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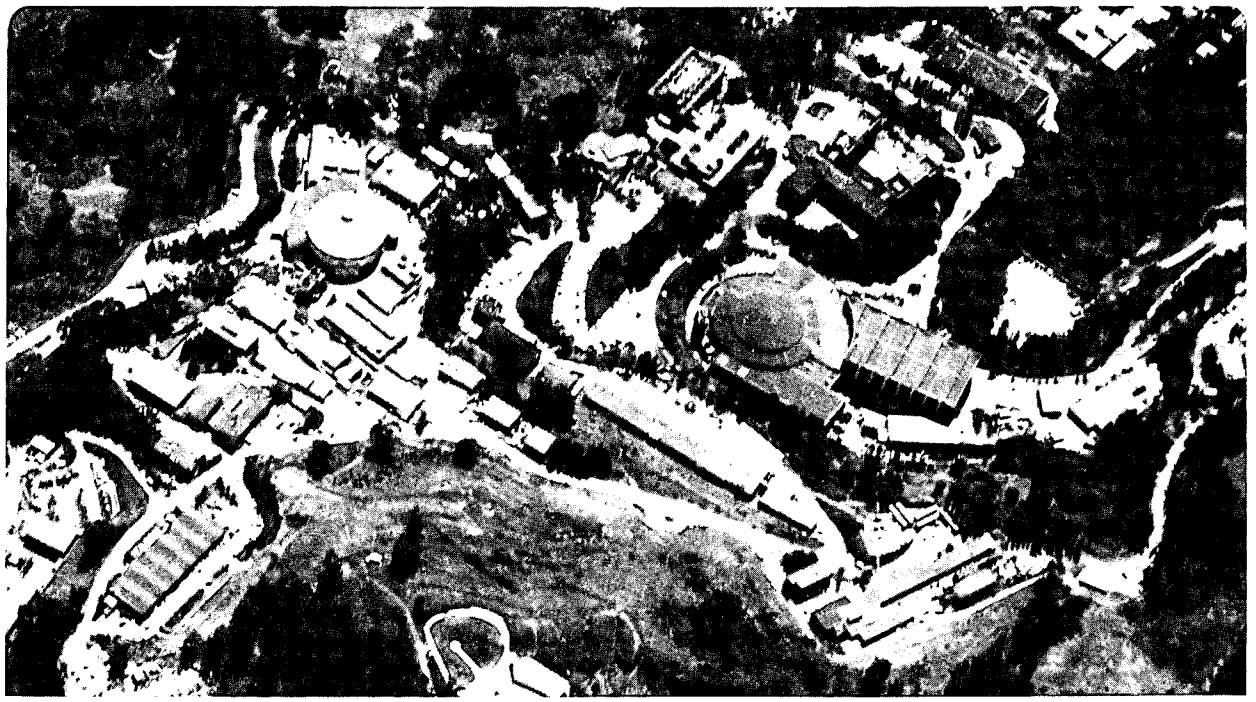
Beyond the Standard Model

L.J. Hall

October 1988



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Beyond The Standard Model. *†

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Abstract

A variety of TeV scale Higgs and flavor sectors are discussed. Key questions are addressed: how can we tell if there is a light Higgs boson or if the Higgs sector is strongly interacting? What new signatures can be used to search for supersymmetry? Can flavor physics be described at a TeV without Yukawa couplings? Ideas are reviewed and some new developments mentioned.

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(1) Scope And Orientation

The standard model [1] has three sectors: the quarks and leptons and their gauge interactions, the Higgs potential which induces electroweak symmetry breaking and the Yukawa interactions which describe fermion masses and mixing. We know that the $SU(3) \times SU(2) \times U(1)$ gauge interactions are correct: the determination of the gauge quantum number assignments of the quarks and leptons has been one of the dominant experimental preoccupations of the last decade, and they have now been verified time and again [2]. Any more complete theory at short distances will yield an effective theory containing these gauge interactions at the weak scale.

The only other knowledge we have about physics at the weak scale is that the W and Z bosons are heavy, and that most of the quarks and leptons also have a mass. This implies that there must be a sector of the theory which breaks $SU(2) \times U(1)$, which I will call the "Higgs" sector, even though it may have nothing to do with the Higgs boson of the standard model; and there must be a sector of the theory which breaks the chiral symmetries of the fermions allowing them to become massive, which I will call the "flavor" sector. Although we know that these sectors exist, and we know the size of various parameters that must exist in these sectors (determined by W, Z and fermion masses) that is all we actually know: we do not know the particles, the interactions or even the nature of the parameters of these sectors.

I have chosen to stress this point because the Higgs potential and the Yukawa interactions of the standard model are so simple that it is tempting to just accept them, assume that the entire standard model is correct, and use it as a launching pad for thinking about physics at a TeV and much beyond. This is a very pervasive view; it is even implicit in the title of this talk: presumably the standard model is known to be correct before going beyond it.

It is well known that the Higgs sector must show up, at least in some of its aspects, by a TeV: without a Higgs sector the gauge theory with a heavy W and Z will break partial wave unitarity [3]. We must come across some new physics to make certain high energy scattering amplitudes (such as $W_L W_L \to W_L W_L$) well-behaved. We also have theoretical prejudices about the nature of

this sector: it should give us some understanding as to the origin of the weak scale and why it is so much less than the Planck scale. These are the best arguments we have for prospective discoveries in the next decade, and hence one focus of this talk will be to review the known theoretical possibilities for the Higgs sector: a fundamental Higgs boson on its own or augmented with supersymmetry, electroweak symmetries broken dynamically by a new strong force, or finally a new strong force which binds a composite Higgs boson. These ideas are not new, but I offer no apology for concentrating on them. They are the only ideas we have about the only new physics which we are sure we will find. There are recent developments, and some of these will be discussed; however, attention will focus on both conventional and exotic signatures for the various Higgs sectors.

The flavor sector is, in many respects, more puzzling and exciting than the Higgs sector. There are more mass scales, more parameters and many more issues. Exotic flavor schemes inevitably have to address questions of neutrino masses, CP violation, approximate flavor conservation, axions and many others issues. Typically they make little, no or negative progress in really understanding the pattern of observed fermion masses; it is amazing how elusive progress in this direction is. The frustrating thing abut the flavor problem is that we have no idea about the relevant mass scale for an understanding of flavor: a TeV, the Planck scale or anywhere in between. This should not deter us: regardless of where a true understanding comes from, the information about fermion masses and mixing must be carried in a flavor sector at all scales. There must be a TeV flavor sector even if it does not answer all of our questions; and at the very least, we must find out what it is.

Many would argue that there is bound to be a Higgs boson and the TeV scale flavor sector will be just the Yukawa coupling of the standard model. To them the question is: how should these Yukawas be derived from a more complete theory? I have much sympathy with this view. Our inability to find any flavor physics beyond the standard model suggests that physics at a TeV may be very minimal: just the standard model. We continue to push up limits on rare μ, τ, K, D and B decays, not to mention limits on ν masses and mixings. Model building at a TeV is becoming a treacherous art form: TeV flavor physics

beyond the standard model, if it is there, is certainly well hidden. However, it is the physics of the TeV world which experiments are now entering; we should be "flavor optimists" and hope that we will be able to probe experimentally at least some of the puzzles of flavor. I do not think we should second guess nature: if the $SU(2) \times U(1)$ particle assignments had been the mass eigenstates (u, b), (c, s) and (t, d), when would hadronic weak interactions have been discovered? Nature is easily capable of hiding things just out of view!

This article does not do justice to the title: it is very incomplete. I have chosen to concentrate entirely on physics at the TeV scale. It is a very personal selection of TeV scale physics, and I apologize in advance to many people whose work I have not mentioned. Most grievous is the omission of any discussion of rare K and heavy flavor decay and of CP violation. These topics are crucial in directing us beyond the standard model and are omitted only for reasons of time. Material has frequently been chosen simply because it illustrates my theme: that there are many fascinating possibilities for the TeV scale Higgs and flavor sectors; and these sectors may have exotic experimental signatures other than the conventional ones currently being searched for.

II. The Higgs Sector

II.1 Overview

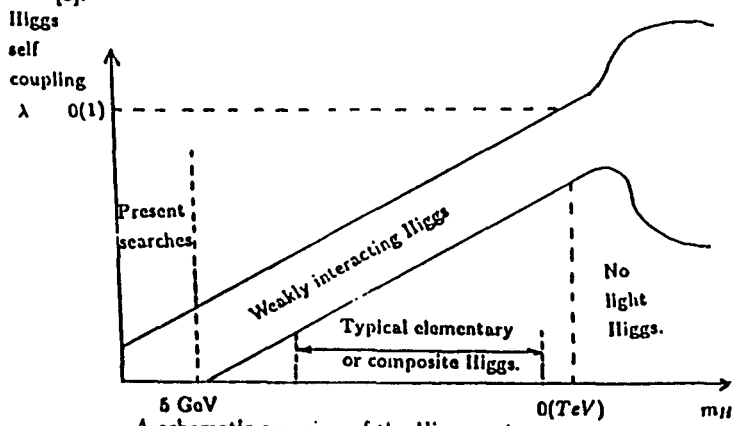
In the standard model [1] it is obvious what the Higgs particle is, it is the only fundamental scalar in the physical spectrum. In more complicated theories there may be several scalars. Are these all to be called Higgs bosons? I will call a scalar a Higgs boson only if its couplings are responsible for restoring unitarity in scattering amplitudes such as $W_L W_L \to W_L W_L$. With this definition, a Higgs is inextricably linked with electroweak symmetry breaking. A theory may have several such Higgs particles, but they need not have direct Yukawa interactions with fermions.

The mass of the Higgs in the standard model, m_H, is a free parameter. It is related to the Higgs self coupling λ by

$$m_H^2 = \lambda v^2 \tag{2.1}$$

where v is the vacuum expectation value. A feature of this result which is generic to any model with a Higgs boson is that the Higgs boson becomes more strongly

self–coupled as it becomes heavier. In fact, a narrow Higgs resonance can only be identified if its mass is less than about 1 TeV. In a strongly coupled Higgs sector, such as technicolor [4], there may not be a Higgs boson. On the other hand, if there is a light Higgs scalar it need not be fundamental, it could be a composite [5].



A schematic overview of the Higgs sector.

Figure 1

II.2 Higgs Searches.

Is there a Higgs? This is one of the most important questions which physics beneath the TeV scale can answer. The direct way of answering is to search for a scalar state with properties that we expect a Higgs boson to have. Since a Higgs boson might not couple to fermions this direct approach may not be successful even if a Higgs exists. On the other hand the Higgs boson of the standard model does couple to all fermions proportional to their mass, and searches for a light Higgs via these couplings have been performed in a variety of processes [6,7].

The Higgs boson of the standard model is very well hidden at low and moderate energies. We have accurate data on light particles (π, μ, K) to which the Higgs couples weakly, and less statistics in decays of particles to which the Higgs couplings more strongly (ψ, B, Υ). The absence of a ZHH coupling means that the Higgs appears only as a rare mode in Z decays [8,9]. In the future we

hope that the large coupling of the Higgs to W pairs will make the Higgs easily visible [10, 11, 12].

All searches to date for the Higgs of the standard model are restricted to Higgs masses less than 10 GeV. The only universally accepted limit is that of 14 MeV from nuclear transitions. There have been many high statistics searches in hadron decays, but in each of these cases theoretical uncertainties cast doubts on quoted limits. For example, there are unknown strong interaction corrections to $K \to H\pi$ [13]. One of the most troublesome problems is the strong interaction corrections to $H \to \pi^+\pi^-$, [14, 15] which affects the theoretical estimate for the branching ratio $H \to \mu^+\mu^-$, which is frequently the signature being sought.

For the Higgs boson of the three generation standard model to be lighter than 7 GeV, a one loop analysis of the Higgs potential requires the top mass to be close to 80 GeV [16]. Since $b \to Hs$ has a rate proportional to m_t^4, such a heavy top results in the exclusion of $m_H < 700$ MeV and 2 GeV $< m_H <$ 3.7 GeV [13, 14, 17].

At this conference an enlarged data set for searches for monochromatic photons, expected from $\Upsilon \to H\gamma$ for a light Higgs, were reported by the CUSB collaboration [18]. This search has the advantage that it covers a wide range of Higgs mass and it is independent of the Higgs decay modes. Comparing the data with the theoretical prediction in the standard model, including one loop QCD corrections, they conclude that at 95% confidence level the Higgs mass is excluded from the region .21 - 4.8 GeV. Again, the difficulty lies with being fully confident with the theoretical prediction: the first order QCD correction was almost a factor of two, so perhaps the next order correction is also large; also there are relativistic corrections to the Υ bound state system.

The mass of the Higgs boson is a crucial and fundamental parameter. We should probably take the cautious and conservative viewpoint that although there is mounting evidence against it being less than 5 GeV, it is still possible that it could be there. It should also be pointed out that if there is a Higgs lighter than, say, 10 GeV, then we will be saddled with another fine tuning problem. Supersymmetric models or models with a composite Higgs allow us to understand why the Higgs can be much lighter than the grand unified or Planck

scale, but we have no theoretical justification for a Higgs much lighter than the weak scale.

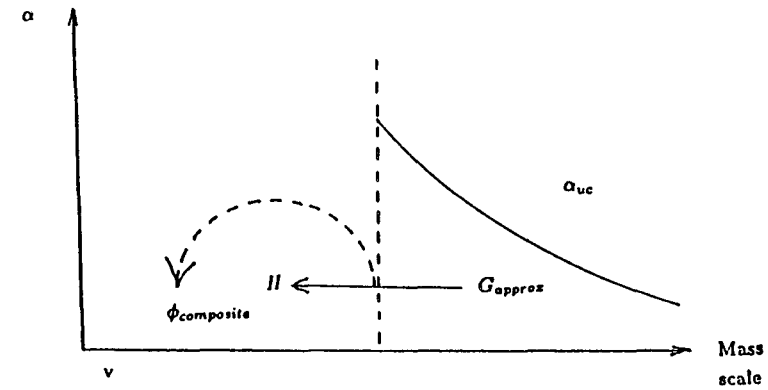
II.3. Origin of the Higgs Mass.

Suppose the Higgs boson exists. Why should its mass be much less than the Planck scale? A scalar mass parameter typically receives quadratically divergent radiative corrections, so that one would expect it to be dragged up to the fundamental mass scale of the theory, even if it were made small at tree level. There are a variety of situations where this is prevented, allowing a light Higgs to be natural. Such cases imply special properties for the Higgs and its mass parameter.

i)Higgs as a pseudo-Goldstone boson: the composite Higgs.

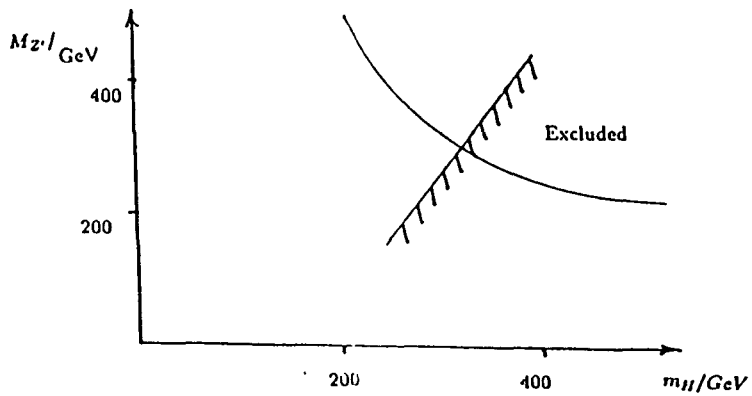
Suppose that a continuous global symmetry G is spontaneously broken to a subgroup H at scale M. This will result in the appearance of massless Goldstone bosons. Although this is a very efficient mechanism for creating light scalars, they cannot be identified with the Higgs: they are absolutely massless and they have no scalar potential. Furthermore, they do not allow a simple description of flavor as they have no Yukawa couplings. You could try to fix this up by adding terms to the theory which explicitly break the symmetry G, with small dimensionless coefficients ζ. The scalars become pseudo-Goldstone bosons (PGB). The problem with this idea is that if M is to be the Planck scale, then we require ζ to be very small: $\zeta \simeq M_W/M_p \simeq 10^{-17}$. This is unnaturally small, and the Yukawa couplings that are generated are proportional to ζ, far too small to be relevant for fermion masses.

It is possible to interpret the Higgs boson as a PGB provided it is a composite state [5]. The idea is sketched in Fig. 2. The standard model is augmented by an



The composite Higgs as a pseudo-Goldstone boson.
Figure 2

additional gauge sector, ultracolor, which has a confining scale Λ_{uc} of a few TeV. When viewed at this scale, the strong ultracolor condensates spontaneously break an approximate global symmetry $G \to H$, producing a PGB which is a composite of ultraquarks. Note that, unlike technicolor, these condensates do not break $SU(2)_L \times U(1)$. The explicit breaking of G is provided by non-ultracolor gauge interactions, and these generate a potential (and possibly Yukawa interactions) which forces the PGB to acquire a vacuum expectation value (vev). While it is not difficult to make the PGB an $SU(2)_L$ doublet, it is not easy to give it a vev. In fact, this appears to require augmenting the gauge interactions of the standard model to include an additional $U(1)$, which acts axially on fermions [19]. Hence, one predicts not only a new confining gauge interaction in the (perhaps multi-) TeV domain, but also a Z' with a sub-TeV mass. Since this Z' generates the potential of the PGB, in any given model there is a relation between the Z' mass and the Higgs mass. This is shown for a particular model [20] in Fig. 3.



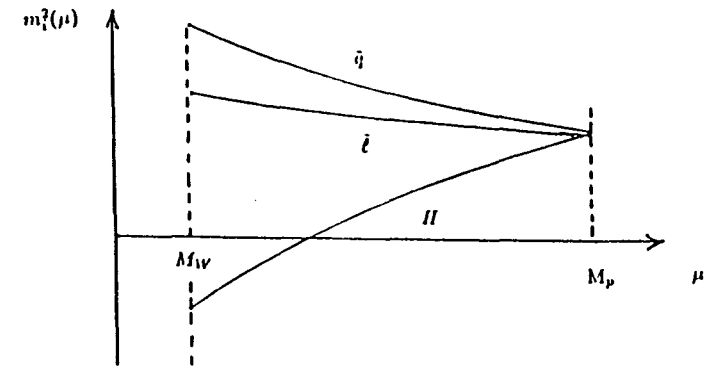
A graph of the Higgs mass against the Z' mass in a particular composite Higgs model.
Figure 3

ii). The Higgs as a superpartner.

Supersymmetric theories can break $SU(2)_L \times U(1)$ in a variety of ways. A superpotential

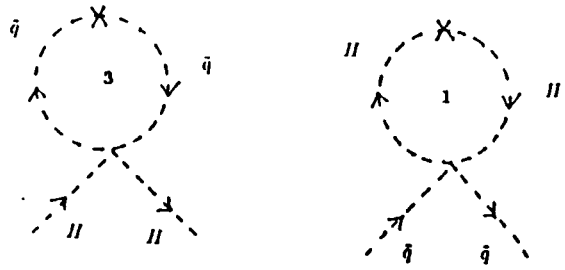
$$f = \lambda X(HH' - v^2) \tag{2.2}$$

where X, H, H' are chiral superfields, causes electroweak breakdown without breaking supersymmetry. However, it is more appealing to have $SU(2)_L \times U(1)$ breaking associated with supersymmetry breaking, since their scales are expected to be close. This preferred possibility occurs very elegantly in the minimal low energy supergravity model [21]. At the Planck scale all scalars of the theory acquire a small, soft, supersymmetry breaking mass $m_i^2(M_p) = m^2$. These various masses scale according to renormalization group equations beneath the Planck scale, as shown schematically in Fig. 4. There are several effects: gauge



The renormalization group scaling of scalarmasses: squarks sleptons and Higgs.
Figure 4

interactions increase scalar masses in the infrared making squarks heavier than sleptons. However, Yukawa interactions typically decrease the mass squared of that scalar in the interaction which has fewest gauge interactions. Thus a large top Yukawa causes the Higgs mass squared to scale negative; while the top squarks maintain positive squared masses. This effect is due to the simple loop counting factor illustrated in Figure 5.



The color loop factor causes Higgs mass to scale negative.
Figure 5

It appears that this mechanism allows for an understanding of why M_W/M_p is so small in the same way that technicolor does: weak interaction breaking occurs at the critical point in the logarithmic evolution of a coupling parameter. However, there is an important difference. For technicolor the coupling is a dimensionless gauge coupling, while for supersymmetry it is a scalar mass parameter. In supersymmetry there is the additional question of why the magnitude of the soft mass parameter is so small to begin with. It has been argued that this can also be understood as a dimensional transmutation which fixes the scale of supersymmetry breaking [22].

Suppose that a Higgs boson is found; can we tell if it comes from a supersymmetric theory? I think the answer is no: the Higgs itself and its couplings to ordinary particles cannot tell you whether the theory is supersymmetric. This will require discovering particles with interactions which fit those of the superpartners. However, it is worth pointing out that in the minimal low energy supergravity model, the lightest Higgs boson must be lighter than the Z [23].

Searches for superpartners are crucial in guiding our speculations on physics beyond the standard model. How do you recognize a superpartner when you

see one? This is a very non-trivial question. It is conventionally argued that missing energy is the prime signature for superpartner production. Although this is a very good signature for the minimal supersymmetric model, it is completely removed by quite minor charges in the theory, as I will discuss in the next section. You will only be sure that you have seen a superpartner when you measure its interactions and find that they have the form and magnitude dictated by supersymmetry. Assuming that the lightest superpartner (LSP) is stable and that missing energy is a good signature, some of the best limits on superpartner masses come from searches at e^+e^- and $\bar{p}p$ colliders for events with missing (transverse) energy. The process $e^+e^- \to \tilde{\gamma}\tilde{\gamma}\gamma$ [24, 25, 26] which proceeds via the diagram of fig. 6, leads to bounds on the selectron mass of about 60 GeV for a massless photino. In fig. 7 I give

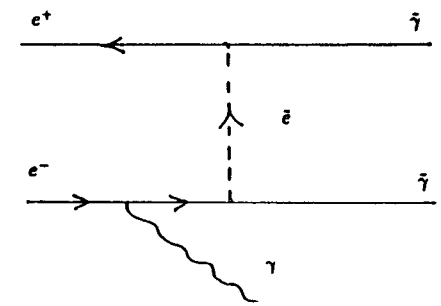
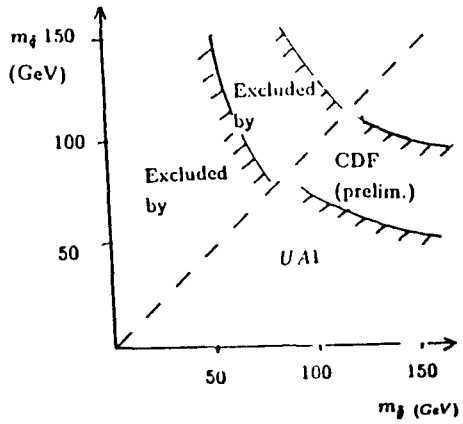


Diagram for $e^+e^- \to \tilde{\gamma}\tilde{\gamma}\gamma$
Figure 6

limits on $m_{\tilde{g}}$ and $m_{\tilde{q}}$ from UA1 [27] and from preliminary data of CDF [28]. It must be stressed that these limits are based on assumptions: they only apply if the LSP is neutral and long lived. Furthermore, the $\tilde{g}$ and $\tilde{q}$ limits assume that the LSP is produced directly in the decay, rather than having superpartner cascade decays degrade the missing energy [29].



Excluded regions of $\tilde{q}$ and $\tilde{g}$ masses from hadron colliders
(see text for assumptions).
Figure 7

Suppose that a Higgs boson is discovered, but no superpartners can be found; at what point should we admit that the scale of weak interaction breaking is not related to that of supersymmetry breaking? If $M_{susy} >> M_W$, then the quadratic divergences will drag the Higgs mass back up to M_{susy}, so that a light Higgs then requires a fine tune. In the minimal low energy supergravity theory the Higgs mass squared parameter is directly related to the soft supersymmetry breaking parameters, and the Higgs quartic interaction is just a gauge coupling. Hence the Higgs vev, which is basically the Z mass, is directly related to the spectrum of superpartners: $M_Z^2 = M_Z^2(a_i)$, where a_i are a set of parameters which include scalar squared masses, m_i^2. While this function is not particularly simple, its schematic form can be thought of as $m_i^2 - m_j^2 + \ldots$, so that heavy superpartners require a cancellation to maintain the Z light. A detailed study of the sensitivity of M_Z to changes in the parameters a_i has recently been performed [30]. The degree of fine tuning may be made quantitative by requiring

$$\left| \frac{a_i}{M_Z^2} \frac{\partial M_Z^2}{\partial a_i} \right| < \Delta \qquad (2.3)$$

for all i. The results are very important. Even for $\Delta = 10$ (*i.e.* a fine tuning of one part in ten) the lightest neutralino and the lightest chargino are both lighter than 350 GeV. Furthermore, because the natural parameters are squared masses, the superpartner spectrum only increases as $\Delta^{1/2}$ as the fine tuning is increased.

iii)The Higgs: supersymmetric or composite?

We know of only two ways to tame the quadratic divergences of the Higgs boson mass. They can be cutoff at some scale Λ_{uc} at which the composite Higgs ceases to be local object, or they can be cutoff at a scale M_{susy} at which the power law radiative divergences are cancelled by others involving superpartners.

It is worth pointing out that these two mechanisms are not mutually exclusive. A supersymmetric preon model [31] can have a low energy supersymmetry but explain M_W/M_p from unusual properties of supersymmetric condensates, rather than from the renormalization group scaling of the Higgs mass.

At the moment, supersymmetry appears to be the favored viewpoint. Apart from the technical reason that it is easier to do calculations with weakly coupled theories, I can offer two reasons for this. Supersymmetry incorporates by far the simplest description of flavor which is known: Yukawa couplings. Secondly, supersymmetric theories at the Planck scale allow for the inclusion of quantum gravity. On second thoughts, it is not obvious that this justifies the prominence that low energy supersymmetry has received: supersymmetric theories of quantum gravity do not require that one supersymmetry survives to low scales. Also, a strongly coupled TeV flavor sector might exist in a yet to be discovered, elegant form; or perhaps experiment will find a very complicated TeV flavor sector. There is no reason why it is the TeV scale at which everything should suddenly be simple and elegant.

II.4. A Strongly Coupled Higgs Sector

Possibly the simplest and most natural explanation of electroweak symmetry breakdown at a scale $<< M_p$ is given by the technicolor idea [4, 32]. Consider an $SU(2)_L$ doublet T_L and singlets T_R of techniquarks transforming as fundamentals under a new $SU(N)$ technicolor gauge force which gets strong at

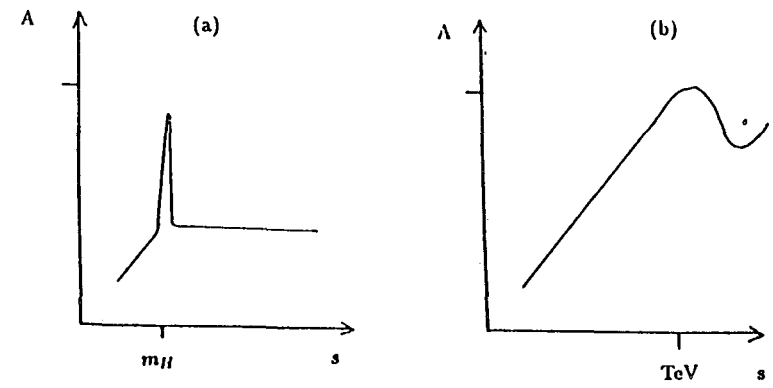
a scale $\Lambda_{TC} \simeq 0(TeV)$. The resulting condensate:

$$\langle \bar{T}_L T_R \rangle = 0(\Lambda_{TC}^3) \tag{2.4}$$

spontaneously breaks $SU(2)_L \times U(1)$ at the dynamically generated scale Λ_{TC}. This is similar to, and motivated by, the QCD quark condensate $\langle \bar{q}_L q_R \rangle = 0(\Lambda_{QCD}^3)$ which spontaneously breaks $SU(2)_A$ producing the PGBs $\pi^\pm$ and π^0. In the technicolor case the $SU(2)$ symmetry is gauged, so that the would-be Goldstone modes are eaten to become $W_L^\pm$ and Z_L. Furthermore, a custodial symmetry preserves the usual relation for M_W/M_Z.

This scaled up version of QCD breaks $SU(2)_L \times U(1)$ so elegantly that it leaves virtually no trace of itself beneath the techni-hadron states lying around a TeV; certainly there is no light Higgs boson. In more elaborate versions of the idea, there may be other light PGBs which are not eaten [32], but this is not guaranteed. It has also been argued that there might be light vector states which mix with the W and Z, but this is not guaranteed either [33].

It seems that beneath the scale of the techni-hadrons there is just a single guaranteed signature [7, 34]. The only light states of the techni-sector are the $W_L^\pm$ and Z_L and these have properties which differ from those in a model with a light Higgs. In QCD the pion self-interactions are due to residual strong forces and can be calculated at low energy using chiral perturbation theory. Similarly, the self-interactions of W_L and Z_L will manifest residual strong techni-forces, and these can also be calculated from arguments of broken symmetry [35]. In a model with a light Higgs the $W_L W_L \to W_L W_L$ amplitude is cut off by the Higgs contribution as shown in Fig. 8a). However, in a strongly coupled Higgs sector there is no light Higgs and the amplitude continues to grow until reaching the resonance region, as shown in Fig. 8b). Thus in high energy hadron



Behavior of $A(W_L W_L \to W_L W_L)$ for models with
(a) a light Higgs bosons (b) a strongly coupled Higgs sector.
Figure 8

collisions it should be possible to discriminate between a light Higgs and a strongly coupled Higgs sector by the number of events induced by the $W_L W_L \to W_L W_L$ sub-process. For example, with pp collisions at $\sqrt{s} = 40TeV$ and an integrated luminosity of $10^{40} cm^{-2}$, one finds a signature, with suitable cuts, in same sign dileptons ($\ell^-\ell^-, \ell^+\ell^+$ for $\ell = e, \mu$) of a few tens of events with a strongly coupled Higgs sector, whereas a model with a 100 GeV Higgs gives fewer than .1 such events [36].

III. The Flavor Sector.

The Yukawa interactions of the standard model are the only simple description of fermion masses and mixing that we have. They are so simple and successful that they may well be correct, in which case an understanding of flavor may be postponed to an arbitrarily high mass scale. In this section I make a few comments about recent alternative ideas on flavor, which either extend or eliminate the Yukawa interactions, and which lead to exciting new physics at the TeV scale.

III.1 Flavor in technicolor theories

i) General approach [37]

The standard model gauge interactions have a flavor symmetry which includes $(SU(3)\times U(1))^5$ with one $SU(3)\times U(1)$ factor for each of $q_L, u_R, d_R, \ell_L, e_R$. Consider an effective theory at the TeV scale with these fermions and the techniquarks T_L, T_{u_R}, T_{d_R}. Then operators which lead to the known fermion masses include

$$\mathcal{L} = \frac{1}{f^2}\left[\bar{q}_L\zeta_U u_R\overline{T}_L T_{u_R} + \bar{q}_L\zeta_D d_R\overline{T}_L T_{d_R} + \bar{\ell}_L\zeta_E e_R\overline{T}_L T_{d_R}\right] \tag{3.1}$$

where f is the scale of the new physics at which flavor is to be understood, and ζ_U, ζ_D and ζ_E are 3×3 matrices of dimensionless parameters which break the flavor symmetries of the gauge interactions and lead to quark and lepton masses $m_{U,D,E} = \zeta_{U,D,E}v^3/f^2$. Except for $(\zeta_U)_{33}$ the parameters in $\zeta_{U,D,E}$ are small.

A minimal assumption is that $\zeta_{U,D,E}$, with their particular transformation properties under the flavor groups, are the only flavor symmetry breaking parameters that survive beneath the scale f. This assumption has the advantage that it automatically incorporates the GIM mechanism [38]. In this case one can write down the most general low energy effective Lagrangian to be expected in a power series in $1/f^2$ and ζ, and study its physical implications [39]. While fermion masses are generated at $0(\zeta/f^2)$, the dominant corrections to the dimension four gauge interactions are $O(1/f^2)$: $\frac{1}{f^2}(\bar{\ell}_L\gamma^\mu\ell_L)^2, \frac{1}{f^2}(\bar{\ell}_L\gamma^\mu\ell_L)(\bar{q}_L\gamma_\mu q_L)$, ... providing corrections to all neutral current processes. $\Delta S = 1$ contributions to rare K decays occur at order ζ^2/f^2 and $\Delta S = 2$ operators appear at ζ^4/f^2.

If the quarks and techniquarks are composites with common constituents, then the above interactions are generated by the residual strong forces amongst the composites. In this case it is plausible that $\zeta_{U,D,E}f$ can be understood to be the preon mass matrices, as was the original idea of the composite technicolor standard models [37]. However, it seems to me that this approach is more widely applicable: it is the most general low energy effective theory of flavor which a theory incorporating technicolor could give, subject to the single assumption that the sole parameters describing flavor in the low energy theory are matrices $\zeta_{U,D,E}$ with the flavor transformation properties of the mass matrices. Any further generalization, to include more parameters with different flavor

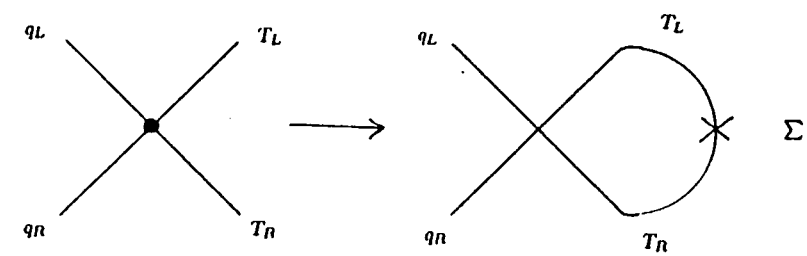
transformation properties, is not fruitful, since, unlike $\zeta_{U,D,E}$, they will not be determined by quark masses and mixing angles, so that the theory will not have any predictive power. It is an interesting question as to what could produce the $\zeta_{U,D,E}$ other than preon masses.

ii) Extended Technicolor Approach

In the extended technicolor (ETC) approach [40], quarks and techniquarks appear in the same irreducible representation of the ETC gauge group:

$$\begin{pmatrix} d' \\ s' \\ T \end{pmatrix} \tag{3.2}$$

This generates four fermion operators of the form $\bar{q}_L q_R\overline{T}_L T_R$. The technicolor condensate then induces fermion masses of sizes $m_f \simeq \Lambda_{TC}^3/M_{ETC}^2$, where M_{ETC} is the mass of the ETC gauge boson. This mechanism is illustrated in Figure 9. Note that the fermion mass hierarchy



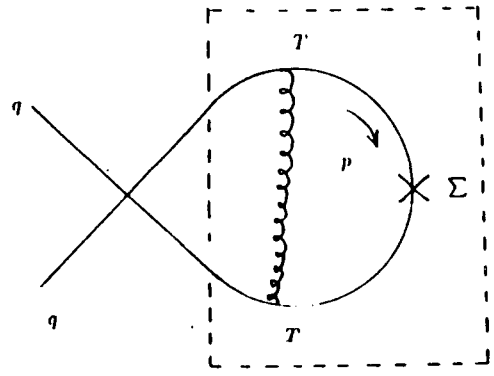
Fermion mass generation via technicolor condensates.

Figure 9

$MeV \lesssim m_f \lesssim 100GeV$ implies a hierarchy in ETC masses $3TeV \lesssim M \lesssim 1000TeV$. A typical difficulty in this approach is that of inducing large flavor changing neutral current effects. The representation of equation 3.2 will lead to $\bar{s}'d'\overline{T}T$ as needed to generate a Cabibbo angle, but it also leads to $(\bar{s}'d')^2$ which contributes

to Δm_K. There is a conflict between the two: for large values of M_{ETC} the Cabibbo angle is too small, while for small M_{ETC}, Δm_K is too large. Actually, while this is a generic problem it is not a theorem. It is possible to construct existence proof models which circumvent this difficulty, but only at the expense of a considerable expansion of the flavor sector [41].

Recently there has been considerable effort in studying whether the strong dynamics of the technicolor forces themselves may alleviate this flavor changing problem [42]. The idea is to calculate the leading log technigluon corrections to the mass generation diagram, as shown in Figure 10. The dashed box encloses



Technigluon radiative corrections to the technicolor condensate and the fermion mass.

Figure 10

the techniquark self energy Σ, and it is the behavior of Σ which is studied. Since there is a loop integral the corrections depend on the p dependence of $\alpha_{TC}(p)$. There are two classes of results: if α_{TC} is easily asymptotically free so that $\alpha_{TC}(p)$ drops quickly with large p then the corrections to Σ, and therefore to m_f, are small. On the other hand, suppose that $\alpha_{TC}(p)$ runs extremely slowly in the region $\Lambda_{TC} < \mu < M_I$. In this case m_f receives power law enhancements

$$m_f \simeq \frac{\Lambda_{TC}^3}{M_{ETC}^2}\left(\frac{M_I}{\Lambda_{TC}}\right)^p \tag{3.3}$$

where p is a positive number which is $0(1)$. Hence to obtain a given m_f, M_{ETC} can now be increased thus suppressing the $\Delta S = 2$ operators which clearly do not have any such technicolor enhancement.

More detailed recent studies of Σ conclude that suppression of $\Delta S = 2$ operator is indeed possible, but the restrictions on the behavior of $\alpha_{TC}(p)$ are quite tight [43]. These ideas have become known as "walking", or "stagnant" technicolor. Models incorporating these ideas have been built [44], but they leave open the question of what breaks the ETC gauge group. It is important to remember that we do not have a complete "standard" model of dynamical electroweak symmetry breaking.

III.2 Supersymmetric Flavor

The standard viewpoint is that models of low energy supersymmetry should have the same number of couplings (not including those which break SUSY) as the standard model: each has three gauge couplings and three Yukawa coupling matrices. While the standard model has a quartic scalar interaction the supersymmetric model has a mass parameter coupling the two Higgs doublets. Thus from the viewpoint of flavor apparently there is no change with the addition of supersymmetry. This viewpoint is a profound mistake, for two reasons.

While supersymmetry helps with the gauge hierarchy problem it makes negative progress elsewhere. In the standard model with 15-plet families we understand why baryon and lepton number violation are suppressed. There are no gauge invariant renormalizable operators which violate B or L which can be constructed. Such violations, if they occur at all, occur at dimension 5, for L, and dimension 6, for B. In supersymmetry this simple understanding is lost. Indeed, the most straightforward expectation is that the proton lifetime is comparable to that of the kaon. To make the theory sensible a discrete symmetry is imposed to remove the unwanted interactions. This is very ugly. However, it is a clue to further flavor physics. This discrete symmetry must come form somewhere. From the viewpoint of superstrings it is interesting to consider compactification physics for the origin of flavor physics and of these discrete symmetries. From the viewpoint of TeV scale physics, it is interesting to study alternative discrete symmetries to the one that is usually chosen. I will do that in this sub-section.

The second error in thinking that supersymmetry does not have anything new to offer for flavor physics at a TeV lies in the enhanced scalar sector of supersymmetric theories. Soft scalar operators can break chiral symmetries; it may not be necessary for some Yukawa couplings to be present [45, 46]. Also there is the possibilities that sneutrinos can acquire vacuum expectation values. This certainly enlarges the physics signatures of flavor; it also introduces a new possibility for fermion masses [47]. These new aspects of the flavor of supersymmetric models will be discussed in the next sub-section.

(A) With Yukawa Couplings.

I would like to mention the flavor physics which can occur even in supersymmetric models of minimal field content when lepton or baryon number violation is not excluded by hand. I think this physics is very important since it causes the LSP to decay visibly; the signature of missing energy is lost, and the assumptions underlying most searches for supersymmetry are not valid.

I will call a supersymmetric model minimal if it has the fewest number of particles possible. These particles are those of the standard model, an extra Higgs doublet and all the superpartners. Further constraining the model to have the fewest Yukawa-type interactions consistent with the known fermion masses produces the standard minimal model. It has the property that the LSP is stable. It is crucial to realize that this further constraint is non-trivial: it excludes, by fiat, those gauge invariant Yukawa interactions which allow the LSP to decay. These other couplings fall into two types; those which violate lepton number ($\ell\ell\bar{e}, q\bar{d}\ell$) and those which violate baryon number ($\bar{u}\bar{d}\bar{d}$).

Here q and ℓ are superfields for $SU(2)$ doublet quarks and leptons, while $\bar{u}, \bar{d}$ and $\bar{e}$ are superfields for the $SU(2)$ singlet quarks and leptons. Yukawa interactions arise when two of the three fields are taken to be fermions and the third a scalar, $\ell\bar{e}\tilde{\ell}$ for example, where we use the same symbol to refer to a field of the standard model and its superfield, and a tilde over that symbol refers to the superpartner.

Why should these possible couplings be absent? There is no firm theoretical answer. The four logical possibilities are shown in table 1. Most theorists might

Table 1. The four supersymmetric models with minimal field content

	Excluded	Standard	$\Delta L \neq 0$	$\Delta B \neq 0$
$\ell\ell\bar{e}$	✓	×	✓	×
$q\bar{d}\ell$	✓	×	✓	×
$\bar{u}\bar{d}\bar{d}$	✓	×	×	✓

have opted for the first column since all gauge invariant couplings are allowed and this seems the most natural. However, this case is excluded because the proton decays with a weak decay rate. The next simplest version is to assume that neither type of coupling is present. This produces the standard minimal model shown in column 2. One motivation for doing this is that grand unification treats quarks and leptons on equal footings and hence suggests that either both or neither type of new coupling should appear. While this argument is correct for certain simple models, it is certainly not always true.

It is high time the remaining two possibilities were taken seriously. In the "$\Delta L \neq 0$" model only lepton number is broken and in the "$\Delta B \neq 0$" model only baryon number. This latter case maintains proton stability since it must decay into an odd number of leptons (we take all superpartners heavier than the proton) and this is forbidden by lepton-number conservation. These models are worth pursuing because their experimental signatures are spectacular and are quite unlike the standard missing energy signatures for supersymmetry.

Work on these models include studies of rare K, π, ν decays [48], neutron oscillations and heavy nucleus decay [49] and baryogenesis [50]. Here I will mention a few of the exotic signatures to be expected at e^+e^- and hadron colliders [51].

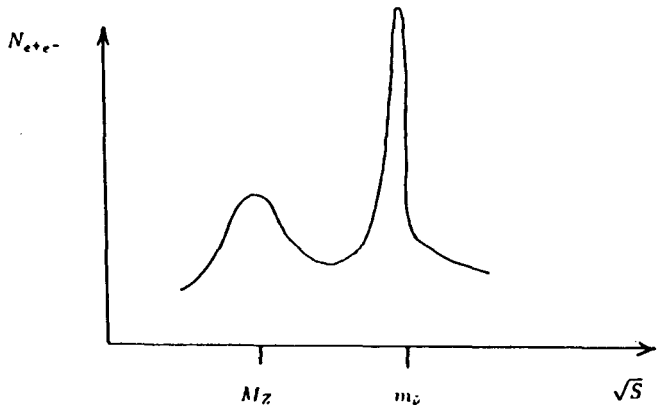
In the $\Delta L \neq 0$ model the LSP decays into leptons. I will assume the LSP is $\tilde{\gamma}$ or $\tilde{\nu}$, although there is no longer any cosmological or astrophysical argument that this should be so. Typical LSP decays would then be $\tilde{\gamma} \to \bar{\mu}e\nu$ or $\tilde{\nu} \to \mu e$. Photino pair production in e^+e^- no longer requires a radiated photon to make it visible. The event rate $N_{\tilde{\gamma}\tilde{\gamma}}$ is very large

$$N_{\tilde{\gamma}\tilde{\gamma}} \simeq 100 \left(\frac{100GeV}{m_{\tilde{e}}}\right)^4 \frac{s}{(40GeV)^2} \left(\frac{\mathcal{L}}{100pb^{-1}}\right) \beta^3 \tag{3.4}$$

This suggests that we have yet to get above the threshold; when we do there

will be events such as $e^+e^- \to \mu^+\mu^+e^-e^-$ together with some missing energy.

Most dramatic of all is the $\tilde{\nu}$ peak shown in Figure 11. The event rate to e^+e^- final states on the $\tilde{\nu}$ peak could be



event rates at the Z and sneutrino peaks.

Figure 11

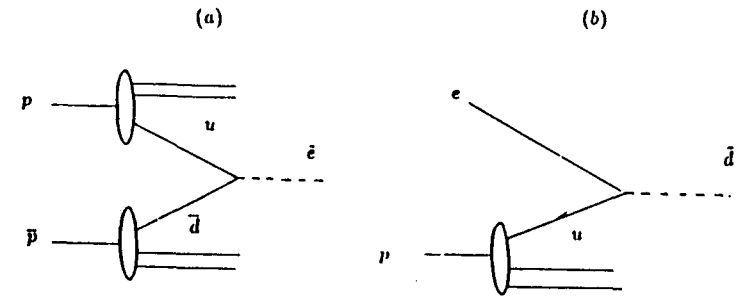
much larger than on the Z peak for comparables luminosities:

$$\frac{N_{e^+e^-}(m_{\tilde{\nu}})}{N_{e^+e^-}(M_Z)} \simeq 100\left(\frac{100GeV}{m_{\tilde{\nu}}}\right)\left(\frac{250MeV}{\Delta E}\right)\left(\frac{\lambda}{.2}\right)^2 \qquad (3.5)$$

where λ is the size of the $\ell\ell\tilde{e}$ vertex and the resonance will probably be narrower than the beam spread ΔE.

In hadronic collisions the best signature for the $\Delta L \neq 0$ model is isolated multicharged leptons. For example, squark or gluino pair production and subsequent decay produces two photinos. When these decay, events with jets + 4 isolated charged leptons will be produced. Missing energy will not be a very good signature unless the photino lives long enough to escape from the detector. Perhaps the most exciting possibility is that the $\Delta L \neq 0$ vertex will have a strength such that photino decays within the active volume of the detector will occur.

If the $q\bar{d}\ell$ operator dominates, and if the sneutrino is light enough, there could be events with up to four jets from $Z \to \tilde{\nu}\tilde{\nu}^* \to (\bar{q}q)^2$. This branching mode could be as large as 1%. This operator also allows for Drell–Yan fusion of a $\tilde{e}$ in $\bar{p}p$ collisions or a $\tilde{q}$ in ep collisions, as shown by the diagrams in Fig. 12. In the former case the signature comes from



Drell–Yan fusion mechanism for (a) sleptons in $\bar{p}p$ collisions, and (b) squarks in ep collisions. Each vertex violates lepton number.

Figure 12

$\tilde{e} \to q\bar{q}$ (two jet bump) or $\tilde{e} \to e\tilde{\gamma}, \tilde{\gamma} \to q\bar{q}\ell$ (isolated lepton) and in the latter case from $\tilde{d} \to eu$ or $\tilde{d} \to d\tilde{\gamma}, \tilde{\gamma} \to q\bar{q}\ell$.

The $\Delta B \neq 0$ model is not quite so easy to discover as it does not give leptonic signatures. However, it is very exciting since it offers the possibility that baryon number violation could be observed at e^+e^- colliders. For example, imagine $e^+e^- \to \tilde{\gamma}\tilde{\gamma}$ followed by each photon decaying as: $\tilde{\gamma} \to bcs$. In this case there would be a multi jet event rate given by equation (3.4). Furthermore, one would expect to see just as many events with dilambdas $\Lambda\Lambda$ or $\overline{\Lambda\Lambda}$ as with the zero baryon number $\Lambda\bar{\Lambda}$ final state.

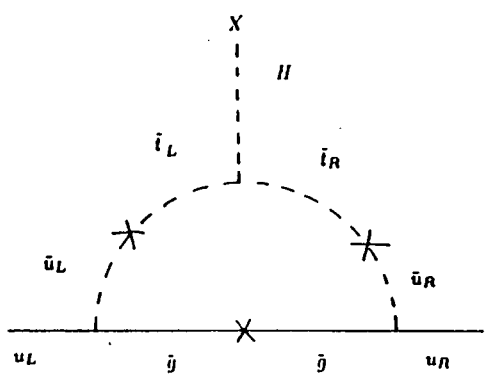
(B) Without Yukawa Coupling

In passing from ordinary to supersymmetric gauge theories, there are so many new couplings introduced that it seems a shame not to try to use them

to understand the flavor problem. This is seldom done; usually supersymmetric models have the same set of Yukawa coupling matrices as the standard model. However, there are alternatives.

From the viewpoint of alternative origins for quark and lepton masses at the TeV scale, perhaps the most important feature of the fermion spectrum is the heaviness of the top. The top is the only "known" fermion with mass $O(M_W)$. Without a strongly coupled flavor sector the top quark mass must occur at tree level. This can occur in two ways in supersymmetry: either there is a Yukawa coupling in the superpotential or the top quark is eaten; it marries a colored gaugino in the supersymmetric extension of the Higgs effect.

One interesting possibility that supersymmetry allows is for the third generation to have tree level masses, while those of the first two generations are radiative [45, 46]. At first sight this is impossible: if u, d, c, s, e, μ do not appear in the superpotential then how are their chiral symmetries to be broken? The answer is that because the chiral symmetries also act on the scalar partners, they can be broken by soft scalar mass terms such as $\tilde{u}^*\tilde{t}, \tilde{d}^*\tilde{b}, \tilde{e}^*\tilde{\tau}$. The light fermion masses are then generated from diagrams such as in Figure 13. Notice that most of the information about



Radiative contribution to the up quark mass. The chiral symmetry of the up quark is linked to that of the top quark via the off-diagonal squark masses.

Figure 13

flavor is now carried by scalar mass terms. Where do these come from? They could come from the remnants of a heavy sector or of a spontaneously broken family symmetry at scale $M_{f\ell}$. In either case, there is no restriction an $M_{f\ell}$; it could be anywhere from a TeV to M_p. If flavor does reside in the soft operators, there should be observable consequences: new contributions to heavy flavor decays and to neutral meson mixing. Indeed the simplest scheme as first proposed [45] is now excluded as it predicted 100% mixing in the B_d system coming from the $\tilde{b}^*\tilde{d}$ operator necessary for m_d. The simplest such model is now with four generations and the large flavor mixing being $\tilde{b}'^*\tilde{d}$.

Most exciting of all is the possibility that none of the fermion masses arise from Yukawa interactions in the superpotential. The top quark mass must then come from the gauge Yukawa interactions. This is impossible with a 15-plet family, but with a 16-plet family having a TeV gauge group $SU(4) \times SU(2)_L \times SU(2)_R$ it is possible for the top quark to be eaten by the color triplet gauginos of the Pati-Salam $SU(4)$ group [47]. This is by far the minimal extension of the standard representation and gauge group which allow a quark to be eaten; it explains why there is only one flavor of quark with mass $0(M_W)$ and why it has charge 2/3. These successes are sufficiently striking that it is worth mentioning the wealth of experimental consequences which immediately follow from this unique minimal mechanism for eating the top.

Supersymmetry and $SU(4)$ breaking are both broken at the same scale, which should not be much more than a TeV. Rare K_L decays ($K_L \to \mu e, \mu\mu, ee, \pi^0 ee$) then require an usual multiplet assignment under the $SU(4) \times SU(2)_L \times SU(2)_R$ gauge group. There are three $\psi(4,2,1)$:

$$\begin{pmatrix} u & d \\ \nu_\mu & \mu \end{pmatrix} \begin{pmatrix} c & s \\ \nu_\tau & \tau \end{pmatrix} \begin{pmatrix} t & b \\ \nu_E & E \end{pmatrix} \tag{3.6}$$

and similarly three $\psi^c(\bar{4},1,2)$ which have the same pattern of anti-fermion fields. The top is eaten by $< \tilde{\nu}_E > \neq 0$ as is the E. The second and third lepton generations couple to the first and second quark generations so that $K_L \to \mu e$ is absent. This introduces the interesting signature of $t \to K\mu$. There is an obvious question: where is the electron! It has an usual assignment as an $SU(4)$

singlet. An L_τ parity forbids $\mu \to e\gamma, \mu \to ee$ as well as $K_L \to \mu e$, but the charge assignments require the introduction of a new $U(1)$: there is a Z' of a few hundred GeV which couples to $T_{3R} + (B + L_e - L_\mu - L_\tau)/2$! The mass matrices for down quarks and charged leptons should be proportional to each other; otherwise relative mixing angles will induce $K_L \to \mu\mu$. This gives a mass relation

$$\frac{m_d}{m_s} = \frac{m_\mu}{m_\tau} \tag{3.7}$$

which is highly successful.

While this is an exciting possibility for the origin of the top quark mass, whether or not all the other fermion masses can arise as radiative corrections is an open and challenging question. The issues are plentiful: can the desired soft operators be generated from physics at a higher scale in a compelling way; m_b is also quite large, can it occur radiatively; what about neutrino masses, in left-right symmetric models they are typically problematic, can ν_R be eaten to implement a see-saw mechanism? Perhaps the biggest question of all is how the alignment of the sneutrino vacuum will take place. These issues are not easy to deal with, they are in some ways similar to the difficult issues in ETC. However, it is a new avenue of exploration, and I see some advantages over the ETC scheme: we know exactly how gauge symmetries get broken and we can calculate the breaking scale; we just renormalization group scale soft mass squared parameters until they go negative. The top quark mass is already dealt with by the eating mechanism within $SU(4)$; in ETC the top implies a very low mass ETC boson. Finally, the origin of the flavor physics of the soft operators can be postponed to a high energy scale, we do not have the flavor-changing difficulties which have led to contortions such as the "stagnant"TC models mentioned earlier.

IV. Conclusions.

In this talk I have argued that there are three questions about physics at the TeV scale which we must answer:

1) Does a light Higgs exist? We could answer this by a direct search or by observing a continued rise in the $W_L W_L \to W_L W_L$ amplitude. If there is a light Higgs, we can be confident that some form of the perturbative Higgs mechanism

is at work in $SU(2) \times U(1)$ breaking. If we find a Higgs the next question is: what is it? Is it composite or supersymmetric, for example, and how does it couple to fermions?

2) Is there a TeV scale supersymmetry? We should explore both conventional, missing energy signatures, and exotic B and L violating signatures. A TeV supersymmetric world would keep us very busy measuring masses and couplings, and after verifying the supersymmetric nature of some of the couplings, the most important further step would be to discover the form of supersymmetry breaking.

3) What is the TeV flavor sector? If there is a Higgs, it could be that this sector will be just the usual Yukawa interactions, but this need not be the case. There could be a wealth of flavor physics in new gauge interactions, sneutrino vevs and in soft supersymmetry breaking operators. If there is no light Higgs, things are bound to be interesting since there certainly will not be any Yukawas. Presumably flavor will reside in four fermion operators, but are these generated by weakly coupled gauge bosons, ETC for example, or by strong forces which are also responsible for another level of substructure?

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