Probing Physics at Short Distances
with Supersymmetry

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Abstract

We discuss the prospect of studying physics at short distances, such as Planck length or GUT scale, using supersymmetry as a probe. Supersymmetry breaking parameters contain information on all physics below the scale where they are induced. We will gain insights into grand unification (or in some cases string theory) and its symmetry breaking pattern combining measurements of gauge coupling constants, gaugino masses and scalar masses. Once the superparticle masses are known, it removes the main uncertainty in the analysis of proton decay, flavor violation and electric dipole moments. We will be able to discuss the consequence of flavor physics at short distances quantitatively.

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Probing Physics at Short Distances with Supersymmetry

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We discuss the prospect of studying physics at short distances, such as Planck length or GUT scale, using supersymmetry as a probe. Supersymmetry breaking parameters contain information on all physics below the scale where they are induced. We will gain insights into grand unification (or in some cases string theory) and its symmetry breaking pattern combining measurements of gauge coupling constants, gaugino masses and scalar masses. Once the superparticle masses are known, it removes the main uncertainty in the analysis of proton decay, flavor violation and electric dipole moments. We will be able to discuss the consequence of flavor physics at short distances quantitatively.

1 Introduction

The aim of particle physics is very simple: to understand the structure of matter and their interactions at as short distant scale as possible. This is the ultimate form of the reductionist approach of physics. This approach has revealed several layers of distance scales in nature, bulk (1 cm), atomic ($10^{-8}$ cm), nuclear ($10^{-13}$ cm). Understanding physics at shorter distance scales always gave us better understanding of physics at a previously known longer distance scale. Knowing the structure of atoms, we can deduce the chemical properties of atoms and molecules. Knowing the statistics of nuclei, we understand the levels of molecular excitations. The quest continues to understand the origin of the known distance scales, such as the electroweak scale of $10^{-16}$ cm, and to discover new layer of physics at shorter distance scales.

The main motivation of supersymmetry is to stabilize the electroweak scale against radiative corrections, which tend to make it much shorter (as short as the Planck length) or much longer (no electroweak symmetry breaking). Whenever we speculate about physics at shorter distance scales, we cannot go around the problem of the stability of the electroweak scale. However, once we accept supersymmetry as the stabilization mechanism, we are allowed to speculate physics at much shorter distances, and ask questions about the origin of gauge forces, fermion masses, and even cosmological issues such as baryon asymmetry.

The aim of this short article is to further elaborate on this point. Not only that supersymmetry allows us to speculate physics as the shortest distance scales, it actually provides probes of it. One can even dream about exploring physics at the GUT scale ($10^{-30}$ cm) or the Planck scale ($10^{-33}$ cm) once we
see superparticles. We will present various possibilities how we may be able to probe physics at such short distance scales using supersymmetry as a probe (see, e.g., a review on this point\textsuperscript{1}). We therefore assume that we will find and can study superparticles at collider experiments.\textsuperscript{2}

Of course, such a dream scenario cannot be discussed without certain assumptions. For each examples of such probes in the following sections, we try to make explicit what the underlying assumptions are. One of the main assumptions in any of these discussions is that the layers in distance scales are exponentially apart from each other. This is not an unreasonable assumption from the historic perspective. All the layers of physics came at very disparate distance scales. There appears to be nothing new between the characteristics scales: deserts. We do not know if this is the way nature is organized; we can only assume that the shorter distance scales also come with exponential hierarchy and discuss its consequences in our further exploration of physics at yet shorter distance scales.

2 Grand Unification

The simplest example in our approach is the grand unification. Needless to explain, a grand unified theory intends to explain the rather baroque pattern of quark, lepton quantum numbers in the Standard Model by embedding its gauge groups into a simple gauge group. Such a theory would resolve bizarre puzzles in the Standard Model. The fact that the matter is neutral (at the level of at least $10^{-21}$) requires a cancellation of electric charges between and electron, two up-quarks and one down-quark. The cancellation of anomalies in the Standard Model appears miraculous. And probably most importantly, the grand unification explains why the strong interaction is stronger than the electromagnetism as a simple consequence of the difference in the size of the gauge groups.

2.1 Gauge Coupling Constants

The supersymmetric grand unification received a strong attention in the past seven years after the precise measurement of the weak mixing angle at LEP in 1991. Many took the agreement of the observed and predicted value of the weak mixing angle as an experimental support for supersymmetry because the prediction was quite off if the minimal non-supersymmetric Standard Model was used to predict the weak mixing angle. The situation has not changed

\textsuperscript{a}Needless to say, it is important to confirm experimentally that the discovered new particles have properties consistent with supersymmetry.
qualitatively since. The detailed discussions on the dependence on superparticle spectrum, GUT-scale threshold, and its correlation to the proton decay are all important issues, and we refer to another chapter on gauge unification.

Since our question is what we will learn by using supersymmetry as a probe, let us suppose that we already have found the superparticles. Then the question on grand unification changes dramatically. First of all, we do not need to motivate supersymmetry assuming grand unification. The question goes the other way around. Since we know that the supersymmetry is there, we will rather ask if the gauge coupling constants unify given the particle content seen at the electroweak scale. More importantly, we measure all the masses of superparticles, which give us quantitative inputs on the supersymmetry thresholds in the renormalization group analysis. Knowing the particle content at the TeV scale and their masses will completely change the rule of the research. Note, however, that this analysis assumes a desert between the electroweak scale under experimental study and the GUT scale.

If they do unify within a certain accuracy, say within a few percents, we will begin asking what the origin of the small mismatch is (if any). For each of the GUT models we construct, we calculate the GUT-scale threshold corrections and compare them to the data. Such an analysis would certainly exclude parts of the parameter space in each model, and in some cases, the model itself. Especially the correlation to the proton decay becomes important, since the GUT-scale threshold corrections contain information about the mass of the color-triplet Higgsino which mediates the dimension-five proton decay such as $p \rightarrow e^+ K^0$ and the mass of the GUT gauge bosons which mediate the dimension-six proton decay such as $p \rightarrow e^+ \pi^0$. We will come back to this point in the section on proton decay. Another origin of a small mismatch may be a higher dimension operator in the gauge kinetic function which depends on the GUT-Higgs field, such as $\int d^2 \theta \text{Tr}(\Sigma \frac{M}{\tilde{M}} W^\alpha W^\alpha)$, where $\Sigma$ is the adjoint Higgs in $SU(5)$ GUT [3, 4]. If we assume that $\tilde{M}$ is the reduced Planck scale and $\Sigma$ has a VEV at the conventional GUT-scale of order $10^{16}$ GeV, such an operator gives an order percent correction to the gauge coupling unification. This possibility may unfortunately contaminate the information on GUT-scale threshold. In any case, however, it is clear that the rule of the game changes from motivating supersymmetry using GUT to making selection of GUT models from observed supersymmetry spectrum.

Unfortunately, the fact that the observed gauge couplings appear consistent with the $SU(5)$ unification does not rule out other possibilities, such as intermediate gauge groups with certain matter content. Fig. 6 shows two patterns of gauge coupling unifications, one with grand-desert $SU(5)$ and the other with intermediate Pati–Salam symmetry. The latter model is intended...
to be a comparison toy model to make the points clear for the later discussions, how the study of superparticles would help us to sort out the physics at shorter distance scales.

2.2 Gaugino Masses

Another aspect of the grand unification is the unification of superparticle masses. This discussion assumes that the supersymmetry breaking parameters are generated at a scale higher than the GUT-scale, such as the string or Planck scales, and hence respect grand-unified symmetry. Under this assumption, we will see if the superparticle masses unify at the same scale as where the gauge coupling constants unify. In fact, the gaugino mass unification,

\[ \frac{M_1}{\alpha_1} = \frac{M_2}{\alpha_2} = \frac{M_3}{\alpha_3} \]  

holds even when the GUT-group breaks to the Standard Model gauge group in several steps, i.e., with intermediate gauge groups such as Pati–Salam $SU(4) \times SU(2) \times SU(2)$ or its subgroup starting from $SO(10)$ or $E_6$, as long as the Standard Model gauge groups are embedded in a simple group with a single gaugino mass. At low-energy, the gaugino masses run in the exactly the same way as the gauge coupling constants squared do, which can be read off from Fig. in these two examples. Therefore, the gaugino mass unification tests the idea of grand unification in a highly model-independent manner.

Experimental strategies have been discussed how to disentangle the mixing in the neutralino-chargino sector to measure the supersymmetry breaking
masses for $SU(2)$ and $U(1)$ gauginos, $M_2$ for wino and $M_1$ for bino, at future $e^+e^-$ linear colliders. At the LHC, mass differences can be measured well by identifying the end points in the decay distributions. Upon the assumption that the first and second neutralinos are close to pure bino and wino eigenstates (which may be cross-checked by other analyses of the data), the mass differences also test the gaugino unification. Putting information from both types of experiments, we will have three new numbers to deal with. This will provide us two more independent test if they unify at the scale determined from the gauge coupling unification.

It is an important question if the gaugino mass unification can be spoiled even within GUT models. One possible effect is the threshold correction at the GUT scale. Fortunately, there is no logarithmic threshold corrections unlike to the gauge coupling constants and hence the gaugino mass unification is a better prediction of GUT than the gauge coupling unification. The only way to spoil the gaugino unification from the threshold corrections is to have extremely large representations under the grand unified group with highly non-universal trilinear and bilinear supersymmetry breaking parameters. Another possible effect is a higher dimension operator in the gauge kinetic function which depends on the GUT-Higgs field, such as $\int d^2\theta \text{Tr}(\Sigma W_\alpha W^{\alpha})$ we discussed before, with an $F$-component VEV of the $\Sigma$ field. If we take $M$ at the reduced Planck scale and $F_\Sigma \simeq m_{\text{SUSY}} M_{\text{GUT}}$, this operator generates an order percent correction to the gaugino mass unification. Note, however, that the size of the $F$-component VEV tends to be only of $m_{\text{SUSY}}^2$ in a wide class of supersymmetry breaking.

On the other hand, there is a case where we may be fooled by the apparent gaugino mass unification. In the models of gauge mediated supersymmetry breaking, the gaugino masses may satisfy the same relation as the case with grand unification even though the supersymmetry breaking gaugino masses have nothing to do with the physics of grand unification. However, this happens only when the messenger fields fall into full $SU(5)$ multiplets and when they acquire masses from the same field which has both $A$- and $F$-component VEVs. This is naturally expected in the GUT models, even when the supersymmetry breaking is induced well below the GUT-scale. On the other hand, there is no reason for the messengers to fall into full $SU(5)$ multiplets and acquire masses from the same field if the theory is not grand unified. Even though this remains as a logically possibility that the data can fool us, the apparent gaugino mass unification still strongly suggests grand unification. One case which probably cannot be distinguished on the bases of gaugino masses is the dilaton-dominant supersymmetry breaking in superstring models. In this case, however, there is a specific prediction on the ratio
Figure 2: Experimental tests of gaugino mass unification at a future $e^+e^-$ collider and the LHC.
of the scalar masses (universal to all scalars) and gaugino masses (universal to all gauginos) and can be confronted to the data.

There are GUT-like models which do not lead to unified gaugino masses, when the unified group is not simple. One example is flipped $SU(5)$. In this case, we expect to see the unification of $M_2$ and $M_3$ at the scale where the gauge coupling constants $\alpha_2$ and $\alpha_3$ meet, while $M_1$ may not. This is an interesting discriminator. The other is the model of dynamical GUT-breaking based on $SU(5) \times SU(3) \times U(1)$. Here the gaugino masses do not appear unified at all.

### 2.3 Scalar Masses

Under $SU(5)$ grand unified group, quarks and leptons belong to either $10$ or $5^*$ multiplets. Under the same assumption that the supersymmetry breaking masses respect grand-unified symmetry, we can extrapolate the observed scalar masses to higher energies and see if they unify at the same scale where the gauge couplings and gaugino masses unify (if they do).

The scalar mass unification will be an independent useful piece of information beyond that from gauge couplings and gaugino masses. One probably very convincing case for grand-desert $SU(5)$ unification is when both the gauge couplings and gaugino masses all unify at the same scale, and also scalars in $10$ and $5^*$ unify there but with different masses. On the other hand, the dilaton-dominated supersymmetry breaking predicts the universal scalar mass, not separate for $10$ and $5^*$, and a definite ratio of the scalar mass to the gaugino mass ratio.

There is a possibility that we get fooled by an apparent unification of gauge couplings and gaugino masses. This happens, for instance, when the supersymmetry breaking is induced by gauge mediation with messengers in full $SU(5)$ multiplets which acquire both supersymmetric and supersymmetry-breaking masses from a single field. We argued in the previous section that such a case already strongly suggests grand unification, but there remains a possibility that it is not. In this case, the scalar masses do not appear grand unified, and provide a way to differentiate the conventional grand-desert $SU(5)$ GUT from gauge-mediated supersymmetry breaking.

Unlike the gaugino masses, the scalar mass spectrum is sensitive to the pattern of GUT symmetry breaking. Many different patterns of scalar masses were discussed from GUT models. An important effect is that the scalars which originally resided in the same GUT multiplet may acquire different contributions from the $D$-term VEV when the rank of the gauge group is reduced. For instance, all quarks and lepton fields live in the same $16$ multiplet.
under $SO(10)$; but the breaking of $SO(10)$ to $SU(5)$ generally splits the $10$ and $5^*$ masses because of the $D$-term. The $D$-term contributions are determined solely by the gauge quantum numbers under the broken gauge group and hence generation blind, and are safe from the point of view of flavor-changing effects. A complicated superpotential interactions may modify the scalar masses as well. An extreme case is when the quarks and leptons in the Standard Model come from different GUT multiplets; then their scalar masses can be totally unrelated. However, the constraint from the flavor-changing neutral currents and smallness of Yukawa couplings for the first, second generation give us a prejudice that the superpotential couplings, which can potentially split the mass of first and second generation scalars in a non-universal manner, are small. Therefore, it is likely that the scalar masses unify according to the patterns of the GUT symmetry breaking, at least for the first and second generations.

By allowing $D$-term contributions but not $F$-term contributions motivated by the above argument, one can try to fit the observed scalar mass spectrum as a function of the symmetry breaking scales and original supersymmetry breaking parameters. For more complicated symmetry breaking patterns, there are less relations and hence the model is harder to test. But still in many interesting symmetry breaking patterns, there remain non-trivial relations among scalar masses which can be confronted to data. The Fig. 3 shows how the scalar masses acquire different patterns at the electroweak scale between the grand-desert $SU(5)$ and the toy Pati–Salam model, which could not be distinguished based on the gauge coupling constants and the gaugino masses. Therefore the role of gaugino mass unification and scalar mass unification are complimentary; the former gives a model-independent test of the grand unification, while the latter selects out particular symmetry breaking patterns and their energy scales.

Even in the case of non-simple GUT groups, the scalar masses still give us useful information. In flipped $SU(5)$, different sets of fields $(Q,d^c,\nu)$ belong to $10$ and $(L,u^c)$ to $5^*$, and $e^c$ is a singlet by itself. Therefore, there is still the scalar-mass unification of $Q$ and $d^c$, and $L$ and $u^c$ separately. In the model of dynamical GUT-breaking based on $SU(5) \times SU(3) \times U(1)$, still all matter fields belong to the ordinary $10+5^*$ multiplets and the pattern of scalar masses is the same as in the grand-desert $SU(5)$.

Experimentally, measurement of scalar masses is also feasible. At an $e^+e^-$ collider, a well-defined kinematics allows a simple kinematic fit to the decay distributions to extract the mass of the parent scalar particle. This comment applies both to the sleptons and squarks using beam polarizations, as long as they are within the kinematic reach. At the LHC, the mass differences
Figure 3: The renormalization group evolution of the scalar masses in two models, SU(5) GUT with grand desert and SO(10) GUT with an intermediate Pati–Salam symmetry used in Fig. 1, which cannot be discriminated based on gauge coupling constants and gaugino masses.

can be measured well as before; especially when the second neutralino decays into on-shell sleptons, one has a high rate and the mass difference between the slepton and the lightest neutralino is measured very well. Many other mass patterns also allow certain mass differences to be measured accurately at a few percent level. It is quite imaginable that the spectroscopy of superparticles will be the main experimental project of the next decade.

3 Proton Decay

Proton decay has been virtually the only direct probe of the physics at the GUT-scale and discussed extensively in the literature. The original idea is that the gauge bosons in SU(5) GUT cause transitions between quarks and leptons in the same SU(5) multiplets and hence allow proton to decay. Assuming that the quarks and leptons of the first generation belong to the same SU(5) multiplets, the exchange of the heavy SU(5) gauge boson generates an operator

\[ \mathcal{L} = \frac{1}{M_V^2} u d e, \]  

which gives rise to a decay \( p \to e^+ \pi^0 \). The current experimental bound excludes the process for heavy gauge bosons approximately up to \( 1.5 \times 10^{15} \) GeV, where we estimated the bound conservatively. Because the operator has a suppression by two powers of a high mass scale, the proton decay rate is suppressed by the fourth power in the mass scale \( \Gamma_p \propto m_p^5 / M^4 \). It is not easy to extend the experimental reach on \( M \). SuperKamiokande will probably extend the limit on the lifetime by a factor of 30 beyond the current one, which translates to a
modest improvement by a factor of 2.3 on the GUT-scale. ICARUS may reach the mass scale of the supersymmetric GUT or \(10^{16}\) GeV.

3.1 \( D = 5 \) operators

The important and novel feature in supersymmetric models is that there are operators of dimension-five which violate baryon and lepton numbers and hence can cause proton decay.\(^{[10]}\) For instance, the following operator is possible in the superpotential:

\[
W = \frac{\lambda}{M} (Q_1 Q_1)(Q_2 L_i),
\]

where the subscript refers to the generation, and \(\lambda\) is a coupling constant. The operator involves squarks and sleptons, which need to be converted to quarks or leptons by a loop diagram. The proton decay rate therefore scales as \(\Gamma_p \propto \frac{m_p \lambda^2}{(16\pi^2)^2 M^2 m_{SUSY}^2}\) where \(m_{SUSY}\) is the mass scale of superparticles. As a result, the reach in the energy scale is drastically improved. The current experimental limit does not allow \(M\) below \(10^{24}\) GeV if \(\lambda \sim 1\). Therefore, we are sensitive to even Planck-scale suppressed operators which, actually, are excluded with \(O(1)\) couplings.

It is interesting that the dimension-five operators necessarily involve quark superfields of different generations (at least two). This is a simple consequence of the Bose symmetry among superfields and the Standard Model gauge invariance. The interesting consequence of this fact is that the proton (or neutron) decay modes preferentially involve kaons in the final state, such as \(p \rightarrow K^+ \bar{\nu}_\mu\) as predicted to be dominant in the minimal \(SU(5)\) GUT. If the dominant proton decay mode will be seen to involve kaons, it is likely to be a consequence of dimension-five operators possible only in supersymmetric theories.

3.2 Minimal SUSY \( SU(5) \)

In the minimal SUSY \( SU(5) \) GUT\(^{[2]}\) the dimension-five operators are generated by the exchange of color-triplet Higgs supermultiplet. The Yukawa couplings of quarks to the Higgs doublets are known from the quark masses, and the couplings to the color-triplet Higgs (\(SU(5)\) partner of the doublets) are the those to the Higgs doublets at the GUT-scale because of the \(SU(5)\) invariance. Therefore, there is little freedom in this model and the size of the dimension-five operators is completely fixed except possible relative phases which become unobservable below the GUT-scale.\(^{[2]}\) The mass of the color-triplet Higgs at

\(^{[10]}\)In this discussion, we assume that there is no dimension-four operators which violate baryon and lepton numbers. Such operators are conveniently forbidden by imposing the \(R\)-parity.
first appears to be a free parameter. However, the gauge unification constrains its mass through its threshold correction. At the one-loop level of the renormalization group equations, one obtains

\[
(3\alpha_2^{-1} - 2\alpha_3^{-1} - \alpha_1^{-1})(m_Z) = \frac{1}{2\pi} \left\{ \frac{12}{5} \ln \frac{M_H}{m_Z} - 2 \frac{m_{SUSY}}{m_Z} \right\}, \quad (4)
\]

\[
(5\alpha_1^{-1} - 3\alpha_2^{-1} - 2\alpha_3^{-1})(m_Z) = \frac{1}{2\pi} \left\{ 12 \ln \frac{M_V^2 M_{\Sigma}}{m_Z^2} + 8 \ln \frac{m_{SUSY}}{m_Z} \right\}. \quad (5)
\]

Here, \( m_{SUSY} \) stands for some weighted average of the superparticle masses. Once the mass spectrum of the superparticle is measured, one can determine \( m_{SUSY} \) in the above formulae, and then extract the mass of the colored Higgs \( M_{HC} \), and a combination of \( M_V \) and \( M_{\Sigma} \) from the renormalization group equations. In fact, the measured \( \alpha_s \) is smaller than the preferred value from the GUT and as a result prefers a low-value of color-triplet Higgs mass. Since various \( \alpha_s \) measurements basically converged recently to \( \alpha_s(m_Z) = 0.118 \pm 0.003 \), the minimal \( SU(5) \) model is almost excluded unless extreme parameters are chosen for gauginos (preferentially light) and squarks (preferentially heavy). Given the uncertainties in the superparticle spectrum, it is hard to announce the definite exclusion of the model. Once the superparticles are found, however, we will be able to make a final word on the model, assuming the current value of \( \alpha_s \) persists and the SuperKamiokande will not find proton decay.

### 3.3 Non-minimal SUSY-GUT

There are many good reasons to discuss extensions of the minimal SUSY SU(5) GUT. Among them, there are two points directly relevant to the nucleon decay. (1) The triplet-doublet splitting problem. In minimal SUSY SU(5) GUT, one needs to fine-tune independent parameters at the level of \( 10^{-14} \) to keep Higgs doublets light while making the color-triplet Higgs heavy. (2) The wrong fermion mass relations. It predicts \( m_s = m_\mu \) and \( m_d = m_e \) at the GUT-scale, which are off from the phenomenologically preferred Georgi–Jarlskog relations \( m_s = m_\mu/3, \ m_d = 3m_e \).

Solutions to the above-mentioned problems modify the predicted rate and branching ratios of the nucleon decay. One possible attempt to obtain Georgi–Jarlskog relations is to use the SU(5)-adjoint Higgs to construct an effective Higgs doublets as composites of ordinary Higgs doublets in 5 and the adjoint. This modification leads to a factor-of-two enhancement in the amplitude; a factor of four in the rate. The relative branching ratios can be also different.
Figure 4: Excluded region on \((\tan \beta, \alpha_s(m_Z))\) plane from nucleon decay based on very conservative assumptions using constraints on superparticle spectrum from LEP-1 only. Expected improvements from SuperKamiokande and LEP-2 are also shown. The range shown for \(\alpha_s(m_Z)\) from PDG96 is two-sigma range. The preferred region from \(b-\tau\) unification is also shown for \(m_t = 176\) GeV as a crescent-shaped region.

It remains true that the \(K^{+0}\bar{\nu}_\mu\) modes are the dominant ones, while the \(K^0\mu^+\) mode may be much less suppressed than in the minimal SU(5) the proton decay in \(SO(10)\) models with realistic fermion mass texture has been also discussed extensively.

There are various proposals to solve the triplet-doublet splitting problem, which lead to different nucleon decay phenomenology. I discuss three of them here. (1) The missing partner model, (2) Dimopoulos–Wilczek–Srednicki mechanism, and (3) flipped SU(5) model.

In the missing partner model, one employs 75 representation to break SU(5) instead of the adjoint 24, and further introduces 50 and 50* representations which mix with the color-triplet Higgs to make them massive. Since the model involves such large representations, the size of the GUT-scale threshold corrections are significantly larger than that in the minimal model. And the correction changes the determination of the color-triplet Higgs mass as done in Eq. (4), and the measured values of the gauge coupling constants prefer larger \(M_{HC}\) than in the minimal model. In this case the proton decay rates are much more suppressed, by a few orders of magnitudes. One drawback of the model is that it becomes non-perturbative well below the Planck scale due to large representations and one needs to complicate the model further to keep it perturbative. It is worth to recall that the minimal SU(5) model is marginally allowed only with very conservative assumptions. Even though there is an ad-
ditional suppression to the proton decay rate in this class of models, the decay rate may still well be within the reach of superKamiokande experiment.

The mechanism proposed by Dimopoulos, Wilczek and further by Srednicki employs SO(10) unification with Higgs fields in adjoint and symmetric tensor representations which naturally keep Higgs doublets light. However, their model breaks SO(10) only to $SU(3) \times SU(2)_L \times SU(2)_R$ and has to be extended to achieve the desired symmetry breaking down to the standard model gauge group. One of such extensions by Babu and Barr eliminates $D = 5$ entirely; but it involves rather complicated Higgs sector, and one needs to forbid some allowed interactions in the superpotential arbitrarily. A later attempt to guarantee the special form of the superpotential by symmetries did not eliminate the $D = 5$ operators entirely, but resulted in a weak suppression of the operators. Again in view of the very marginal situation in the minimal model, the decay rate could be within the reach of the superKamiokande.

The flipped SU(5) model solves the triplet-doublet splitting problem in a way that it also eliminates the $D = 5$ operators entirely. A possible problem with this model is that the gauge unification becomes more or less an accident rather than a prediction. On the other hand, the elimination of the $D = 5$ operator is a natural consequence of the structure of the Higgs sector, and is rather a robust prediction of the model except the Planck-scale effects which will be discussed below. An interesting feature of the model is that the GUT-scale is determined by $\alpha_2$ and $\alpha_3$ and hence can be lower than the scale in the minimal SU(5) which is determined by $\alpha_2$ and $\alpha_1$. Since the model does not predict the relation between $\alpha_1$ and $\alpha_{SU(5)}$, $\alpha_1$ does not need to meet with the other coupling constants at the same scale. Therefore, the GUT-scale can be as low as $M_{GUT} = 4 - 20 \times 10^{15}$ GeV. If the $M_{GUT}$ is at the low side within this range, the $D = 6$ operator may be observable in the $\pi^0 e^+$ mode since the superKamiokande is expected to extend the reach by a factor of 30.

Certain models of direct gauge mediation also have $SU(5)$ group broken below the typical GUT-scale and can lead to dimension-six proton decay at a rate observable at superKamiokande. There is also a variant of the missing-partner model with dimension-six proton decay within the reach of superKamiokande.

3.4 Planck-scale Operators

Planck-scale physics may generate $D = 5$ operators suppressed by the reduced Planck scale $M_p = 2 \times 10^{18}$ GeV. Even when there is no color-triplet Higgs, such as in string compactifications which breaks the gauge group down to the standard model (with possible U(1) factors) directly, the higher string
excitations may give rise to effective non-renormalizable $D = 5$ operators which break baryon- and/or lepton-number symmetries. For $D = 5$ operators which involve first- and second-generation fields, $1/M_*$ suppression is far from enough: one needs a coupling constant of order $10^{-7}$ to keep the nucleons stable enough as required by experiments.

It is a serious question in supersymmetry phenomenology why the Planck-scale $D = 5$ operators are so much suppressed. One possibility is to forbid them by employing discrete gauge symmetries which are believed to be respected by quantum gravitational effects unlike global symmetries. In this case, there is no baryon-number-violating $D = 5$ operator from Planck-scale physics and we do not have any handle on it. A different type of solution is probably more interesting: the $D = 5$ operators are suppressed because of the same reason why the Yukawa couplings of light generations are suppressed. One way to understand why the Yukawa couplings are so small, such as $10^{-6}$ for the case of the electron, may be a natural consequence of an approximate flavor symmetry. If a flavor symmetry exists and is only weakly broken to explain smallness of the Yukawa couplings, the same flavor symmetry can well suppress the $D = 5$ operators at the Planck-scale. The $D = 5$ operators with such a flavor origin may have very different flavor structure from those in the GUT models, and may lead to quite different decay modes like $p \rightarrow K^0 e^+$.

Suppression of Planck-scale $D = 5$ operators based on certain flavor symmetries were discussed. For instance the $S_3^3$ model explains the hierarchical Yukawa matrices as a consequence of sequential breaking of the flavor symmetry while the symmetry preserves sufficient degeneracy among the squarks and sleptons to suppress flavor-changing neutral currents. It happens that the flavor symmetry in this model also suppresses $D = 5$ operators to the level of about $1/9$ of the minimal $SU(5)$ model, so that it can well be within the reach of superKamiokande. What is particularly interesting in this model is that it predicts $p \rightarrow K^0 e^+$ as the dominant mode over the $K^+ \bar{\nu}$, while $n \rightarrow K^0 \bar{\nu}_e$ is the dominant mode in neutron decay with a comparable rate. In general, decay modes of proton, if observed, will provide interesting information on the flavor physics.

4 Flavor Physics

Another interesting topic is how well we will be able to understand the origin of flavor, fermion masses and mixing based on the study of supersymmetry possible at the electroweak scale. Unlike the case of grand unification and proton decay, the answer to this question depends heavily on what the true story is.
4.1 Neutrino Physics

An analogue of proton decay discussed in the previous section is a consequence of flavor physics suppressed by powers of the mass scale, such as the neutrino mass via the seesaw mechanism. The neutrino masses are generated from their Dirac masses $m_D$ with right-handed neutrinos neutral under the Standard Model gauge groups and their Majorana masses $M$ which violate the lepton numbers by two units. The one-generation case is given by a two-by-two mass matrix

$$\mathcal{L}_{mass} = \frac{1}{2} (\nu, N^c) \begin{pmatrix} 0 & m_D \\ m_D & M \end{pmatrix} \begin{pmatrix} \nu \\ N^c \end{pmatrix},$$

where $\nu$ is the Weyl field of the left-handed neutrino, and $N$ the right-handed neutrino. The Lagrangian is written in terms of the charge-conjugated Weyl spinor $N^c$ with the same chirality as the left-handed field $\nu$. After diagonalization of the mass matrix, one obtains a mass for the left-handed neutrino of $m_D^2 / M$, which is power suppressed. This mechanism naturally explains why the neutrino masses are so small, if finite, and leaves imprint of short-distance physics in the pattern of neutrino masses and mixings.

Let us emphasize that this is the area where a dramatic progress is likely to be made in the next few years, from SuperKamiokande (together with neutrino beam from KEK), CHORUS, NOMAD, KARMEN, SNO, BOREXINO, MINOS, and more. Even though supersymmetry does not necessarily help to study the physics at the scale of right-handed neutrinos, many flavor models in supersymmetry predict interesting patterns on neutrino masses. We will certainly be making selections on different flavor models based on neutrino physics if finite neutrino masses and their mixings will be established.

4.2 Flavor-Changing Neutral Currents

As discussed in other chapters, there are severe constraints on the superparticle masses and mixings from the flavor-changing neutral currents (FCNC). There are broadly three categories of models which naturally suppress the FCNC. (1) Flavor symmetry enforces the squarks, sleptons to be degenerate or aligns their mass basis to that of the down-quark, charged-lepton mass basis. (2) The string theory generates universal scalar mass. (3) The supersymmetry breaking is generated in a flavor-blind fashion below the scale of flavor physics.

In cases (1) and (2), there may be interesting imprints of flavor physics in the small mixing between squarks and sleptons. The case of GUT also belongs to the category: the large top Yukawa coupling above the GUT-scale may affect the slepton masses-squared with small mixing between, for instance, selectron.
Figure 5: Contours of constant $\sigma(e^+e^-\rightarrow e^\pm\mu^\mp\tilde{\chi}^0\tilde{\chi}^0)$ (solid) and $\sigma(e^-e^-\rightarrow e^-\mu^-\tilde{\chi}^0\tilde{\chi}^0)$ (solid) in fb for $e^+e^-$ or $e^-e^-$ linear colliders, with $\sqrt{s} = 500$ GeV, $m_{\tilde{\chi}_R}, m_{\mu_R} \approx 200$ GeV, and $M_1 = 100$ GeV (solid). The thick gray contour represents the experimental reach in one year. Constant contours of $B(\mu\rightarrow e\gamma) = 4.9 \times 10^{-11}$ and $2.5 \times 10^{-12}$ are also plotted for degenerate left-handed sleptons with mass 120 GeV and $\tilde{t} \equiv -(A + \mu \tan \beta)/m_R = 0$ (dotted), 2 (dashed), and 50 (dot-dashed), with left-handed sleptons degenerate at 350 GeV.

The search for rare decays such as $\mu\rightarrow e\gamma$ or electric dipole moments of electron or neutron may reveal the imprints of flavor physics in scalar masses. At present, the main uncertainty in the quantitative analysis is the mass of superparticles. Once they are measured, however, we can try to extract the mixing effects in the scalar mass matrices from the FCNC data.

An interesting case where the flavor-mixing effects in scalar masses can be probed at colliders was discussed. Analogous to neutrino oscillation, a selectron produced from an $e^+e^-$ or $e^-e^-$ collider can oscillate to a smuon as a result of the flavor mixing and is detected as $e\mu$ final state (see Fig. 5).

In the case (3) where the scalar masses are generated in a flavor-blind fashion below the flavor physics scale, such as in the models of low-energy gauge mediation, we unfortunately may not learn about the origin of flavor from the study of flavor signatures at the electroweak scale.

5 Conclusion
Experiments at the electroweak scale will remove the cloud which masks the physics at yet shorter distance scales. If supersymmetry turns out to be the mechanism of stabilizing the electroweak scale, we will have a wealth of new
data on superparticle spectroscopy. Combined with data on proton decay, neutrino physics, and FCNC, we will obtain useful information on physics such as grand unification, string, flavor physics. At this point it is just a dream; but we may be able to glimpse the physics at the shortest possible distance scales by this program, which is nothing but the goal of particle physics after all.

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1. H. Murayama, Invited Talk at the ICEPP Symposium “From LEP to the Planck World,” University of Tokyo, Dec 17–18, 1992. In Proceedings of the ICEPP Symposium “From LEP to the Planck World,” eds. K. Kawagoe and T. Kobayashi, UT-ICEPP 93-12, TU-451, 11pp.
2. J.L. Feng, M.E. Peskin, H. Murayama, and X. Tata, Phys. Rev. D 52, 1418 (1995).
3. M. M. Nojiri, K. Fujii, and T. Tsukamoto, Phys. Rev. D 54, 6756 (1996).
4. J. Hisano, H. Murayama, and T. Yanagida, Phys. Rev. Lett. 69, 1014, (1992); Nucl. Phys. B402, 46 (1993).
5. L. J. Hall and U. Sarid, Phys. Rev. Lett. 70, 2673 (1993).
6. T. Dasgupta, P. Mamales, and P. Nath, Phys. Rev. D 52, 5366 (1995).
7. N.G. Deshpande, E. Keith, and T.G. Rizzo, Phys. Rev. Lett. 70, 3189 (1993).
8. Y. Kawamura, H. Murayama, and M. Yamaguchi, Phys. Lett. B324, 52 (1994).
9. T. Tsukamoto, K. Fujii, H. Murayama, M. Yamaguchi, and Y. Okada, Phys. Rev. D 51, 3153 (1995).
10. J. L. Feng and M. J. Strassler, Phys. Rev. D 51, 4661 (1995); Phys. Rev. D 55, 1326 (1997); H. Baer, R. Munroe, and X. Tata, Phys. Rev. D 54, 6735 (1996); Erratum-ibid 56, 4424 (1997).
11. I. Hinchliffe, F.E. Paige, M.D. Shapiro, J. Soderqvist, and W. Yao, Phys. Rev. D 55, 5520 (1997).
12. J. Hisano, H. Murayama, and T. Goto, Phys. Rev. D49, 1446 (1994).
13. Y. Kawamura, H. Murayama, and M. Yamaguchi, Phys. Rev. D51, 1337 (1995).
14. A. Pomarol and S. Dimopoulos, Nucl. Phys. B 453, 83 (1995); R. Rattazzi, Phys. Lett. B 375, 181 (1996).
15. M. Dine and A.E. Nelson, Phys. Rev. D48, 1277 (1993).
16. M. Dine, A.E. Nelson and Y. Shirman, Phys. Rev. D51, 1362 (1995).
17. M. Dine, A.E. Nelson, Y. Nir and Y. Shirman, Phys. Rev. D53, 2658 (1996).
18. S. M. Barr, Phys. Lett. B 112, 219 (1982); J.P. Derendinger, J. E. Kim, and D.V. Nanopoulos, Phys. Lett. B 139, 170 (1984); I. Antoniadis, J. Ellis, J.S. Hagelin, and D.V. Nanopoulos, Phys. Lett. B 194, 231 (1987).
19. T. Yanagida, Phys. Lett. B 344, 211 (1995);
20. N. Arkani-Hamed, H.-C. Cheng, and T. Moroi, Phys. Lett. B 387, 529 (1996).
21. V. S. Kaplunovsky and J. Louis, Phys. Lett. B 306, 269 (1993); A. Brignole, L.E. Ibanez, and C. Munoz, Nucl. Phys. B 422, 125 (1994), Erratum-ibid. 436, 747 (1995); A. Brignole, L.E. Ibanez, C. Munoz, and C. Scheich, Z. Phys. C 74, 157 (1997).
22. C. Kolda, and S. P. Martin, Phys. Rev. D 53, 3871 (1996).
23. H.C. Cheng, and L.J. Hall, Phys. Rev. D 51, 5289 (1995).
24. M. Drees, Phys. Lett. B 181B, 279 (1986); J.S. Hagelin and S. Kelley, Nucl. Phys. B342, 95 (1990); A.E. Faraggi, J.S. Hagelin, S. Kelley, and D.V. Nanopoulos, Phys. Rev. D45, 3272 (1992).
25. H. Murayama, Invited plenary talk given at 4th International Conference on “Physics Beyond the Standard Model,” Lake Tahoe, CA, 13-18 Dec 1994. Proceedings, eds. by J. Gunion, T. Han, J. Ohnemus, World Scientific, 1995.
26. S. Dimopoulos and A. Pomarol, Phys. Lett. B 353, 222 (1995).
27. J. L. Feng and D. E. Finnell, Phys. Rev. D 49, 2369 (1994).
28. J. Hisano, T. Moroi, K. Tobe, and T. Yanagida, Mod. Phys. Lett. A 10, 2267 (1995); J. Bagger, K. Matchev, and D. Pierce, Phys. Lett. B 348, 443 (1995).
29. H. Murayama, Invited talk presented at 28th International Conference on High-energy Physics (ICHEP 96), Warsaw, Poland, 25-31 Jul 1996. Published in the Proceedings of the 28th International Conference on High Energy Physics, eds., Z. Ajduk and A. K. Wroblewski, World Scientific, 1997, pp. 1377–1382.
30. N. Sakai and Tsutomu Yanagida, Nucl. Phys. B 197, 533 (1982); S. Weinberg, Phys. Rev. D 26, 287 (1982).
31. S. Dimopoulos and H. Georgi, Nucl. Phys. B 193, 150 (1981); N. Sakai, Z. Phys. C 11, 153 (1981).
32. P. Nath and R. Arnowitt, Phys. Rev. D 38, 1479 (1988).
33. H. Murayama, Phys. Rev. Lett. 79, 18 (1997); S. Dimopoulos, G. Dvali, R. Rattazzi, and G.F. Giudice, CERN-TH/97-98, hep-ph/9705307.
34. J. Hisano, Y. Nomura, and T. Yanagida, KEK-TH-547, hep-ph/9710279.
35. H. Murayama and D. B. Kaplan, Phys. Lett. B336, 221 (1994).
36. H. Murayama, Invited talk presented at the 22nd INS International Symposium on Physics with High Energy Colliders, Tokyo, Japan, March 8–10, 1994, published in Proceedings of INS Symposium, World Scientific, 1994.

37. A. Antaramian, LBL-36819, Feb 1995, Ph.D. Thesis.

38. K.S. Babu and S.M. Barr, Phys. Lett. B 381, 137 (1996).

39. V. Lucas and S. Raby, Phys. Rev. D 55, 6986 (1997).

40. K.S. Babu, J. C. Pati, and F. Wilczek, IASSNS-HEP-97-136, hep-ph/9712307.

41. A. Masiero, D.V. Nanopoulos, K. Tamvakis, and T. Yanagida, Phys. Lett. B 115, 380 (1982); B. Grinstein, Nucl. Phys. B 206, 387 (1982).

42. S. Dimopoulos and F. Wilczek, NSF-ITP-82-07 (unpublished); M. Srednicki, Nucl. Phys. B 202, 327 (1982).

43. K. Hagiwara and Y. Yamada, Phys. Rev. Lett. 70, 709 (1993); Y. Yamada, Z. Phys. C 60, 83 (1993).

44. J. Hisano, T. Moroi, K. Tobe, and T. Yanagida, Phys. Lett. B 342, 138 (1995).

45. K.S. Babu and S.M. Barr, Phys. Rev. D 48, 5354 (1993).

46. J. Ellis, J. L. Lopez, and D.V. Nanopoulos, Phys. Lett. B 371, 65 (1996).

47. L. E. Ibanez and G. G. Ross, Nucl. Phys. B368, 3 (1992).

48. V. Ben-Hamo and Y. Nir, Phys. Lett. B 339, 77 (1994).

49. L. J. Hall and H. Murayama, Phys. Rev. Lett. 75, 3985 (1995).

50. C. D. Carone, L. J. Hall, and H. Murayama, Phys. Rev. D53, 6282 (1996).

51. T. Yanagida, in Proceedings of Workshop on the Unified Theory and the Baryon Number in the Universe, Tsukuba, Japan, 1979, edited by A. Sawada and A. Sugamoto (KEK, Tsukuba, 1979), p. 95; M. Gell-Mann, P. Ramond and R. Slansky, in Supergavity, proceedings of the Workshop, Stony Brook, New York, 1979, edited by P. Van Nieuwenhuizen and D.Z. Freedman (North-Holland, Amsterdam, 1979), p. 315.

52. Y. Grossman and Y. Nir, Nucl. Phys. B 448, 30 (1995); M. Schmaltz, Phys. Rev. D 52, 1643 (1995); C. D. Carone and L. J. Hall, Phys. Rev. D 56, 4198 (1997); P. Binetruy, S. Lavignac, S. Petcov, and P. Ramond, Nucl. Phys. B 496, 3 (1997).

53. M. Dine, A. Kagan, and R. Leigh, Phys. Rev. D 48, 4269 (1993); P. Pouliot and N. Seiberg, Phys. Lett. B 318, 169 (1993); D.B. Kaplan and M. Schmaltz, Phys. Rev. D 49, 3741 (1994); A. Pomarol and D. Tommasini, Nucl. Phys. B 466, 3 (1996); R. Barbieri, G. Dvali, and L. J. Hall, Phys. Lett. B 377, 76 (1996); R. Barbieri, L. J. Hall, S. Raby, and A. Romanino, Nucl. Phys. B 493, 3 (1997).
54. Y. Nir and N. Seiberg, *Phys. Lett.* B 309, 337 (1993).
55. L. J. Hall, V. A. Kostelecky, and S. Raby, *Nucl. Phys.* B 267, 415 (1986).
56. R. Barbieri and L.J. Hall, *Phys. Lett.* B 338, 212 (1994).
57. R. Barbieri, L.J. Hall, and A. Strumia, *Nucl. Phys.* B 445, 219 (1995); J. Hisano, T. Moroi, K. Tobe, M. Yamaguchi, and T. Yanagida, *Phys. Lett.* B 357, 579 (1995).
58. R. Barbieri, L.J. Hall, and A. Strumia, *Nucl. Phys.* B 449, 437 (1995).
59. S. Dimopoulos and L.J. Hall, *Phys. Lett.* B 344, 185 (1995).
60. N. Arkani-Hamed, H.-C. Cheng, J. L. Feng, and L. J. Hall, *Phys. Rev. Lett.* 77, 1937 (1996).