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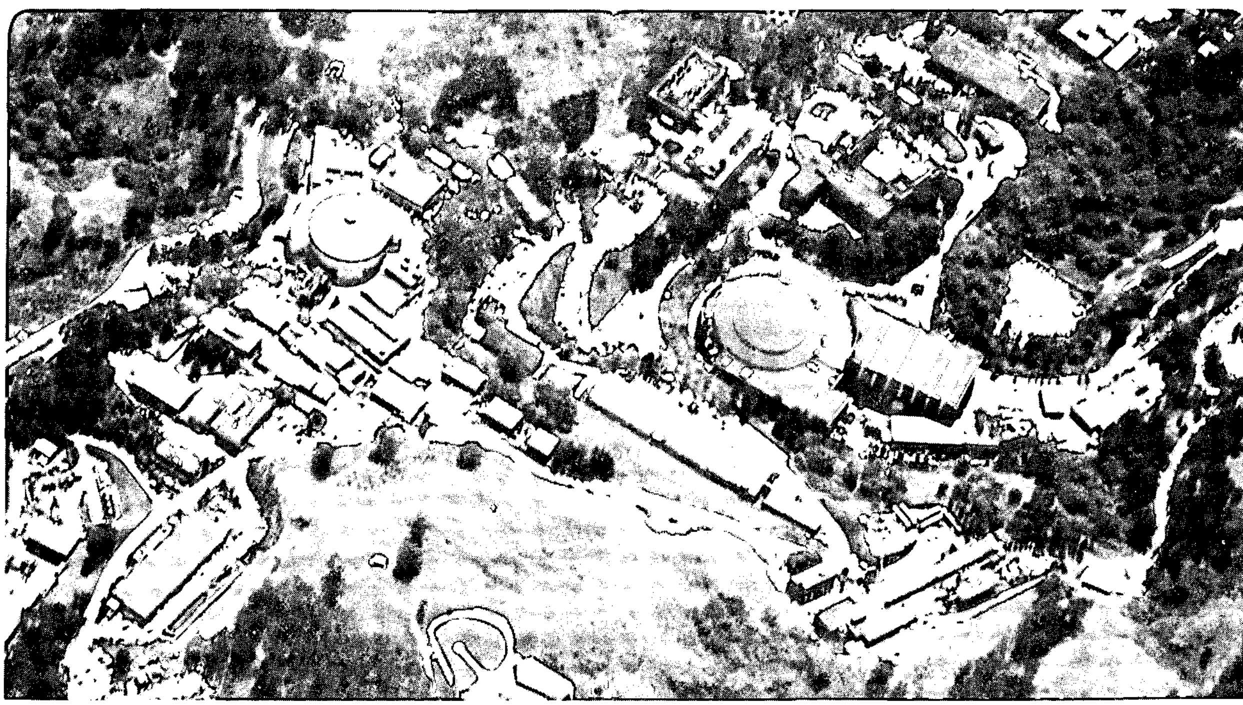

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Supersymmetry*†

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Abstract

I review phenomenologically interesting aspects of supersymmetry. First I point out that the discovery of the positron can be regarded as a historic analogue to the would-be discovery of supersymmetry. Second I review the recent topics on the unification of the gauge coupling constants, m_b–m_τ relation, proton decay, and baryogenesis. I also briefly discuss the recent proposals to solve the problem of flavor changing neutral currents. Finally I argue that the measurements of supersymmetry parameters may probe the physics at the Planck scale.

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1 INTRODUCTION

Low-energy supersymmetry was introduced into particle physics in early 80's mainly because of a theoretical motivation, namely stabilizing

$$m_W^2 \ll M_P^2$$

against radiative corrections [1], where M_P is the Planck mass. This is still the main motivation for the theorists (at least for me) to regard low-energy supersymmetry as a serious candidate of physics beyond the standard model. If this idea is true, we expect the discovery of supersymmetry in near-future collider experiments. This would be really exciting and a historic event in particle physics.

Recently there has been a revival of interest in supersymmetry because LEP measurements on $\sin^2\theta_W$ [2] have shown that the gauge coupling constants extrapolated up to very high energies beautifully meet at a single point using the supersymmetric particle content [3]. More attention is now being paid to phenomenological success of supersymmetry.

We should not believe in supersymmetry in a religious way, even after the LEP measurement on $\sin^2\theta_W$. Only experiments can decide whether the world is supersymmetric. The aim of this talk is not to convince the audience to *believe* in supersymmetry. I am trying to demonstrate how *interesting* supersymmetry is, from the theoretical, phenomenological, and cosmological aspects. In particular, I wish to emphasize another virtue of supersymmetry in this talk, namely, supersymmetry may offer us the unique possibility that the measurement of supersymmetry breaking parameters at collider experiments supplies us with a "window to the Planck world" [4].

2 MOTIVATION FOR SUPERSYMMETRY

It is often stated that the progress of elementary particle physics has revealed an "onion-like" structure of microscopic substances. If history repeats itself, then we may find another substructure of the particles which appear to be elementary, especially in the Higgs sector of the standard model. This understanding of the history leads to speculations like preons or subquarks, or technicolor type scenarios.

However, I wish to remind the audience first that there is another history which is often overlooked. A repetition of that history would lead to a discovery of supersymmetry. It is the discovery of the positron.

At the end of 19th century, there was a problem in electrodynamics that the self-energy of the electron diverges. To see this, let us suppose that the electron is a uniformly charged sphere with radius r_e. Then a simple calculation shows that the electron has a self-energy due to the Coulomb potential generated by itself,

$$E_{self} = \frac{3}{5}\frac{1}{4\pi\epsilon_0}\frac{e^2}{r_e}. \tag{1}$$

The self-energy is linearly divergent in the point-like limit $r_e \to 0$. This contribution can be depicted by a diagram where the electron emits a photon and then re-absorbs it.

The observed mass of the electron is then a sum of the "bare" mass and its self-energy, *i.e.*,

$$m_e c^2 = (m_e)^0 c^2 + E_{self}. \tag{2}$$

As we reduce the "size" of the electron, the smaller we should take its "bare" mass $(m_e)^0$, maybe down to a negative value. It requires increasing fine-tuning to reproduce the observed electron mass. Such a theory cannot be true down to a small distance $r \lesssim \frac{1}{4\pi\epsilon_0}e^2/m_e c^2 \simeq 4$ fm [5].[1] But of course we now know the "size" of the electron is smaller than 10^{-3} fm!

The cure to this problem was supplied by the discovery of the positron. The existence of the positron suggests that there is a fluctuation in the vacuum where an electron positron pair is created and then annihilated. Then another process is possible that an electron "hits" a positron created by the vacuum fluctuation and annihilates it, while the other electron in the vacuum fluctuation remains and pretends it were the original electron coming in. This gives us an intrinsic uncertainty in the position of the electron of the order of the Compton length $r_{Compton} \equiv \hbar/m_e c \simeq 400$ fm. Indeed, the self-energy is cut-off at this scale due to a cancelation between two processes (re-absorption and vacuum fluctuation) down to mild logarithmic divergence [6],

$$E_{self} = \frac{3}{4\pi}\frac{1}{4\pi\epsilon_0}\frac{e^2}{r_{Compton}}\ln\frac{m_e c r_e}{\hbar} \tag{3}$$

[1] Let me quote some sentences by Landau and Lifshitz [5].

"Since the occurrence of the physically meaningless infinite self-energy of the elementary particle is related to the fact that such a particle must be considered as point-like, we can conclude that electrodynamics as a logically closed physical theory presents internal contradictions when we go to sufficiently small distances. We can pose the question as to the order of magnitude of such distances. We can answer this question by noting that for the electromagnetic self-energy of the electron we should obtain a value of the order of the rest energy mc^2. If, on the other hand, we consider an electron as possessing a certain radius R_0, then its self-potential energy would be of order e^2/R_0. From the requirement that these two quantities be of the same order, $e^2/R_0 \sim mc^2$, we find

$$R_0 \sim \frac{e^2}{mc^2}.$$

"This dimension (the "radius" of the electron) determines the limit of applicability of electrodynamics of the electron, and follows already from its fundamental principles. We must, however, keep in mind that actually the limits of applicability of the classical electrodynamics which is presented here lie much higher, because of the occurrence of quantum phenomena."

where the "size" of the electron appears only in the log. Even for a size equal to the Planck length, the self-energy is only about 10 % correction to the "bare" mass.

The cancelation of the linear divergence is a consequence of a new symmetry, (softly-broken) *chiral symmetry*.[2] This symmetry transform electron to positron.[3] The discovery of a new symmetry lead to the cure of the problem of the linear divergence in the electron self-energy.

The motivation for supersymmetry is very similar to the above situation. If we consider the Higgs potential of the standard model and calculate the self-energy of the Higgs field, it turns out to be quadratically divergent. Therefore, the standard model with a naive Higgs potential cannot be a true theory applicable to a length scale much smaller than 10^{-17} cm. Though a fermion mass can be protected by chiral symmetry as we have seen above, a scalar mass cannot be protected by any symmetry of the scalar field alone. The way supersymmetry cures this problem is as follows. First, chiral symmetry protects fermion masses against quadratic and linear divergences. Second, supersymmetry relates the scalar mass to the mass of its superpartner, the fermion. The combination of these leads to only a mild logarithmic divergence in the scalar mass. In diagrammatic language, there is a cancelation between the diagrams of particles and their superpartners.

To be more realistic, supersymmetry should be (softly) broken because we have never observed a superparticle degenerate with the particles which we already know. The effects of breaking can be characterized by the supersymmetry breaking scale, m_{SUSY}. Then the self-energy of the Higgs boson is roughly[4]

$$\delta m_H^2 \sim \frac{\alpha}{4\pi} m_{SUSY}^2 \ln(m_H^2 r^2) \tag{4}$$

where r is the "size" of the Higgs boson. Since the Higgs mass parameter is supposed to be around the m_Z scale, we also expect[5]

$$m_{SUSY} \sim m_Z. \tag{5}$$

All these arguments are highly heuristic. For instance, another possible cure for the problem of the quadratic divergence in the Higgs mass is the replacement of the Higgs boson

[2]The chiral symmetry is exact only in the limit where the electron is massless. Though explicitly broken by non-zero electron mass, the chiral symmetry prohibits the appearance of a huge self-energy, because the breaking parameter m_e is dimensionful and does not change the short-distance behavior of the theory. Explicit breaking of a symmetry which does not change the short-distance behavior of the theory is called "soft breaking."

[3]More precisely, it transforms a positive energy solution of the Dirac equation to a negative energy solution, and the *absence* of an electron in the negative energy state corresponds to a positron.

[4]This is in exact analog with chiral symmetry. Chiral symmetry is softly broken by the non-vanishing electron mass. The self-energy is proportional to the explicit breaking.

[5]Note that the suppression factor by coupling constants is roughly compensated by a large log, if we take r at the Planck length.

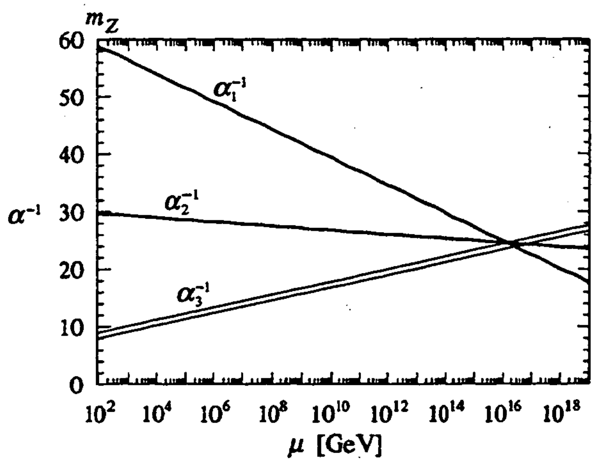

Figure 1: The renormalization group evolution of the gauge coupling constants with $m_{SUSY} = m_Z$.

by a fermion pair bound state. This is the idea of technicolor theories [7]. Though this idea itself is very beautiful and intuitive, it is, unfortunately, hard to implement fermion masses into this scenario; there is a competition between realizing large enough top quark mass and small enough flavor changing neutral current. There is the further logical possibility that the "bare" mass of Higgs and the self-energy cancel to many digits to give a Higgs mass at the 100 GeV scale. But I do not think we can rely on such a theory as a framework to explain physics at its most fundamental level. This is the point where individual taste comes into the discussion. I do not think one can argue that supersymmetry is the only possibility to solve this problem. Supersymmetry is one good candidate which is known to be consistent with the present phenomenology. I will concentrate on more phenomenological aspects of supersymmetry in the next sections.

3 Unification of gauge coupling constants

Three gauge coupling constants are measured precisely at LEP/SLC [2, 8]. If we extrapolate them to very high energies assuming the particle content of minimal supersymmetric version of the standard model, they beautifully meet at a scale $\simeq 2 \times 10^{16}$ GeV, as everybody knows. This observation revived strong interest in supersymmetric grand unified theories. The real excitement of this observation is that the naive choice

$$m_{SUSY} \sim m_Z$$

is consistent with the unification of the gauge coupling constants (see Fig. 1).

After this simple but exciting observation, many people refined the renormalization-group analysis to include threshold corrections both at the GUT- and SUSY-scales [9, 10, 11, 12, 13, 14, 15, 16].

The questions asked in these works are the following.

1. How robust is the success of unification? Can it be destroyed by threshold corrections?
2. Can we constrain m_{SUSY} if we require unification?
3. Is the particle content of the MSSM unique to achieve the unification?
4. What can we learn about GUT-scale physics?

Concerning the first question, it was found that the unification is in general not destroyed even with threshold corrections. The threshold corrections at the SUSY-scale are in general small. A simple parametrization of m_{SUSY} in terms of superparticle masses was found, and in general m_{SUSY} appearing in the renormalization group analysis may be very different from the actual superparticle masses [14]. Also, one has to take care of the difference of definitions in $\overline{MS}$ and $\overline{DR}$ schemes where the latter definition is "more supersymmetric." Even the effects of the $SU(2) \times U(1)$ breaking in superparticle masses have been discussed [16]. In the end, SUSY-scale threshold corrections do not do any harm to the unification of the gauge coupling constants.

Threshold corrections at the GUT-scale can in general be large. For example, while the minimal $SU(5)$ model [17] does not give us large corrections, the missing partner model [18] which includes large $SU(5)$ representations leads to a systematic difference in the α_s prediction [15]. More complicated $SO(10)$ models also may lead to big threshold corrections [19]. These corrections may somewhat weaken the beauty of unification, again they do not destroy the unification.

The second question is a natural question especially in view of near-future collider experiments. Unfortunately, the renormalization-group analysis is only weakly sensitive to m_{SUSY}. The dependence is only logarithmic with a small coefficient, and a naive analysis shows that any m_{SUSY} between m_Z and 10 TeV equally well leads to unification. Furthermore, it was pointed out that the GUT-scale threshold correction can completely destroy any more precise prediction of m_{SUSY} [11]. This is sad, but *c'est la vie.*

The correct particle content is crucial in achieving unification of gauge coupling constants. The MSSM particle content is usually assumed in renormalization group analyses. It is interesting that the minimal supersymmetric extension of the standard model leads to unification, without the addition of arbitrary new fields into the model. On the other hand, the particle content below the GUT-scale is now severely constrained by the unification condition. For instance, the MSSM has two Higgs doublets, while models with four doublets are in contradiction with unification. In general, addition of gauge non-singlet fields destroy unification. There is still room for adding $SU(5)$ complete multiplets to the MSSM particle content, since they change the slope of the renormalization group running of the three gauge coupling constants by the same amount. But again there is a severe constraint from the requirement that the gauge coupling constants do not blow up below the GUT-scale. Introduction of one family and anti-family at the TeV scale leads to the gauge coupling constants blowing up exactly at the GUT-scale [20].[6] Therefore, the particle content one can introduce at the TeV scale is completely classified to be (i) $\mathbf{5}^* + \mathbf{10}$ (fourth generation), (ii) up to three pairs of $\mathbf{5} + \mathbf{5}^*$, and (iii) $\mathbf{10} + \mathbf{10}^*$. Of course, the introduction of singlets is completely harmless as far as the gauge coupling constants are concerned.[7]

An interesting possibility is that there is an enhanced gauge symmetry below the GUT-scale, so that the contribution from new gauge multiplets and new matter multiplets sum up to achieve the correct unification [23, 24]. Though the unification is somewhat accidental in this case, there are interesting aspects to these models. There is a right-handed neutrino at an intermediate scale in accordance with the MSW solution to the solar neutrino deficit as well as τ-neutrino hot dark matter.

The last question is a very ambitious question, and one has to completely specify one particular GUT model to answer it. The simplest example is the minimal $SU(5)$ model [17]. The new particle content at the GUT-scale is an adjoint Higgs $\Sigma(\mathbf{24})$ and Higgs quintets $H(\mathbf{5})$, $\bar{H}(\mathbf{5}^*)$, where the doublet components of H, $\bar{H}$ are contained in the MSSM. Then the mass spectrum at the GUT-scale can be parametrized by three quantities, M_V, the mass of the heavy gauge fields corresponding to broken generators of $SU(5)$, M_Σ, the mass of the adjoint Higgs, and M_{H_C}, the mass of the color-triplet partner of the Higgs doublets in H and $\bar{H}$. What is interesting in this simple model is that the current precision of LEP/SLC data can already constrain the mass spectrum at the GUT-scale, on one combination $(M_V^2 M_\Sigma)^{1/3}$ and M_{H_C} separately [12].

Using LEP data in 1992, one obtains [12]

$$0.95 \times 10^{16}\ \mathrm{GeV} < (M_V^2 M_\Sigma)^{1/3} < 3.3 \times 10^{16}\ \mathrm{GeV} \tag{6}$$

$$2.2 \times 10^{13}\ \mathrm{GeV} < M_{H_C} < 2.3 \times 10^{17}\ \mathrm{GeV} \tag{7}$$

at a 90 % confidence level. SLD data on $\sin^2\theta_W$ prefers a lower value of M_{H_C} with central value down by a factor of 37, while new $\alpha_s(m_Z) = 0.127 \pm 0.005$ from LEP prefers a larger value of M_{H_C}. The constraint on M_{H_C} is crucial to the prediction of the proton decay rate, and it will be interesting to see what values of α_s and $\sin^2\theta_W$ the data converge to.

[6]In this case, $\sin^2\theta_W$ turns out to be at an "infrared fixed point," so that its value at m_Z does not depend on the initial (large) values of gauge coupling constants. Then one obtains the correct $\sin^2\theta_W$ even without grand unification.

[7]It should be noted that the introduction of singlets has a potential danger to destroy the hierarchy [21, 22].

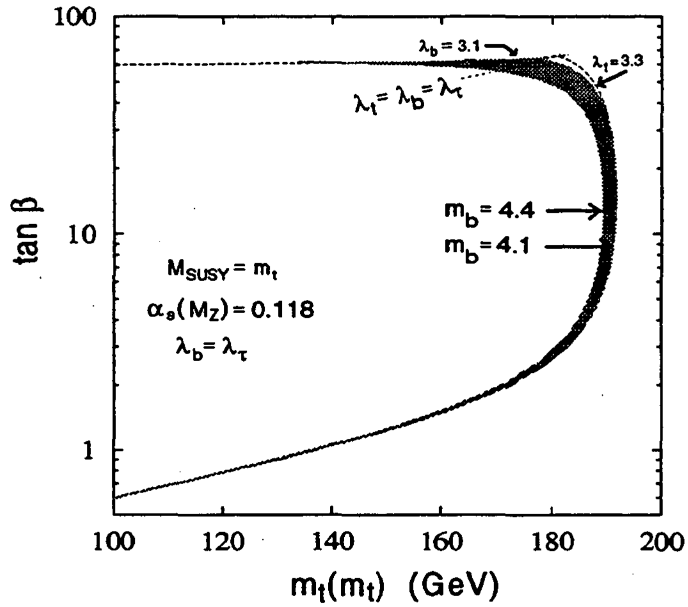

Figure 2: Constant m_b contour in $(m_t, \tan\beta)$ plane assuming $m_b = m_\tau$ at the GUT-scale [27].

4 m_b–m_τ RELATION

The phenomenological success of SUSY-GUT is now not only the gauge coupling constant unification, but also the m_b–m_τ relation. Simple GUT models predict that m_b and m_τ are the same at the GUT-scale. The disparity between their observed masses is supposed to arise from renormalization effects as we scale down from M_{GUT} to their mass shells [25], in a similar manner to the gauge coupling constants.

Recently there have been active discussions of the m_b–m_τ relation in the literature [26]. The first observation is that the successful unification of m_b and m_τ at the GUT-scale requires heavy top, or more precisely, large top quark Yukawa coupling to the Higgs boson, namely $\simeq 1.1$ at the m_Z scale. This does not directly lead to a prediction of the top mass since there are two Higgses in the MSSM, and we should introduce a parameter $\tan\beta$ to relate the Yukawa coupling constant to the physical top mass. An interesting point is that for any given $\tan\beta$ the top mass should be nearly at its maximum possible value allowed by perturbation theory.

Further speculation is that even m_t may be unified with m_b and m_τ at the GUT-scale. This is possible when $\tan\beta \sim 55$, the smallness of m_b/m_t being explained by the smallness of one Higgs expectation value to the other. This led to a prediction of the top quark mass since we know $\tan\beta$ in this assumption; typically $m_t \simeq 180$ GeV [28]. This looks consistent with the top mass "evidence" found by CDF [29]! However, it was found later that one needs to include a finite threshold correction for m_b due to superparticle loops [28], that weakened the predictive power.

To test whether m_b and m_τ really unify at the GUT-scale, we need to know m_t more precisely, and also need to measure $\tan\beta$ by studying the Higgs bosons in the MSSM. Also in the large $\tan\beta$ region, we need to know (at least roughly) the superparticle spectrum.

5 Proton Decay

It is often stated that the non-SUSY $SU(5)$ GUT was excluded by the proton decay experiments. It is true historically, but now it is an empty statement because the gauge coupling constants do not meet in non-SUSY $SU(5)$, so that one cannot predict where the GUT-scale is. Therefore one first has to build models where the coupling constants are unified to make predictions in non-SUSY GUT models [30]. I will not go into this direction in this talk.

There is another claim that SUSY GUT's do not have the same problem with proton decay because the GUT-scale turns out to be much larger than that of non-SUSY theories. This is true for proton decay induced by exchange of heavy gauge bosons, leading to $p \to \pi^0 e^+$ as the dominant mode. However, in most of the SUSY GUT models there is another potentially more dangerous problem, namely proton decay via dimension-five operators [31].[8] Let us restrict ourselves to the minimal SUSY $SU(5)$ model for a while.

Dimension-five operators are caused by the exchange of the color-triplet Higgs, which is the $SU(5)$ partner of the doublet Higgs in the MSSM. Because of supersymmetry, there is also a color-triplet Higgsino which is a fermion. While exchange of heavy bosons induces operators suppressed by their mass squared $\propto 1/M_{GUT}^2$, exchange of heavy fermions induces operators suppressed only linearly in their masses, $\propto 1/M_{GUT}$. Therefore, exchange of the color-triplet Higgsino is potentially a very dangerous mechanism of rapid proton decay.

Fortunately, dimension-five operators have very small Yukawa coupling constants in front, since the color-triplet Higgsino couples to the quark/lepton fields with the same strength as that of the doublet Higgs in the MSSM because of $SU(5)$ symmetry. Furthermore, exchange of a fermion cannot directly induce four-fermi operators which cause proton decay. One needs to "dress" the dimension-five operators with a loop of superparticles to obtain operators which are directly responsible for proton decay. This gives us another small factor of $\alpha_2/4\pi$ or so. In the end, the proton decay rate in the minimal SUSY $SU(5)$ model turns out to be marginally allowed by the experiments, mainly Kamiokande [32]. It is noteworthy that proton decay prefers a light chargino (maybe observable at LEP200?) and heavy squarks (at the margin of LHC reach).

[8] There is yet another problem on proton decay by dimension-four operators, but they can be forbidden by a discrete symmetry called R-parity.

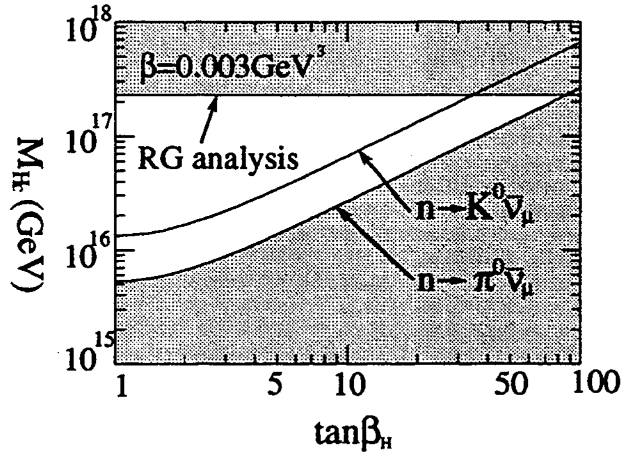

Figure 3: Constraints on M_{H_C} from the renormalization group analysis (from above) and the present nucleon decay experiments (from below) in the minimal $SU(5)$ SUSY-GUT [32]. The shaded regions are excluded at 90 % C.L. The horizontal axis is $\tan\beta$ of the MSSM Higgs sector.

As I explained in section 3, LEP/SLC data are already sensitive to the GUT-scale particle spectrum, and we have an *upper bound* on the mass of the color-triplet Higgs, $M_{H_C} \leq 2.3 \times 10^{17}$ GeV. On the other hand, proton decay (actually, neutron decay can give better bounds in some cases) puts a *lower bound* on M_{H_C}, and it is interesting to see whether there is a remaining region. The present lower bound is (see Fig. 3) [32]

$$M_{H_C} \geq 5.3 \times 10^{15} \text{ GeV}, \tag{8}$$

when taking the most conservative parameter set of hadron matrix element *etc.*

As seen so far, proton decay in the minimal SUSY $SU(5)$ GUT is predictive and can be probably tested by the superKamiokande experiment. However, one may be suspicious of this model because it gives us the wrong predictions $m_\mu = m_s$ and $m_e = m_d$ at the GUT-scale. If we take the current quark masses determined by chiral perturbation theory, more appropriate relations are $m_\mu = 3m_s$, $m_e = m_d/3$ [33]. We have to modify the GUT-model to reproduce the correct fermion mass spectrum. Then the prediction of proton decay is also affected by the same modification, since the dimension-five operators are induced by Yukawa interactions. For example, the relation $m_\mu = 3m_s$ would increase the proton decay rate by a factor of four compared to the minimal model. On the other hand, one can make a complicated Higgs sector to kill dimension-five operators completely [34]. Therefore, an improved lower bound on the proton lifetime imposes stronger constraints on the SUSY-GUT models, but cannot exclude them completely.

It is noteworthy that proton decay puts constraints even on the physics at the Planck scale. If we allow arbitrary baryon-number violating interactions at the Planck scale (as in some of the superstring theories), they cause too-rapid proton decay. These dangerous dimension-five operators can be forbidden using "B-parity" [35], but this discrete symmetry is not consistent with grand unification. One needs to resort on a flavor symmetry to ensure the absence of dangerous operators. This puts strong constraint on flavor physics model building [36]. In this way, an improvement in proton decay experiments may lead to a deeper understanding of flavor physics in the future.

6 Neutralino Dark Matter

One of the virtues of low-energy supersymmetry is that it leads to a natural candidate for cold dark matter particle, the lightest neutralino.

As explained in the previous section, one usually imposes a discrete symmetry, R-parity, in the MSSM. This is a pure assumption, but seems a plausible method to kill the dangerous dimension-four operators which may cause too-rapid proton decay (with a life time of $O(10^{-12})$ sec!). R-parity is a very simple symmetry which assigns even parity to all the fields of the standard model while odd parity to their superpartners. An immediate consequence of this symmetry is that the lightest superparticle (LSP) is stable. Then any LSP produced in the early universe (at a temperature of $T \sim 10^2$ GeV) may possibly remain until the present.

Actually, cold dark matter is needed anyway to explain structure formation in the universe in a way consistent with the fluctuation in the cosmic microwave background measured by COBE [37]. For a typical particle χ with a perturbative annihilation cross section, $\sigma \simeq \pi\alpha^2/m_\chi^2$, its contribution to the present energy density of the universe is $\Omega_\chi \sim 10^{-5}(m_\chi/\mathrm{GeV})^2$. The coincidence between the typical SUSY breaking scale $m_{SUSY} \sim m_Z$ and $\Omega_\chi \sim 1$ from this rough estimation is suggestive.

The precise calculation of the cosmic abundance depends on the details of the superparticle mass spectrum. What is usually done is to assume the "minimal supergravity" boundary condition on the SUSY parameters at the GUT-scale, which introduces five additional parameters to the standard model. In this framework it turns out that the LSP is almost a pure bino, the superpartner of the hypercharge gauge boson [38, 39]. A rough formula for the cosmic abundance is [40]

$$\Omega_{\tilde{B}} \sim \left(\frac{m_{\tilde{B}}}{60\ \mathrm{GeV}}\right)^{-2} \left(\frac{m_{\tilde{l}}}{140\ \mathrm{GeV}}\right)^{4}, \tag{9}$$

and the requirement to avoid a charged dark matter $m_{\tilde{B}} < m_{\tilde{l}}$ leads to an upper bound on superparticle masses, $m_{\tilde{B}}, m_{\tilde{l}} \lesssim 500$ GeV, or $m_{\tilde{g}} < 3$ TeV. This is an interesting constraint on the superparticle masses without resorting on the naturalness arguments, though heavily relying on the assumption of a bino-like LSP.

There are continuing active discussions/experiments relating to the detection of the neutralinos in the halo in our galaxy [41]. Unfortunately, neutralinos have a very weak mainly spin-dependent interaction with nuclei (weaker than the neutrino). We still have only a very weak limit on neutralino dark matter from direct search experiments, while heavy neutrino dark matter is almost excluded [42]. A recent topic is the discussion of constraints from an indirect search at Kamiokande [43]. Neutralinos in the galactic halo may be accumulated in the Sun or Earth, and they may annihilate with each other to produce energetic neutrinos. The neutrinos may then convert to energetic muons before entering the neutrino Čerenkov detectors to leave an "upgoing muon" signature. Here the constraints are very sensitive to the SUSY parameters, especially to the masses of the Higgs bosons.

7 Baryogenesis

Baryogenesis has been regarded as one of the main virtues of grand unified theories [44]. There are many new particles at the GUT-scale whose interaction violates baryon number. When they decay in the very early universe, the decay may generate an asymmetry in the baryon number by CP-violation. Unfortunately, the minimal $SU(5)$ SUSY-GUT does not share this virtue. The predicted baryon asymmetry turns out to be too small because CP-violation in the decay of the color-triplet Higgs appears only at two-loop order. One needs to extend the Higgs sector of the model to begin with [45].

Furthermore, there are general constraints on baryogenesis in the SUSY-GUT models which prefer new physics at an intermediate scale, rather than generating the baryon asymmetry by a decay of a GUT-scale particle. The first constraint is the monopole problem. Monopoles necessarily exist in any grand unified theory based on a simple group [46]. If the universe starts at a very high temperature where the GUT-symmetry is restored, monopoles are produced basically one for each horizon [47]. Usually the annihilation of monopoles is negligible [48]. Then the energy density of the universe is completely dominated by the monopoles, and the universe is very short-lived; it turns back to a "Big Collapse." One can avoid this problem if there is inflation [49], but the temperature after inflation (reheating temperature T_{RH}) cannot be beyond the GUT phase transition temperature so as not to produce monopoles again. Therefore it is not easy to produce GUT-scale particles after inflation. The situation becomes even more severe in supersymmetric theories. In supergravity, there exists a spin 3/2 partner of the graviton, the gravitino, whose mass is supposed to be of the same order as the other superparticles. The lifetime of the gravitino is very long, ~ 10 minutes for $m_{3/2} \sim 1$ TeV. Then it decays *after* the nucleosynthesis producing many high-energy γ's, destroying the light elements. To avoid this, the number density of gravitinos should be small, and this requires the reheating temperature T_{RH} to be smaller than $T_{RH} \lesssim 10^{10}$ GeV [50]. If we take this constraint seriously, GUT-scale particles cannot play any role in generating the baryon asymmetry.

On the other hand, it is now widely accepted that the standard model itself violates baryon (B) and lepton (L) numbers at high temperature due to the anomaly effect [51]. Even if the decay of a GUT-scale particle could generate a baryon asymmetry, it would be wiped out except as far as there is a non-vanishing $B-L$ asymmetry. Therefore it is necessary to generate a $B-L$ asymmetry at some stage in the early universe.[9] However, simple $SU(5)$ GUT models preserve $B-L$ symmetry and hence no asymmetry can be generated.

These problems can be cured just by adding right-handed neutrinos N to the MSSM particle content (MSSM+N) [54]. Then baryon asymmetry is generated "automatically" irrespective of details in the inflationary scenario.

We introduce the right-handed neutrino supermultiplets to the MSSM. The superpotential of this model is

$$W_{\mathrm{MSSM+N}} = W_{\mathrm{MSSM}} + h_{ij} N_i L_j H_u + \frac{1}{2} M_{ij} N_i N_j, \tag{10}$$

and the mass term breaks L and $B-L$ invariance. I wish to remind the audience that the existence of a right-handed neutrino is also preferred in explaining the small neutrino mass in the MSW solution [55] to the solar neutrino problem via the seesaw mechanism [56]. The neutrino mass is characterized by $\Delta m^2 \sim 10^{-5}\ \mathrm{eV}^2$ [57]. If we assume that the MSW effect is due to the ν_e–ν_μ oscillation and take an $SO(10)$-like ansatz for the neutrino Yukawa matrix h_{ij}, one obtains $M \simeq 10^{10}$–10^{13} GeV. This mass range has also a cosmological interest since the τ-neutrino mass turns out to be around $m_{\nu_\tau} \sim 1$–100 eV, and can contribute to the hot dark matter density of the present universe. Therefore, the MSSM+N can naturally lead to a co-existence of the neutralino cold dark matter and ν_τ hot dark matter (mixed dark matter), which is favored now by the observed spectrum of the density fluctuation from COBE to galaxy clusters [58].

The important point is that the MSSM+N "automatically" generates a lepton asymmetry if the mass M of the right-handed neutrino is smaller than the expansion rate during the inflation [54]. The scalar component of the right-handed neutrino supermultiplet is driven to large values at the end of inflation due to the quantum fluctuations in the de Sitter spacetime [59]. It oscillates after inflation, and decays. The decay generates a lepton asymmetry via CP violation in the neutrino Yukawa matrix. Finally electroweak sphaleron effects partially convert the lepton asymmetry to a baryon asymmetry [60]. This scenario depends on the assumption that there is an inflationary period, but does not depend on the details of the inflationary models. Then everything occurs without any further assumptions; hence an "automatic" scenario.

Actually one may be even more ambitious to expect that the scalar component of the right-handed supermultiplet itself can drive a chaotic inflation [61]. Then the simplest exten-

[9] There are two other possibilities discussed in the literature. One is to generate baryon asymmetry at the electroweak phase transition [52]. The other is to employ some mechanisms to keep baryon asymmetry even when $B-L=0$ [53].

sion of the MSSM, namely the MSSM+N can accommodate chaotic inflation, baryogenesis, mixed dark matter, and solve monopole and gravitino problems. The only price one has to pay is to assume that the quadratic potential $V = M^2|\tilde{N}|^2$ persists beyond M_P [62].

8 FCNC Problem

Though I have been presenting various phenomenological virtues of supersymmetry in the previous sections, there is a big embarrassment due to the existence of the superparticles below 1 TeV. The exchange of squarks and gluino may lead to an unacceptably large flavor-changing neutral current, especially in K–$\overline{K}^0$ oscillations. This requires the squark masses to be highly degenerate at least for $\tilde{d}$ and $\tilde{s}$ [63],

$$\frac{m^2_{\tilde{d}} - m^2_{\tilde{s}}}{m^2_{\tilde{d}}} \lesssim 6 \times 10^{-3} \left(\frac{m_{\tilde{d}}}{\text{TeV}}\right). \tag{11}$$

We need an understanding of why squarks are so degenerate in mass.

The standard lore to explain the degeneracy is the following. Suppose supersymmetry is broken in a "hidden sector," which interacts with our "observable sector" of quarks and leptons only through gravitational interactions. Let us denote the scale of supersymmetry breaking in the hidden sector as Λ_{SUSY}. The masses of squarks and sleptons are generated by gravitional interactions with the hidden sector, at the order of Λ^2_{SUSY}/M_P or Λ^3_{SUSY}/M_P^2 depending on the hidden sector models, and turn out to be the same for any scalar particles because gravity is flavor-blind. Though this may sound plausible, the supergravity Lagrangian does not have this feature in general. No symmetry principle restricts the interactions between the hidden and observable sectors to be flavor-blind.

Recently, there have appeared several interesting proposals to ensure the absence of FCNC processes due to superparticle loops. I'll briefly describe the basic ideas in the paragraphs below.

The first one is based on superstring theory (don't be afraid; I myself am an amateur). In four-dimensional superstring models, there is a so-called "dilaton" field which plays a unique role in superstring theory. It has a completely flat potential to any finite order in perturbation theory, but is supposed to have an expectation value due to non-perturbative effects. This expectation value determines the gauge coupling constants dynamically. The point is that the dilaton field has a universal coupling to all fields. If the dilaton plays another role in breaking supersymmetry, the squarks and sleptons acquire the same masses at the Planck scale [64]. The assumption is that the dilaton has two expectation values, one is the usual scalar expectation value which determines the gauge coupling constants and the other is the so-called F-component which breaks supersymmetry. Though this scenario is based on relatively firm theoretical grounds, so far no concrete model of the hidden sector is known which gives rise to an expectation value of the F-component of the dilaton field.

Another possibility that has been pointed out is that the supersymmetry breaking effects are fed to quarks and leptons by the gauge rather than the "gravitational" interactions. Since the gauge interactions are flavor-blind, the generated supersymmetry breaking terms would be also flavor-blind. Though the idea sounds very simple, it actually requires a drastic modification of supersymmetric models. First of all, the gauge interactions are characterized by dimensionless coupling constants, and the supersymmetry breaking masses of the squarks and sleptons are not suppressed by $1/M_P$, rather $m_{SUSY} \sim (\alpha/4\pi)^n \Lambda_{SUSY}$ with n being model-dependent. Therefore this scenario requires many new particles at multi-TeV energies. An explicit realization was worked out recently [65], which has a rather complex structure with a symmetry group $SU(7) \times SU(2) \times SU(3)_L \times SU(3)_R$ in addition to the standard model gauge group. Attractive features are that there is no gravitino problem, supersymmetry is broken dynamically, and the model can be embedded into $SU(5)$ unification. The most unattractive feature is that the vacuum is only a local minimum.

The two proposals which I have described above deal with the physics of supersymmetry breaking to explain the smallness of the flavor changing neutral currents. Other scenarios below take a different attitude with emphasis on the flavor dynamics, the physics of the Yukawa coupling.

A simple solution to the FCNC problem is to assume a symmetry among squarks of different generations to ensure the degeneracy of squarks. This sounds a very natural idea, but it is in apparent contradiction with the non-degeneracy of quarks. First of all, one needs a non-abelian flavor symmetry so that (at least the first two generations of) the quarks lie in an irreducible representation to ensure $m_{\tilde{d}} = m_{\tilde{s}}$. Then it also restricts the form of the Yukawa couplings. Therefore such a scenario should explain both the observed structure of the Yukawa coupling matrix and the degeneracy of squarks at the same time. One such example is based on a non-abelian discrete group $\Delta(75)$ [66]. It is noteworthy that this symmetry also prohibits dangerous dimension-five operators at the $1/M_P$ level [36]. However one needs a relatively complicated Higgs sector to break the flavor symmetry in the desired pattern. No explicit Higgs potential has been presented so far.

A completely different direction is to give up the degeneracy of the squarks, and try to explain the smallness of the flavor-changing neutral current by yet another flavor symmetry. One of the reasons why an exchange of the squarks can lead to large flavor-changing neutral current processes is that the mass eigenstates of quarks and squarks may be completely unrelated in the flavor space. One possible way to suppress the flavor changing neutral currents is to align the mass eigenstates of quarks and squarks [67]. This sounds miraculous, but it was shown that a certain abelian horizontal symmetry can restrict the form of both the Yukawa matrix and the squark mass matrix in such a way that the diagonalization of both matrices can be done with almost the same rotation. An unattractive point is that the assignment of horizontal charges is rather ad hoc, fitted to explain the hierarchical structure of the Yukawa matrix. Once one accepts it, however, the alignment of the quark and squark bases is automatic.

Currently, there is no reason to choose one scenario over the others except from the aesthetic point of view. We will, however, be able to discriminate among them experimentally after we find the superparticles. This will be the topic of the last part of my talk.

9 Future Prospects

First of all, supersymmetry predicts a relatively light Higgs boson. In the MSSM, it has long been known that the lightest Higgs boson should be lighter than the Z^0 [68]. Thus LEP-200 could have made a definitive test of the MSSM. Unfortunately, a large top quark Yukawa coupling induces an important correction to this prediction at the one-loop level. The upper bound is pushed up to around 130 GeV, well beyond the reach of LEP-200 [69]. We should recall that the current LEP bound on the MSSM Higgs is not so tight: $m_h \gtrsim 45$ GeV for standard $b\bar{b}$ mode and $\gtrsim 25$ GeV in the worst case when it decays invisibly into a neutralino pair. There is still a wide parameter space which LEP-200 will explore.

It is very nice that the proposed hadron facility, LHC, will cover most of the parameter space of the MSSM Higgs sector.[10] Recently there has been the interesting suggestion that the Tevatron can search for a MSSM Higgs up to 120 GeV if a luminosity upgrade is performed [72].

What is more encouraging is that the upper bound on the lightest Higgs mass does not change much even if we consider more complicated models, with singlets, fourth generation, *etc.* The mass of the Higgs boson is proportional to (the square root of) the strength of the self-interaction among the Higgs bosons, just like the mass of a fermion is proportional to its Yukawa coupling. The self-coupling tends to become stronger at higher energies due to the renormalization-group running. If the self-coupling is too strong, it can happen that it becomes infinite at some scale below the GUT-scale. Requiring that the theory remains within the validity of the perturbation theory up to the GUT-scale, we can put an upper bound on the strength of the self-coupling, leading to an upper bound on the Higgs mass. At tree-level, models with singlets have an upper bound of 150 GeV [73]. This bound is pushed further up by one-loop effects, but never beyond 180 GeV [74]. The same is true for models with heavy fermions beyond the top quark [75]. An e^+e^- collider of $\sqrt{s} = 300$ GeV can definitely exclude supersymmetric theories based on the GUT idea [76].

Concerning the superparticles, the following characteristic of the mass spectrum is important to future searches. Colored particles tend to be heavy, while colorless particles are light. This is because the form of the renormalization group equations implies that the gauge interactions make the superparticle masses heavier. Colored particles have the strongest gauge

[10] It has been known that there is a region in the parameter space of the MSSM Higgs sector ("ozone hole") which is very hard to be covered at LHC [70]. However, continuous efforts are being devoted to ensure the covering of the whole MSSM parameter space at LHC [71].

interactions and hence become the heaviest among the superparticles. Typical mass ratio of the gluino to the lighter chargino is about a factor of 4. Therefore, supersymmetry is an ideal target of the high-energy experiments that e^+e^- and hadron colliders literally play complementary roles. I refer further discussions to the talk by Michael Peskin [77].

Currently the most stringent bounds on supersymmetry come from LEP and Tevatron experiments. LEP has excluded charginos below 45 GeV in a model-independent way, while CDF excluded gluino masses below 100 GeV.[11] These two bounds are comparable because of the above-mentioned characteristics of the superparticle mass spectrum. A Tevatron with Main Injector would extend the reach up to 300 GeV or so with like-sign dilepton and $\not{E}_T$ when $m_{\tilde{g}} \simeq m_{\tilde{q}}$ [78], while LEP 200 will find charginos below 90 GeV. So far constraints from hadron and e^+e^- colliders will improve more or less in parallel.

Meanwhile, HERA can search for the superparticles produced singly if R-parity is broken. For instance, a squark below 270 GeV may be found if it is produced as an s-channel resonance in eq collision [79].

To go beyond the reach of the present machines, we must wait until LHC begins operation. Depending on parameters and analyses, LHC will extend the reach of the gluino search up to 1.2–1.8 TeV [80, 81]. This would basically cover the whole "natural" region of the superparticle mass spectrum.

In view of these numerous experimental programs in the near future, I strongly expect the discovery of supersymmetry in a few to ten years.

10 Window to the Planck World

Once supersymmetry is discovered, one may worry about losing one's job. Actually it will be exactly the opposite. Since there are so many superparticles which await detailed studies, it will be just the beginning of a whole series of experiments. Indeed, measurement of superparticle masses and mixings will give us information crucial to distinguish between various GUT-models, scenarios of flavor physics and supersymmetry breaking. In this sense, low-energy supersymmetry is a messenger of the physics of the very high-energy scale to the scale which is accessible by experiments.

First of all, I should emphasize that precision measurements of supersymmetry parameters to several percent are possible at an e^+e^- linear collider with a high beam polarization [83]. I refer to the talk by K. Fujii [84] concerning this point on which we worked together. We have demonstrated that the measurements of the superparticle masses and cross sections at a few percent level enable us to extract parameters in the supersymmetric Lagrangian, for example, the mass parameters of $SU(2)_L$ and $U(1)_Y$ gauginos.[12]

[11] This bound depends weakly on an assumption of the MSSM parameters.

[12] It was also pointed out that one can measure the gluino mass at a *pp* supercollider with a precision better than 10 % [82]. One has to assume the low-lying superparticle spectrum to do the analysis. Presumably the

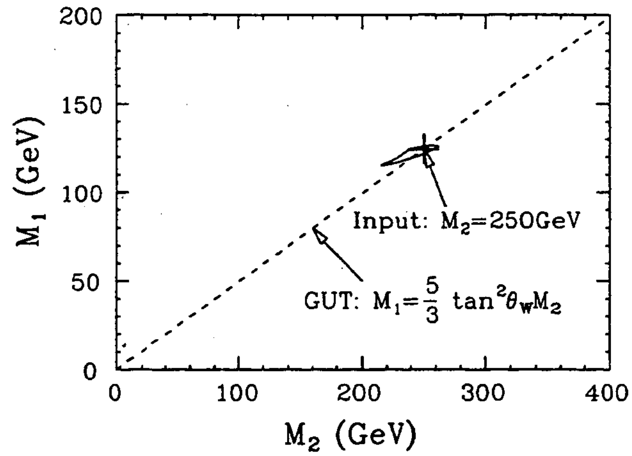

Figure 4: A $\chi^2 = 1$ contour in (M_1, M_2) plane extracted from chargino and selectron pair production processes at an e^+e^- linear collider after detector simulation [83]. One can test whether the gaugino masses unify at the GUT-scale with a 5 % precision.

There is a simple prediction in the SUSY-GUT models that the mass of the gauginos should satisfy the relation

$$\frac{M_1}{\alpha_1} = \frac{M_2}{\alpha_2} = \frac{M_3}{\alpha_3} \tag{12}$$

at any scale. We already know α_1, α_2 and α_3 precisely thanks to LEP/SLC. They turned out to be consistent with SUSY-GUT as we all know. The measurements of M_1 and M_2 would allow us to make another test on SUSY-GUT models at the few percent level, namely a test whether the gaugino masses unify at the same scale where the gauge coupling constants unify (see Fig. 4). It would be even more exciting if the gluino mass measured at a hadron collider also fits into the same relation. If the gaugino masses will be consistent with the SUSY-GUT prediction, it would leave little doubt about unification, at least of some kind.

The GUT-relation of the gaugino masses holds as far as the standard model gauge group is embedded into a simple group, irrespective of the symmetry breaking pattern [24]. For instance, a model based on the chain breaking (Pati-Salam)

$$SO(10) \to SU(4)_{PS} \times SU(2)_L \times SU(2)_R \to SU(3)_C \times SU(2)_L \times U(1)_Y \tag{13}$$

predicts the same relation among the gaugino masses. Therefore, this relation, if confirmed experimentally, would suggest a unification based on a simple group,[13] but does not tell us

combination of data from e^+e^- on low-lying superparticle masses and from hadron supercolliders on missing $\not{E}_T$ and multi-lepton signals would allow us a measurement of the gluino mass at this precision.

[13]It is noteworthy that the superstring with dilaton F-term also leads to the same relation [85]. This is

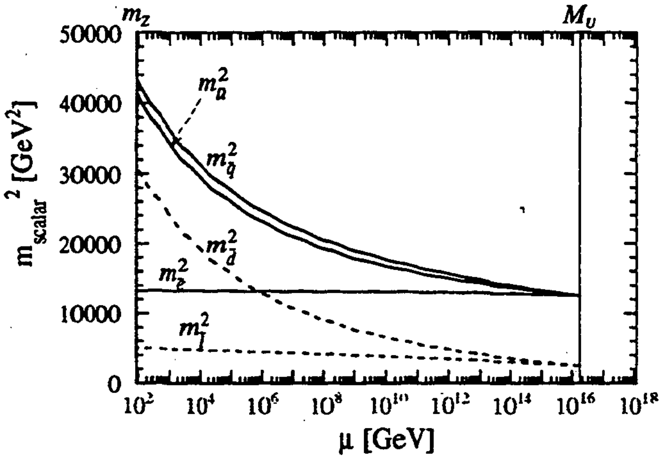

Figure 5: The renormalization group running of the scalar masses in $SU(5)$ SUSY-GUT. They should unify at the same scale where the gauge coupling constants unify [4].

about the existence of an intermediate scale or the symmetry breaking pattern. On the other hand, models not based on a simple group (*e.g.*, $SU(5) \times U(1)$ or superstring models with moduli F-term) does not predict this relation.

Though all the GUT-models give the same relations for the gaugino masses, the scalar mass spectrum distinguishes different models. Let me discuss the "vertical" direction first, the spectrum within the same generation. I will discuss the "horizontal" direction later, comparing the masses of different generations.

The difference of the masses for different quantum numbers reflects the physics of the gauge symmetry. The first and the second generations have tiny Yukawa interactions which can be neglected to a good approximation. Then the splitting of the scalar masses within the same generation arises only by their renormalization-group running due to the gauge interactions. The pattern of the scalar mass spectrum will tell us the pattern of the symmetry breaking. For instance, $SU(5)$ SUSY-GUT predicts the scalar masses of $\tilde{e}_R$, $\tilde{u}_L$, $\tilde{d}_L$ and $\tilde{u}_R$ unify at the GUT-scale in the same manner as the gauge coupling constants unify (see Fig. 5). Similarly for $\tilde{e}_L$, $\tilde{\nu}_L$ and $\tilde{d}_R$. On the other hand, the cases with an intermediate symmetry, such as Pati-Salam, predict certain "sum rules" of the scalar masses [24],

$$m_{\tilde{q}}^2(M_{PS}) - m_{\tilde{l}}^2(M_{PS}) = m_{\tilde{e}}^2(M_{PS}) - m_{\tilde{d}}^2(M_{PS}), \tag{14}$$

$$g_{2R}^2(M_{PS})(m_{\tilde{q}}^2 - m_{\tilde{l}}^2)(M_{PS}) = g_4^2(M_{PS})(m_{\tilde{u}}^2 - m_{\tilde{d}}^2)(M_{PS}). \tag{15}$$

amusing because one needs rather big threshold corrections for the gauge coupling constants to reconcile the difference between the apparent GUT-scale and the string scale [86]. Exactly the same correction appears both for the gauge coupling constants and the gaugino masses to give the same relation as in the (field theoretical) GUT models.

	α_i	M_i/α_i	m_i^2
$SU(5) \to G_{SM}$	natural	common	testable
$SO(10) \to G_{SM}$	natural	common	testable
$SO(10) \to G_{PS} \to G_{SM}$	adjustable	common	testable
$SO(10) \to G_{3221} \to G_{SM}$	adjustable	common	not testable
$SO(10) \to G_{3211} \to G_{SM}$	adjustable	common	not testable
$SU(5) \times U(1) \to G_{SM}$	adjustable	common only for $i = 2, 3$	testable
superstring with dilaton F-term	adjustable	common	testable
superstring with moduli F-term	adjustable	not common	not testable

Table 1: The "score sheet" of how well we can distinguish between various models [22]. The intermediate groups are defined as $G_{PS} = SU(4)_{PS} \times SU(2)_L \times SU(2)_R$, $G_{3221} = SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, $G_{3211} = SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{B-L}$. The row α_i refers to the unification of the gauge coupling constants, where "natural" means that the unification is automatic, while "adjustable" employs either particular particle content or threshold corrections to reproduce observed gauge coupling constants. The row M_i/α_i refers to the gaugino masses. The row m_i^2 states whether the model predicts a definite pattern which is testable using the low-energy scalar mass spectrum.

at the scale M_{PS} where Pati-Salam symmetry is broken. Though one of the relations should be used to determine M_{PS}, the other relation can be used to test the prediction. For models with lower symmetry, predictivity becomes lower, so that they may not be able to be tested. In any case, the scalar mass spectrum in the "vertical" direction reflects the pattern of gauge symmetries at high energies. Table 1 summarizes the "score sheet," of how different observables can distinguish various models.

The "horizontal" direction in the scalar mass spectrum carries information on flavor physics at high scales. For instance, non-abelian flavor symmetry gives us the degeneracy between the scalar masses of different generations, for given gauge quantum numbers. Superstring theories based on the dilaton F-term breaking also give the same degeneracy, but with a further prediction of the ratio of the scalar mass to the gaugino mass. On the other hand, a flavor group which leads to an alignment of the quark and squark bases does not need a degeneracy among the squarks, and we expect a baroque spectrum in the horizontal direction.

In any case, the point is that we will be able to play the same game with the superparticle masses as we play now with the gauge coupling constants. We can make numerous checks whether experimentally independent masses unify at higher energies. This is what I mean by "window to the Planck world" [4]. This is a truly ambitious program, but would become possible if supersymmetry were true.

11 Conclusion

I reviewed interesting aspects of supersymmetry with emphasis on recent topics. There are many more virtues and problems in supersymmetric models which were not covered in this talk.[14] Though supersymmetry is certainly a very interesting candidate of physics beyond the standard model, we theorists do not know whether the nature is supersymmetric; only experiments can decide it. And once supersymmetry is found, we will gain many hints on the physics at ultra-high energy scales.

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