The Supersymmetric Top-Ten Lists

HOWARD E. HABER

Santa Cruz Institute for Particle Physics
University of California, Santa Cruz, CA 95064

Abstract

Ten reasons are given why supersymmetry is the leading candidate for physics beyond the Standard Model. Ultimately, the experimental discovery of supersymmetric particles at future colliders will determine whether supersymmetry is relevant for TeV scale physics. The grand hope of supersymmetry enthusiasts is to connect TeV scale supersymmetry with Planck scale physics. The ten most pressing theoretical problems standing in the way of this goal are briefly described.

Invited Talk presented at the Workshop on Recent Advances in the Superworld
Houston Advanced Research Center, April 14–16, 1993.

* Work supported in part by the U.S. Department of Energy.
1. Introduction

The organizers of this workshop asked me to give a talk entitled “Low-Energy Supersymmetry Basics”. I assume that the organizers intended this to serve as an introduction to subsequent talks on supersymmetry in particle physics. In this regard, I thought that it would be useful to present a collection of the main reasons we all have attended this workshop. Namely, why is it that we are all partisans of supersymmetry, and what are the main theoretical problems that keep us interested.

Seven months prior to this workshop, I was attending another workshop, not all that dissimilar to this one, held in Erice. During a sumptuous Italian meal, Gordon Kane made the provocative statement that there are already at least seven strong experimental indications that low-energy supersymmetry is a correct description of nature. A number of us joined in on the discussion, sometimes challenging Gordy’s assertions and sometimes trying to augment or modify his list. Since that evening in Erice, Gordy’s list has increased to nine phenomenological indications for low-energy supersymmetry which now appear in ref. 1.

However, this talk differs from ref. 1 in a number of important respects. First, any list that describes the reasons for pursuing supersymmetry should contain both theoretical arguments as well as phenomenological hints. In particular, the phenomenological hints are often ambiguous or open to a variety of interpretations. That such hints lend their support to a supersymmetric interpretation is strengthened by the theoretical motivations for supersymmetry. Second, I believe that it is important to highlight the main theoretical challenges for supersymmetric theories. The supersymmetric framework is a very ambitious one: it attempts to connect physics at low energies (the TeV scale and below) with the ultimate energy scale of fundamental physics—the Planck scale. A list of the main unsolved problems of the supersymmetric approach will illustrate how far we are from achieving this ultimate goal.

The plan of this talk is as follows. I will provide two “top-ten” lists (hence the use of the plural in the title). In section 2, I will give ten reasons why I (and many supersymmetry enthusiasts) favor supersymmetry. Here, there is substantial overlap with Gordy’s list; although I do not subscribe to all the phenomenological hints of ref. 1. In section 3, I provide ten challenges for theorists who attempt to use supersymmetry to connect low-energy physics and the Planck scale. In some sense, the second list is the more important one. The question of whether low-energy supersymmetry exists (or more specifically, whether supersymmetry is responsible for the scale of electroweak symmetry breaking) is ultimately an experimental one. Yet even if explicit evidence for low-energy supersymmetry is
eventually revealed, the issues addressed in the second top-ten list will remain. Of course, one hopes that additional information such as a detailed supersymmetric particle spectrum might significantly improve our chances to successfully address some of the items on the list in section 3.

An addiction of “top-ten” lists may be indicative of a night person (which is true in this case). However, unlike the lists of David Letterman, I will not be overly dramatic and save the best bits for the end. The order of the points on each list is somewhat arbitrary. I invite the reader to re-order the items on the list, as well as to subtract and/or add various items to them. In the end, I hope that these lists serve as an overview to the status of supersymmetry in particle theories today, as well as providing some impetus as to where we must focus our attention in the future.

2. The First Top-Ten List: Why Believe in Supersymmetry?

Low-energy supersymmetry is the leading candidate for physics beyond the Standard Model. The simplest model of this type is the Minimal Supersymmetric extension of the Standard Model (MSSM), which is reviewed in ref. 2. Of course, there are other candidates as well, including technicolor,\(^3\) composite models,\(^4\) models based on effective four-fermi Lagrangians (\(e.g.,\) top-mode condensate models\(^5\)), and perhaps the true model of nature based on physical principles not yet invented. But, an informal poll, taken with the help of SPIRES, indicates that supersymmetry has attracted the most attention of both theorists and experimentalists. Here are ten reasons why.

1. **Supersymmetry is elegant.**

   Supersymmetry is a symmetry that associates fermionic and bosonic degrees of freedom. It allows one to evade the famous Coleman-Mandula theorem\(^6\) which asserted the impossibility of putting together space-time symmetries and internal symmetries in a non-trivial way. The twentieth century has seen the triumph of gauge symmetries as the underlying structure of all theories of fundamental forces and particles. Supersymmetry is a beautiful generalization of the concept of continuous symmetries; it would be surprising if nature did not make use of it.

2. **Gravity exists.**

   Supersymmetry may be the link between theories of elementary particles and a fundamental theory of gravity. Local supersymmetric theories necessarily contain gravity.\(^7\) Moreover, the only consistent quantum theories that incorporate gravity are superstring theories which possess supersymmetry at some stage in the theory.\(^8\) It is important to note that this argument by itself does not set the energy scale
at which supersymmetry breaks. In particular, it is theoretically conceivable that supersymmetry is relevant only at the Planck scale ($M_P \simeq 10^{19}$ GeV) in which case it would never affect the physics we see at present or future colliders.

3. **The gauge hierarchy problem**\(^{[9]}\)

It is very puzzling how the ratio $m_W^2/M_P^2 \simeq 10^{-34}$ emerges from the fundamental Planck scale theory. Low-energy supersymmetric models have the potential to solve this problem. In such theories, the effective scale of supersymmetry breaking lies below 1 TeV and provides the connection to the scale of electroweak symmetry breaking. The most successful model of this type is the radiative symmetry breaking scenario of minimal supergravity models.\(^{[10]}\) In this picture, the effective theory at the Planck scale is a globally supersymmetric model broken by soft-supersymmetry breaking mass terms of order 1 TeV. The low-energy consequences of such a theory is revealed by using renormalization group equations (RGEs) with Planck scale boundary conditions. One of the Higgs scalar squared-masses, which is positive at the Planck scale, is driven negative due to the effects of the large top-quark Yukawa coupling; this triggers electroweak symmetry breaking at the required scale. Note that because the renormalization group evolution is logarithmic in nature, an exponential hierarchy between the electroweak scale and the Planck scale can develop.\(^{[11]}\)

4. **Naturalness**\(^{[12]}\)

Despite the simplicity of the Higgs mechanism in the Standard Model, the existence of fundamental scalars in field theory is problematical. If the electroweak model is embedded in a more fundamental structure characterized by a much larger energy scale (e.g., the Planck scale), the Higgs boson would tend to acquire mass of order the large scale due to radiative corrections. Only by adjusting (i.e., “fine-tuning”) the parameters of the Higgs potential “unnaturally” can one arrange a large hierarchy between the Planck scale and the scale of electroweak symmetry breaking. This requires new physics beyond the Standard Model. The virtue of the supersymmetric solution to the hierarchy problem is that it is natural. That is, the no-renormalization theorems of supersymmetry\(^{[13]}\) guarantee that if the hierarchy is established at tree-level, it is not upset when radiative corrections are included.

5. **Unification of gauge couplings.**

Consider the three SU(3)×SU(2)×U(1) gauge couplings ($g_1$, $g_2$, and $g_3$) evaluated at $m_Z$. Now, run these couplings up to the Planck scale. In a low-energy supersymmetric model, where the supersymmetry breaking scale is characterized by $M_{\text{SUSY}}$, one uses the RGEs of the Standard Model for energies between $m_Z$ and $M_{\text{SUSY}}$ and the RGEs of the supersymmetric model for energies between $M_{\text{SUSY}}$ and $M_P$. Do the three gauge coupling constants meet at a single point (call it
$M_X$? In the Standard Model, the answer is no. In the MSSM, the answer is yes! Unification occurs at $M_X \simeq 10^{16}$ GeV for $M_{SUSY} \sim 1$ TeV. This result has been known for some time,\cite{14,15} although a re-analysis by Amaldi and co-workers a few years ago based on LEP data caused a great stir in the particle physics community.\cite{16} Is this result truly significant?

Before being carried away by all the hype, consider the reactions of the optimist, the pessimist, and the cynic. The optimist says: the unification of coupling constants is the first experimental verification of the low-energy supersymmetric scenario. The pessimist says: the unification of coupling constants only rules out the simplest GUT extensions of the Standard Model. It may imply new physics at any scale between the weak scale and the GUT scale and says nothing about TeV scale physics. The cynic says: in the GUT extension of low-energy supersymmetry, there are three unknown parameters: $g_U$ (the unified coupling constant at the GUT scale), $M_X$ (the GUT scale or unification point) and $M_{SUSY}$. Thus the RGEs for $g_1$, $g_2$ and $g_3$ provide three equations and three unknowns. A unique solution is essentially guaranteed, so the unification of coupling constants is no surprise at all. The optimist clearly overstates the (experimental) case for supersymmetry. On the other hand, the pessimist admits that the unification of couplings implies that the desert hypothesis of no new physics between the electroweak scale and the GUT scale is incorrect. New physics must enter somewhere between $m_Z$ and $M_X$. This is an exciting result! Clearly, low-energy supersymmetry is one possible model for such new physics. Although there is no guarantee that the new physics is associated with the TeV scale, the arguments based on the hierarchy and naturalness problems of the Standard Model strongly suggest that new TeV scale physics must exist. The simplest possible scenario would be one in which this TeV scale physics also accounts for the unification of couplings. Finally, the cynic’s remarks that the unification of couplings is guaranteed is technically true (if we ignore the effects of supersymmetric thresholds). However, in solving the RGEs for $M_{SUSY}$ and $M_X$, there was no guarantee that the coupling constant unification that emerges would be consistent with sensible values for these parameters. The fact that such values correspond precisely to the expected range of a successful grand unified extension of low-energy supersymmetry may be more than coincidental and should not be simply dismissed.

Of course, there are numerous complications to the conclusion that coupling constant unification is a hint for low-energy supersymmetry.\cite{17} One must consider the effects of thresholds, both at the low-energy scale (e.g., the various MSSM particle masses) and at the high-energy scale (e.g., the superheavy grand unified particle masses). Non-renormalizable operators induced at the Planck scale can also affect the unification of couplings.\cite{18} Nevertheless, I believe that the message
is clear. The unification of couplings is a strong hint for grand unification at scales near $M_P$. The failure of coupling constant unification in the Standard Model means that there is no desert between $m_Z$ and $M_X$—new physics at some intermediate scale must exist. The fact that coupling constant unification does occur in the MSSM presents an intriguing clue that the physics of the desert may have been identified!

6. Proton decay has not yet been observed.\textsuperscript{[19]}

If one accepts the grand unification scenario just discussed, then one must consider carefully the predictions of proton decay. It is interesting to note that non-supersymmetric grand unified models tend to predict proton decay rates that are incompatible with current experimental bounds (primarily because $M_X$ tends to lie below $10^{15}$ GeV). In contrast, in supersymmetric grand unified models, $M_X$ turns out to be significantly larger, and the conventional proton decay modes are unobservable. One must still check that other decay modes that are induced by new (dimension-five) operators particular to supersymmetric models are consistent with present experimental bounds. This imposes interesting constraints in some cases but does not rule out supersymmetric grand unified models.\textsuperscript{[20]}

7. Relations between third generation quark and lepton masses.

If one uses Standard Model RGEs and assumes no new physics between the electroweak scale and $M_X$, then the prediction of $m_b = m_\tau$ at the unification scale is not compatible with low-energy data.\textsuperscript{[21]} However, the relation $m_b = m_\tau$ at $M_X$ is still viable in supersymmetric grand unified models.\textsuperscript{[22]} Relations among other quark and lepton masses require a more complicated structure at the grand unification scale. Supersymmetric models can accommodate such a structure, although it is not clear whether this constitutes a real hint for low-energy supersymmetry.\textsuperscript{[23]}

The result $m_b = m_\tau$ at $M_X$ may be a less compelling clue than the unification of gauge coupling constants. For example, in some string models, unification of gauge couplings can occur without grand unification, whereas the Yukawa couplings depend in part on the structure of the compactification manifold and do not necessarily satisfy standard unification relations.\textsuperscript{[24,25]}

8. The existence of cold dark matter.

Most theoretical cosmologists believe that the ratio of the matter density in the universe to the critical density, $\Omega \equiv \rho/\rho_c = 1$. This result follows from theories of inflation, and there are some observational hints that also support this conclusion.\textsuperscript{[26]} But, the baryonic matter density cannot contribute more than $0.2\rho_c$, which strongly suggests the existence of dark matter making up a significant portion of the total matter density of the universe. The precise nature of the dark matter has
been much debated in the astrophysical community. Evidence based on theories of galaxy formation and the fluctuations of the microwave background radiation suggest that a substantial fraction of the dark matter is likely to be “cold”. The lightest supersymmetric particle (LSP) is an ideal candidate for cold dark matter. Ranges of MSSM parameter space exist where the primordial abundance of the LSP provides exactly the right amount of “missing mass” to reach the critical closure density.\(^\text{[27]}\) See ref. 28 for further details.

9. **Precision electroweak measurements at LEP show no deviation from the Standard Model.**

Suppose that the origin of the electroweak scale lies with new physics beyond the Standard Model. By the naturalness arguments of section 2.4, the energy scale at which this new physics enters must not lie much above 1 TeV. One must then check that the effects of virtual new heavy particle exchange is not in conflict with the precision electroweak measurements at LEP, which at present show no evidence of departures from the Standard Model. One might be tempted to conclude that the effects of new heavy physics should decouple from LEP observables. However, this is not always true, since violation of decoupling can occur in spontaneously broken gauge theories. In many cases, the effects of new heavy particles on electroweak radiative corrections can be neatly summarized by three different combinations of vector boson self-energies.\(^\text{[29,30]}\) (These are called oblique radiative corrections.) \(T\) is proportional to the shift in the \(\rho\)-parameter and \(S\) counts the number of very massive degenerate chiral weak multiplets. A third parameter, \(U\), also enters although it is typically smaller than \(S\) and \(T\).

Recent LEP measurements show that the contributions to \(S\), \(T\) and \(U\) from new physics beyond the Standard Model must be less than 1.\(^\text{[31]}\) This is not a trivial constraint. For example, it has been shown that \(S\) can be reliably estimated in a class of technicolor models.\(^\text{[29,32]}\) New heavy technifermion doublets do not decouple from \(S\), so precision electroweak measurements can potentially rule out such models. In contrast, supersymmetric models have the property that their contributions to \(S\), \(T\) and \(U\) precisely decouple in the limit of large supersymmetry breaking scale, \(M_{\text{SUSY}}\). Still, one must check the coefficient of the leading terms to determine the numerical importance of the supersymmetric contributions.\(^\text{[33,34]}\) I have computed the contributions of the various supersymmetric sectors to \(S\), \(T\) and \(U\) as a function of the MSSM parameters.\(^\text{[34]}\) An example of these results is shown in fig. 1. I find that once supersymmetric particle masses all become larger than about 150 GeV, the effects of supersymmetry on the oblique radiative corrections become negligible. Thus, the non-observation of deviations from the Standard Model at LEP is easily compatible with supersymmetric extensions of the Standard Model, in contrast to other excursions beyond the Standard Model.
mentioned above.

To be complete, it is important to note that some non-oblique radiative corrections (e.g., vertex corrections) can arise in supersymmetric models that are observable. Perhaps the most interesting example of this type is $b \rightarrow s \gamma$ which is induced in the Standard Model at one-loop.\cite{30} In the MSSM, new contributions enter which could alter the Standard Model prediction. Recent bounds from CLEO\cite{36} can already place interesting limits on the MSSM parameter space. Nevertheless, in the limit of large supersymmetric masses, these contributions vanish as well. Thus, it is likely that supersymmetric particles, if they exist, will be discovered by direct production at future colliders before their virtual effects are uncovered.

10. *LEP has not discovered the Higgs boson.*\cite{37}

The tenth reason is admittedly given with a little tongue in cheek. (I am sure that technicolor enthusiasts would place this point prominently on their top-ten list.) Nevertheless, low-energy supersymmetry leads to important constraints on the Higgs sector. The experimental discovery or absence of the Higgs boson in future experiments will have a significant impact on the validity of low-energy supersymmetry. Consider first the MSSM. The Higgs sector of the MSSM is a constrained two-Higgs-doublet model, in which all quartic Higgs self-couplings are given in terms of the gauge couplings.\cite{38} This means that at least one physical Higgs boson of the model cannot be arbitrarily heavy. It is easy to show that the lightest CP-even Higgs scalar satisfies the following tree-level bound: $0 \leq m_{h^0} \leq m_Z |\cos 2\beta|$, where $\tan \beta$ is the ratio of Higgs vacuum expectation values. In light of this result, it seems that the LEP non-discovery of the Higgs is somewhat of an embarrassment for the MSSM.

This perception changed dramatically a few years ago when it was realized that both the lower and upper bounds on the lightest Higgs mass in the MSSM are significantly affected by radiative corrections. This is mainly a result of the incomplete cancellation between top quark and top-squark loop corrections to the Higgs two-point function. For example, the most significant one-loop radiative correction to the neutral CP-even Higgs masses grows as $m_t^4 \ln(M_\tilde{t}^2/m_t^2)$. Since these effects were first uncovered independently by three groups,\cite{39} there have been many papers in the literature examining the impact of the radiative corrections on the MSSM Higgs sector.\cite{40–50} Here, I shall quote only two results. First, if $\tan \beta = 1$, then the tree-level prediction for the light CP-even Higgs mass is $m_{h^0} = 0$. In this case, the physical mass of the $h^0$ is entirely due to radiative corrections. Marco Diaz and I have performed an exact one-loop calculation of the neutral Higgs masses for values of $\tan \beta$ near 1.\cite{49} We confirmed that for $\tan \beta$ near 1, large radiative corrections to the light Higgs mass can easily push $m_{h^0}$ to values
beyond the current LEP experimental lower bound of 60 GeV. Thus, the possibility that the light Higgs mass is due entirely to radiative corrections is not ruled out! Second, in fig. 2, I show the predicted value of $m_{h^0}$ vs. $\tan \beta$ for various values of the CP-odd Higgs mass, $m_{A^0}$. Two graphs are shown corresponding to $m_t = 150$ GeV and 200 GeV, respectively. For values of $m_{A^0} > 150$ GeV (assuming characteristic supersymmetric particle masses of order 1 TeV), one sees that the predicted value for $m_{h^0}$ lies above the current LEP Higgs mass bound.

The fact that a significant region of MSSM parameter space leads to a predicted value of $m_{h^0} > 60$ GeV is a consequence of enhanced radiative corrections driven by a large top-quark mass. To put it another way, given the large value of $m_t$ suggested by LEP precision electroweak experiments, the most probable MSSM parameters would put the light Higgs boson outside the reach of LEP-I.

So far, the above discussion has focused on the MSSM. What about non-minimal models of low-energy supersymmetry? In such models, the main new feature is the possibility of Higgs-self couplings that are not related to the gauge couplings. Top-quark mass enhanced radiative corrections would still tend to raise the tree-level value of the lightest Higgs boson. But, the upper Higgs mass limit is seemingly unconstrained, since it depends on a new unknown parameter. In this case, it is tempting to make use of the observation of gauge coupling constant unification to conclude that no new physics enters between the TeV scale and $M_X$. If this is true, then it seems likely that all couplings of the model remain perturbative below $M_X$. In the case of the Standard Model, the Higgs squared-mass is proportional to the Higgs self-coupling. The renormalization group scaling of this coupling indicates that the Landau pole would be reached below the Planck scale (indicating that new physics must enter) if the Higgs self-coupling at the electroweak scale lies above a certain value. This result translates into the Higgs mass bound: $m_H \lesssim 175$ GeV. In non-minimal low-energy supersymmetric models, new Higgs self-coupling parameters enter which must be bounded in the same way as in the Standard Model. This yields a similar Higgs mass upper bound, which according to ref. 54 is around 150 GeV.

Thus, if no Higgs boson is found below 150 GeV (and assuming that such a Higgs scalar is not unexpectedly difficult to detect by the standard experimental techniques), one would have to be prepared to either give up on low-energy supersymmetry or abandon the concept of the desert between 1 TeV and the Planck scale. Is there any experimental hint that the Higgs mass might be light (i.e., of order $m_Z$ rather than, say, 1 TeV)? Without a good measurement of the top-quark mass, the precision electroweak data from LEP is not accurate enough for one to reach any conclusion on the Higgs mass. Nevertheless, there are intriguing hints from some of the theoretical analyses of LEP data that give a weak preference for
light Higgs mass values. If such an indication is confirmed, it could provide yet another confirmation of the expectations of low-energy supersymmetry. Of course, the discovery of the Higgs boson at LEP-II would give a much larger boost to low-energy supersymmetry. Based on the Higgs mass calculations quoted above, I give LEP-II about a 50-50 chance for a Higgs discovery if they can extend their search up to $m_Z$. This may be the best hint for low-energy supersymmetry prior to turning on the supercolliders.

3. The Second Top-Ten List: Challenges to a Supersymmetric Theory of Particle Physics

The primary motivation for supersymmetry is that it has the potential for providing a consistent, natural embedding of the Standard Model of particle physics in a more fundamental theory whose natural scale is $M_P$. The unification of coupling constants discussed in section 2.5 provides the strongest hint that one can extrapolate from the TeV scale all the way up to energies near $M_P$. However, a fundamental supersymmetric theory of particles remains an elusive goal. Many theorists insist that the fundamental supersymmetric theory can be truly understood only in the context of superstring theory. This is a very ambitious point of view which proposes that superstring theory is the “theory of everything”. For example, in such a framework, one could in principle derive the effective low-energy broken-supergravity model that emerges at the Planck scale. If this is your point of view, then perhaps you should replace the following list with the top-ten list of the outstanding problems in string theory and string model building. Such a list would contain questions such as: (i) what is the correct string vacuum? (ii) how does one compute the effective Planck-scale broken-supergravity model parameters from string theory? (iii) etc. These questions lie beyond the scope of this talk (and my expertise), and I refer you to the string talks of this workshop. Nevertheless, it is certainly worthwhile to contemplate the solutions to the questions posed in the list below in the context of string theory.

Here are ten theoretical problems that must be overcome on the way to constructing a successful supersymmetric theory of particle physics from the TeV scale to the Planck scale.

1. The origin of supersymmetry breaking.

The origin of supersymmetry breaking is one of the most pressing theoretical problem in fundamental theories of supersymmetry. Here, I shall only briefly outline the most common scenario for producing low-energy supersymmetry from a more fundamental broken supergravity model. This scenario has been called the
hidden sector scenario.\textsuperscript{[56]} In this scenario, one posits two sectors of fields. One sector (called the “visible” sector) contains all the fields of the Standard Model (and perhaps additional heavy fields in a grand unified model of the strong and electroweak forces). A second “hidden” sector contains fields which lead to the breaking of supersymmetry at some large scale $\Lambda_{\text{SUSY}}$. One assumes that none of the fields in the hidden sector carry quantum numbers of the visible sector. Thus, the two sectors are nearly decoupled; they communicate only by weak gravitational interactions. Thus, the visible sector only finds out about supersymmetry breaking through its very weak gravitational couplings to the hidden sector. In the visible sector, the effective scale of supersymmetry breaking (denoted by $M_{\text{SUSY}}$) is therefore much smaller than $\Lambda_{\text{SUSY}}$. A typical result is

$$M_{\text{SUSY}} \simeq \frac{\Lambda_{\text{SUSY}}}{M_P^{n-1}},$$ (3.1)

depending on the mechanism for supersymmetry breaking in the hidden sector. Two popular models for the breaking mechanism are the Polonyi model (where $n = 2$) based on $F$-type breaking in the hidden sector, and gaugino condensate models\textsuperscript{[57]} (where $n = 3$). In both cases, $\Lambda_{\text{SUSY}}$ can be quite large, above $10^{10}$ GeV, while still producing $M_{\text{SUSY}}$ of order 1 TeV or less.

In these scenarios, supersymmetry has the potential for solving the hierarchy and naturalness problems described sections in 2.3 and 2.4. However, at this point, we have only scenarios rather than realistic models. The gaugino condensate model is indicative of the difficulty in constructing a realistic and viable fundamental model of supersymmetry breaking. It suggests that the origin of supersymmetry breaking is probably nonperturbative. As a result, reliable calculations are difficult. This situation is somewhat reminiscent of the status of technicolor approaches to electroweak symmetry breaking; \textit{i.e.}, a number of scenarios have been advanced, but no standard model of technicolor exists. In this regard, I would like to make a plea to the advocates of technicolor and related strong interaction approaches to electroweak symmetry breaking. Lend us your skills of nonperturbative analysis and help us to unravel the secrets of the fundamental origin of supersymmetry breaking!

In string models, supersymmetry breaking should emerge as a consequence of the dynamics of the model. If one were able to successfully solve the string theory and determine the correct vacuum, one would in principle have the tools for determining the “low-energy” effective broken supergravity model at the Planck scale. The soft-supersymmetry breaking parameters would then be computable, and would serve as boundary conditions for renormalization group evolution down to the electroweak scale. Recently, there have been some attempts to explore
model-independent features of the soft-supersymmetry breaking terms that emerge from string theory. Kaplunovsky will summarize some of these results later in this workshop.

2. The cosmological constant problem.

Even if one is successful in making use of supersymmetry to solve the gauge hierarchy problem with no fine-tuning, there is one unsolved fine-tuning problem which remains, called the cosmological constant problem. Once we take gravity into account in particle theory, the vacuum energy density \( \Lambda_0 \) is a physical quantity which in principle is calculable in a fundamental theory of gravity. Theoretically, the vacuum energy density is naively expected to be of order \( M_P^4 \), but this is not our universe. A universe with such a large vacuum energy density would have a lifetime of order the Planck time, \( \hbar/M_P c^2 \simeq 10^{-43} \) sec! Thus, based on the fact that the universe has endured over 10 billion years (and looks very flat at large scales), \( \Lambda_0/M_P^4 < 10^{-121} \). This is the mother of all fine-tuning and naturalness problems!

The extent of this fine-tuning problem is slightly alleviated in broken supersymmetric models. But at best, \( M_P^4 \) is replaced by \( M_{\text{SUSY}}^4 \). Most theorists simply put this question aside, perhaps to be solved at an undetermined future time. However, it is not clear that this is justified. One could ask whether one should accept the theoretical motivation of low-energy supersymmetry to solve the gauge hierarchy and fine-tuning problems while ignoring the most severe fine-tuning problem of them all. Perhaps a solution of the cosmological constant problem will automatically solve all other fine-tuning problems by some presently unknown theoretical mechanism, thus rendering supersymmetry unnecessary. Although it is difficult to argue against such a proposition, the fact that supersymmetry seems to significantly reduce the severity of the cosmological constant problem (as indicated above) may be a hint that supersymmetry will play a key role in its eventual solution.

In the absence of a solution to the cosmological constant problem, one must simply be prepared to accept for now the required fine-tuning in models that incorporates both particle physics and gravity. In models of spontaneously-broken supergravity, there is some freedom that allows a fine-tuning of parameters to set the cosmological constant to zero. All model builders must do this in order to have a theoretically consistent framework. In some string models, the cosmological constant cannot be adjusted by hand, rather it is fixed by the theory. For example, in the models of ref. 61, the cosmological constant is of order \( M_{\text{SUSY}}^4 \). While this is certainly an improvement over the natural value of \( M_P^4 \), it is difficult to ascertain whether one can make sense of such models by pretending that the cosmological constant is zero.
3. The origin of the unification scale ($M_X$).

In section 2.5, we saw that the unification of gauge coupling constants takes place at $M_X \approx 10^{16}$ GeV. This implies that $M_X \neq M_P$. How is the scale $M_X$ generated? Perhaps threshold effects of super-heavy particles are sufficient to push $M_X$ toward $M_P$ such that the two scales are not distinct. In some superstring theories, the unification of coupling constants is predicted to occur at $M_P$, so the fact that $M_X \neq M_P$ is somewhat problematical. One possible way around this problem is to add extra multiplets to the theory to delay the unification to $M_P$; some examples can be found in ref. 62.

4. The gauge hierarchy and tree-level fine-tuning problem.

One of the theoretical motivations of supersymmetry is to solve the gauge hierarchy and naturalness problems, as discussed in sections 2.3 and 2.4. In supersymmetric grand unified models, the ratio $m_W^2 / M_X^2 \approx 10^{-28}$ is stable under radiative corrections. This follows from the supersymmetric no-renormalization theorems, which imply that the parameters of the superpotential (where the above ratio is set) are not renormalized. Still, one can ask: where does such a small number arise in the first place? For example, in supersymmetric SU(5), this small number must be inserted into the theory (i.e., in the superpotential) by hand at tree-level. Specifically, one must fine-tune tree-level parameters of the theory to an accuracy of one part in $10^{28}$, in order that the uncolored doublet Higgs fields remain light (of order $m_W$) while the color triplet Higgs fields are superheavy (with masses of order $M_X$). This is the famous doublet-triplet mass splitting problem which is shared by most grand unified theories. In ordinary SU(5), this hierarchy is unstable under radiative corrections. In supersymmetric SU(5), one can “set it and forget it”, but this is not a desirable attribute of a fundamental theory. Possible solutions to the tree-level fine-tuning problem do exist. One such example is the missing-partner mechanism; other solutions have also been proposed. However, having eliminated tree-level fine-tuning does not guarantee that the gauge hierarchy problem has been solved. One must check that higher dimensional operators (which are suppressed by inverse powers of the Planck scale) do not re-introduce fine-tuning (which may be less severe than the original fine-tuning, but may still be too large for comfort) in order to maintain the required gauge hierarchy. String theory also provides another mechanism for generating (approximately) massless Higgs doublets, in models where no formal grand unification occurs.

5. The $\mu$-problem.

A low-energy supersymmetric model is specified by its superpotential and collection of soft-supersymmetry breaking terms. The latter are dimension two or three terms with coefficients with units of mass to the appropriate power. The
scale of the soft-supersymmetry-breaking terms is the electroweak scale; the origin of this scale is tied to the mechanism of supersymmetry-breaking, discussed in section 3.1. But what about the terms in the superpotential with units of mass? For example, the superpotential of the MSSM contains the term $\mu \hat{H}_1 \hat{H}_2$ (where the $\hat{H}_i$ are the two Higgs superfields). The parameter $\mu$ has dimensions of mass, which must be no larger than about 1 TeV in order to preserve the naturalness of the electroweak theory. What is the origin of the scale $\mu$? There is danger that in the fundamental theory at the Planck scale, $\mu$ could be generated with a value of order $M_P$. Perhaps the most natural solution is to demand that only dimensionless parameters in the superpotential can be nonzero. However, this is not acceptable in the case of the MSSM, since by setting $\mu = 0$, the theory would have a Peccei-Quinn symmetry, leading to a weak scale axion which is experimentally untenable. Other solutions to the $\mu$-problem have been proposed. One solution is to add a singlet superfield $\hat{N}$ to the theory and eliminate the $\mu$-term in the superpotential in favor of the term $\hat{H}_1 \hat{H}_2 \hat{N}$. Then, $\mu$ would be generated when $\hat{N}$ acquires a vacuum expectation value. However, singlet superfields are dangerous in that they can destroy the gauge hierarchy. A more natural mechanism is one in which $\mu = 0$ initially, but a non-zero value of $\mu$ of order 1 TeV (or less) is generated when supersymmetry-breaking effects are taken into account.

6. The gravitino problem.

In low-energy supersymmetry, one typically expects the gravitino to possess a mass of order $M_{\text{SUSY}}$. The gravitino may or may not be the lightest supersymmetric particle.† Since its interactions with ordinary matter are gravitational in strength, its lifetime would exceed the lifetime of the universe by many orders of magnitude. Thus, the gravitino is another candidate for dark matter (in addition to the LSP mentioned in section 2.8). This is problematical, since as shown in ref. 68, a gravitino whose mass is of order $M_{\text{SUSY}} \simeq 100 \text{ GeV} - 1 \text{ TeV}$ would lead to a mass density of the universe significantly larger than the critical density. That is, the number of primordial gravitinos is predicted to be too large. One solution to this problem is to suppose that the universe reheats only up to about $10^{10} \text{ GeV}$ after inflation. Inflation dilutes the primordial gravitinos sufficiently and a low reheating temperature would insure that gravitinos are not regenerated in significant numbers. In the early days of supersymmetry model building, such a solution was disfavored, since a successful model of baryogenesis at the GUT scale implied that the baryons were generated after inflation, which required a reheating

---

* In the standard low-energy supergravity approach, if $\mu = 0$, then so is the soft-supersymmetry-breaking term $\mu^2 \equiv B \mu$ which is the coefficient of $H_1 H_2$ in the scalar potential.
† I shall stick to the notation in which the LSP is the lightest supersymmetric particle, excluding the gravitino.
temperature substantially above $10^{10}$ GeV. Recently, there has been much theoretical work which indicates that baryogenesis at the electroweak scale \cite{70} is possible.\footnote{In low-energy supersymmetric models, electroweak baryogenesis is possible in the context of the MSSM.\cite{71}} In this case, inflation and a low reheating temperature can be a viable solution to the gravitino problem. (See also ref. 72 for an alternative suggestion.)
7. Flavor changing neutral current (FCNC) problems.

One of the great successes of the Standard Model is that FCNCs are very suppressed, as required by experimental bounds on FCNC processes. The suppression of FCNCs is a consequence of the GIM-mechanism. On the other hand, the origin of flavor in the Standard Model is a complete mystery. Extended technicolor models attempt to solve both the origin of electroweak symmetry breaking and the origin of flavor with new physics in the energy range between 1 TeV and 1000 TeV. Perhaps it is not surprising that such an ambitious program is generally plagued with FCNCs which violate the strict experimental bounds.

Supersymmetry is often touted as being superior in that there is no FCNC problem. This is only partially correct. To avoid FCNCs, it must be true that the squark mass matrices are approximately diagonal in the same basis that the corresponding quark mass matrices are diagonal. In addition, since the dominant contributions to the squark masses arise from soft-supersymmetry breaking terms, one finds that the squarks must be roughly degenerate in mass.\(^\text{[73]}\) In supergravity model building, a standard assumption is that all soft-supersymmetry-breaking scalar masses at the Planck scale are universal. Of course, flavor information does enter the renormalization group evolution, so that the low-energy squark mass parameters will not be exactly flavor independent. Nevertheless, it is easy to show that the assumption of universal soft scalar masses at the Planck scale is sufficient to keep FCNCs below their experimental upper limits.

Are universal soft scalar mass terms at the Planck scale natural? The answer appears to be model-dependent. Such a result appears automatically in supergravity models with canonical kinetic energy terms, although there is no fundamental reason why a theory of supergravity should only possess the simplest kinetic energy terms. In superstring models (at string tree-level), universal scalar masses appear in models in which the supersymmetry-breaking arises solely from the dilaton $F$-term.\(^\text{[74]}\) More general models of low-energy supergravity generate non-universal soft scalar masses, although the corrections to universality is calculable and in some cases may be sufficiently small.\(^\text{[75]}\) Another possibility, where the required squark degeneracy is obtained by exploiting flavor symmetries at the Planck scale, is explored in refs. 76 and 77.

There are other dangers lurking if one begins to allow for new physics at intermediate scales (between the TeV scale and the Planck scale). As shown in ref. 78,

\(^\star\) Phenomenological requirements impose strong constraints only on the first two generations of squarks. In the $\tilde{q}_L - \tilde{q}_R$ basis, significant off-diagonal mixing in the bottom and top-squarks sector (which splits the corresponding squark masses from the common diagonal soft-supersymmetry-breaking mass) cannot be ruled out at present.
integrating out the effects of physics at an intermediate scale can produce effective non-universal scalar mass terms at that scale. Evolving the parameters of the effective Lagrangian down to the electroweak scale can generate FCNCs larger than the allowed bounds. This is an important constraint on models that attempt to attribute the origin of flavor to an intermediate scale.

8. The flavor puzzle.

As mentioned in the previous item, the Standard Model and the MSSM treat the fermion generations in the same way. Neither provide any insight into the origin of quark and lepton masses and mixing angles. The discussion of section 3.7 suggests that in a fundamental supersymmetric model, the origin of flavor probably lies at the Planck scale. At least two different scenarios are possible. In the first scenario, the physics of flavor is imprinted on the fermion-Higgs Yukawa couplings that arise from the underlying superstring theory. Examples are known in which the Yukawa couplings are computable and depend on topological properties of the compactified space that defines the string vacuum. In the second scenario, quark and lepton mass matrices are generated from the dynamics at the grand unified scale. Examples of this approach have recently appeared in refs. 23 and 80. In this approach, supersymmetry does not play a fundamental role in the generation of the quark and lepton matrices; rather it is required in order to have consistent unification of couplings. One particularly elegant scenario suggests that the three third-generation Yukawa couplings ($h_t$, $h_b$ and $h_\tau$) all unify at some large scale $M_X$. Such models predict that $m_t \sim 180$ GeV and $\tan \beta \sim m_t/m_b$. Large $\tan \beta$ models have a number of interesting phenomenological implications including enhanced radiative corrections to the light Higgs mass (if $m_{A^0} > m_{h^0}$) and enhanced Higgs couplings to the $b$-quark (and $\tau$-lepton).

9. The CP-violation puzzle.

In the Standard Model, CP-violation arises from a complex phase of the CKM-matrix. This complex phase is also a source of CP-violation in the MSSM; in both cases, there is no clue to the fundamental origin of CP-violation, or the relevant energy scale involved. The solution to this problem may be intimately connected to the flavor puzzle discussed above, since the CKM-phase arises after diagonalizing the quark mass matrix. However, supersymmetric theories introduce new complex phases. For example, in the MSSM complex phases can appear in the gaugino Majorana mass terms and the $A$-parameters. If these phases were $O(1)$, one

---

* Flavor could arise from intermediate scale physics. But realistic models of this kind (that satisfy, e.g., the FCNC constraints) are difficult to construct. In addition, the presence of intermediate scales could seriously disrupt the successful unification of gauge coupling constants discussed in section 2.5.
would compute an electric dipole moment for the neutron which is 2 to 3 orders of magnitude larger than the present experimental bounds. Thus, one must conclude that the new supersymmetric phases are no larger that $10^{-3}$--$10^{-2}$. Can such a result arise in a natural way?

In the MSSM, one typically sets the unwanted phases to zero. But in a more fundamental supersymmetric theory, the question of these phases must be addressed. If these phases are zero at the Planck scale, then their magnitudes at the electroweak scale (driven by renormalization group evolution) are certainly small enough to avoid potential phenomenological problems. But, this then shifts the question to the Planck scale. What sets the Planck scale phases to zero?

String theory may provide a hint at a solution to this problem. It turns out that in string theory, the CP-transformation is in fact a gauge transformation in higher-dimensional space-time. Since gauge symmetries cannot be explicitly broken, it follows that at a fundamental level, CP-violation must arise from a spontaneous breaking of CP at some higher energy scale (perhaps the Planck scale). This could provide a theoretical motivation for setting Planck scale phases to zero. Remarkably, in a model of spontaneous CP-breaking at a very high scale, upon integrating out the physics at the high scale, the effective low-energy CP-violation has precisely the form of a single phase in the CKM-matrix. This scenario requires new intermediate scale (or Planck scale) physics associated with the scale of CP-violation. It appears impossible to construct a viable model of spontaneous CP-violation solely in the context of the MSSM.

Although the above scenario sounds compelling, there is a potential problem associated with the strong CP phase. Based on the present limits on the electric dipole moment of the neutron, it is known that $\bar{\theta} < 10^{-9}$. But in models of spontaneously broken CP, $\bar{\theta}$ is calculable. One can think of the mechanism as follows: set $\bar{\theta} = 0$ at the high scale where CP is a good symmetry. Below the scale of CP-breaking, a non-zero (finite) value of $\bar{\theta}$ is generated. If only Standard Model particles remained after integrating out the physics above the CP-breaking scale, one would find an incredibly small value for $\bar{\theta}$ well below the experimental limits. This result follows because one needs to go to a high order in perturbation theory before the first nontrivial correction to $\bar{\theta}$ arises. But, in models with low-energy supersymmetry, contributions to $\bar{\theta}$ may arise at one-loop and yield a value for $\bar{\theta}$ larger than the experimental bound. Thus, it is a challenge to model-builders to construct a viable supersymmetric model of spontaneously broken CP-violation which does not generate a value of $\bar{\theta}$ that is incompatible with the bound on the neutron electric dipole moment.

10. The origin of low-energy discrete symmetries.
One of the great triumphs of the Standard Model is that \( SU(3) \times SU(2) \times U(1) \) gauge invariance is sufficient to eliminate the possibility of baryon number \( (B) \) and lepton number \( (L) \) violating operators of dimension four or less.\(^{[93]}\) This provides a natural explanation why the Standard Model conserves \( B \) and \( L \) to such great accuracy.\(^*\) Unfortunately, this elegant result of the Standard Model is lost in the MSSM.\(^{[92]}\) To see why, recall that the supersymmetric interactions are fixed once one specifies the superpotential. The most general gauge-invariant superpotential of the MSSM has the following form:

\[
W = W_R + W_{NR}.
\]

First, \( W_R \) is given by

\[
W_R = \epsilon_{ij} \left[ h_\tau \hat{H}_1^i \hat{L}_j \hat{E} + h_b \hat{H}_1^i \hat{Q}_j \hat{D} - h_t \hat{H}_2^i \hat{Q}_j \hat{U} - \mu \hat{H}_1^i \hat{H}_2^j \right].
\]

In eq. (3.3), \( \epsilon_{ij} \) is used to combine two \( SU(2) \) doublets [where \( \epsilon_{ij} = -\epsilon_{ji} \) with \( \epsilon_{12} = 1 \)]. The parameters introduced above are the Yukawa coupling matrices \( h_\tau, h_b \) and \( h_t \) (generation labels are suppressed) and the Higgs superfield mass parameter, \( \mu \). Second, \( W_{NR} \) is given by

\[
W_{NR} = \epsilon_{ij} \left[ \lambda_L \hat{L}_i \hat{L}_j \hat{E} + \lambda_L' \hat{L}_i \hat{Q}_j \hat{D} - \mu' \hat{L}_i \hat{H}_2^j \right] + \lambda_B \hat{U} \hat{D} \hat{D},
\]

where generation labels are again suppressed. One quickly observes that the terms in \( W_{NR} \) violate either baryon number \( (B) \) or lepton number \( (L) \). Specifically,

\[
\hat{L} \hat{L} \hat{E}, \hat{L} \hat{Q} \hat{D}, \hat{L} \hat{H}, \Delta L \neq 0, \\
\hat{U} \hat{D} \hat{D}, \Delta B \neq 0.
\]

In the MSSM, one sets \( W_{NR} = 0 \) in order to recover \( B \) and \( L \) symmetry. This can be implemented by introducing a discrete symmetry. There are two equivalent descriptions:

(i) Matter parity\(^{[93]}\)

The MSSM does not distinguish between Higgs and quark/lepton superfields. One can define a discrete matter parity under which all quark/lepton superfields are odd while the Higgs superfields are even.

\(^*\) In fact, \( B \) and \( L \) are not exact in the Standard Model but are violated due to the electroweak anomaly. But the size of such violations is exponentially suppressed and not relevant to the discussion here.
(ii) R-parity$^{[94]}$

In the supersymmetric limit, one can show that the theory possesses a continuous $U(1)_R$ symmetry if $R = 2$ for all terms in the superpotential $W$. It follows that in order to set $W_{NR} = 0$, one may choose

$$R = 1 \quad \text{for } \hat{H}_1, \hat{H}_2,$$

$$R = \frac{1}{2} \quad \text{for } \hat{L}, \hat{E}, \hat{Q}, \hat{U}, \hat{D}. \quad (3.6)$$

The full continuous $U(1)_R$ symmetry does not survive when the soft-supersymmetry-breaking terms are included; the $U(1)_R$ symmetry breaks down to a discrete $Z_2$ symmetry called $R$-parity. It is easy to check that the $R$-parity quantum number is given by

$$R = (-1)^{3(B-L)+2S} \quad (3.7)$$

for particles of spin $S$.‡

If low-energy supersymmetry is correct, will it be $R$-parity invariant? From a purely phenomenological point of view, one cannot rule out the possibility that some of the operators listed in eq. (3.5) are present. One can set bounds on $\lambda_L, \lambda_L', \mu'$ and $\lambda_B$ [see eq. (3.4)], based on $B$ and $L$ violation limits in a variety of Standard Model processes.$^{[99,100]}$ There is one weak argument in favor of a conserved $R$-parity. If $R$-parity is violated, then the LSP is no longer stable and therefore cannot be the dark matter. If one regards the existence of dark matter as one of the selling points for low-energy supersymmetry (see section 2.8), then one must demand that $R$-parity is a good symmetry.

Whether one believes in $R$-parity or not, it is clear that phenomenological requirements prevent the simultaneous appearance of all the operators listed in eq. (3.5). Thus, it seems inevitable that some sort of discrete symmetry will be required to remove the unwanted operators. Any fundamental theory of supersymmetry must address the origin of such discrete symmetries. Examples of superstring models are known in which the discrete symmetries necessary for $R$-parity invariance emerge naturally.$^{[95]}$ This may provide a clue as to the origin of discrete symmetries in low-energy supersymmetry.

---

† My normalization of the $R$-quantum number differs by a factor of two from that of ref. 98.

‡ Interesting alternative supersymmetric models exist in which the $R$-parity symmetry described above is modified. Among such models are $R$-parity-violating models, and models that promote the $Z_2$ $R$-parity of the MSSM to a larger discrete symmetry group$^{[95]}$ or even to the full continuous $U(1)_R$ symmetry.$^{[96]}$ In the latter case, one must introduce new color octet fermions to mix with the gluinos,$^{[97]}$ in which case $U(1)_R$-symmetric massive color-octet Majorana fermions are permitted. Such models represent interesting alternatives to the MSSM.
4. Conclusions

With the discovery of the $W$ and $Z$ gauge bosons in the early 1980s, particle physics entered the era of electroweak symmetry breaking. Nevertheless, the fundamental origin of electroweak symmetry breaking is still unknown. Theoretical arguments have made a convincing case that the dynamics underlying electroweak symmetry breaking must be associated with physics at the TeV scale. Low-energy supersymmetry is a leading candidate for this dynamics, for the reasons presented in section 2. However, these ten reasons can never carry the weight of one reason—the discovery of supersymmetric particles at some future collider. It is in this regard that the next generation of supercolliders—the LHC and SSC—are indispensable. These machines are designed to have the capability of determining whether low-energy supersymmetry or some other dynamics is responsible for the masses of the gauge bosons. Without the supercolliders, progress at the forefront of particle physics will stop cold.

In the meantime, the top-ten list of section 3 provides a useful menu of theoretical problems that supersymmetric theorists must address. Perhaps some progress will occur as we await the deliberations of the politicians. But, I suspect that major theoretical breakthroughs will elude us until we have more experimental input. It will be instructive to see how the top-ten lists presented here are viewed ten and twenty years from now. If supersymmetry survives, its main promise will be that it provides a window to the Planck scale. May we be so fortunate!

Acknowledgements

The invitation and hospitality of Jorge Lopez and his colleagues is much appreciated. I would like to thank Gordy Kane whose advocacy of the phenomenological hints for supersymmetry inspired this talk. In addition, I gratefully acknowledge conversations with Michael Dine, Robert Leigh and Pierre Ramond. Finally, I thank the hospitality of the Aspen Center for Physics which provided me with another opportunity to give this talk, and where the written version of this work was completed. This work was supported in part by the Department of Energy and in part by the Texas National Research Laboratory Commission grant #RGFY93-330.
REFERENCES

1. G.L. Kane, UM-TH-93-10 (1993), presented at the Coral Gables Conference on “Unified Symmetry in the Small and in the Large” (January, 1993), and at the XVIII Rencontre de Moriond, Les Arcs, France (March, 1993).

2. H.E. Haber, SCIPP 92/33 (1993), to appear in The Proceedings of the 1992 Theoretical Advanced Study Institute, Boulder, CO, June 1992.

3. E. Farhi and L. Susskind, Phys. Rep. 74 (1981) 277; R.K. Kaul, Rev. Mod. Phys. 55 (1983) 449.

4. I.A. D’Souza and C.S. Kalman, Preons (World Scientific, Singapore, 1992).

5. For a recent review, see M. Lindner, Int. J. Mod. Phys. A8 (1993) 2167.

6. See, e.g., P. West, Introduction to Supersymmetry and Supergravity (World Scientific, Singapore, 1990).

7. P. van Nieuwenhuizen, Phys. Rep. 68 (1981) 189.

8. M.B. Green, J.H. Schwarz and E. Witten, Superstring Theory (Cambridge University Press, Cambridge, England, 1987).

9. E. Gildener, Phys. Rev. D14 (1976) 1667; S. Weinberg, Phys. Lett. 82B (1979) 387; L. Susskind, Phys. Rep. 104 (1984) 181.

10. H.P. Nilles, Phys. Rep. 110 (1984) 1.

11. For a recent review, see L.E. Ibáñez and G.G. Ross, in Perspectives on Higgs Physics, edited by G.L. Kane (World Scientific, Singapore, 1993) p. 229.

12. S. Weinberg, Phys. Rev. D13 (1976) 974; D19 (1979) 1277; L. Susskind, Phys. Rev. D20 (1979) 2619; G. ’t Hooft, in Recent Developments in Gauge Theories, Proceedings of the NATO Advanced Summer Institute, Cargese, 1979, edited by G. ’t Hooft et al. (Plenum, New York, 1980) p. 135.

13. M.T. Grisaru, W. Siegel and M. Roček, Nucl. Phys. B159 (1979) 429; I. Jack, D.R.T. Jones and P. West, Phys. Lett. B258 (1991) 382.

14. M.B. Einhorn and D.R.T. Jones, Nucl. Phys. B196 (1982) 475; W.J. Marciano and G. Senjanovic, Phys. Rev. D25 (1982) 3092.

15. U. Amaldi et al., Phys. Rev. D36 (1987) 1385.

16. U. Amaldi, W. de Boer and H. Furstenau, Phys. Lett. B260 (1991) 447; U. Amaldi et al., Phys. Lett. B281 (1992) 374.

17. R. Barbieri and L.J. Hall, Phys. Rev. Lett. 68 (1992) 752; L.J. Hall and U. Sarid, Phys. Rev. Lett. 70 (1993) 2673; P. Langacker and N. Polonsky, Phys. Rev. D47 (1993) 4028; A.E. Faraggi, B. Grinstein, and S. Meshkov, Phys. Rev. D47 (1993) 5018.
18. L.J. Hall and U. Sarid, LBL-32905 (1992).
19. P. Langacker, UPR-0539-T (1992), invited talk given at The Benjamin Franklin Symposium in Celebration of the Discovery of the Neutrino, Philadelphia, PA, April 29—May 1, 1992.
20. R. Arnowitt and P. Nath, Phys. Rev. Lett. 69 (1992) 725; Phys. Lett. B287 (1992) 89; NUB-TH-3056/92 (1992); J. Hisano, H. Murayama and T. Yanagida, Tohoku preprint TU-400 (1992).
21. H. Arason et al., Phys. Rev. D46 (1992) 3945.
22. V. Barger, M.S. Berger and P. Ohmann, Phys. Rev. D47 (1993) 1093.
23. S. Dimopoulos, L.J. Hall and S. Raby, Phys. Rev. Lett. 68 (1992) 1984; Phys. Rev. D45 (1992) 4192; D46 (1992) 4793.
24. E. Witten, Nucl. Phys. B258 (1985) 75.
25. See, e.g., G.G. Ross, in The Santa Fe TASI-87, Proceedings of the 1987 Theoretical Advanced Study Institute, Santa Fe, NM, edited by R. Slansky and G. West (World Scientific, Singapore, 1988) p. 628; in Particles and Fields—3, Proceedings of the 3rd Banff Summer Institute on Particles and Fields, Banff, Alberta, August 14—27, 1988, edited by A.N. Kamalan and F.C. Khanna (World Scientific, Singapore, 1989) p. 223.
26. E.W. Kolb and M.S. Turner, The Early Universe (Addison-Wesley Publishing Company, Reading, MA, 1990); T. Padmanabhan, Structure Formation in the Universe (Cambridge University Press, Cambridge, England, 1993).
27. S. Kelley, J.L. Lopez, D.V. Nanopoulos, H. Pois and K. Yuan, Phys. Rev. D47 (1993) 2461; R.G. Roberts and L. Roszkowski, Phys. Lett. B309 (1993) 337.
28. K. Yuan, contribution to these Proceedings.
29. M.E. Peskin and T. Takeuchi, Phys. Rev. Lett. 65 (1990) 964; Phys. Rev. D46 (1992) 381.
30. G. Altarelli and R. Barbieri, Phys. Lett. B253 (1990) 161.
31. P. Langacker, in Electroweak Physics Beyond the Standard Model, Proceedings of the International Workshop on Electroweak Interactions Beyond the Standard Model, Valencia, Spain, October 2—5, 1991, edited by J.F.W. Valle and J. Velasco (World Scientific, Singapore, 1992) p. 75.
32. B. Holdom and J. Terning, Phys. Lett. B247 (1990) 88; M. Golden and L. Randall, Nucl. Phys. B361 (1991) 3.
33. R. Barbieri, M. Frigeni, F. Giuliani, and H.E. Haber, Nucl. Phys. B341, 309 (1990).
34. H.E. Haber, SCIPP 93/06 (1993), to appear in the Proceedings of the 23rd Workshop of the INFN Eloisatron Project, “Properties of Supersymmetric Particles”, Erice, Italy, September 28—October 4, 1992.

35. A. Buras, P. Krawczyk, M.E. Lautenbacher and C. Salazar, Nucl. Phys. B337 (1990) 284; S. Bertolini, F. Borzumati, A. Masiero and G. Ridolfi, Nucl. Phys. B353 (1991) 591; J.L. Hewett, Phys. Rev. Lett. 70 (1993) 1045; V. Barger, M.S. Berger, and R.J.N. Phillips, Phys. Rev. Lett. 70 (1993) 1368; R. Barbieri and G.F. Giudice, Phys. Lett. B309 (1993) 86.

36. R. Ammar et al. [CLEO Collaboration] CLNS-93-1212 (1993).

37. See, e.g., D. Decamp et al. [ALEPH Collaboration], Phys. Rep. 216 (1992) 253.

38. See chapter 4 of J.F. Gunion, H.E. Haber, G.L. Kane and S. Dawson, The Higgs Hunter’s Guide (Addison-Wesley Publishing Company, Reading, MA, 1990).

39. H.E. Haber and R. Hempfling, Phys. Rev. Lett. 66 (1991) 1815; Y. Okada, M. Yamaguchi and T. Yanagida, Prog. Theor. Phys. 85 (1991) 1; J. Ellis, G. Ridolfi and F. Zwirner, Phys. Lett. B257 (1991) 83; B262 (1991) 477.

40. Y. Okada, M. Yamaguchi and T. Yanagida, Phys. Lett. B262 (1991) 54; R. Barbieri, M. Frigeni, and F. Caravaglios, Phys. Lett. B258 (1991) 167; J.R. Espinosa and M. Quiros, Phys. Lett. B267 (1991) 27.

41. R. Barbieri and M. Frigeni, Phys. Lett. B258 (1991) 395; A. Yamada, Phys. Lett. B263 (1991) 233.

42. A. Brignole, J. Ellis, G. Ridolfi and F. Zwirner, Phys. Lett. B271 (1991) 123 [E: B273 (1991) 550].

43. P.H. Chankowski, S. Pokorski and J. Rosiek, Phys. Lett. B274 (1992) 191; B281 (1992) 100; MPI-Ph/92-117 and DFPD 93/TH/13 (1993).

44. A. Brignole, Phys. Lett. B277 (1992) 313.

45. M.A. Diaz and H.E. Haber, Phys. Rev. D45 (1992) 4246.

46. D.M. Pierce, A. Papadopoulos, and S. Johnson, Phys. Rev. Lett. 68 (1992) 3678.

47. K. Sasaki, M. Carena and C.E.M. Wagner, Nucl. Phys. B381 (1992) 66.

48. A. Brignole, Phys. Lett. B281 (1992) 284.

49. M.A. Diaz and H.E. Haber, Phys. Rev. D46 (1992) 3086.

50. H.E. Haber and R. Hempfling, SCIPP 91/33 (1992), Phys. Rev. D49 (1993) in press.
51. D. Buskulic et al. [ALEPH Collaboration], CERN-PPE/93-40 (1993).
52. N. Cabibbo, L. Maiani, G. Parisi, R. Petronzio, Nucl. Phys. B158 (1979) 295.
53. H.E. Haber and M. Sher, Phys. Rev. D35 (1987) 2206; M. Drees, Phys. Rev. D35 (1987) 2910; Int. J. Mod. Phys. A4 (1989) 3635; J. Ellis, J.F. Gunion, H.E. Haber, L. Roszkowski, and F. Zwirner, Phys. Rev. D39 (1989) 844; J.R. Espinosa and M. Quiros, Phys. Lett. B279 (1992) 92; B302 (1993) 51.
54. G.L. Kane, C. Kolda and J.D. Wells, Phys. Rev. Lett. 70 (1993) 268.
55. J. Ellis, G.L. Fogli and E. Lisi, Phys. Lett. B279 (1992) 169; Phys. Lett. B285 (1992) 238; Phys. Lett. B286 (1992) 85; Nucl. Phys. B393 (1993) 3; F. del Aguila, M. Martinez and M. Quiros Nucl. Phys. B381 (1992) 451.
56. See, e.g., L. Hall, J. Lykken and S. Weinberg, Phys. Rev. D27 (1983) 2359; S.K. Soni and H.A. Weldon, Phys. Lett. 126B (1983) 215.
57. H.-P. Nilles, Int. J. Mod. Phys. A5 (1990) 4199.
58. L.E. Ibáñez and D. Lust, Nucl. Phys. B382 (1992) 305; V.S. Kaplunovsky and J. Louis, Phys. Lett. B306 (1993) 269.
59. V. Kaplunovsky, contribution to these Proceedings.
60. S. Weinberg, Rev. Mod. Phys. 61 (1989) 1.
61. B. de Carlos, J.A. Casas and C. Muñoz, Nucl. Phys. B399 (1993) 623.
62. I. Antoniadis, J. Ellis, R. Lacaze and D.V. Nanopoulos, Phys. Lett. B268 (1991) 188; I. Antoniadis, J. Ellis, S. Kelley and D.V. Nanopoulos, Phys. Lett. B272 (1991) 31; S. Kelley, J.L. Lopez and D.V. Nanopoulos, Phys. Lett. B278 (1992) 140.
63. H. Georgi, Phys. Lett. B108 (1982) 283; A. Masiero, D.V. Nanopoulos, K. Tamvakis, and T. Yanagida, Phys. Lett. 115B (1982) 380; B. Grinstein, Nucl. Phys. B206 (1982) 387; S.-C. Chao, Nucl. Phys. B256 (1985) 705.
64. A.A. Anselm and A.A. Johansen, Phys. Lett. B200 (1988) 331; G.R. Dvali, Phys. Lett. B287 (1992) 101.
65. I thank Michael Dine for stressing this point to me.
66. For a recent discussion, see J.A. Casas and C. Muñoz, Phys. Lett. B306 (1993) 288.
67. This point has recently been re-examined in J. Bagger and E. Poppitz, Johns Hopkins preprint (1993).
68. S. Weinberg, Phys. Rev. Lett. 48 (1982) 1303.
69. J. Ellis, A.D. Linde and D.V. Nanopoulos, *Phys. Lett.* **118B** (1982) 59; S. Dimopoulos and S. Raby, *Nucl. Phys.* **B219** (1983) 479; L.M. Krauss, *Nucl. Phys.* **B227** (1983) 556; J. Ellis, J.E. Kim and D.V. Nanopoulos, *Phys. Lett.* **145B** (1984) 181; B.A. Ovrut, *Phys. Lett.* **147B** (1984) 263.

70. For a recent review, see A.G. Cohen, D.B. Kaplan and A.E. Nelson, UCSD-PTH-93-02 (1993).

71. A.G. Cohen and A.E. Nelson, *Phys. Lett.* **B297** (1992) 111.

72. J. Cline and S. Raby, *Phys. Rev.* **D43** (1991) 1781.

73. See, e.g., F. Gabbiani and A. Masiero, *Nucl. Phys.* **B322** (1989) 235.

74. V. Kaplunovsky and J. Louis, ref. 58; R. Barbieri, J. Louis and M. Moretti, CERN-TH.6856/93 (1993).

75. B. de Carlos, J.A. Casas and C. Muñoz, *Phys. Lett.* **B299** (1993) 234.

76. For a recent discussion, see Y. Nir and N. Seiberg, *Phys. Lett.* **B309** (1993) 337.

77. M. Dine, R.G. Leigh and A. Kagan, SLAC-PUB-6147 (1993) and SCIPP 93/04 (1993).

78. L.J. Hall, V.A. Kostelecky and S. Raby, *Nucl. Phys.* **B267** (1986) 415.

79. M. Dine and A.E. Nelson, SCIPP 93/03 (1993).

80. G. Anderson, S. Raby, S. Dimopoulos, L.J. Hall and G. Starkman, LBL-33531 (1993).

81. B. Ananthanarayan, G. Lazarides and Q. Shafi, *Phys. Rev.* **D44** (1991) 1613; *Phys. Lett.* **B300** (1993) 245; L.J. Hall, R. Rattazzi and U. Sarid, LBL-33997 (1993).

82. For example, see M. Dugan, B. Grinstein, and L.J. Hall, *Nucl. Phys.* **B255** (1985) 413.

83. M. Dine, R.G. Leigh and D.A. MacIntire, *Phys. Rev. Lett.* **69** (1992) 2030; K. Choi, D.B. Kaplan and A.E. Nelson, *Nucl. Phys.* **B391** (1993) 515.

84. H.E. Haber and Y. Nir, *Nucl. Phys.* **B335** (1990) 363; L. Bento and G.C. Branco, *Phys. Lett.* **B245** (1990) 599.

85. A. Nelson, *Phys. Lett.* **136B** (1984) 387; S.M. Barr, *Phys. Rev. Lett.* **53** (1984) 329.

86. S.M. Barr and A. Masiero, *Phys. Rev.* **D38** (1988) 366; S.M. Barr and G. Segre, *Phys. Rev.* **D48** (1993) 302; M. Dine, R.G. Leigh and A. Kagan, SLAC-PUB-6090 and SCIPP 93/05 (1993), *Phys. Rev.* **D49** (1993) in press.

87. A. Dannenberg, L. Hall and L. Randall, *Nucl. Phys.* **B271** (1986) 574.
88. A. Pomarol, Phys. Lett. B287 (1992) 331; Phys. Rev. D47 (1992) 273.
89. J. Ellis and M.K. Gaillard, Nucl. Phys. B150 (1979) 141.
90. R. Akhouri, I. Bigi and H.E. Haber, Phys. Lett. 135B (1984) 113; R. Akhouri and I. Bigi, Nucl. Phys. B234 (1984) 459.
91. S. Weinberg, Phys. Rev. Lett. 43 (1979) 1566; F. Wilczek, Phys. Rev. Lett. 43 (1979) 1571.
92. S. Weinberg, Phys. Rev. D26 (1982) 287.
93. S. Dimopoulos, S. Raby and F. Wilczek, Phys. Lett. 112B (1982) 133.
94. P. Fayet, Nucl. Phys. B90 (1975) 104; Phys. Lett. 69B (1977) 489.
95. L.J. Hall, Mod. Phys. Lett. A5 (1990) 467.
96. L.J. Hall and L. Randall, Nucl. Phys. B352 (1991) 289.
97. P. Fayet, Phys. Lett. 78B (1978) 417.
98. J. Wess and J. Bagger, Supersymmetry and Supergravity (Princeton University Press, Princeton, NJ, 1992).
99. S. Dimopoulos and L.J. Hall, Phys. Lett. B207 (1987) 210; S. Dimopoulos, R. Esmailzadeh, L.J. Hall, J.-P. Merlo and G.D. Starkman, Phys. Rev. D41 (1990) 2099.
100. H. Dreiner and G.G. Ross, Nucl. Phys. B365 (1991) 597.
101. B.R. Green, K.H. Kirklin, P.J. Miron and G.G. Ross, Nucl. Phys. B278 (1986) 667; B292 (1987) 606; R. Arnowitt and P. Nath, Phys. Rev. Lett. 62 (1989) 2225; Phys. Rev. D40 (1989) 191; D42 (1990) 2948.
FIGURE CAPTIONS

1) The contribution to the $S$ and $T$ parameters from the neutralino and chargino sector of the MSSM as a function of $\mu$ for $\tan \beta = 2$. The four curves shown correspond to $M = 50, 250, 500$ and $1000$ GeV [with $M_2 \equiv M$ and $M_1 = (5g'^2/3g^2)M$]. In (a), curves in the region of $|\mu| \leq 100$ GeV are not shown, since in this region of parameter space the light chargino mass is less than of order $m_Z$. In (b), $T$ is related to the $\rho$ parameter via $\delta \rho = \alpha \delta T$ which is an experimental observable over the entire mass parameter region. Taken from ref. 34.

2) RGE-improved Higgs mass $m_{h^0}$ as a function of $\tan \beta$ for (a) $m_t = 150$ GeV and (b) $m_t = 200$ GeV. Various curves correspond to $m_{A^0} = 0, 20, 50, 100$ and $300$ GeV as labeled in the figure. All $A$-parameters and $\mu$ are set equal to zero. The light CP-even Higgs mass varies very weakly with $m_{A^0}$ for $m_{A^0} > 300$ GeV. Taken from ref. 50.