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by

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PHENOMENOLOGICAL ASPECTS OF UNIFIED THEORIES

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After some preliminary observations concerning attempts to go beyond the standard model, I briefly discuss two new phenomena of recent interest: the 5^{th} force and variant axions. The former, for its elucidation, will require further gravitational experiments, but I conclude that variant axions are now definitly ruled out experimentally. Various aspects of superstring phenomenology are then addressed, including some of the generic predictions of superstrings and some of its generic problems. In particular, I discuss some of the phenomenological consequences of having an extra Z^o boson and the circumstances under which this excitation is a genuine prediction of superstrings. Since it is likely that a more reliable relic of superstrings will be provided by the presence of superpartners at low energy ($\lesssim$ TeV), I discuss some of the bounds for squarks and gluinos obtained at the SppS collider and the expectations for their production at the Tevatron. As a final topic, I touch upon some of the consequences that would result from having the Fermi scale arise from an underlying theory. Some aspects of the composite Higgs model and of the strongly coupled standard model are briefly reviewed.

I. PRELIMINARY OBSERVATIONS

The standard SU(3)xSU(2)xU(1) model of the strong and electroweak interactions works exceedingly well phenomenologically. This has been amply demonstrated again at this conference in the review talks by Altarelli[1] and Scott[2]. Yet theorists remain unhappy, even in the face of success, because they do not really understand the deep reasons that lie behind certain structural aspects of the standard model. Putting it rather succintly, theorists would like to understand three main points:

i) Why are these the forces we see, and where does gravity fit into this picture?

ii) Why does the matter we see, the quarks and leptons, have the peculiar transformation properties it has in the standard model?

iii) What fixes the dynamics which generates all masses?

This last point is associated with a further problem, that of the hierarchy of scales. In the standard model all masses are proportional to the scale of the SU(2)xU(1) breakdown, the Fermi scale: $\Lambda_F = (\sqrt{2}G_F)^{-1/2} \simeq$ 250 GeV. However, the constants of proportionality for fermions and the Higgs boson are unknown. Only for gauge bosons are their masses predictable, since they are related to Λ_F via the gauge coupling constant. Schematically one has

$$m \sim \left\{ \begin{array}{c} \Gamma_f \\ e \\ \lambda^{1/2} \end{array} \right\} \Lambda_F \qquad \begin{array}{l} \text{Fermions} \\ W^\pm \\ \text{Higgs} \end{array} \qquad (I.1)$$

The hierarchy problem is two-fold. There is a "small" hierarchy problem connected with what physics forces the Yukawa couplings Γ_f to be so varied, so as to give the rather spread out quark and lepton mass spectrum we observe .However, the real hierarchy problem is why the Fermi scale itself is so different from the scale associated with gravitational phenomena, the Planck scale: $M_P = (G_N)^{-1/2} \sim 10^{19}$ GeV, where G_N is the Newtonian coupling. The Fermi scale, in the standard model, is not a dynamical scale like that associated with the running of the strong coupling constant, Λ_{QCD}. Thus its only natural value, in a world where the high energy cutoff is provided by gravity, is the Planck scale. Why is then $\Lambda_F \ll M_P$?

The theoretical landscape, in which one attempts to answer the open deep questions of the standard model, is quite vast. As Fig I.1 shows, one can adopt a bottom up approach, to try to answer the questions of forces, matter and masses, by focusing on issues connected with the dynamics which generates the Fermi scale. On the other end of the scale, however, one can adopt a top down approach and start with physics at the Planck scale; connected with gravity and superstrings, and attack in this way the standard model fundamental problems

The experimental landscape, shown in Fig I.2, is necessarily much more restricted. The next generation of accelerators, SLC, LEP, Tevatron and HERA will probe the 100 GeV region deeply, but we will have to wait for the SSC or a possible TeV e^+e^- collider (2TLC) to get into the TeV region. Non accelerator experiments (NAC) can provide crucial information about very high energy scales, but in rather restricted regions. Thus, most of the territory will never get direct experimental probing.

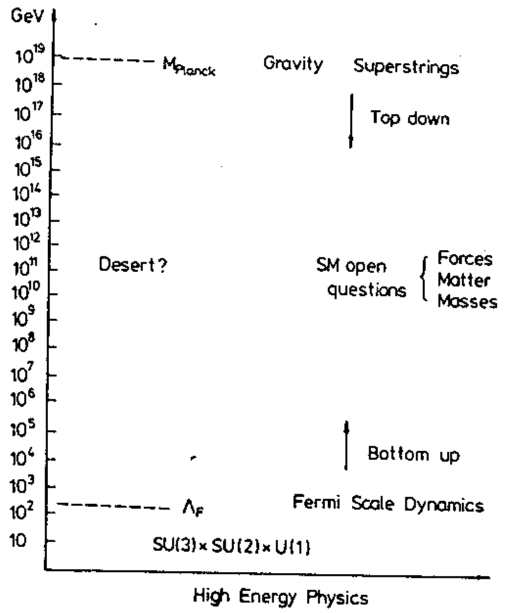

Fig I.1 Theoretical landscape to address the standard model open questions

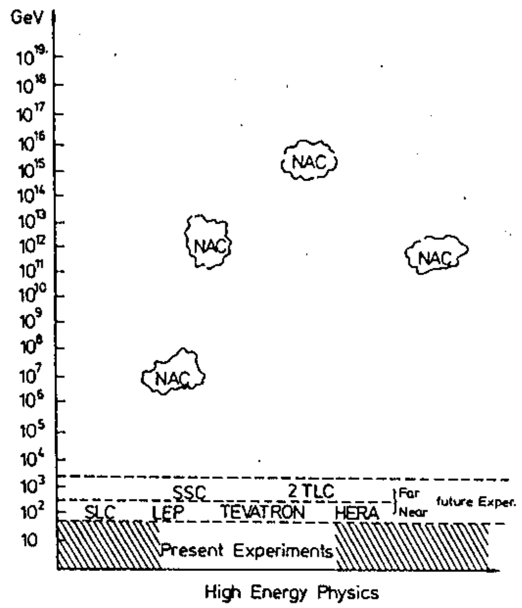

Fig I.2 Experimentally accessible landscape

Given this state of affairs, it is obvious that theories beyond the standard model, to make contact with reality, must predict some subTeV phenomena and/or some phenomena which will be accessible to non accelerator experiments. After all, physics is an experimental science! It is here that phenomenologists can play a useful role in trying to translate theoretical predictions of Planck scale physics, or coming out of Fermi scale dynamics, into possible signals of subTeV phenomena, which might be experimentally detectable. Before discussing some of the speculative novel effects studied by the Desert Trekkers and Moose Herders of Fig I.3 (which, if detected in future experiments, will point to physics beyond the standard model), let me first discuss two phenomena which already could signal new physics.

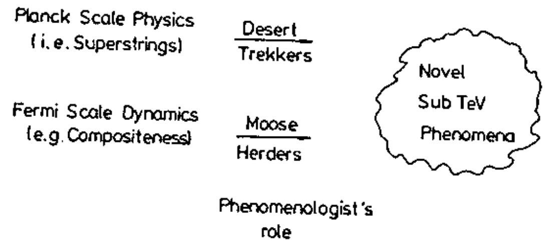

Fig I.3 Phenomenologist's role in modern day high energy physics

II. NEW PHENOMENA

a) The Fifth Force

Considerable excitement, and a certain amount of controversy and confusion, has surrounded a recent reanalysis by Fischbach, Sudarsky, Szafer, Talmadge and Aronson[3] (FSSTA) of the classic Eötvos, Pekár and Fekete[4] experiment on the equality of gravitational and inertial mass. What FSSTA found was a correlation in the data of the Eötvos experiment between the discrepancy in the torque measurements ΔK and the relative baryon number in the sample used $\Delta(B/\mu)$. This correlation is shown in Fig II.1, which is a slightly updated version[5] of the figure presented in the original FSSTA paper. Bearing in mind that the data plotted was obtained almost three quarters of a century ago, and so is subject to difficult to quantify systematic uncertainties, the correlation found should be treated with some skepticism. However, optically, Fig II.1 looks quite impressive. FSSTA took the correlation they found seriously and suggested that it was an indication for a new, extremely weak, force in nature which coupled to baryon number or strong hypercharge Y - a fifth force. Since this force, if it existed, would give rise to a typical Yukawa potential between two bodies:

$$V_5 = f^2 \frac{B_1 B_2}{r} e^{-r/\lambda} \qquad \text{(II.1)}$$

with B_i here standing for either baryon number of hypercharge and λ being the range of the force, FSSTA tried to relate the Eötvos reanalysis to discrepancies found in geophy-

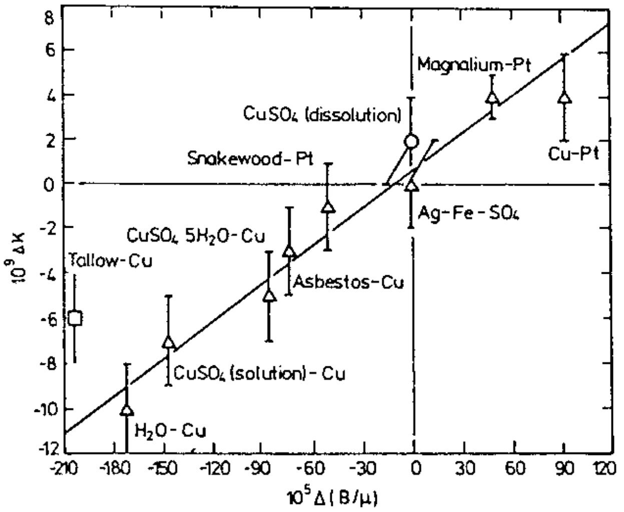

Fig II.1 Updated plot (from Ref 5) of the measured values of ΔK in the Eötvos experiment plotted against $\Delta(B/\mu) = B_1/\mu_1 - B_2/\mu_2$. Here μ_i is the mass of the sample expressed in terms of m(,H') and B_i is its baryon number.

sical determinations of the Newtonian constant G_N. It has been known for some time[6] that the value of G_N, inferred from measurements of the gravitational acceleration g in mines and boreholes, is about 1% larger than that determined from the classical Cavendish type laboratory measurement[7]: $G_N = 6.6720 \pm 0.0041 \times 10^{-11} m^3 kg^{-1} sec^{-2}$. For example, a particularly careful recent investigation in the Hilton mine in Queensland gives a value[8] (II.2)

$$(G_N)^{\text{Hilton Mine}} = (6.720 \pm 0.002 \pm 0.024) \times 10^{-11} m^3 kg^{-1} sec^{-2}$$

where the last error is an estimate of the possible systematic error, arising from a lack of precision in the density determination of the area surrounding the mine. This discrepancy (which, however, is only at the 2σ level) could be attributed to a non gravitational addition to the potential, so that

$$V = -G_N \frac{m_1 m_2}{r} \left[1 + \alpha e^{-r/\lambda} \right] \qquad \text{(II.3)}$$

where the last term above comes from the fifth force. The geophysical data is actually only sensitive to the combination of $\alpha\lambda$ and the discrepancy given in (II.2) gives the bounds[8] $0.004 \leq -\alpha\lambda \leq 10$m. However, satelite data[9] severely restrict large values of λ and one deduces[8]

$$0.035 \lesssim -\alpha \lesssim 0.15 \qquad \text{(II.4)}$$
$$1 \leq \lambda \leq 10^3 m$$

These results imply a value $f^2/4\pi \sim 10^{-40}$ for the fifth force coupling and a superlight mass, $2\times10^{-10} \leq m_5 \leq 2\times10^{-7}$eV, for the boson associated with this force.

FSSTA tried to connect the parameters inferred from the geological anomalies with the slope of Fig II.1. However, in fact one cannot really establish a definite correlation between the two phenomena. One can show[10] that, if the fifth force exists, the torque that it produces in the Eötvos experiment is mainly a function of the local topography. This is easy to understand, since there is no residual torque if the fifth force is parallel to the direction of the effective gravity (the direction of the vector sum of the gravitational and of the centrifugal acceleration). Thus average matter beneath the experiment is of little importance for the torque on the wire. But the results are crucially dependent on what nearby large buildings (and basements) existed at the time of the Eötvos experiment! Thus it seems pointless to compare the predictions of the fifth force - with the parameters fixed by the geological observations - with the slope of Fig II.1, although an attempt to do this is presented in Ref 5. Rather, the message one draws from these considerations is that the Eötvos experiment should be repeated, under careful controlled conditions, near the side of a mountain, so as to maximize the effect.

Although FSSTA put forth the possibility that the fifth force could couple to hypercharge, this suggestion can be ruled out by using the strong bound[11]

$$B(K^+ \to \pi^+ \text{Nothing}) < 3.8 \times 10^{-8} \qquad \text{(II.5)}$$

obtained some time ago at KEK. If the fifth force coupled to hypercharge a K^+ could decay into a π^+ emitting a hyperphoton - the vector boson associated with this force, of mass m_5. Since the coupling f is so small, one might think that the rate for this process would be negligible. However, as pointed out by Weinberg[12] more than 20 years ago, the fact that hypercharge is not conserved gives a contribution from longitudinally polarized hyperphotons proportional to f/m_5. Thus the process $K^+ \to \pi^+\gamma_5$ is not negligible and one can bound this ratio. This has been studied recently by a number of authors[13] and a typical bound, taken from the paper of Suzuki[13] is

$$\frac{f^2}{4\pi m_5^2} < 7 \times 10^{-26}\ (\mathrm{eV})^{-2} \qquad (\mathrm{II}.6)$$

This is the conflict with the values of f and m_5 given earlier. However, the bound is trivially avoided by supposing that the fifth force couples to baryon number only. In fact, this is much more reasonable, since strong hypercharge has only a meaning neglecting weak interactions.

My conclusions on the fifth force are twofold:

i) The whole subject of departures from gravity is very interesting and FSSTA should be given credit for having stimulated a variety of new experiments of the Eötvos and Galilei type, whose results should be soon forthcoming. However, until these results are in, very little can be settled. It remains an open question if there is a connection between the residual torque correlation, obtained by FSSTA, and the geological anomaly and indeed if either or both of these phenomena are real.

ii) Theoretical attempts to fit the fifth force in a grand picture[14], although useful and clever exercises, are probably premature.

b) Variant Axions

A second phenomena which elicited a great deal of attention this year was the sharp positron peaks[15] and the correlated e^+e^- signals seen in heavy ion collisions at GSI. As Fig II.2 shows, the spectrum of positrons in U-Cm collisions exhibits a narrow line at $T_{e^+} \sim 350$ KeV, above the continuum spectrum

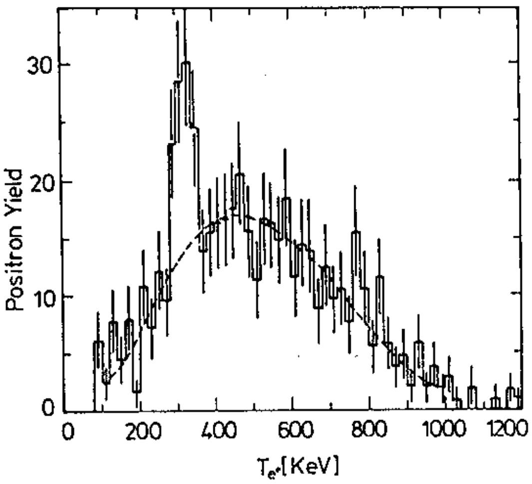

Fig II.2 Positron energy spectrum for U-Cm collisions at the Coulomb barrier. From Schweppe et al, Ref 15.

expected from spontaneous positron creation at the Coulomb barrier[17]. Similar sharp lines are seen[18] in other heavy ion collisions, for sufficiently large total Z: $Z_1 + Z_2 > 180$. What made this phenomena particularly intriguing was the report[16] that the positron peaks were correlated with analogous peaks in the spectrum of emitted electrons. These correlated signals are shown in Fig II.3.

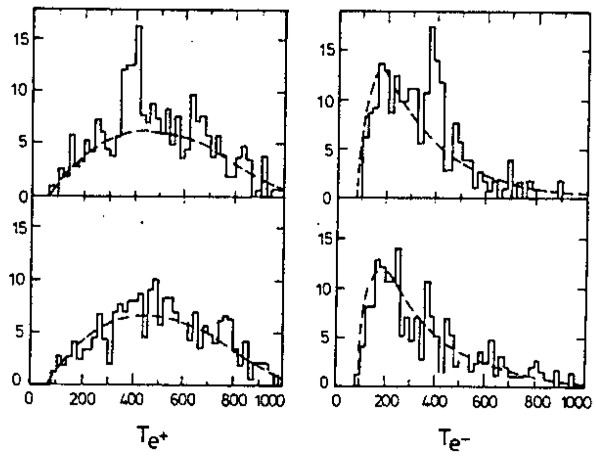

Fig II.3 Peaks in the positron (electron) spectum for electrons (positrons) with $340 < T_e < 420$ KeV. The bottom curves show these spectrums when the energy cuts correspond to the adjacent bins. From Ref 16.

These observations have a "trivial" kinematical explanation if in the heavy ion collision one produces, essentially at rest, a particle of mass $M_a \approx 1.7$ MeV, which decays rapidly into e^+e^- pairs. Although no convincing mechanism has been invented to justify why such a particle should be produced at rest[19], the fact that M_a is so light naturally leads one to the speculation that this particle might be an axion. On the other hand, there is no way to associate a particle of mass as heavy as 1.7 MeV with the standard axion[20], since then it would have very enhanced couplings to either charm or bottom quarks, and so would be in violent conflict with the existing bounds on $\psi \to \gamma a$ or $\Upsilon \to \gamma a$[21]. It is, however, possible to construct variant axion models[22], where axions with masses above the e^+e^- threshold can exist, without being in direct contradiction with the quarkonia bounds.

Axions arise out of an attempt to solve the strong CP problem by imposing an additional global symmetry[20], $U(1)_{PQ}$, on the standard model Lagrangian. To achieve this it is necessary to have at least two Higgs doublets, Φ_1 and Φ_2, in the theory. Although the

$U(1)_{PQ}$ symmetry allows one to set to zero the overall coefficient of the CP violating $F_a^{\mu\nu}\tilde{F}_{a\mu\nu}$ term, where $F_a^{\mu\nu}$ is the gluon field strength, the fact that the fields Φ_i have non zero vacuum expectation value implies that this symmetry is spontaneously broken. The axion is the resultant Goldstone boson which, however, obtains a slight mass since the $U(1)_{PQ}$ symmetry is anomalous. For the standard axion model of Ref 20, one finds

$$M_a = \frac{m_\pi f_\pi}{\Lambda_F}\left\{\frac{m_u m_d}{(m_u + m_d)^2}\right\}^{1/2} N_F\left(x + \frac{1}{x}\right)$$

$$\simeq 25\ N_F\ \left(x + \frac{1}{x}\right)\ \mathrm{KeV} \qquad \text{(II.7)}$$

where N_f is the number of families and $x = \langle\Phi_2\rangle/\langle\Phi_1\rangle$ is the ratio of the Higgs vacuum expectation values. Clearly for $M_a \simeq 1.7$ MeV x or x^{-1} must be very large. Since the coupling of the standard axion to charge 2/3 quarks (charge-1/3 quarks) is proportional to $x(x^{-1})$, one of the branching ratios: $B(\psi \to a\gamma) \sim x^2$ or $B(\Upsilon \to a\gamma) \sim x^{-2}$ is predicted to be very large, in contradiction with experiment[21].

Variant axion models[22] treat quarks asymetrically under $U(1)_{PQ}$. Thus in the simplest model, for example, only the u quark and the electron have a coupling to the axion proportional to x, while all other fermion couplings are proportional to x^{-1}. If x is large, then <u>both</u> the ψ and Υ decays into γ axion will be suppressed. Furthermore, since the electron coupling is enhanced, the decay lifetime for the process $a \to e^+e^-$ will also be very short, making it possible for these kind of variant axions to escape some previous beam dump bounds, which assumed that axions were long lived.

Although variant axion models appeared for a while this year to be viable, and could be adduced as an "explanation" for the GSI positrons[22], their existence has proven ephemeral. Three distinct factors contributed to their demise:

i) Very recent experiments at GSI[24], although confirming the existence of the e^+e^- correlations, appear to see two distinct correlated e^+e^- signals whose origins, obviously, are difficult to reconcile with a single axion.

ii) New electron beam dump experiments, discussed by Davier[25] at this conference, are sensitive to the production and decay of variant axions. However, no signal for these excitations is seen in the data and all values for the ae^+e^- coupling, not in violation of g-2 bounds, are excluded[25].

iii) New experiments measuring axion deexcitations in hadronic transitions also rule out these models entirely[26].

Let me briefly discuss this last point here.

Variant axion models[22] are characterized by the number of quark doublets N_{PQ} which are active under the $U(1)_{PQ}$ transformation. This parameter replaces N_f in Eq(II.7) for the axion mass. Hence, for x large, and $M_a \simeq 1.7$ MeV, the combination $N_{PQ}x \simeq 70$ is fixed. In hadronic decays it is important to know to which extent the axion acts an an isovector or an isoscalar excitation. This is detailed in variant axion models by the axion's isovector and isoscalar mixing parameters, which depend again on x and N_{PQ}. One finds[27]

$$\lambda_3 \simeq \frac{x}{2}\left[1 - N_{PQ}\frac{(m_d - m_u)}{(m_d + m_u)}\right] \simeq \frac{x}{8}\,(4 - N_{PQ}) \qquad \text{(II.8a)}$$

$$\lambda_s \simeq \frac{x}{2}\,(1 - N_{PQ}) \qquad \text{(II.8b)}$$

Note that it is not possible for both of these parameters to be small since

$$\lambda_3 - \lambda_8 \simeq \frac{3}{8}\,(xN_{PQ}) \simeq 25 \qquad \text{(II.9)}$$

The strongest bound on λ_3 comes from a recent experiment at SIN on π^+ decay, where the rare process $\pi^+ \to e^+e^-e^+\nu_e$ was measured, thereby allowing a bound to be set on the process $\pi^+ \to ae^+\nu_e$[28]:

$$B(\pi^+ \to ae^+\nu_e) \leq (1\text{-}2)\times 10^{-10} \qquad \text{(II.10)}$$

where the range given above depends on the precise value of the axion lifetime. Theoretically one computes[26]

$$B(\pi^+ \to ae^+\nu_e) \simeq 3\times 10^{-9}(\lambda_3)^2 \qquad \text{(II.11)}$$

yielding $|\lambda_3| \leq 0.25$, in contradiction to Eq(II.8a) unless $N_{PQ} = 4$. However, such a value of N_{PQ} is not compatible with the recent result of an isoscalar, axion induced, nuclear deexcitation experiment in ^{10}B[29]. The axion to photon rate for the 3.59 MeV $2^+0 \to 3^+0$ rate is predicted to be[27)30]

$$\frac{\Gamma_a}{\Gamma_\gamma} \simeq 7.9\times 10^{-4}(\lambda_s)^2 \qquad \text{(II.12)}$$

while experimentally one has the bound[29]

$$\frac{\Gamma_a}{\Gamma_\gamma} \leq 7.2\times 10^{-3} \qquad \text{(II.13)}$$

so that $|\lambda_s| \lesssim 3$. For $N_{PQ} = 4$, however, one expects $\lambda_s \simeq 25$.

Given the above state of affairs, my conclusions are easily drawn:

i) The GSI phenomena, although very interesting by itself, has nothing to do with axions.

ii) Since neither variant or standard axions exist, if one insists in solving the strong CP problem by imposing a $U(1)_{PQ}$ symmetry, this symmetry must be broken at very high scales, leading to the invisible axion scenario[31].

III. SUPERSTRING INSPIRED PHENOMENOLOGY

Superstring theories were very much on the backburner at the time of the last International Conference of High Energy Physics in Leipzig. Indeed, I could find only one reference to them in the Conference proceedings, and that in the last paragraph of the summary talk of C. Callan[32]! However, after the publication of the anomaly cancellation paper of M. Green and J. Schwarz[33], superstrings have become a major theoretical industry. J. Schwarz[34], in this meeting, has thoroughly discussed the motivation and structure of these elegant theories. My job here is to summarize the status of the phenomenology which superstrings have inspired.

Superstrings are unfortunately not directly amenable to phenomenological study, since they are only consistent theories in a D = 10 dimensional space-time[34]. Therefore, if superstrings are to connect at all to reality, six of these ten dimensions must spontaneously compactify. Since these theories contain gravity, the scale associated with the compact dimensions is related to the Planck mass. Physics in four dimensional space-time, according to these theories, is set at the scale of compactification, $M_{comp} \sim M_P$, and depends on the geometry of the manifold K which compactified. At present there is no proof that this compactification actually takes place. Indeed, it is not even clear if one can expect this to happen for a unique space K, or if for a given superstring theory there is an infinity of such compact spaces.

These uncertainties notwithstanding, it has been argued by Candelas, Horowitz, Strominger and Witten[35] that particularly interesting manifolds for superstring compactification are provided by manifolds of SU(3) holonomy, which are known as Calabi Yau spaces[36]. Recall that the holonomy group is the group of all rotations generated when a vector, or spinor, is parallel transported around a closed loop in K. The fact that the holonomy group is not the full O(6) group, but only SU(3), means that there is at least one spinor which is not rotated under parallel transport. The existence of a covariantly constant spinor in the compact space assures that in the 4 dimensional theory an N = 1 supersymmetry is retained. Furthermore, one can show that the existence of the SU(3) holonomy allows a complex structure to exist in K, leading to a Kähler manifold whose metric is Ricci flat. This last circumstance is important because one can argue[37] that precisely such metrics will provide solutions to the string theory, since they preserve the conformal invariance of the two dimensional σ model on the string world sheet[38]. Since superstrings are endowed with a fixed gauge group[34], one must also specify the background gauge field in a consistent manner to guarantee these solutions. As Candelas et al show[35], the simplest consistent specification is to identify the background gauge fields, in an SU(3) subgroup of the superstring Yang Mills group, with the spin connection, ω, which transforms as an SU(3) matrix ($A = \omega$).

My discussion of superstring phenomenology will be restricted to the case when the compactification occurs in a Calabi-Yau space, in which the identification $A = \omega$ is made. Furthermore, I shall only consider the heterotic $E_8 \times E_8$ superstring[42], since the SO(32) superstring does not lead to realistic models[35]. There are other compactification possibilities which have been explored, including orbifolds[39] - which can be thought as limiting cases of Calabi-Yau spaces - and manifolds where the identification $A = \omega$ is not made[40]. Orbifolds are known to provide solutions to the string equations[39], but their phenomenology is largely unexplored. Manifolds with $A \neq \omega$, in general, can be shown not to provide solutions to the string equations, due to non perturbative effects[41]. These same effects, however, do not affect manifolds where $A = \omega$.

Let me begin by detailing the main features of the four dimensional theory which emerges from the $E_8 \times E_8$ superstring, after Calabi-Yau compactification[35]:

i) *The theory possesses an $N = 1$ supersymmetry.*

This is a very nice feature since supersymmetry allows naturally two different scales, like Λ_F and M_P, to coexist. So hierarchies are not unnatural, except that for the moment the only scale of the theory is that of the compact dimensions: $M_{comp} \sim M_P$.

ii) *The full $\mathcal{E}_8 \times \mathcal{E}_8$ gauge group of the ten dimensional theory is reduced to $\mathcal{G} = g \times \mathcal{E}_8$ where $g \subset \mathcal{E}_6$.*

One can understand this reduction by decomposing E_8 in terms of its maximal subgroup $E_6 \times SU(3)$. Since, in the compactification, the gauge fields associated with an SU(3) subgroup of E_8 were identified with the spin connection, clearly only an E_6 symmetry remains. In fact, if the manifold K is not simply connected, non trivial gauge configurations (Wilson loops) can be trapped in the manifold, even though the E_6 gauge field strength vanishes[43]. This can lead to a further breakdown of the E_6 group, with the flux trapping mechanism acting analogously to a breakdown induced by a Higgs field in the adjoint representation. It is important for phenomenology that at M_{comp} the remaining symmetry group g be not a GUT group, because otherwise one risks having a low energy group with unacceptable baryon number violations. Note that the second E_8 group is left untouched in the compactification and it provides a shadow matter world, which interacts only gravitationally with ordinary matter.

iii) *The matter representations are fixed by the properties of the manifold K.*

The four dimensional massless fermions which emerge correspond to chiral zero modes of the Dirac operator in K, and therefore have non trivial SU(3) properties in D = 10. Since the adjoint representation of E_8 decomposes under $SU(3) \times E_6$ as

$$248 = (1,78) + (3,27) + (\bar{3},\overline{27}) + (8,1) \tag{III.1}$$

one expects fermions in the 27 and $\overline{27}$ representations of E_6. Let me write these fermions as $n_f\, 27 + \delta(27 + \overline{27})$. The number n_f, which details the difference between 27's and $\overline{27}$'s, is entirely fixed by the topology of K[35] and it turns out to be one half of the Euler number χ of K:

$$n_f = \frac{1}{2}\,|\chi| \tag{III.2}$$

Since, as I will detail below, the ordinary quarks and leptons fit in the 27 of E_6, we see that n_f is just the number of families and this number is topologically determined. If there is flux breaking, the number δ of $(27 + \overline{27})$ fermions need not remain in complete E_6 representations, but the n_f 27's must remain, since their number is fixed by (III.2). The 27 of E_6 can be decomposed with respect to its SO(10) and SU(5) subgroups, respectively, as

$$27 = 16+10+1 = (10+\bar{5}+1) + (5+\bar{5}) + 1 \tag{III.3}$$

The ordinary quarks and leptons fit in the 16 of SO(10), including a right-handed neutrino. In addition there appears a new charge −1/3 quark and its antiquark plus a new electroweak doublet and its antidoublet, along with a total SU(5) and SO(10) singlet state. The new matter in the 27 is vector like, since ψ and ψ^c have conjugate transformation properties under the group. All the states in the 27 are displayed in Table III.1.

Table III.1 States in the 27 dimensional representation of E_6

SO(10)	SU(5)	states
16	10	$\begin{pmatrix} u \\ d \end{pmatrix}_L$ u^c_L e^c_L
	$\bar{5}$	d^c_L $\begin{pmatrix} \nu \\ e \end{pmatrix}_L$
	1	ν^c_L
10	5	g_L $\begin{pmatrix} N^c \\ E^c \end{pmatrix}_L$
	$\bar{5}$	g^c_L $\begin{pmatrix} N \\ E \end{pmatrix}_L$
1	1	S_L

iv) *The Yukawa couplings in the four dimensional theory, at the scale of M_{comp}, are in principle computable.*

This point is simple to understand. A gauge fermion fermion coupling in D = 10 contains, when all fields are expanded in terms of four dimensional fields, also a scalar fermion fermion coupling, with the scalar fields corresponding to components of the gauge fields in the compact directions. This correspondence is shown schematically in Fig III.1.

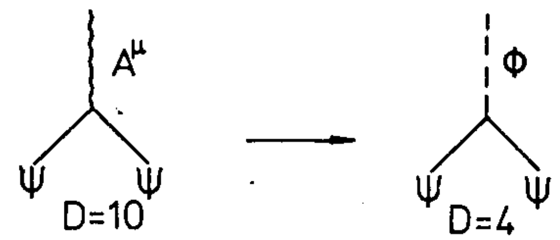

Fig III.1 Generation of Yukawa couplings on compactification.

It is impossible at the moment to really go ahead and compute all Yukawa couplings, although there have been some very interesting suggestions of how one might actually be able to achieve this by using topological considerations[44]. For some manifolds, however, even now one is able to infer, from the existence of certain discrete symmetries, that certain of the Yukawa couplings vanish[45].

In my opinion, properties i) - iv), along with the promise that superstrings may indeed provide one with a consistent quantum theory of gravity, are responsible for the extreme interest that these, otherwise rather remote, theories have stirred up in the high energy physics community. The above four results provide a reasonable starting point for answering the questions about hierarchy, forces, matter and mass dynamics, which I raised in my introductory remarks. However, one should remember that one must still transit from an energy scale of order $M_{comp} \sim M_P \sim 10^{19}$ GeV down to present accessible energies. What, if anything, survives of these beautiful patterns at 100 GeV?

The answer to this question depends a bit on how optimistic or pessimistic one is about the issue of supersymmetry breaking. Although the N=1 supersymmetry at the compactification scale makes hierarchies natural, it is clear that this supersymmetry must break down, allowing to split ordinary matter from its superpartners. If supersymmetry is to stabilize the low energy theory, so that a parameter like Λ_F is naturally 250 GeV and not 10^{19} GeV, the difference in mass between superpartners and ordinary matter must also be of this order of magnitude:

$$\tilde{m} - m \lesssim O(\Lambda_F) \sim \text{TeV} \qquad \text{(III.3)}$$

Superstring phenomenology is based on the idea that the shadow sector of the theory, that corresponding to the other E_8, triggers this breakdown and then transmits it to the visible sector. In this respect, this scenario is precisely analogous to that used a few years ago in the, so called, low energy N = 1 supergravity models[46]. There also supersymmetry breaking takes place in a hidden sector, which is coupled only gravitationally to the observable world. In these schemes, SU(2) x U(1) breaking occurs in a natural way as a radiative effect of supersymmetry breaking[47]. Thus the Fermi scale is intimately connected with the way one breaks supersymmetry.

There are, however, some important differences between the case of shadow sector supersymmetry breaking and that which occurs in the N = 1 supergravity models[48]. Even though supersymmetry is broken in both cases in a hidden sector, which is coupled to ordinary matter only gravitationally, for the superstring case the transmission of this breaking to the ordinary sector is not so straightforward[49]. For instance, it is usually assumed that the hidden sector breaking occurs through the formation of gluino condensates of the shadow E_8, $\langle\chi\chi\rangle$, and/or condensates involving the field strength $F_{\alpha\beta\gamma}$ of the second rank field $a_{\alpha\beta}$[50], crucial for the anomaly cancellation[33]. Although these condensates break the supersymmetry, leading to a gravitino mass $m_{3/2} \sim \langle\chi\chi\rangle/M_P^2$, their contributions cancel in the scalar potential. Thus, at tree level, even though supersymmetry is broken in the shadow sector, the observable world remains supersymmetric [49][50]. It turns out that also one loop radiative effects do not change this situation for the scalar fields[51], although gauginos can obtain a mass radiatively[52]. More generally, one can argue that quantum corrections involving heavy string modes[53] can be used to transmit the supersymmetry breaking from the hidden sector to the observable sector. However, one is then confronted with terms which destabilize the vacuum[54], unless one can find a mechanism to cancel the cosmological constant.

Clearly the present situation regarding supersymmetry breaking in superstring theories is unsatisfactory. In these circumstances a pessimist would argue that it is impossible to extrapolate down from the compactification scale, since there is no way to reliably generate any other scales in the theory, including the Fermi scale. An optimist, on the other hand, would argue that, in time, the matter of supersymmetry breaking will be resolved and that, for practical purposes, one can just assume that the supersymmetry breaking scenario will turn out to be just like that of the N = 1 supergravity theory [46]. Obviously, superstring inspired phenomenology is pursued by optimists!

The matter of supersymmetry breaking is not the only source of uncertainty in trying to connect superstrings to reality. Since one does not know precisely what the compact space K is, it is also necessary to make some assumptions on what the resulting four dimensional group g is. As I mentioned earlier, a necessary assumption of desert trekkers is that g must be smaller than SU(5), to avoid immediate problems with proton decay. Thus flux breaking must be allowed in the manifold K[43]. The pattern of sensible g's obtained after flux breaking has been studied by many people[55]. The results obtained depend crucially on whether one has or does not have an intermediate scale of symmetry breaking between M_{comp} and Λ_F. If there is no intermediate symmetry breaking then, due to flux breaking, E_6 breaks at M_{comp} to g and this is the surviving low energy group. It turns out, as I will demonstrate below, that g necessarily is <u>bigger</u> than the standard model SU(3) x SU(2) x U(1). If there is one, or more, intermediate scales of symmetry breaking, then the low energy group g obtained could be bigger than the standard model group, but it could also be precisely the standard model. Since flux tube breaking is equivalent to adjoint breaking and since matter and therefore also the Higgs fields are in 27's, one can characterize these two possibilities as:

i) Direct breaking (III.4)

$$E_6 \xrightarrow[M_{comp}]{\langle 78\rangle} g \qquad g > SU(3)\times SU(2)\times U(1)$$

ii) Intermediate scale breaking (III.5)

$$E_6 \xrightarrow[M_{comp}]{\langle 78\rangle} g' \xrightarrow[M_{int}]{\langle 27\rangle} g \qquad g \geq SU(3)\times SU(2)\times U(1)$$

Let me first discuss the case of direct breaking, Eq(III.4). If this happens, as I indicated above, g is necessarily bigger than the standard model group[55]:

$$g = SU(3) \times SU(2) \times U(1) \times \mathfrak{g}$$

$$\mathfrak{g} = \begin{cases} U(1) \\ U(1)^2 \\ SU(2) \times U(1) \end{cases} \qquad \text{(III.6)}$$

Thus one is lead to expect at least one extra neutral gauge boson, which survives at low energies, providing a characteristic signal for these patterns of compactification. The necessary existence of $\mathfrak{g}$ is rather easy to see by decomposing the 78 and 27 representations of E_6 in terms of the $SU(3)_c\times SU(3)_L\times SU(3)_R$ maximal subgroup[56]:

$$78 = (3,\bar{3},\bar{3})+(\bar{3},3,3)+(8,1,1)+(1,8,1)+(1,1,8) \qquad \text{(III.7a)}$$

$$27 = (3,3,1)+(1,\bar{3},3)+(\bar{3},1,\bar{3}) \qquad \text{(III.7b)}$$

If one does not want to break color, then only the last two terms in (III.7a) can contribute to the vacuum expectation of the 78. Now $\langle(1,8,1)\rangle \neq 0$ will break $SU(3)_L$ to $SU(2)_L \times U(1)_L$. However, it is not possible to identify $U(1)_L$ with the U(1) of the standard model since, according to (III.7b), the antiquarks which are singlets of $SU(3)_L$ would then have also no hypercharge. So the U(1) of the standard model gets contributions both from $U(1)_L$ and some $U(1)\subset SU(3)_R$. Thus it is not possible to break down $SU(3)_R$ completely and a non trivial $\mathfrak{g}$ ensues, which could be as large as SU(2) x U(1).

Besides having an extra Z^0, a further general property of direct breaking is that n_f 27's survive at low energy. This number, since it is a topological property of the compact space K, is not affected by flux breaking. Hence, in the case of direct breaking, all the exotic fermions of Table III.1 survive at low energy. Although the discovery of these states would provide evidence for superstring ideas, their presence at low energy, as I will discuss below, is far from being an unmitigated blessing.

The phenomenology of models with extra Z^0's has been studied by a great many authors in the last year and several papers on this topic have been submitted to this conference[57]. To discuss these models, it is particularly convenient to characterize the two extra U(1)'s in E_6, orthogonal to the usual hypercharge, via the decomposition[58]:

$$E_6 \to SO(10)\times U(1)_\psi \to SU(5)\times U(1)_\psi\times U(1)_\chi$$

with the standard model group being embedded in the SO(10) group. The charges corresponding to $U(1)_\psi$ and $U(1)_\chi$ for the 27 dimensional E_6 representation can be read off from the tables in the review of Slansky[56] and are displayed in Table III.2. If E_6 is broken down to $[SU(3)\times SU(2)\times U(1)]\times U(1)$, the additional U(1) charge, in general, will be a linear superposition of these charges.

$$Q' = Q_\psi \sin\theta + Q_\chi \cos\theta \qquad \text{(III.9)}$$

However, if this breakdown is induced by flux breaking then the angle θ is fixed to be[59]: $\theta = \cos^{-1}\sqrt{3/8}$. This can be seen as follows. Instead of Q' and Y, the two U(1)'s can be described in terms of $U(1)_L$ and $U(1)_R$ in the $SU(3)_c\times SU(3)_L\times SU(3)_R$ decomposition of E_6. Their properly normalized charges are

$$Y_L = \sqrt{3/5}\begin{pmatrix} 1/6 & & \\ & 1/6 & \\ & & -1/3 \end{pmatrix}_L ;$$

$$Y_R = \sqrt{3/5}\begin{pmatrix} 1/3 & & \\ & -1/6 & \\ & & -1/6 \end{pmatrix}_R \qquad \text{(III.10)}$$

acting on the appropriate SU(3) group. The hypercharge Y is simply expressible in terms of Y_L and Y_R

$$Y = Y_L + 2Y_R \qquad \text{(III.11)}$$

and Q' is the orthogonal combination

$$Q' = -2Y_L + Y_R \qquad (III.12)$$

Comparing (III.12), with (III.9) for an u quark, for example, establishes $\theta = \cos^{-1}\sqrt{3/8}$. The value of Q' is also given in Table III.2.

Table III.2 Charge values for the 27 of E_6.

States	Q_ψ	Q_χ	Q'
10: $\begin{pmatrix} u \\ d \end{pmatrix}_L u^c_L e^c_L$	$-\frac{1}{2\sqrt{6}}$	$-\frac{1}{2\sqrt{10}}$	$-\frac{1}{\sqrt{15}}$
$\bar{5}$: $d^c_L \begin{pmatrix} \nu \\ e \end{pmatrix}_L$	$-\frac{1}{2\sqrt{6}}$	$\frac{3}{2\sqrt{10}}$	$\frac{1}{2\sqrt{15}}$
1: ν^c_L	$-\frac{1}{2\sqrt{6}}$	$-\frac{5}{2\sqrt{10}}$	$-\frac{5}{2\sqrt{15}}$
5: $\begin{pmatrix} N^c \\ E^c \end{pmatrix}_L g_L$	$\frac{1}{\sqrt{6}}$	$\frac{1}{\sqrt{10}}$	$\frac{2}{\sqrt{15}}$
$\bar{5}$: $\begin{pmatrix} N \\ E \end{pmatrix}_L g^c_L$	$\frac{1}{\sqrt{6}}$	$-\frac{1}{\sqrt{10}}$	$\frac{1}{2\sqrt{15}}$
1: S_L	$-\frac{2}{\sqrt{6}}$	0	$-\frac{5}{2\sqrt{15}}$

An extra Z^0, if sufficiently light, would give rise to departures of neutral current experiments from the predictions of the standard model. Thus one can use neutral current data to put bounds on the mass of this particle[60]. In particular for the $Z^{0\prime}$, which couples to Q', since its coupling to quarks and leptons is not as strong as that of the ordinary Z^0, the bounds obtained are rather weak, hovering around $M_{Z^{0\prime}} \gtrsim 100$ GeV. These bounds have been reviewed by Deshpande[61] in this conference. An illustration is provided by Fig III.2, where the limits on $M_{Z^{0\prime}}$, from its non observation at the CERN collider, are detailed. Clearly if the $Z^{0\prime}$ cannot decay into exotic matter, these bounds are stronger. Note that even for a $Z^{0\prime}$ with the same mass as the Z^0, $\sigma B(Z^{0\prime} \to e^+e^-)$ is much smaller than for the Z^0. Analyses for other possible $Z^{0\prime}$s, like $Z(\theta)$, which couples to $Q_\psi \sin\theta + Q_\chi \cos\theta$, give similar bounds, although for particular θ-values these bounds can be rather weak[60].

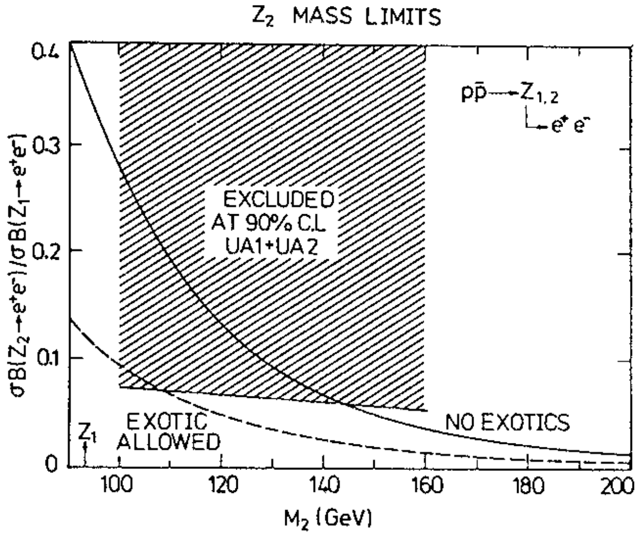

Fig III.2 Mass limits on $M_{Z^{0\prime}}$ from the CERN collider, taken from Barger et al, Ref 60.

Even though the $Z^{0\prime}$ is harder to produce than the ordinary Z^0, the Tevatron can push the mass limits for this excitation to near 300 GeV. For instance, London and Rosner[57] estimate a cross section times branching ratio into e^+e^-, for a $Z^{0\prime}$ of 200 GeV, of 1 pb at $\sqrt{s} = 1.8$ TeV. Such a signal should be detectable in a high luminosity run. Perhaps more favorable is the situation regarding the $Z^{0\prime}$ in e^+e^- collisions, since the presence of the $Z^{0\prime}$ can give rise to dramatic effects in various asymmetries, which will be measured at LEP and the SLC[57][62][63]. I illustrate this in Fig III.3, taken from Belanger and Godfrey[57], which shows the shift in the forward backward asymmetry and the left-right asymmetry expected at $\sqrt{s} = M_Z$ for various values of $M_{Z(\theta)}$. One sees that for the superstring $Z^{0\prime}$, corresponding to $\theta = \cos^{-1}\sqrt{3/8}$, the shifts in A_{FB} and A_{LR} are not as large as those for other values of θ. Nevertheless, these shifts are of a comparable order of magnitude to the expected effect in the standard model and much above the hoped for accuracy in these measurements[54]. If the $Z^{0\prime}$ is really as low as 150 - 200 GeV it will be visible directly in experiments at LEP 200, leading to large departures in the forward backward asymmetry for energies well below the resonance behaviour, as illustrated in Fig III.4. Indeed one should be sensitive to $Z^{0\prime}$ effects up to $M_{Z^{0\prime}} \lesssim 1$ TeV[63].

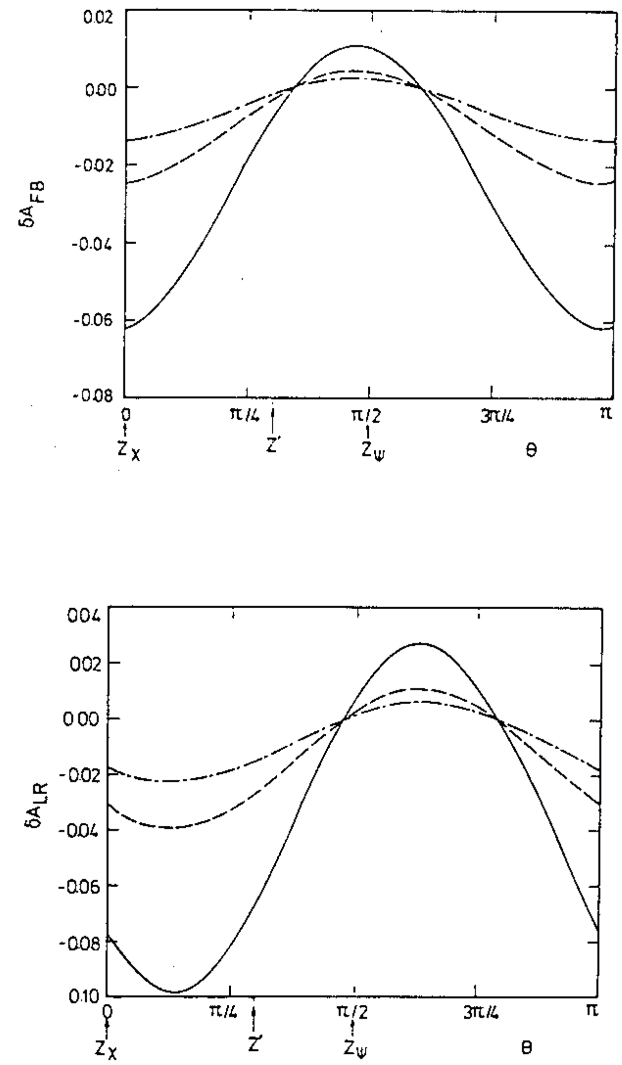

Fig III.3 Shifts in A_{FB} and A_{LR} caused by a new U(1) gauge boson $Z(\theta) = Z_\psi \sin\theta + Z_\chi \cos\theta$. The superstring gauge boson has $\theta = \cos^{-1}\sqrt{3/8}$. From Belanger and Godfrey, Ref 57.

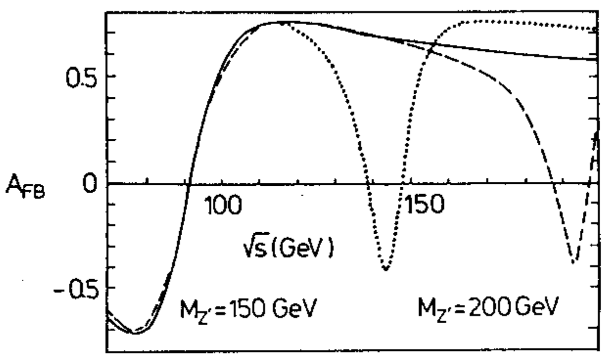

Fig III.4 Departures of A_{FB} from standard model expectations in the LEP 200 range due to a $Z^{0\prime}$. From Matsuoka et al, Ref 57.

Since there is no plausible motivation, except superstrings, for having an extra Z^0 at low energy, the detection of such a signal would provide strong corroboration for the superstring ideas. However, recall that if there is direct breaking, one also requires the presence at low energy of n_f full multiplets of 27's. The extra fermions in the 27, although providing further evidence for superstrings, are in themselves quite problematic[49]. Let me briefly discuss some of these potential difficulties:

i) Having both supersymmetry and 27, rather than 15 (or 16), fermions per family makes the SU(3) and SU(2) coupling constants run very differently from the case where one has only the usual quarks and leptons. The relevant β-functions, for scales above possible low energy thresholds read, in this case:

$$\beta_3 = -\frac{3g_3^3}{8\pi^2}(3 - n_f)$$
$$\beta_2 = -\frac{3g_2^3}{8\pi^2}(2 - n_f) \qquad \text{(III.13)}$$

Hence for $n_f > 3$ both α_3 and α_2 grow at high energy. Indeed, one can convince oneself from (III.13) that if $n_f = 4$ $\alpha_3 > 1$ already at $q^2 \simeq 10^8$ GeV! Since α_3, α_2 and α_1 must unify at M_{comp}, in these schemes, and $M_{comp} \sim M_P$, it is clear that *direct breaking is only possible if $n_f = 3$*. For $n_f = 3$, α_3 does not run at high energy. Nevertheless it is possible to have SU(3)xSU(2)xU(1)xU(1) unified at scales of order 4×10^{17} GeV[65], with a resulting $\sin^2\theta_W \simeq 0.21$, which is adequate, if not spectacular.

ii) If the g quarks survive at low energy, there can be universality violations, and flavor changing neutral currents, induced by mixings of g with the other charge -1/3 quarks. Similarly the charged lepton E can mix with $e\mu\tau$. Both of these circumstances are catastrophic, unless one can effectively suppress, to an arbitrary low level, all possible mixings. Although one can argue for the absence of these mixing effects[66], by allowing only certain Higgses to have vacuum expectation values, these arguments are not particularly compelling to me.

iii) The presence of the exotic fields in the 27 can lead to very rapid proton decay, through the presence of dimension four terms in the superpotential. On general grounds, one knows that the superpotential itself contains terms involving the $(27)^3$ of E_6. However, the coefficients of these terms are not fixed by E_6[43], since they get altered by flux breaking. The dangerous terms for proton decay are the ones connecting g to quarks and leptons:

$$g^cLQ \; ; \; ge^cu^c; \; gd^c\nu^c; \; gQQ; \; g^cu^cd^c \qquad (III.14)$$

Here $Q=\binom{u}{d}_L$ and $L=\binom{\nu}{e}_L$. One may argue[66] that only a subset of these couplings (the first three, or the last two) are nonvanishing because of topological reasons. In this case one can define a baryon number and there is no problem with proton decay. However, to my knowledge, no specific example has been found of a Calabi Yau manifold where the above can be demonstrated.

iv) Similar problems exist with neutrino masses. In principle, a coupling in the superpotential involving ν^c, L and the extra doublet $L^c = \binom{N^c}{E^c}$ exists. Since the scalar partners of L^c play the role of Higgs fields in the theory one is lead naturally to a Dirac mass for neutrinos, unless this coupling itself is absent.

In view of the above difficulties it is perhaps easier envisaging models with an intermediate scale breaking, since in these models one can, in principle, push some of the unwanted fermions in the 27 to high scales. The possibility of intermediate scale breaking requires more dynamical assumptions and obtains only if one has a certain number, δ, of chirally paired states $(\delta(27+\overline{27}))$. As pointed out by Dine et al[65] and Witten[40)43)], it is only consistent to have some of the scalar fields in the 27 acquire large vacuum expectation values, if there are directions in the scalar potential which are not affected by these expectation values (flat directions). In general, the scalar potential can be written as

$$V = |F|^2 + D^2 + \text{soft Susy breaking terms} + \text{non renorm. terms} \qquad (III.15)$$

Ignoring for the moment the last two terms, it is clear that intermediate scale breaking only will obtain if V is both D and F flat. D flat directions occur if the vacuum expectation value of a 27 component can be cancelled by that of a $\overline{27}$. Hence δ must be non vanishing. Furthermore, for example, since in the $(27)^3$ terms in the superpotential, no factors containing S^2 or S^3 appear one sees that $\langle S\rangle = \langle\bar{S}\rangle \neq 0$ is also an F-flat direction[67]. Hence the value of $\langle S\rangle = \langle\bar{S}\rangle$ is totally determined by the supersymmetry breaking and non renormalizable terms in V. A non zero and large value of $\langle S\rangle$ obtains if the soft supersymmetric breaking terms for S are, in fact, <u>negative</u> and act against the non renormalizable pieces in V. That is

$$V \simeq -\Lambda^2|S|^2 + \frac{|S|^6}{M^2_{comp}} \qquad (III.16)$$

where, presumably $\Lambda \simeq \Lambda_F$, and the scale which typifies the non renormalizable terms is M_{comp}. From (III.6) it follows that

$$\langle S\rangle = \langle\bar{S}\rangle = M_{int} \sim \sqrt{\Lambda_F M_{comp}} \sim 10^{10}\ \text{GeV} \qquad (III.17)$$

Note that although it is necessary to have D and F flat directions to be able to generate an intermediate scale, the existence of these directions is not enough to guarantee that (III.17) obtains. For this, it is necessary that the soft supersymmetry breaking terms really be driven to have a negative coefficient - something which is not so easy to demonstrate in practice. Superstring partisans, in general, are content to find F and D flat directions and optimistically assume that if these flat directions exist, the conspiracy of Eq(III.16) will follow[65].

With intermediate scale breaking, as indicated in Eq(III.5), the final group g can be bigger or equal to the standard model group. What pattern ensues depends both on what the manifold K is, which causes the first stage of breakdown: $E_6 \xrightarrow{\langle 78\rangle} g'$, and on which components of the 27 cause the further breakdown. Apart from $\langle S\rangle \neq 0$, one can check that the only other realistic possibility is $\langle\nu^c\rangle \neq 0$. If the first non renormalizable terms in the superpotential are of the form of $(27)^2(\overline{27})^2$ then the scalar potential along the flat directions will be as in (III.16) and the estimate (III.17) for M_{int} follows. It may be, however, that only higher terms in the superpotential are allowed, like $(27)^3(\overline{27})^3$. In this case M_{int} is higher[68]

$$M_{int} \sim (\Lambda_F M^3_{comp})^{1/4} \sim 10^{14}\ \text{GeV} \qquad (III.18)$$

Intermediate scales as high as this are interesting since if $m_g \sim 10^{14}$ GeV, then g mediated proton decay is sufficiently suppressed. However, one must make sure that there remain in the theory light doublets to allow for a low energy breakdown of the standard model. The splitting of the triplets from the doublets is rather natural if the doublets arise from the fields in the $\delta(27+\overline{27})$ and flux breaking has already removed the triplet fields from these components at compactification[69].

In the case of intermediate scale breaking, as I have indicated, a breakdown pattern to the standard model is possible. A nice example of this possibility has been considered by Greene, Kirklin, Miron and Ross[70], who studied one of the few three generation Calabi Yau manifolds known[71]. This manifold has a first homotopy group $\pi_1(K)=Z_3$, so it admits flux breaking. Before flux breaking, the model has 3(27) and $6(27 + \overline{27})$ multiplets. The two non trivial embeddings of Z_3 in E_6 give SU(3) x SU(3) x SU(3) and SU(6) x U(1), respectively, as the resulting group after flux breaking. Greene et al[70] concentrate on the first possibility, since it can lead to a realistic theory. After flux breaking all the existing $SU(3)_c$ singlet fields in the chirally paired $(27 + \overline{27})$ representations survive, as well as 4 out of 6 of the $SU(3)_c$ triplet and antitriplet fields in these representations. This very unpleasant phenomenologial situation is remedied by a sequence of two intermediate scale breakings, triggered by ν^c and S vacuum expectation values:

$$SU(3)^3 \xrightarrow{\langle\nu^c\rangle} SU(3) \times SU(2) \times SU(2) \times U(1)$$
$$\xrightarrow{\langle S\rangle} SU(3) \times SU(2) \times U(1) \tag{III.19}$$

Greene et al[70], by studying the manifold's discrete symmetries show that the superpotential is F-flat to $O(27^3\overline{27}^3)$ for the first breaking. Thus $\langle\nu^c\rangle \sim 10^{14}$ GeV, while $\langle S\rangle \sim 10^{10}$ GeV, provided, of course, that the appropriate supersymmetry breakdown to trigger this sequence exists. This multiple breakdown gives high mass to almost all the remaining vector-like states in the theory. Remarkably, however, even though all g quarks are heavy, two doublet superfields stay light. These are precisely the Higgs multiplets necessary for a supersymmetric extension of the standard model[46]. In addition, the coupling of these multiplets to the quarks, possesses certain discrete symmetries which yield a reasonable structure for the Kobayashi-Maskawa matrix[70].

The results of Greene et al[70] are both encouraging and discouraging. I find it encouraging that there exist Calabi Yau manifolds with topological properties which can lead one to a model at low energy with many of the characteristics of the standard model. It is discouraging, however, that there is so little to show from the superstring superstructure, except certain interrelations among Yukawa couplings. Furthermore to get from the superstrings to the standard model one has had to make many assumptions, each of them hard to justify. So one is left in the ambivalent position of not being able to decide whether defects in the resulting theory are due to poor intermediate assumptions or are really signals of some profound sickness in the scheme. A case in point is provided by SU(3) x SU(2) x U(1) unification in the manifold studied by Greene et al[70]. Although one can get unification into $(SU(3))^3$ at 10^{14} GeV with a reasonable $\sin^2\theta_W$, the presence of so many matter fields beyond this scale drives the gauge couplings above unity much before M_{comp}. Is this a deadly defect, or can it be conveniently ignored? Questions of this ilk, unfortunately, abound in trying to bring superstring ideas down to laboratory energies.

My conclusions on superstring inspired phenomenology are two fold:

i) Most of the "predictions" of superstrings at low energy are strongly dependent on implicit assumptions made at the compactification scale and at possible intermediate scales. Although certain specific predictions are phenomenologically appealing, like the presence of an extra Z^0, none of these predictions are sure things.

ii) To motivate the existence of superstrings it is important to find evidence for the "super" aspect of these theories. The existence of low mass ($\tilde{m} \lesssim$ TeV) superpartners remains the best "smoking gun" for superstrings.

IV LOOKING FOR SUPERSYMMETRY

Long before superstrings became popular, there was considerable theoretical activity in low energy supersymmetry[72]. The physical motivation for considering supersymmetric extensions of the standard model was related to the hierarchy problem. Although supersymmetry cannot explain the existence of mass hierarchies, it allows hierarchies naturally to exist. Mass shifts in the scalar sector are no longer quadratically divergent since there is a cancellation between bosonic and fermionic contribution, so radiative corrections do not destabilize the theory. Although supersymmetry is not the only way to make the Fermi scale a natural parameter, it is obviously the solution chosen by superstrings. So, for these theories also, it is sensible to expect to have superpartners of the known excitations in a mass range below, say, 1 TeV. Unfortunately, since one does not know precisely how supersymmetry is broken, one cannot really pin down the masses of the superpartners. All that is necessary is that these masses be low enough to provide a credible mechanism for having Λ_F of O(250 GeV).

The usual assumption pursued is that supersymmetry is broken spontaneously. The favored scenario is that discussed earlier, which is based on an N=1 supergravity theory in which the supersymmetry breaking occurs in a hidden sector[46)48]. In the low energy theory the manifestation of this breakdown is the appearance of soft breaking terms which give masses to the scalars and the gauginos and provide corrections to scalar vertices. The schematic structure of these terms is shown in Fig IV.1.

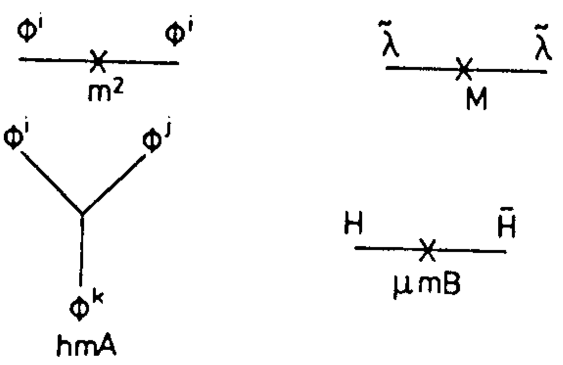

Fig IV.1 Soft supersymmetry breaking terms generated through hidden sector breaking. ϕ_i are scalars, $\tilde\lambda$ are gauginos and H,$\bar H$ are the Higgs doublets.

The resulting superpartner spectrum depends on detailed assumptions one makes on M, m, A and B at some high scale (M_{comp} or M_p), plus the renormalization group evolution of these parameters down from this scale. In general one takes the gaugino and scalar masses, M and m, to be universal at the high scale and one gets SU(2) x U(1) to break down when the soft Higgs mass squared is driven negative during the renormalization group evolution[47].

In the absence of precise knowledge of the superpartner spectrum, the most important quantity to know for phenomenology is which is the lightest supersymmetric particle (LSP). Supersymmetric theories are invariant under an additional symmetry, R parity[73]. R parity is a multiplicatively conserved quantity, given by $R = (-1)^{2J+3B+L}$, so that R = +1 for all the known particles and R = -1 for their superpartners. The existence of R parity implies that:
i) Sparticles are produced in pairs.
ii) The lightest sparticle, LSP, is absolutely stable.
Because of ii) all sparticle decay chains end in an LSP and so for supersymmetric searches one needs to know which is this particle. There are good astrophysical arguments[74] that an LSP cannot be charged or have strong interactions, if not it would have condensed in galaxies and planets and would have been detected in searches of matter with anomalous values of e/m.

The usual assumption made is that the LSP is the photino, $\tilde\gamma$. One can adduce a variety of astrophysical, cosmological and particle physics arguments in favor of this hypothesis [74][75]. For instance, if the gaugino soft masses M are universal, then the different renormalization behaviour for gluinos and photinos implies $m_{\tilde g} \simeq 7m_{\tilde\gamma}$. In certain models[76] it is possible that the sneutrino is the LSP. However, I shall not consider this possibility here. I shall also not discuss in detail the present status of bounds on supersymmetric particles, since Davier[25] has discussed this topic in his rapporteur talk here. Rather, I'll consider only one example of a supersymmetric particle search: that of squarks and gluinos at the CERN SppS collider, and how this will be extended at the Tevatron collider. If the photino is the LSP, the produced squarks and gluinos will decay to photinos through the chains

$$m_{\tilde g} > m_{\tilde q} \qquad \tilde g \to q\bar q\tilde\gamma \; ; \; \tilde q \to q\tilde\gamma$$
$$m_{\tilde g} < m_{\tilde q} \qquad \tilde g \to q\bar q\tilde\gamma \; ; \; \tilde q \to q\tilde g \to qq\bar q\tilde\gamma \qquad \text{(IV.1)}$$

Since the produced photino interacts weakly, it provides a missing energy signal. Hence the well known experimental signature to expect, in the case of gluino or squark production, is missing energy plus (multi) jets. How many jets, however, is a sensitive issue that depends crucially on experimental cuts.

The famous (infamous?) UA_1 monojet signal and its relation to supersymmetry were the hot subject in 1984-1985[77]. At this conference the UA_1 collaboration has presented[78] a very complete analysis of their missing energy signal, which consists of 53 monojets and 3 dijets. Already the preponderance of monojets, even with the UA_1 cuts, is a bad signal for squark or gluino production, since one would expect from these decays a sizeable fraction of multijet events[79]. As Honma reported[78], the missing energy events are essentially accounted for by standard model backgrounds. This allows the UA_1 collaboration to set rather strong bounds on the masses of squarks and gluinos. These bounds are shown in Fig IV.2. One sees from this figure that $m_{\tilde q} \lesssim 80$ GeV and $m_{\tilde g} \lesssim 60$ GeV are excluded. Actually, the UA_1 collaboration, cannot also exclude a light gluino window. However, this light gluino scenario has fallen in theoretical disrepute[80], and I have taken the liberty of removing this window from Fig IV.2.

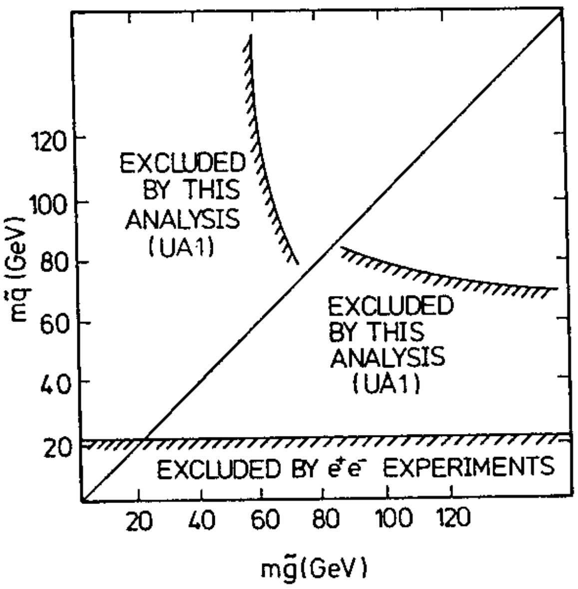

Fig IV.2 Limits on gluino and squark masses obtained by the UA_1 collaboration[78].

It would have been nice if the UA_1 collaboration could have given a bound also on the Wino. Although the mass matrix for the charged gauginos is model dependent, for models with a small supersymmetry breaking gaugino mass, one of the eigenstates, $\tilde W$, is lighter than the W[81]. So sequential decays like $W \to \tilde W \tilde\gamma$; $\tilde W \to q\bar q\tilde\gamma$ are possible. These processes give rise mostly to monojets, when one takes into account of the UA_1 cuts. So the preponderance of monojets in the missing energy signal is at least not unfavorable to Winos. In this conference, Arnowitt[32] estimated that a $\tilde W$ mass $M_{\tilde W} \gtrsim 40$ GeV is still compatible with the UA_1 signal. However, it is clear that a reliable analysis can only be done by the UA_1 collaboration itself.

Although the bounds on gluinos and squarks obtained at the CERN collider are impressive, it is important to emphasize that if these excitations exist near these bounds then they will be rather easily seen at the Tevatron. This point has been forcefully made by Baer and Berger[83] and Reya and Roy [84] and has been discussed at this conference by Reya[85]. For given sparticle masses the production cross section grows rapidly with energy in the range from $\sqrt{s}$ = 620 GeV to $\sqrt{s}$ = 2 TeV and the signal relative to standard model background also increases. As an example, I show in Fig IV.3, the event rate for producing gluinos at $\sqrt{s}$ = 1700 GeV for various cuts on the missing energy, for the case $m_{\tilde q} = 2m_{\tilde g}$. One sees that even for the very safe cut of p_T^{miss} > 60 GeV one gets nearly 5 events per 100 nb^{-1}, for gluinos as heavy as 100 GeV.

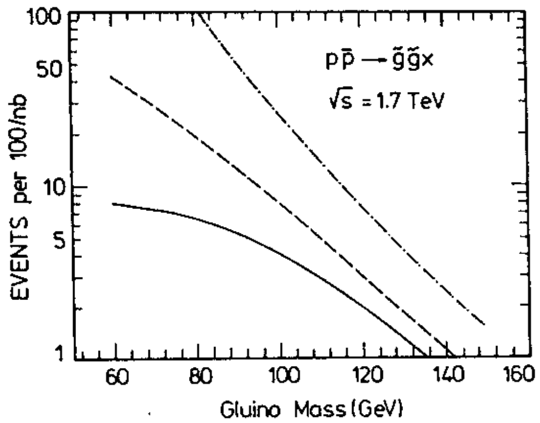

Fig IV.3 Gluino production at $\sqrt{s}$ = 1700 GeV, for $m_{\tilde q} = 2\, m_{\tilde g}$, as a function of various p_T^{miss} cuts (Dot dash: total; dashed > 40 GeV; solid > 60 GeV) from Ref 86. For similar curves see also Ref 83 and 84.

Secondly, at the Tevatron multijet events will be a dominant feature of squark and gluino production[83][84]. Not only dijets will dominate over monojets but also, for sufficiently large gluino masses, trijets become quite important, as shown in Fig IV.4.

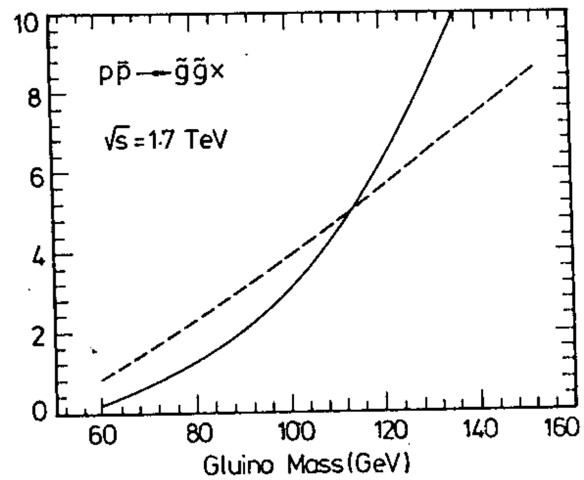

Fig IV.4. Fraction of dijets (dashes) and trijets (solid) to monojets for $\sqrt{s}$ = 1.7 TeV. From Ref 86.

Operating the Tevatron at $\sqrt{s}$ = 1.6 TeV one should be able to discover[83,85] squarks and gluinos if $m_{\tilde g}$, $m_{\tilde q} \lesssim 150$ GeV. This discovery range can be pushed to near 200 GeV when $\sqrt{s}$ is raised to 2 TeV[85].

V FERMI SCALE PHYSICS - CHALLENGES AND HOPES

There is a minority of theorists, to which I belong, who contend that the origin of the Fermi scale is not directly related to phenomena occurring at energies much above a TeV. These present day heretics believe that Λ_F is a dynamical scale, related to the presence of condensates of an underlying strong interaction theory. The Higgs picture gives only an approximate description of the true theory, just like the Landau Ginzburg model was an approximation for the fundamental BCS theory[87]. Thus questions of stability and naturalness are irrelevant for the Higgs sector and it is unnecessary to appeal to supersymmetry to stabilize the theory. (Supersymmetry might well exist, however, for deeper reasons).

One can imagine that the role of the underlying theory is just to provide a mechanism for generating the Fermi scale, as was the case for technicolor[88]. However, it is perhaps more reasonable to suppose that this strongly interacting theory is also responsible for producing quarks and leptons as composite bound states of more fundamental objects - preons[89]. The status of these theories has been summarized by Harari at this conference[90] in a very kind way, by pointing out that although the motivation for compositeness remains as good as ever, the major problems are largely unchanged! Particularly troublesome to me is the absence of any model which can really serve as a paradigm, so one is left only with a collection of disconnected dynamical ideas.

One of the principal difficulties of composite models is related to the fact that one is asking the theory to do two separate things, which are hard to reconcile. On the one hand, one would like the theory to provide the spontaneous breakdown of SU(2) x U(1). Hence, it is natural to imagine that the dynamical scale of the theory, Λ, is of the order of Λ_F: $\Lambda \simeq \Lambda_F$. On the other hand, one would like these theories to provide a mechanism for generating family replications and small fermion masses, while at the same time avoiding large flavor violations. This last point seems to demand that $\Lambda >> \Lambda_F$.

The physics of the underlying theory will generate, in general, effective non renormalizable terms which violate flavor and which must be added to the standard model Lagrangian[91]:

$$\mathcal{L}_{eff} = \mathcal{L}_{SM} + \sum_i \frac{\lambda_i}{\Lambda^2} O_i \qquad (V.1)$$

where the O_i are SU(3)xSU(2)xU(1) invariant operators. Taking $\lambda_1 = 1$, an analysis of a variety of flavor changing processes, reported by Wyler at this conference[92], gives typical bounds $\Lambda \gtrsim 10^2 - 10^3$ TeV. It is difficult for a model which has Λ of this order of magnitude to generate $\Lambda_F \sim 0.1$ TeV. Even admitting two different scales Λ_F and Λ, if Λ is so large it is also difficult to get fermion masses large enough. An example being ETC[93], where $m_f \sim \Lambda_F^3 / \Lambda^2$.

There has been really no theoretical progress on the flavor issue. However, skirting the flavor problem, some theoretical advances has been made, which could have some phenomenological implications. I would like to briefly discuss two recent examples, whose lessons perhaps can be useful, even though they do not illuminate the role of flavor. One should keep in mind, in this respect, that it is perfectly possible for the compositeness scale of electrons and muons individually to be of O(TeV), and yet only to be able to probe the e-μ difference at distances of order $(10^3 \text{TeV})^{-1}$ [90].

The first example which I will discuss is the, so called, composite Higgs model[94]. This model separates the scales Λ and Λ_F by making the Higgs bosons essentially Goldstone bosons. One imagines that, in the limit in which the electroweak interactions are turned off, the underlying theory posesses a global symmetry G which is broken down to another group H, producing certain bound state Goldstone bosons Φ. Turning on the SU(2)xU(1) couplings causes a realignement of the vacuum and Φ acquires a non zero vacuum expectation value

$$\langle\Phi\rangle = \Lambda_F = f(\alpha)\Lambda \qquad (V.2)$$

The function $f(\alpha)$ is dynamically determined and it is possible that $\Lambda >> \Lambda_F$. For the vacuum realignement to actually take place, it is necessary that the low energy group be bigger than SU(2)xU(1)[95]. Although H > SU(2)xU(1), the unbroken group H should not contain the full gauge theory, if one wants the vacuum to reorient itself when the gauge couplings are turned on. Amusingly enough, therefore, the simplest realization of these theories contains also an extra Z^0, although with characteristics quite different from the superstring Z^0's.

Dugan, Georgi and Kaplan[96] studied a toy model where the weak group is H_W = SU(2)xU(1)xU(1)$_A$ and G = SU(5) while H = O(5). In this model the effective Higgs potential depends on a function of the ratio of the SU(2)xU(1) and U(1)$_A$ couplings:

$$c_0 = \frac{3g^2 + g'^2}{g_A^2} \qquad \text{(V.3)}$$

and one can establish that Λ_F vanishes in the limit as $c_0 \to 1$, so that

$$f(\alpha) \sim \sqrt{1-c_0} \qquad \text{(V.4)}$$

The Higgs effective potential is calculable in the model, and therefore the Higgs mass and that of the $\tilde{Z}$ boson are given as functions of c_0. The dependence of these quantities on c_0 is displayed in Fig V.1. To be in agreement with neutral current data $M_{\tilde{Z}}$ must be high enough, which in this case implies $c_0 > 0.6$[96]. For values of c_0 in the allowed range, it is easy to see from Fig V.1 that the Higgs mass lies in a band near 200 GeV.

The model discusses in Ref 96 is not quite realistic since, for example, one needs to introduce spectator fermions to cancel some of the anomalies associated with the extra axial U(1)$_A$ gauge interactions. Trying to make the composite Higgs model more realistic has engendered a growing set of bizzare and baroque, but clever, models: the moose models[97]. Unfortunately, these models are not yet ripe for phenomenology.

The second example I want to discuss is the strongly coupled standard model, which was reported on, in this conference, by B. Schrempp[98]. This model was proposed originally by Abbott and Farhi[99] and has been recently reexamined in some detail by Claudson, Farhi and Jaffe[100]. The Lagrangian for the model is precisely that of the standard model, except that all left handed fields ψ_L and the Higgs doublet Φ are considered as preons. Most importantly, the SU(2) gauge group is supposed to confine and not suffer spontaneous breakdown. In the limit of vanishing g', and neglecting all Yukawa couplings, the theory possesses a global SU(4n_f)xSU(2)$_W$ symmetry, where the last symmetry group arises from the O(4) symmetry properties of the Higgs potential. This symmetry can be preserved in the binding if one can find a set of massless composite fermions to match the global symmetry anomalies [101]. This set is trivially provided by the left handed quarks and leptons constructed as

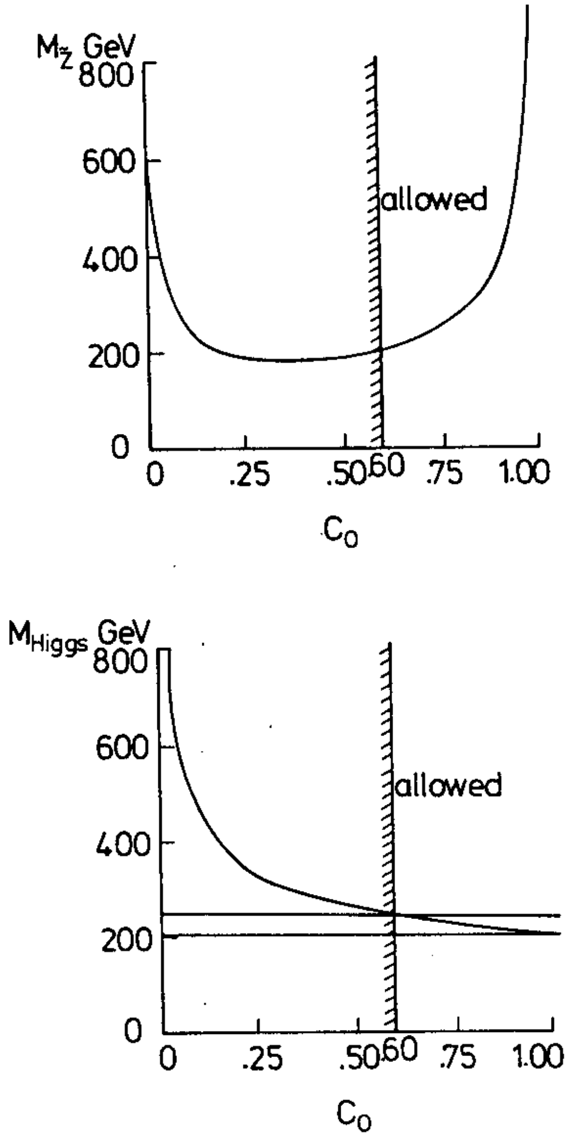

Fig V.1 Masses of the $\tilde{Z}$ and composite Higgs boson as a function of c_0, for the model of Ref 96. Phenomenologically $c_0 > 0.6$

$l_L, q_L \sim (\Psi_L\ \Phi)$ (V.5)

which obviously transforms as $(4n_f, 2)$ under the global group.

Since the SU(2) confines, one has to assume that the weak interactions are mediated by composite $W^{\pm}$ and Z bosons, which are bound states of the Higgs field Φ. The crucial difficulty in this model is to demonstrate that these vector bosons are well separated from other J=1 states in the spectrum. If this is so, then, to a very good approximation, the strongly coupled standard model is analogous in content to the spontaneously broken standard model. The existence of a large gap in the J=1, T=1 channel can be seen[100] to be equivalent to having $M_W = M_Z \cos\theta_W$ and to having a small effective coupling $\bar{g}$ for the Wf_Lf_L vertex, where f_L is a left-handed bound state fermion in the theory[100]. Let me focus on this last point. The effective coupling constant $\bar{g}$ can be plotted as a function of $(\Lambda/\Lambda_F)^2$, where Λ is the SU(2) dynamical scale, and Λ_F is now just a parameter in the Higgs potential. For $\Lambda << \Lambda_F$ $\bar{g}$ is precisely the gauge coupling of the standard model in the spontaneously broken phase, which vanishes as $\Lambda \to 0$. For $\Lambda >> \Lambda_F$, on the other hand, one is in the strong coupling phase and following the perturbative evolution of $\bar{g}$ one would expect it to become large. For the strongly coupled standard model to make sense, however, $\bar{g}$, for Λ^2/Λ_F^2 large, must attain again a small value, $\bar{g} \simeq 0.7$, since this is what is dictated by phenomenology. The required dynamical behaviour of $\bar{g}$ is plotted in Fig V.2, taken from Ref 100.

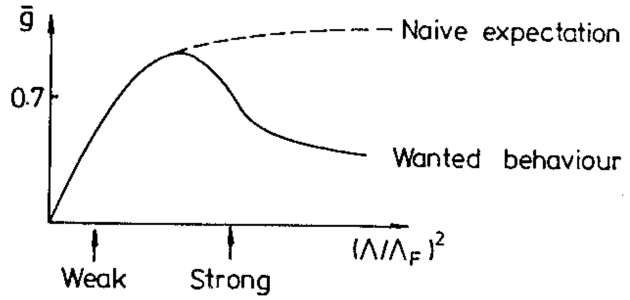

Fig V.2 Behaviour of the effective Wf_Lf_L coupling $\bar{g}$ needed for the strongly coupled standard model to be consistent.

It is an unsolved problem whether dynamically the strongly coupled standard model makes sense; i.e. if the behaviour shown in Fig V.2 for $\bar{g}$ really obtains. If this happens, however, one has a model which for most purposes is equivalent to the standard model but which, in addition, has a rich spectrum of composite bosons and fermions[100]. Some of the phenomenology of these states has been discussed in this conference by B. Schrempp[98]. Here I will only comment on a noteworthy class of candidates for the model: spin zero difermions. These states are bound states of two Ψ_L which are antisymmetric in the flavor indices a,b of the Ψ_L's, with a,b going from 1 to $4n_f$. That is

$$S^{[a,b]} \sim \Psi_L^a\ \Psi_L^b \quad (V.6)$$

The $S^{[a,b]}$ contain both charge -1 dileptons, charge -1/3 leptoquarks and charge +1/3 antitriplet diquarks. Hence their coupling to the ordinary quark and lepton left handed doublets $D^a = \begin{pmatrix} q_L, & l_L \end{pmatrix}$ is given by

$$\mathcal{L}_{eff} = \frac{\lambda}{2} S^+_{ab}\ D^{aT} C D^b + h.c. \quad (V.7)$$

Note that since these scalar excitations are <u>not</u> Goldstone bosons, there is no reason to suppose that λ is proportional to the mass of the bound state fermions. Korpa and Ryzak[102] used data on neutrino nucleon scattering to put a bound on λ/M_S, where M_S is the mass of the difermions:

$$M_S > 275\ \lambda\ \text{GeV} \quad (V.8)$$

Hence if λ is of the same order of magnitude as $\bar{g}$, the leptoquarks in S would be in the discovery range of HERA. Indeed, in this case, as discussed by Wudka[103], the cross sections are very large, since the leptoquark production is a resonance process. Unfortunately the value of λ, like that of $\bar{g}$, is beyond our present computational capacity.

VI CONCLUDING REMARKS

I hope this report has demonstrated that theoretical ideas beyond the standard model are rich and varied, leading to potentially very interesting phenomenology. However, phenomenology needs phenomena! Thus we are all waiting with extreme interest for the results which shall be forthcoming in the coming years from the Tevatron, SLC, LEP and HERA, as well as from non accelerator experiments. Only then we shall know what kind of physics, if any, lies beyond the standard model. From this point of view, the tremendous theoretical investment in D>4 physics could be a bit premature. After all, Fermi scale physics may well turn out to be different!

REFERENCES

(1) G. Altarelli, these proceedings

(2) W. Scott, these proceedings

(3) E. Fischbach, D. Sudarsky, A. Szafer, C. Talmadge and S.H. Aronson, Phys. Rev.Lett. 56(1986)3; E. Fischbach, these proceedings

(4) R.V. Eötvos, D. Pekár and E. Fekete, Ann.Phys. (Leipzig)68(1922)11

(5) E. Fischbach, D. Sudarsky, A. Szafer, C. Talmadge and S.H. Aronson, Proceedings of the 2nd Conference on Intersections between Particle and Nuclear Physics, Lake Louise, May 1986

(6) F.D. Stacey and G.J. Tuck, Nature 292 (1981)230

(7) Review of Particle Properties, Phys. Lett. 170B(1986)1

(8) S.C. Holding, F.D. Stacey and G.J. Tuck, Phys.Rev. D33(1986)3487

(9) G.W. Gibbons and B.F. Whiting, Nature 291(1981)636

(10) M. Milgrom, IAS preprint; S.C. Holding, F.D. Stacey and G.J. Tuck, Ref 8; P. Thieberger, Phys.Rev.Lett. 56(1986)2347; A. Bizzeti, Florence University preprint; D.A. Neufeld, Phys.Rev.Lett. 56(1986) 2344; C. Talmadge, S.H. Aronson and E. Fischbach, in Proceedings of the XXI Rencontre de Moriond, Les Arcs, March 1986

(11) Y. Asano et al, Phys.Lett. 107B(1981) 159; 113B(1982)195

(12) S. Weinberg, Phys.Rev.Lett. 13(1964) 495

(13) M. Suzuki, Phys.Rev.Lett. 56(1986) 1339; C. Bouchiat and J. Iliopoulos, Phys.Lett.169B(1986)447; M. Lusignoli and A. Pugliese, Phys.Lett. 171B(1986) 468;S.H. Aronson, H.-Y. Cheng, E. Fischbach and W. Haxton, Phys.Rev.Lett. 56 (1986)1342

(14) I. Bars and M. Visser, Phys.Rev.Lett. 57(1986)25; I. Bars, these proceedings

(15) J. Schweppe et al, Phys.Rev.Lett. 51 (1983)2261; M. Clemente et al, Phys. Lett. 137B(1984)41;

(16) T. Cowan et al, Phys.Rev.Lett. 56 (1986)446

(17) J. Reinhardt, B. Müller and W. Greiner, Phys.Rev. A24(1981)173

(18) T. Cowan et al, Phys.Rev.Lett. 54 (1985)1761

(19) A. Schäfer et al, J.Phys.G: Nucl.Phys. 11(1985)L69; A. Chodos and L.C.R. Wijewardhana, Phys.Lett. 56(1986)302; J. Reinhardt et al, Phys.Rev. C33 (1986)194; K. Lane, Phys.Lett. 169B (1986)97

(20) R.D. Peccei and H.R. Quinn, Phys.Rev. Lett. 38(1977)1440; Phys.Rev. D16 (1977) 1791; S. Weinberg, Phys.Rev. Lett. 40(1978)223; F. Wilczek, Phys. Rev.Lett. 40(1978)279

(21) For a review of these bounds, see for example, S. Yamada in Proceedings of the International Symposium on Lepton Photon Interactions at High Energies, Cornell University, Ithaca, 1983

(22) R.D. Peccei, T.T. Wu and T. Yanagida, Phys.Lett. 172B(1986)435; L.M. Krauss and F. Wilczek, Phys.Lett. 173B(1986) 189

(23) W.A. Bardeen and S.H. Tye, Phys.Lett. 74B(1978)229

(24) J. Greenberg, Proceedings of the International Symposium on Weak and Electromagnetic Interactions in Nuclei, Heidelberg, July 1986

(25) M. Davier, these Proceedings

(26) W.A. Bardeen, R.D. Peccei and T. Yanagida, Nucl.Phys. B279(1987)401; L.M. Krauss and M. Wise, Phys.Lett. 176B(1986)483; M. Suzuki, Phys.Lett. 175B(1986)364

(27) W.A. Bardeen et al, Ref 26

(28) E. Eichler et al, Phys.Lett. 175B (1986)101

(29) F.W.N. de Boer et al, Groningen preprint, Phys.Lett. B to be published

(30) T.W. Donnelly et al, Phys.Rev. D18 (1978)1607

(31) For a review of this scenario and its astrophysical and cosmological constraints see for example, F. Wilczek in

How far are we from the Gauge Forces, Erice lectures 1983 (Plenum Press, New York 1985)

(32) C. Callan, Proceedings of the XXII International Conference of High Energy Physics, Leipzig 1984

(33) M. Green and J. Schwarz, Phys.Lett. 149B(1984)117

(34) J. Schwarz, these proceedings

(35) P. Candelas, G.T. Horowitz, A. Strominger and E. Witten, Nucl.Phys. B258 (1985)46; Proceedings of the Argonne Symposium on Anomalies, Geometry, and Topology, (World Scientific, Singapore)

(36) E. Calabi, Algebraic Geometry and Topology: a Symposium in Honour of S. Letschetz (Princeton University Press, Princeton, N.Y. 1957); S.T. Yau, Proc. Nat.Acad.Sci. 74(1977)1798

(37) P. Candelas et al, Ref 35; C.G. Callan, E. Martinec, D. Friedan and M. Perry, Nucl.Phys. B262(1983)593; A. Sen, Phys.Rev. D32(1985)2102; C.M. Hull and P.K. Townsend, Nucl.Phys. B274(1986)349

(38) Even though the β function for these σ models does not vanish at four loops (M.T. Grisaru, A. van de Ven and D. Zanon, Nucl.Phys. B277(1986)388, C.N. Pope, M.F. Sohnius and K. Stelie, Imperial College preprint) one can perperturbatively construct such conformally invariant σ models. See D. Nemeschansky and A. Sen, Phys.Lett. 178B (1986)365; D. Gross and E. Witten, Nucl.Phys. B277(1986)1

(39) L. Dixon, J.A. Harvey, C. Vafa and E. Witten, Nucl.Phys. B261(1985)678; B276 (1986)285; C.S. Hamidi and C. Vafa, Caltech preprint; L. Dixon, D. Friedan, E. Martinec and S. Shenker, Chicago preprint

(40) E. Witten, Nucl.Phys. B268(1986)79

(41) M. Dine, N. Seiberg, X.G. Wen and E. Witten, Princeton preprint; J. Ellis, C. Gomez, D.V. Nanopoulos and M. Quiros, Phys.Lett. 173B(1986)59

(42) D.J. Gross, J.A. Harvey, E. Martinec and R. Rohm, Phys.Rev.Lett. 54(1985) 502; Nucl.Phys. B256(1985) 253; B267; (1986)75

(43) E. Witten, Nucl.Phys. B258(1985)75; Y. Hosotani, Phys.Lett. 126B(1983)309; 129B(1983)193

(44) A. Strominger and E. Witten, Comm. Math.Phys. 101(1985)341; A. Strominger, Phys.Rev.Lett.55(1985)2547

(45) B.R. Greene, K.H. Kirklin, P.J. Miron and G.G. Ross, Oxford preprint

(46) For a review see H.P. Nilles, Phys.Rept. 110(1984)1

(47) L.E. Ibañez, Phys.Lett. 118B(1982)73; J. Ellis, D.V. Nanopoulos and K. Tamvakis, Phys.Lett.121B(1983)123; H.P. Nilles, Phys.Lett. 115B(1982)193; K. Inoue, A. Kakuto, H. Komatsu and S. Takeshita, Prog.Theo.Phys. 68(1982) 927; L. Alvarez Gaume, J. Polchinski and M. Wise, Nucl.Phys. B221(1983)495

(48) E. Cremmer et al, Phys.Lett. 79B(1978) 231; Nucl.Phys. B147(1979)105; E. Cremmer et al, Phys.Lett. 116B(1982)231; Nucl.Phys. B212(1983)413; R. Barbieri, S. Ferrara and C.A. Savoy, Phys.Lett. 119B(1982)343; R. Arnowitt, A.H. Chamseddine and P. Nath, Phys.Rev.Lett. 49 (1982)970;50(1983)232

(49) L.E. Ibañez, these proceedings

(50) J.P. Derendinger, L.E. Ibañez and H.P. Nilles, Phys.Lett. 155B(1985)65; M. Dine, R. Rohm, N. Seiberg and E. Witten, Phys.Lett. 156B(1985)55

(51) J. Breit, B. Ovrut and G. Segré, Phys.Lett. 162B(1985)303; P. Binétruy and M.K. Gaillard, Phys.Lett. 168B (1986)347

(52) P. Binétruy, S. Dawson and I. Hinchliffe, LBL-21606; S. Dawson these proceedings

(53) J.P. Derendinger, L.E. Ibañez and H.P. Nilles, Nucl.Phys. B267(1986)365; L.E. Ibañez and H.P. Nilles, Phys.Lett. 169 B(1986)354

(54) M. Dine and N. Seiberg, Phys.Lett 162B (1985)299; and in Unified String Theories (World Scientific, Singapore 1986)

(55) E. Witten, Nucl.Phys. B258(1985)75, M. Dine, V. Kaplunovsky, M. Mangano, C. Nappi and N. Seiberg; Nucl.Phys. B259 (1985)549; G. Cecotti, J.P. Derendinger, S. Ferrara, L. Girardello and M.

Roncadelli, Phys.Lett. 156B(1985)318; J. Breit, B. Ovrut and G. Segre, Phys. Lett. 158B(1985)33; F. del Aguila, G. Blair, M. Daniel and G.G. Ross, Nucl. Phys. B272(1986)413; P. Binétruy, S. Dawson, I. Hinchliffe and M. Sher, Nucl.Phys. B273(1986)501; T. Matsuoka and D. Suematsu, Nagoya preprint; L.E. Ibañez and J. Mas, CERN preprint TH 4426/86; G.G Ross, these proceedings

(56) R. Slansky, Phys.Rept. 79(1981)1

(57) J.P. Ader, S. Narison and J. Wallet, Montpellier preprint (#7889); G. Belanger and S. Godfrey, Phys.Rev. D34(1986) 1309(#7730 and #7749); T. Matsuoka, H. Mino, D. Suematsu and S. Watanabe, Nagoya preprint (#7617); T.G. Rizzo, Iowa State preprint (#7773); D. London and J. Rosner, Phys.Rev. D34 (1986)1530 (#9288)

(58) S. Barr, Phys.Rev.Lett. 55(1985)2778

(59) E. Witten, Ref 43

(60) V. Barger, N.G. Deshpande and K. Wishnant, Phys.Rev.Lett. 56(1986)30; L.S. Durkin and P. Langacker, Phys. Lett. 166B(1986)436; J. Ellis, K. Enqvist, D.V. Nanopoulos and F. Zwirner, Nucl.Phys. B276(1986)14; F. del Aguila, G.A. Blair, M. Daniel and G.G. Ross, CERN TH4376; D. London and J. Rosner, Phys.Rev. D34(1986)1530

(61) N.G. Deshpande, these proceedings

(62) V.D. Angelopoulos, J. Ellis, D.V. Nanopoulos and N.D. Tracas, Phys.Lett. 176B(1986)203

(63) M. Cvetic and B.W. Lynn, SLAC PUB 3900; P.J. Franzini, SLAC PUB 3920; P.J. Franzini and F.J. Gilman, SLAC PUB 3932

(64) For an estimate of the accuracy expected for A_{FB} at LEP, see P. Baillon et al, Physics at LEP CERN 86-02

(65) M. Dine et al, Ref 55

(66) E. Cohen, J. Ellis, K. Enqvist and D.V. Nanopoulos, Phys.Lett. 165B (1985)76; J. Ellis, K. Enqvist, D.V. Nanopoulos and F. Zwirner, Ref 60 and Mod.Phys.Lett. A1(1986)57

(67) I am using here the same notation for scalars as that for the fermions in table III.1. I hope this is not confusing

(68) F. del Aguila et al, Ref 55

(69) E. Witten, Ref 43; J. Breit et al, Ref 55; A. Sen, Phys.Rev.Lett. 55(1985)33

(70) B.R. Greene, K.H. Kirklin, P.J. Miron and G.G. Ross, Ref 45; G.G. Ross, these proceedings

(71) S.T. Yau in Proceedings of the Argonne Symposium on Anomalies Geometry and Topology (World Scientific, Singapore 1985)

(72) See for example, D.V. Nanopoulos, Proceedings of the XXII International Conference of High Energy Physics, Leipzig 1984

(73) P. Fayet, Phys.Lett. 69B(1977)489; G. Farrar and P. Fayet, Phys.Lett. 76B (1978)575

(74) J. Ellis, J.S. Hagelin, D.V. Nanopoulos, K.A. Olive and M. Srednicki, Nucl. Phys. B238(1984)453

(75) L. Ibañez, Phys.Lett. 137B(1984)160; J.S. Hagelin, G.L. Kane and S. Raby, Nucl.Phys. B241(1984)638

(76) For a discussion, see for example, J. Ellis, K. Enqvist, D.V. Nanopoulos and F. Zwirner, Mod.Phys.Lett. A1(1986)57

(77) For a critical discussion see: J. Ellis in Proceedings of the 1985 International Symposium on Lepton and Photon Interactions at High Energies, Kyoto, Japan

(78) A. Homna, these proceedings

(79) R.H. Barnett, H.E. Haber and G.L. Kane, Nucl.Phys. B267(1986)625

(80) See for example, K. Hagiwara in The Quark Structure of Matter (World Scientific, Singapore 1986)

(81) S. Weinberg, Phys.Rev.Lett. 50(1983) 387; for a general discussion, see for example, H.E. Haber and G.L. Kane, Phys.Rept 117(1985)75

(82) R. Arnowitt, these proceedings

(83) H. Baer and E.L. Berger, Phys.Rev. D34 (1986)1361

(84) E. Reya and D.P. Roy, Phys.Lett. 166B

(1986)223 and Zeit.f.Phys.C to be published

(85) E. Reya, these proceedings

(86) R.M. Barnett and H.E. Haber, to be published

(87) V.L. Ginzburg and L. Landau, Zh.Eksp. Teor.Fiz. 20(1950)1064; J. Bardeen, L.N. Cooper and J.R. Schrieffer, Phys. Rev.108(1957)1175

(88) L. Susskind, Phys.Rev. D20(1979)2619; S. Weinberg, Phys.Rev. D19(1979)1277

(89) For a recent review see M.E. Peskin, Proceedings of the 1985 International Symposium on Lepton and Photon Interactions at High Energies, Kyoto, Japan

(90) H. Harari, these proceedings

(91) C. Burges and H. Schnitzer, Nucl.Phys. B228(1983)464; C.N. Leung, S. Love and S. Rao, Zeit.f.Phys. 31C(1986)433; W. Buchmüller and D. Wyler, Nucl.Phys. B268(1986)621

(92) D. Wyler, these proceedings

(93) S. Dimopoulos and L. Susskind, Nucl.Phys. B155(1979)237, E. Eichten and K. Lane, Phys.Lett. 90B(1980)125

(94) D.B. Kaplan, these proceedings; H. Georgi and D.B. Kaplan, Phys.Lett. B145(1984)216; H. Georgi, D.B. Kaplan and B. Galison, Phys.Lett. B143(1984) 152; M.J. Dugan, H. Georgi and D.B. Kaplan, Nucl.Phys. B254(1985)299; for and extensive discussion see H. Georgi, Proceedings of the 1985 Les Houches Summer School

(95) M.E. Peskin, Nucl.Phys. B175(1981)197; J. Preskill, Nucl.Phys. B177(1981)21

(96) M.J. Dugan, H. Georgi and D.B. Kaplan, Ref 94

(97) H. Georgi, Phys.Lett. B151(1985)57; H. Georgi and J. Preskill, Phys.Lett. B156(1985)369; see also H. Georgi, Ref 94

(98) B. Schrempp, these proceedings

(99) L. Abbott and E. Farhi, Phys.Lett. 101B(1981)69; Nucl.Phys. B189(1981)547

(100) M. Claudson, E. Farhi and R.L. Jaffe, Phys.Rev. D34(1986)873

(101) G. 't Hooft in Recent Development in Field Theory (Plenum Press, New York, 1980)

(102) C. Korpa and Z. Ryzak, MIT preprint

(103) J. Wudka, Phys.Lett. 167B(1986)337