (1985).
- [29] Cohen, E., J. Ellis, K. Enqvist and D.V. Nanopoulos, Phys. Lett. **161B**, 85 (1985).
- [30] Derendinger, J.P., L. Ibáñez and H.-P. Nilles, Phys. Lett. **155B**, 65 (1985);
Dine, M., R. Rohm, N. Seiberg and E. Witten, Phys. Lett. **156B**, 55 (1985).
- [31] Ahn, Y.J. and J.D. Breit, Nucl. Phys. **B273**, 75 (1986).
Diamandis, G.A., J. Ellis, A.B. Lahanas and D.V. Nanopoulos, Phys. Lett. **173B**, 303 (1986);
Ahn, Y.J. and G. Segrè, private communication (1986).
- [32] Cremmer, E., B. Julia et al., Phys. Lett. **79B**, 28 (1978) and Nucl. Phys. **B147**, 105 (1979);
Cremmer, E., S. Ferrara, L. Girardello and A. Van Proeyen, Phys. Lett. **116B**, 231 (1982) and
Nucl. Phys. **B212**, 413 (1983);
Arnowitz, R., A.H. Chamseddine and P. Nath, Phys. Rev. Lett. **49**, 970 (1982) and **50**, 232
(1983) and Phys. Lett. **121B**, 33 (1983);
Bagger, J., and E. Witten, Phys. Lett. **115B**, 202 (1982) and **118B**, 103 (1982);
Bagger, J., Nucl. Phys. **B211**, 302 (1983).
- [33] Barbieri, R., S. Ferrara and C. Savoy, Phys. Lett. **119B**, 343 (1982);
Nilles, H.-P., M. Srednicki and D. Wyler, Phys. Lett. **120B**, 345 (1983);
Arnowitz, R., A.H. Chamseddine and P. Nath in Ref. 32.
- [34] Ellis, J., J.S. Hagelin, D.V. Nanopoulos and K. Tamvakis, Phys. Lett. **125B**, 275 (1983).
- [35] Inoue, K., A. Kakuto, H. Komatsu and S. Takeshita, Prog. Theor. Phys. **68**, 927 (1982) and
71, 413 (1984).
- [36] Ellis, J., C. Kounnas and D.V. Nanopoulos, Phys. Lett. **143B**, 410 (1984);
Ellis, J., K. Enqvist and D.V. Nanopoulos, Phys. Lett. **147B**, 99 (1984).
- [37] Wen, X.-G., and E. Witten, Phys. Lett. **166B**, 397 (1986).
- [38] Ellis, J., K. Enqvist, D.V. Nanopoulos and F. Zwirner, Nucl. Phys. **B276**, 14 (1986) and Mod.
Phys. Lett. **A1**, 57 (1986).
- [39] Campbell, B.A., J. Ellis, K. Enqvist, M.K. Gaillard and D.V. Nanopoulos, CERN preprint
TH. 4473/86 (1986).
- [40] Strominger, A., and E. Witten, Commun. Math. Phys. **101**, 341 (1985).
- [41] Bonora, L., and P. Cotta-Ramusino, Phys. Lett. **169B**, 187 (1986).
- [42] Vilenkin, A., Phys. Rep. **121C**, 263 (1985).
- [43] Campbell, B.A., J. Ellis, C. Gómez and D.V. Nanopoulos, CERN TH preprint in preparation
(1986).
- [44] Derendinger, J.-P., L. Ibáñez and H.-P. Nilles, Nucl. Phys. **B267**, 365 (1985);
Ibáñez L., and H.-P. Nilles, Phys. Lett. **169B**, 354 (1986).
- [45] Ibáñez, L. and J. Mas, CERN preprint TH. 4426/86 (1986).

- [46] Wess, J., and B. Zumino, Phys. Lett. **49B**, 52 (1974);
 Iliopoulos J., and B. Zumino, Nucl. Phys. **B76**, 310 (1974);
 Ferrara, S., J. Iliopoulos and B. Zumino, Nucl. Phys. **B77**, 413 (1974).
- [47] Ellis, J., C. Gómez and D.V. Nanopoulos, Phys. Lett. **168B**, 215 (1986).
- [48] Dine, M. and N. Seiberg, Nucl. Phys. **B273**, 109 (1986).
- [49] Cohen, E., J. Ellis, K. Enqvist and D.V. Nanopoulos, Phys. Lett. **165B**, 76 (1985).
- [50] Barger, V., N. Deshpande and K. Whisnant, Phys. Rev. Lett. **56**, 30 (1986).
- [51] Durkin, L.S. and P. Langacker, Phys. Lett. **166B**, 436 (1986).
- [52] Del Aguila, F., G. Blair, M. Daniel and G.G. Ross, CERN preprint TH. 4376/86 (1986).
- [53] CHARM Collaboration, F. Bergsma et al., CERN EP preprint in preparation (1986).
- [54] CHARM II Collaboration, CERN proposal CERN/SPSC/83-24 (1983).
- [55] Yang, J., M.S. Turner, G. Steigman, D. Schramm and K. Olive, Astrophys. J. **281**, 493 (1984).
- [56] Ellis, J., K. Enqvist, D.V. Nanopoulos and S. Sarkar, Phys. Lett. **167B**, 457 (1986).
- [57] Davidson, K. and T.D. Kinman, Astrophys. J. Suppl. **58**, 321 (1985).
- [58] Rood, R.T., T.M. Bania and T.L. Wilson, Astrophys. J. **280**, 629 (1984).
- [59] Bondarenko, L.N., et al., Sov. Phys. – JETP Lett. **28**, 303 (1978).
- [60] Steigman, G., K. Olive, D. Schramm and M.S. Turner, Univ. of Minnesota preprint UMN-TH-562 (1986).
- [61] Goldberg, H., Phys. Rev. Lett. **50**, 1419 (1983).
 Ellis, J., J.S. Hagelin, D.V. Nanopoulos, K. Olive and M. Srednicki, Nucl. Phys. **B238**, 453 (1984).
- [62] Fayet, P., talk at this Symposium.
- [63] Georgi, H. and D.V. Nanopoulos, Phys. Lett. **82B**, 95 (1979).
- [64] Baer, H., et al., Physics at LEP, CERN 86-02, Vol. 1, p. 297 (1986).
- [65] Angelopoulos, V., J. Ellis, H. Kowalski, D.V. Nanopoulos, N. Tracas and F. Zwirner, CERN preprint in preparation (1986).
- [66] Wen, X.-G., and E. Witten, Princeton Univ. preprint (1985).
- [67] Schwarz, J.H., talk at this Symposium.

Table I

Limits on the masses of unseen particles in the minimal low-energy superstring models considered in the text. The first column gives lower limits, using experimental data, the assumption that supersymmetry breaking in the observable sector is triggered by a universal primordial gaugino mass and the result $0.2 \lesssim \bar{v}/v \lesssim 0.6$ of our dynamical calculations. The second column gives more tentative and model-dependent upper limits, obtained by imposing a ‘naturalness’ constraint. All the masses are expressed in GeV units.

Particle	Lower limit	Upper limit
t	–	70
Z_2	185	320
$\tilde{g}$	100	250
$\tilde{\gamma}$	15	35
$\tilde{e}_L$	55	150
$\tilde{e}_R$	90	170
$\tilde{\nu}$	0	150
$\tilde{q}_{1,2}$	180	500

Table II

Summary of past and possible future history

	Past	Future
Beautiful Theory	Yang–Mills	Superstring
Formalism not properly understood	before B.R.S.	before a complete covariant formulation of interacting superstrings
Renormalizable/Finite?	’t Hooft	Green, Schwarz, ...
Full structure only apparent at unreachably high energies	$m_W \sim 80 \text{ GeV}$	$m_P \sim 10^{19} \text{ GeV}$
Ambiguous predictions for low energies	neutral currents? heavy leptons?	more neutral currents? more quarks?
Stimulates experiment	Gargamelle	you?
Evidence found	Yes	?

Figure captions

Fig. 1 Sketch of the strategy for making contact between the superstring and reality.

Fig. 2 a) Contours of $m_{H_L^0}$ in the $(x/v, \bar{v}/v)$ -plane corresponding to the representative values $m_{1/2} = 250 \text{ GeV}$, $\lambda = 0.15$ and $A/m_{1/2} = -0.35$: the dependence on $m_{1/2}$ is very slight.
b) Contours of m_a in the $(x/v, m_{1/2})$ -plane corresponding to the representative values $\bar{v}/v = 0.6$, $\lambda = 0.15$ and $A/m_{1/2} = -0.35$: m_a increases by ~ 1.5 if $\bar{v}/v$ is decreased to 0.2.

Fig. 1

Fig. 2