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ABSTRACT

In these lectures I discuss the issue of the origin of the quark and lepton massés, both in the case in which these objects are elementary and in the case they are composite. Some of the generic predictions and dynamical assumptions of GUTS, family symmetry models and superstrings are detailed. They are contrasted to the dynamics required for composite models of quarks and leptons. In this latter case, the difficulties of protecting dynamically fermion masses and yet still generating intra and interfamily hierarchies is emphasized.

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I. What is the Origin of the Quark and Lepton Masses?

Probably one of the deepest open questions in particle physics today is what fixes the spectrum of quarks and leptons, as shown schematically in Fig. 1. Although there are some overall regularities - leptons are lighter than quarks, in each generation; the average mass increases with generation number; etc. - there is really no clear pattern in the intra- and inter-generational mass splittings. Furthermore, the physical reason for the appearance of the generational replications is itself an open question. In these lectures I shall describe a variety of attempts which have been pursued to explain this mass pattern. Although my own prejudice is that the quark and lepton masses are the result of a deep underlying bound state problem, I have tried to emphasize in this report also alternative ideas, in which quarks and leptons retain their elementarity. As I will try to make clear, composite models of quarks and leptons face very serious dynamical questions, whose answers are largely unknown. Hence, these models are only a sensible alternative, if the fermion mass issue cannot be resolved otherwise. To be able to judge whether this is the case or not, one must also explore thoroughly the "elementary" quark and lepton option.

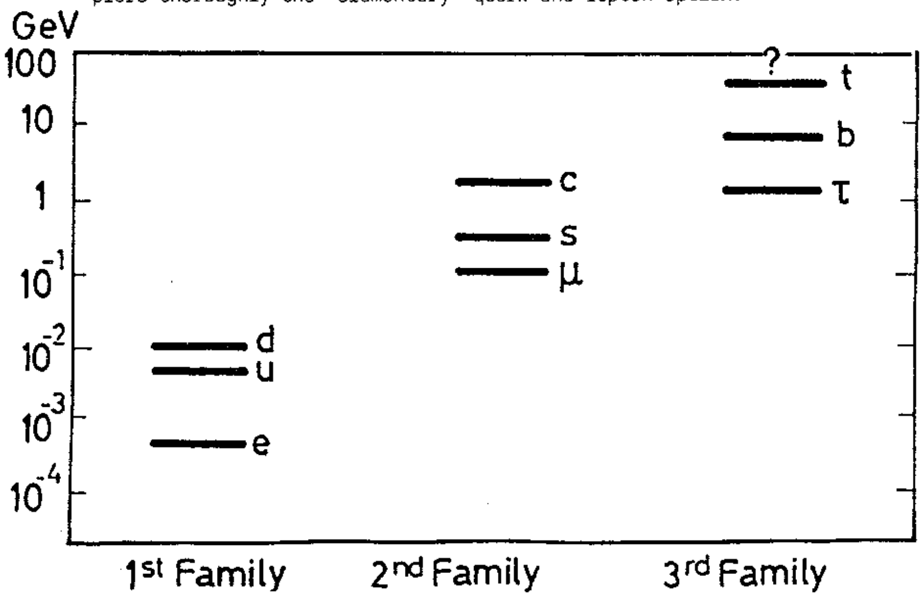

Fig. 1: Schematic quark and lepton mass spectrum

To begin the discussion, it is useful to recall how the quark and lepton

masses appear in the standard electroweak theory /1/. As is well known, fermion masses involve left-right transitions. That is, one has

$$\mathcal{L}_{mass} = -m[\bar{\psi}\psi] = -m[\bar{\psi}_L\psi_R + \bar{\psi}_R\psi_L] \quad \text{(I.1)}$$

where ψ_L, ψ_R are the usual helicity projections:

$$\psi_L = \tfrac{1}{2}[1-\gamma_5]\psi \quad ; \quad \psi_R = \tfrac{1}{2}[1+\gamma_5]\psi \quad \text{(I.2)}$$

Because ψ_L and ψ_R, for both quarks and leptons, have different SU(2) x U(1) transformation properties, no direct SU(2) x U(1) conserving mass term is allowed in the standard model *. Of course, since SU(2) x U(1) is spontaneously broken to $U(1)_{em}$, eventually quark and lepton mass can appear, but only as the result of the spontaneous breakdown!

In the standard model one introduces a Higgs doublet $\Phi = \begin{pmatrix}\phi^0\\ \phi^-\end{pmatrix}$, with an asymmetric self interaction

$$V = \lambda\left(\Phi^\dagger\Phi - \frac{v^2}{2}\right)^2 \quad \text{(I.3)}$$

so as to cause the spontaneous breakdown of SU(2) x U(1). The field Φ has a vacuum expectation value $\langle\Phi\rangle = \frac{v}{\sqrt{2}}\begin{pmatrix}1\\0\end{pmatrix}$ which preserves charge, but breaks all the remaining symmetries in SU(2) x U(1). The existence of a nonvanishing vacuum expectation value for Φ can generate quark and lepton masses. Since Φ is a doublet under SU(2),

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* Right-handed neutrinos, if they existed, could in principle have an SU(2) x U(1) invariant mass term. This term, however, necessarily would violate lepton number.

Yukawa couplings of Φ to the fermions in the theory are allowed. These trilinear fermion-fermion-Higgs terms become mass terms for the quarks and leptons, when Φ is replaced by its vacuum expectation value $\langle\Phi\rangle$. For instance, the SU(2) x U(1) invariant interaction

$$\mathcal{L}_{\text{Yukawa}} = -h\left\{(\bar{u},\bar{d})_L \begin{pmatrix}\phi^0\\ \phi^-\end{pmatrix} u_R + \bar{u}_R\,(\phi^{0*},\phi^+)\begin{pmatrix}u\\ d\end{pmatrix}_L\right\} \tag{I.4}$$

when $\Phi \rightarrow \langle\Phi\rangle$ gives rise to a mass term for the u-quark

$$m_u = \frac{h\,v}{\sqrt{2}} \tag{I.5}$$

Note, however, that although v is known, being related to the Fermi constant characterizing the breakdown:

$$v = (\sqrt{2}\,G_F)^{-1/2} \simeq 250\ \text{GeV} \tag{I.6}$$

the mass m_u is <u>unknown</u> since it depends on the <u>unknown</u> coupling h.

The above considerations are easily generalized to the case of many families. The most general SU(2) x U(1) invariant Yukawa couplings between Φ and the fermions in the theory reads

$$\begin{aligned}\mathcal{L}_{\text{Yukawa}} = &-h^u_{ij}\left\{(\bar{u}_i\bar{d}_i)_L \begin{pmatrix}\phi^0\\ \phi^-\end{pmatrix} u_{jR} + h.c.\right\}\\ &+ h^d_{ij}\left\{(\bar{u}_i\bar{d}_i)_L \begin{pmatrix}\phi^+\\ -\phi^{0*}\end{pmatrix} d_{jR} + h.c.\right\}\\ &+ h^l_{ij}\left\{(\bar{\nu}_{l_i}, l_i)_L \begin{pmatrix}\phi^+\\ -\phi^{0*}\end{pmatrix} l_{jR} + h.c.\right\}\end{aligned} \tag{I.7}$$

where i,j are family indices and I have assumed that there are no

right-handed neutrinos. The coefficients h^u_{ij}, h^d_{ij} and h^l_{ij} - just like h above - are unknown. When Φ gets replaced by its vacuum expectation value $\langle\Phi\rangle$ the Yukawa couplings of Eq. (I.7) generate mass matrices for the charge 2/3 and charge -1/3 quarks and the charge -1 leptons:

$$M^f_{ij} = \frac{1}{\sqrt{2}}\,h^f_{ij}\,v \qquad f = \{u,d,l\} \tag{I.8}$$

The eigenvalues of M^f_{ij} are the quark and lepton masses, while the charged current mixing angles (the Cabibbo angle in the case of two generations) are related to the unitary matrices which diagonalize M^d and M^u.

The standard electroweak model allows for the existence of quark and lepton masses and mixing angles once SU(2) x U(1) $\rightarrow$ $U(1)_{em}$. However it <u>does not predict</u> these parameters since the Yukawa couplings h^f_{ij} are arbitrary. It is clear that if one wants to predict the mass spectrum of quarks and leptons one must go <u>beyond</u> the standard model. Basically two rather distinct options appear to be open. Either

i) One believes in the existence of a Higgs doublet Φ and the associated Yukawa couplings of Eq. (I.7). However, in contrast to what happens in the standard model the h^f_{ij} are no longer arbitrary but are fixed by a theory beyond SU(2) x U(1).

or

ii) One does not believe in the existence of a Higgs doublet and Yukawa couplings. Fermion masses then must arise as solutions of some bound state problem.

The first option above is, in a sense, a perturbative approach to the generation of mass. Since the quark and lepton masses are much smaller than $v \sim 250$ GeV, the Yukawa couplings h^f_{ij} are small numbers. The physics associated with determining these numbers need not have anything to do with a strong coupling bound state problem. Hence option i) is contrary in spirit to the idea that quarks and leptons are composite, although even in these cases one could think of the Higgs field and the Yukawa couplings as being the result of an effective theory.* The second option is clearly a non perturbative option: quarks and leptons are composite of yet more elementary objects - preons - and their peculiar spectrum is the result of the underlying dynamics. Furthermore, at some level, this dynamics must break SU(2) x U(1) since no quark and lepton masses can arise unless this symmetry is spontaneously broken.

II. Anti Composite Ideas and their Generic Predictions

In this section I want to explore some of the ideas put forth to "explain" the quark and lepton mass spectrum, assuming that these objects are elementary. In this instance - which corresponds to option i) above - it is necessary to invent some physics beyond the standard model to fix the values of the Yukawa couplings h^f_{ij} **. A good first

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* This point of view has been recently explored in a series of very interesting papers by H. Georgi and collaborators /2/.

** In some cases, for example for extended technicolor /3/, there may be no real Yukawa coupling, but only effective mass generating interactions.

example is provided by grand unified theories (GUTS) /4/. Actually, although GUTS do predict some interrelations among Yukawa couplings they do not fix all Yukawa couplings. So GUTS can only provide some partial information on the quark and lepton mass pattern. Even so, it is interesting and illustrative to show how this comes about for the simplest of all GUTS, SU(5) /5/.

The idea behind GUTS is that at very high energy there is only one gauge interaction, instead of the separate electroweak and strong interactions. At some very high scale, the GUT symmetry suffers a spontaneous breakdown to SU(3) x SU(2) x U(1). This scale can be estimated by studying the evolution of the SU(3), SU(2) and U(1) coupling constants. If something like GUTS is correct, these coupling constants should approach each other at the unification scale M_X. Extrapolating the behaviour of the standard SU(3) x SU(2) x U(1) model over many orders of magnitude one finds /4/, amazingly enough, that the running coupling constants seem to merge at $M_X \sim 10^{14} - 10^{15}$ GeV, lending support to the whole GUT idea.

The fermions in the theory - plus perhaps a few additional states - for consistency must sit in some GUT representation. It is in fact easy to see that this is the case for the SU(5) GUT. Describing all fermions in terms of their left handed projections, using that $\psi_R \sim \psi^c_L$ *, it is easy to convince oneself that the 15 fermions of one generation comprise a $\bar{5}$ and a 10 representation of SU(5). To wit, one has

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* Here ψ^c is the charge conjugate to ψ.

$$\bar{5} = \left\{ d^c_L \, , \begin{pmatrix} \nu_e \\ e \end{pmatrix}_L \right\} \equiv \psi_a \quad (a=1,\dots,5) \tag{II.1}$$

$$10 = \left\{ \begin{pmatrix} u \\ d \end{pmatrix}_L , u^c_L , e^c_L \right\} \equiv \psi^{ab} = - \psi^{ba} \quad (a,b=1,\dots,5)$$

Also the Higgs fields in an SU(5) GUT must transform irreducibly. For these purposes the doublet Higgs field Φ of the standard model must be augmented by an additional SU(3) triplet field, to form an SU(5) quintet: $H^a \sim 5$. In contrast to the three possible SU(2) x U(1) invariant Yukawa couplings of Eq. (I.7), SU(5) allows only two possible couplings. Thus the SU(5) GUT theory has some non trivial prediction for the quark and lepton masses. The allowed Yukawa couplings are

$$\mathcal{L}_{\bar{5}\bar{5}10} = (h_1)_{ij} \left\{ \psi^T_{ai} \, C \, \psi^{ab}_j \, H^+_b + h.c. \right\} \tag{II.2a}$$

and

$$\mathcal{L}_{51010} = -(h_2)_{ij} \left\{ \psi^{T\,ab}_i C \psi^{de}_j H^c \epsilon_{abcde} + h.c. \right\} \tag{II.2b}$$

where C is a charge conjugation matrix. If the Higgs field H has the vacuum expectation value $\langle H \rangle = \frac{v}{\sqrt{2}} \begin{pmatrix} 0 \\ 0 \\ 0 \\ 0 \\ 1 \end{pmatrix}$, then it is easy to check that the mass matrices are

$$\begin{aligned} M^d_{ij} &= M^l_{ij} = \frac{v}{\sqrt{2}} (h_1)_{ij} , \\ M^u_{ij} &= \frac{v}{\sqrt{2}} (h_2)_{ij} \end{aligned} \tag{II.3}$$

Hence at the scale M_x, where the grand unified theory is presumed to hold, the SU(5) GUT predicts the equality of the charge -1/3 quark and charge -1 lepton masses

$$m_\tau = m_b \Big|_{q^2 = M_x^2} \; ; \; m_\mu = m_s \Big|_{q^2 = M_x^2} \; ; \; m_e = m_d \Big|_{q^2 = M_x^2} \tag{II.4}$$

These equations can be extrapolated back to low q^2 by using the renormalization group. The principal effect comes from the running of the strong coupling constant and one finds, for instance:

$$\frac{m_b(q^2)}{m_\tau(q^2)} \simeq \left[\frac{\alpha_3(q^2)}{\alpha_{GUT}(M_x^2)} \right]^{\frac{4}{11 - \frac{4}{3} n_f}} \simeq 2.8 \tag{II.5}$$

Here $n_f = 3$ is the number of families and the numerical result is that corresponding to $q^2 \sim (5 \text{ GeV})^2$ /6/. This very nice result, which agrees with experiment, however, is also accompanied by a bad prediction. Using Eq. (III.4) one has

$$\frac{m_s(q^2)}{m_d(q^2)} = \frac{m_\mu}{m_e} \simeq 206 \tag{II.6}$$

while the best current algebra prediction gives only about 20 for this ratio! /7/

What can one conclude from these considerations? First and foremost that the assumptions of having an SU(5) GUT are <u>not</u> sufficient to determine the quark and lepton mass spectrum. After all, M^u is arbitrary and the result $M^d = M^l$ at M_x is not totally successful. Similar problems also plague other GUTS, so that in general there is agreement that GUTS per se cannot really illuminate the quark and lepton mass question. Furthermore, GUTS have fallen somewhat in disfavor since their most spectacular prediction - that of proton decay - has not received experimental confirmation /8/. The result (II.5) remains tantalyzing, but it might just be an accidental coincidence.

A less motivated set of theories than GUTS - involving the idea of a family symmetry - also "predict" Yukawa couplings. Basically, by assuming that quarks and leptons of different families transform irreducibly under a family group G_F, one allows only certain Yukawa couplings in the theory. In general, the Yukawa couplings take the form:

$$\mathcal{L}_{\text{Yukawa}} = h_{ij\ell}\, \bar{\psi}_{iL}\, \Phi_\ell\, \psi_{jR} + \text{h.c.}$$

where h_{ijl} are Clebsch Gordan coefficients of G_F and, usually, more than one Higgs field Φ_ℓ transforming irreducibly under G_F is involved. Although there have been a number of suggestions for the family group G_F (G_F = SU(3), U(1), etc.) /9/, none of these suggestions is really deeply motivated. Because of this circumstance, and the fact that detailed predictions of the various models are really crucially dependent on G_F, I will consider only generic predictions of family models, and no particular model in detail.

Perhaps the most characteristic feature of family models is that the symmetry G_F is not exact (after all $m_e \neq m_\mu$!). Although G_F can be assumed to be explicitly broken, this possibility is inelegant. Usually, therefore, one presumes that G_F is spontaneously broken. In fact this is rather natural. The Higgs fields Φ_ℓ carry both SU(2) x U(1) and G_F quantum numbers. Their expectation values $\langle\Phi_\ell\rangle$, besides breaking SU(2) x U(1) spontaneously, therefore also break down G_F. In general, two distinct possibilities are open: either i) G_F is a global symmetry of the theory or ii) G_F is a local symmetry. Both of these options give

rise to interesting physics "beyond" the standard model, which I shall now discuss.

If G_F is a global symmetry, the existence of $\langle\Phi_\ell\rangle \neq 0$ implies that certain Goldstone bosons must arise in the theory. These states are generically known in the literature as Majorons /10/ or Familons /11/. Because they do have zero mass, one can well imagine that they could lead to physical inconsistencies. However, if the scale of the breakdown of G_F - Λ_F - is much greater than the scale of the SU(2) x U(1) breakdown $v = (\sqrt{2}\, G_F)^{-1/2}$, one can avoid any direct physical contradictions /10/ - /12/. Obviously for $\Lambda_F \gg v$ to obtain, one of the Higgs fields Φ_ℓ which acquires vacuum expectation value must carry no SU(2) x U(1) quantum numbers. It is this field whose scale is Λ_F.

To appreciate how it is possible to have real Goldstone bosons in nature, recall that the typical Goldstone boson coupling to fermions is of the form

$$\mathcal{L}_{GB} = \frac{h_{ij}}{\Lambda_F}\, \bar{\psi}_i \gamma_5 \gamma^\mu \psi_j\, \partial_\mu \phi \tag{II.7}$$

where h_{ij} is a coefficient of O(1) and Λ_F is the scale of the breakdown. The derivative coupling above implies that the effective coupling constant of ϕ to fermions is of the order $g_{ij} \sim \frac{(m_i - m_j)}{\Lambda_F}$, which is very small if $\Lambda_F \gg v \approx 250$ GeV. The long range force, due to massless ϕ exchange, runs into no trouble if ϕ is a pseudoscalar (as implicitly assumed here). Because of the γ_5 in (II.7), instead of an r^{-1} potential one gets a spin-spin interaction proportional to r^{-3}:

$$V(r) = \frac{h^2}{4\pi\Lambda_F^2}\left\{\frac{1}{r^3}\left[\vec{\sigma}_1\cdot\vec{\sigma}_2 - 3\,(\vec{\sigma}_1\cdot\hat{r})(\vec{\sigma}_2\cdot\hat{r})\right]\right\} \quad \text{(II.8)}$$

The presence of such non magnetic tensor interactions in nature has been looked for by Ramsey /13/ and by Feinberg and Sucher /14/. Their investigations exclude values of $\Lambda_F \lesssim$ 10-100 GeV. Thus if $\Lambda_F \gg v$ no contradiction ensues from these considerations.

Much more stringent bounds on Λ_F follow, however, from the non observation of decays like $\mu \to e\,\phi$ or $K \to \pi\,\phi$, which a Lagrangian like (II.7) could in principle allow. * From the analysis of Gelmini, Nussinov and Yanagida /12/ these decays necessitate $\Lambda_F \gtrsim 10^9 - 10^{10}$ GeV. The natural theoretical question, therefore, for these family models is why the scale of breakdown of G_F is so much greater than that associated with the $SU(2) \times U(1) \to U(1)_{em}$ breakdown. Perhaps more challenging still is the study of possible experimental tests which might detect the ephemeral effects of the presence of this breakdown. Some very nice suggestions in this respect have been put forward by Moody and Wilczek /15/ and by Sikivie /16/, however their detailed discussion is beyond the scope of these lectures.

The second possibility, which one might envisage, connected to the spontaneous breakdown of the family symmetry G_F, is that G_F is really gauged.

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* Since one is dealing with family symmetries, it is to be expected that ϕ connect between different generations.

In this case no Goldstone bosons will appear in the theory, but one expects family gauge bosons to exist which - for instance - will change a μ into an e. A typical process like $K^0 \to \mu e$ would then be allowed in the theory, with an amplitude of order

$$A \sim \frac{g_{\mu e}\, g_{ds}}{M^2_{W_F}} \sim \frac{1}{\Lambda_F^2} \quad \text{(II.9)}$$

where I have used that the family gauge boson mass is of the order of the gauge coupling constant g times Λ_F. Typical bounds on flavour changing neutral currents then imply that $\Lambda_F \gtrsim$ 10-100 TeV /17/ and again one has to ask what physics forces $\Lambda_F \gg v$?

To summarize: although one can invent certain family symmetries to restrict the form of the Yukawa couplings, family symmetries are really not very predictive as far as quark and lepton masses go. One gets out basically what one puts in. Nevertheless, family symmetries might exist and one should continue active experimental probing of inter-family transitions like $\mu \to e\,\phi$ or $K^0 \to \mu e$. However, no one has really a good idea of what the scale Λ_F is, and what is the connection of this scale to the Yukawa couplings which fix the quark and lepton masses. It is clear that if one wants some insight into this latter problem, one needs a cleverer theory.

A cleverer theory may be at hand, in the form of superstring theories /18/. Although the practical difficulties to get reliable predictions out of these theories are daunting - as I shall explain - their potential and intrinsic interest makes them well worth discussing. In fact, in my opinion, they constitute the only real alternative to composite models, as a possible source for the explanation of the quark and lepton masses *. It is not my intention here to describe superstring theories in any detail (indeed I am far from an expert in these matters). Rather I want to concentrate on the aspects of these theories which may make them relevant for the quark and lepton mass question.

Superstrings are supersymmetric string theories in d-dimensional space time. They are theories which describe extended objects - strings - which are endowed both with a supersymmetry and a gauge group G. Green and Schwarz /20/ showed that these theories possess a number of remarkable properties: They are consistent only in d = 10 dimensions; their massless excitations contain among others a spin 2 graviton and calculations show that the string theory may provide a finite theory of gravity; for the specific gauge groups $G = SO(32)$ or $E_8 \times E_8$ the theory has no chiral or gravitational anomalies, and so only for these groups are superstring theories tenable.

The no anomaly result of Green and Schwarz, for the specific groups above, is particularly significant. The string has both massless and massive modes, where the string tension T sets the scale of the massive modes. Since the superstring theories want to be an alternative theory of gravity - with better short distance behaviour - it follows that

$$M^2_{\substack{\text{Massive}\\ \text{states}}} \sim T \sim M^2_{\text{Planck}} \tag{II.10}$$

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* Certain Kaluza Klein theories, which generically have many of the properties of superstrings, could provide also a viable explanation for quark and lepton masses /19/.

Obviously, for particle physics purposes only the massless modes of the string are therefore relevant. Dropping the massive excitations, the string theory reduces to a d = 10 field theory of the massless modes. This theory is a supergravity theory coupled to a Yang Mills supersymmetric theory based on the group G. These d = 10 field theories contain chiral fermions and were known /21/ to be inconsistent because they suffered from gravitational and chiral anomalies. It turns out, however, that the d = 10 field theory which emerges from the superstring theories has some additional pieces (Wess Zumino terms) which for $G = SO(32)$ or $E_8 \times E_8$ precisely cancel the unwanted anomalies. Therefore these theories are perfectly consistent and these groups are specially selected.

What do these theories have to do with our d = 4 real world? The belief is that the ground state of the superstring theories corresponds to configurations where 6 dimensions compactify, at a scale $M_{\text{Comp}} \sim M_{\text{Planck}}$, leaving one with a d = 4 field theory. This theory, in general, will contain the standard $SU(3) \times SU(2) \times U(1)$ model but all of it parameters are fixed, since the superstring theory and the ground state are unique! In particular, the Yukawa couplings between Higgs field(s) and fermions are not arbitrary but calculable in terms of the Yang Mills coupling g. Therefore quark and lepton masses are predictable, in principle.

There is, as yet, no hope of checking these assertions since no one knows how to find the superstring ground state. One has been able to make some phenomenological headway by trying to guess the compactification pattern and seeing if the guess has pleasing enough properties.

A particularly appealing suggestion, in this respect, has been made by Candelas, Horowitz, Strominger and Witten /24/. What these authors assumed is that the compactification from $d = 10$ to $d = 4$ occurred leaving an $N = 1$ supersymmetry exact. For this to obtain the 10 dimensional space has to be of the form $M_4 \times K$, where M_4 is the usual Minkowski space and K is a compact 6 dimensional space which is Ricci flat and has an SU(3) holonomy. These spaces are known as Calabi-Yau spaces /23/.

The consequences of this compactification, if the group of the superstring is $G = E_8 \times E_8$, are very remarkable at first glance. I enumerate them /22/:

i) Because of the SU(3) holonomy the gauge symmetry in $d = 4$ is reduced. Since $E_8 \supset E_6' \times SU(3)$ the $d = 4$ symmetry is at most $E_6 \times E_8$, but the E_6 group (depending on K) could be further broken down. The appearance of an E_6 is pleasing, since this is one of the natural GUT groups, of the exceptional sequence: $E_4 = SU(5)$; $E_5 = SO(10)$; E_6

ii) The chiral fermions which emerege in $d = 4$ must transform nontrivially under the holonomy group. Since the fermions are in the adjoint of E_8, and this representation in terms of $E_6 \times SU(3)$ reads

$$248 = (78,1) + (1,8) + (27,3) + (\overline{27},\bar{3}), \tag{II.11}$$

the chiral fermions necessarily are in 27's of E_6. This is also quite nice since the 27 does contain in a natural way the 16 quarks and leptons of one SO(10) generation

$$27 = 16+10+1 = (\bar{5}+10+1) + (\bar{5}+5)+1 \tag{II.12}$$

where the second result is the SU(5) decomposition

iii) The actual number of families of 27's and $\overline{27}$ is given purely by topological properties of the manifold K

$$n_{27} - n_{\overline{27}} = \frac{1}{2} \,|\chi(K)| \tag{II.13}$$

where $\chi(K)$ is the Euler characteristic of K.

iv) The Yukawa couplings of fermions to Higgs fields is fixed by harmonic analysis on the manifold K /24/. This is easily understood since the fermion fermion Higgs coupling arises by integrating over the compact directions the fermion fermion gauge couplings in $d = 10$ dimensions. The point is that the coefficients of the gauge fields in the compact direction are just scalar fields. If w denotes the $d = 10$ coordinates, x the $d = 4$ coordinates and y the compact coordinate then one has the expansion

$$A_m(w) = \sum_i A^i_m(x)\, \phi^i(y) + \sum_i \phi^i(x)\, A^i_m(y)$$

$$(m = 1, \ldots 10) \tag{II.14}$$

where $\phi^i(x)$ are 4-dimensional scalar fields. Hence if one could really perform the integral over d^6y and could do the relevant harmonic analysis on the Calabi-Yau spaces K, one could predict the Yukawa couplings at the compactification scale M. This would then fix the mass matrices for quarks and leptons provided one knew how to go from the scale M to the scale of a few GeV's.

Although these four results above are qualitatively quite spectacular, to render them quantitative is a very hard task indeed. First of all, there exist many Calabi Yau manifolds K and thus, even though the number of families is given by $\frac{1}{2}\chi(K)$, there is really no prediction on how many generations exist *. Furthermore, doing harmonic analysis on some of the manifolds K, which lead to just a few families, appears to be almost a prohibitive task. Fortunately, as Strominger has pointed out /25/, the structure of the Yukawa couplings appears to be determined entirely by topological considerations, which may bypass the need for a detailed harmonic analysis. Even so, it appears necessary to look elsewhere for confirmation of the superstring idea, rather than just wait for someone to compute the quark and lepton spectrum!

It behooves us, therefore, to examine the generic predictions of superstring theories, to see if one of these predictions can provide the "smoking gun" evidence for the reality of these beautiful, but rather fanciful, ideas. There appear to me to be three principal superstring "predictions" to date:

1) The existance of shadow matter
2) The necessity of having superpartners
3) A natural enlargement of the standard model symmetry group.

Let me discuss these in turn.

Recall that the $E_8 \times E_8$ symmetry group of the theory, after compactification, was reduced to $E_6 \times E_8$ (or a subgroup therein). Quarks and

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* Indeed, the simplest Calabi Yau manifolds have very large Euler characteristic /24/ ($\chi(K) \sim 10^2$) and are thus useless.

leptons then appear in the 27-dimensional representations of E_6. However, the theory has also other fermions which transform according to the 248 of the other E_8. These other fermions comprise what has been dubbed shadow matter /26/. After compactification, matter and shadow matter interact among themselves only gravitationally, and therefore very weakly. Shadow matter has obviously cosmological consequences. For instance, we know that primordial nucleosynthesis provides the bound $N_\nu \lesssim 4$ /27/. Clearly, if shadow and ordinary matter where in equilibrium and had the same density one would have too much Helium produced in the universe. This example is clearly too naive, but there may well be significant cosmological traces for shadow matter which would establish the validity of superstring ideas.

The shadow sector of the $E_8 \times E_8$ superstrings may have a more theoretical use. Since at compactification, by assumption, an N = 1 supersymmetry remains one must provide some plausible mechanism for how this supersymmetry breaks down. It has been suggested /28/ that this breakdown can perhaps be induced by the shadow matter, with the shadow matter playing the role of the hidden sector in N = 1 supergravity models /29/. That is, the shadow matter precipitates a breakdown of supersymmetry, which in turn induces a spontaneous breakdown of the electroweak SU(2) x U(1) theory. It is an open question whether the shadow matter can really catalyze this sequence of breakdowns, and whether one can follow perturbatively the behaviour of the theory from scales of the order of $M_{\text{compactification}}$ to ordinary electroweak scales. If one must necessarily go through a non perturbative step to go from $M_{\text{compactification}}$ to M_W, as Dine and Seiberg have argued /30/, then the knowledge of the Yukawa couplings after compactification (if they could be computed!) would not directly suffice to predict the quark and lepton mass spectrum. I fear that this may well be the Achilles heel of the whole superstring program.

Since supersymmetry is an essential ingredient of superstring theories, the discovery of supersymmetric partners of quarks and leptons would be of fundamental importance for these theories. Of course, since no one knows the scale of the breakdown of supersymmetry, the absence of supersymmetric partners at present energies is not an argument against superstrings. Nevertheless, since the effective field theory emerging from superstrings has elementary scalar fields, the naturalness of having a small Fermi scale vis a vis $M_{\text{compactification}}$ can only be guaranteed if there is a supersymmetry approximately valid at scales of O ($v \sim 250$ GeV). So it would be unnatural if no supersymmetry partners appeared at scales below, say, a TeV. Thus these ideas would really become rather forced if in the next decade or so no superpartners were found.

The final generic prediction of superstrings, which might be well to bear in mind, concerns a possible enlargement of the standard model. The group which remains after compactification generally can be much smaller than E_6, since the E_6, depending on the topological properties of K, can be further broken down /31/. Typically, one is left with a symmetry group, which is the standard model group times various U(1) factors. These extra factors have corresponding gauge bosons, which presumably end up with masses comparable to M_W. The phenomenology of these extensions of the standard model is being actively investigated /32/. Unfortunately, again, the predictions of superstrings are not

unique since they depend on K, which at the moment is unknown. Thus, although a positive signal would be tremendously exciting, a negative result can always be argued away.

III. Compositeness Dynamics - Mass protection

Having discussed some of the attempts to get at the quark and lepton spectrum, assuming that these objects are elementary, let me now turn to composite model ideas on the subject. Composite models for quarks and leptons, compared to superstring theories, are very much more pedestrian. In their defense, however, one should remark that there exist no evidence whatsoever for:

1. Strings with tension $T \sim 10^{38}$ GeV^2
2. Supersymmetry
3. Space-time dimensions $d > 4$

Imaging that the solution to the quark and lepton mass question is due to these states being composite bound states of preons, rather than elementary excitations, faces in contrast only one experimental problem *: There is no evidence whatsoever for any kind of substructure of quarks and leptons! Indeed as Perrottet and Renard /33/ have discussed in their lectures at this school, the experimental limits on the scale of compositeness are typically of the order $\Lambda_c \gtrsim$ Tev. (Very roughly speaking, the intrinsic size of quarks and leptons is measured by $\langle r \rangle \sim \frac{1}{\Lambda_c}$).

* The reader should please forgive my tongue in check!

The result $\Lambda_c \gtrsim$ TeV has profound significance for the dynamics of the preon theory. It means that the compositeness scale is very much greater than the masses of the quark and lepton bound states:

$$\Lambda_c \gg m_{q,1} \tag{III.1}$$

Hence, to a good approximation the preon dynamics must be able to produce certain bound states - to be identified, a posteriori, with the quarks and leptons - which are essentially massless. Such a dynamics is radically different from the bound states dynamics encountered in physics heretofore. For a weakly bound state object, like positronium, one has a bound state whose size is much greater than its Compton wavelength:

$$\langle r \rangle_{\text{positronium}} \sim \frac{1}{m\alpha} \gg \lambda_{\text{Compton}} \sim \frac{1}{m} \tag{III.2a}$$

For a strongly bound state object, like a proton, one has typically an intrinsic radius comparable to the Compton wavelength of the state

$$\langle r \rangle_{\text{proton}} \sim \lambda_{\text{Compton}} \sim \frac{1}{M_{\text{proton}}} \tag{III.2b}$$

What is required in the preon theory is that the inequality in (III.2a) be reversed. That is, one needs to bind states to a size much less than their Compton's wavelength

$$\langle r \rangle_{q,1} \sim \frac{1}{\Lambda_c} \ll \lambda_{\text{Compton}} \sim \frac{1}{m_{q,1}} \tag{III.2c}$$

Even though this is a novel dynamical requirement, one can devise models, with a built in mass protection mechanism, which give rise to massless spin 1/2 bound states. One can understand why the preon models must

have a mass protection mechanism if one imagines - as it is usually done - that the preon theory is a non Abelian gauge theory *. In such theories one has only one scale Λ_c, which can be taken as the dynamical scale where $\alpha_{\text{preon}}(\Lambda_c^2) \sim 1$. Clearly, all bound state masses will be proportional to this scale ($M \sim \Lambda_c$), unless one has some protection mechanism that forces some states to have $M = 0$. A very nice example, in this context, is provided by QCD. All hadrons (neglecting quark mass effects) except for pions have masses $M_{\text{hadron}} \sim \Lambda_{\text{QCD}}$, this being also the relevant scale of their intrinsic size: $\langle r \rangle_{\text{hadron}} \sim \frac{1}{\Lambda_{\text{QCD}}}$. The pions, however, in the limit of zero quark masses, have precisely $M_\pi = 0$, since they are the dynamical Goldstone bosons reflecting the breakdown of chiral symmetry in QCD.

In the literature, two mass protection mechanisms have been discussed for composite models. They involve either chiral protection /34/ or the quasi Goldstone fermion mechanism /35/. Models in which both protection mechanisms operate simultaneously are particularly interesting /36/. In the case of the chiral protection mechanism, massless bound state fermions arise if chirality, assumed to be a good symmetry at the underlying level, is preserved in the binding. The quasi Goldstone fermion mechanism, on the other hand, assumes that the underlying theory is supersymmetric, with some global symmetry G suffering a spontaneous breakdown to a subgroup H. As a result of this spont-

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* It would be nice to be able to sensibly discuss other alternatives for preon models than NAGT, but there has been very little work on this area.

aneous breakdown, dim G/H Goldstone bosons ensue. Because of the supersymmetry these Goldstone bosons have fermionic m = o partners, which are identified as the quarks and leptons.

Let me discuss both these options in a little more detail. Having a chiral symmetry in the underlying theory can give rise to massless bound states, only if the chirality is not spontaneously broken in the process of binding. QCD provides a well known example where chirality at the quark level does not imply massless fermion bound states. What happens in QCD, in the massless quark limit where a chiral symmetry obtains, is that condensates of the type $\langle \bar{u}u \rangle \neq o$ form. Such condensates break chirality spontaneously and no massless fermionic states bind. Rather, in QCD one has massive protons but massless pions, as a signal of the spontaneous breakdown of chirality. One is interested, at the preon level, in theories which behave completely differently than QCD. That is, chirality must be a good symmetry both at the underlying and bound state level in preon models.

't Hooft /34/ has spelled out necessary conditions which must obtain if chirality is to be preserved in the binding. These conditions require that the coefficients of the chiral anomalies at the preon and overlying level match. To explain the reasons why this anomaly matching is necessary, let me recall that the currents J_f^μ, associated with a given chiral symmetry, in general have Adler Bell Jackiw anomalies /37/. That is, although the theory has a chiral symmetry, the three point Green functions $\Gamma^{\mu\nu\lambda}$ has an anomalous divergence:

$$q_{3\mu} \Gamma^{\mu\nu\lambda} = A_{preon} \epsilon^{\nu\lambda\alpha\rho} q_{1\alpha} q_{2\rho} \qquad \text{(III.3)}$$

The anomaly coefficient A_{preon} can be computed at the underlying level from the triangle graph shown in Fig. 2.

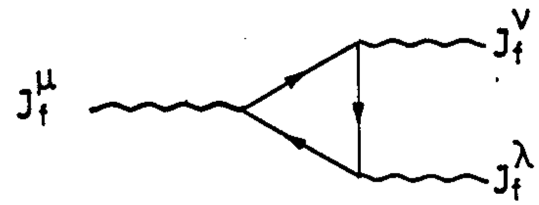

Fig. 2: Triangle graph connected to the chiral anomaly.

't Hooft condition for the preservation of chirality in the binding is that the anomaly coefficient A_{preon} must be matched at the bound state level. That is, for chirality to remain unbroken, it is necessary that the anomaly coefficient $A_{Bound\ State}$ computed by including the presumed m = o bound states in the triangle graph of Fig. 2 be precisely the same as A_{preon}.

$$A_{Bound\ State} = A_{preon} \qquad \text{(III.4)}$$

The necessity of Eq. (III.4) follows because the chiral anomaly implies that the 3 point Green's function $\Gamma^{\mu\nu\lambda}$ at the symmetry point $q_1^2 = q_2^2 = q_3^2 = q^2$ is singular /38/:

$$\Gamma^{\mu\nu\lambda}\Big|_{q_i^2=q^2} = \frac{A_{preon}}{q^2} \left\{ \epsilon^{\mu\nu\alpha\rho} q_{1\alpha} q_{2\rho} q_3^\lambda + \text{cycl. terms} \right\} + \text{non sing. terms} \qquad \text{(III.5)}$$

Such a singularity must also appear in $\Gamma^{\mu\nu\lambda}$ by calculating it at the bound state level. Obviously, if the chiral symmetry is broken, so that m = o bound state Goldstone bosons exist, the graph of Fig. 3 will provide such a singularity. One has in this case

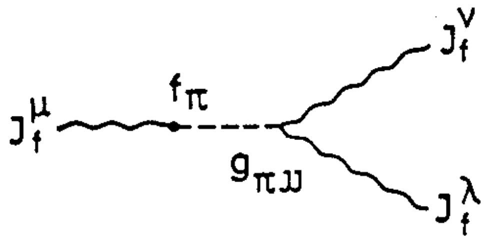

Fig. 3: Singular contribution to $\Gamma^{\mu\nu\lambda}$, if chirality is broken

$$\Gamma^{\mu\nu\lambda} = f_\pi \frac{1}{q^2} g_{\pi JJ} \left\{ \epsilon^{\mu\nu\alpha\beta} q_{1\alpha} q_{2\beta} q_3^\lambda + \text{cyclic terms} \right\} \qquad \text{(III.6)}$$

and the matching of the q^2 singularity gives a low energy theorem for the coupling of the Goldstone boson to the chiral currents

$$g_{\pi JJ} = A_{preon}/f_\pi \qquad \text{(III.7)}$$

The q^2 singularity, however, can also be reproduced if m = o fermion bound states exist. In this case the triangle graph of Fig. 2, with the m = o bound states going around the loop, will produce the q^2 singularity. Obviously, for consistency, Eq. (III.4) must obtain and in this case chirality is preserved in the binding. One can show /39/ that the q^2 singularity in $\Gamma^{\mu\nu\lambda}$ can only be reproduced by one of the above options. That is, by either having Goldstone bosons or by having massless m = o fermions.

There are also some general dynamical results that can be stated for the case of the quasi Goldstone fermion mechanism. In this case, as I mentioned above, one is dealing with a supersymmetric theory at the preon level which has a global-symmetry G, spontaneously broken to another symmetry H. The breakdown causes the appearance of dim G/H bound state Goldstone bosons and the supersymmetry requires one to have also m = o bound state fermionic partners. More precisely, it is necessary that the number of fermionic degrees of freedom for the m = o states match those of the bosonic degrees of freedom. However, it is not necessary that all the m = o bosonic degrees of freedom be associated with actual Goldstone bosons. Just as there must be quasi Goldstone fermion partners of the Goldstone bosons (QGF) there may well be bosonic m = o partners of the Goldstone bosons - quasi Goldstone bosons (QGB). Supersymmetry requires that the number of QGF, QGB and Goldstone bosons (GB) obey

$$2n_{QGF} = n_{GB} + n_{QGB} \qquad \text{(III.8)}$$

Although n_{GB} is fixed by the breakdown $G \to H$, n_{QGB} and therefore n_{QGF} is a dynamical issue. One can show /40/ that n_{QGB} ranges from

$$1 \leq n_{QGB} \leq n_{GB} \qquad \text{(III.9)}$$

Two remarks ought to be added with respect to these mass protection mechanisms. First, it may well be in the case of the quasi Goldstone mechanism that, in the breakdown $G \to H$, the remaining unbroken group is chiral. If this is the case, then it is necessary to check that the anomalies of the group H at the preon level and at the bound state level match. If they do not, then H cannot remain as an unbroken chiral group. In some very nice examples - one of which I shall discuss in more detail below - it turns out that the bound state fermions needed to match the H anomalies are precisely the m = 0 fermions required to bind as supersymmetric partners of the G/H Goldstone bosons. In these cases one has a double protection /36/. The m = o fermions are present in the theory both because they are required by supersymmetry and because they are required by chirality.

The second remark concerns many of the simpler popular composite models in the literature (Rishons, Haplons, etc.) in which one constructs quarks and leptons rather straightforwardly algebraically. These models, in general, do not have a dynamical protection mechanism which guarantees that their simple bound states (quarks and leptons) are light - i.e. have m = o in some appropriate limit. These models, to my mind, are dynamically suspect and probably ought to be discarded *.

I want to illustrate these ideas with a model /41/ - the Novino model - which has both chirality protection and the quasi Goldstone fermion mechanism. The underlying preon theory is a supersymmetric SU(2) gauge theory and one has six preon supermultiplets described by chiral superfields ϕ_α^p, with α = 1,2 being an SU(2) index and p = 1,...6 being a "flavor" index. The superfields ϕ_α^p, of course, describe both a

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* One should keep in mind, however, that new dynamical principles may one day render these models again viable. The quark model of hadrons is a salutary example, in this respect.

left handed fermion and two scalars for each p and α. The global symmetry of this model is SU(6) x $U(1)_x$ where x is a linear combination of preon number and R-symmetry, which has no SU(2) anomalies /41/. This theory possesses a natural condensate

$$v = \langle \epsilon^{\alpha\beta} \phi_\alpha^1 \phi_\beta^2 \rangle \qquad \text{(III.10)}$$

which, if it is nonvanishing, breaks SU(6) x $U(1)_x$ to SU(4)xSU(2)x$U(1)_{x'}$. Hence this model will have some m = o fermions as partners of the Goldstone bosons arising from the breakdown

$$SU(6) \times U(1)_x \rightarrow SU(4) \times SU(2) \times U(1)_{x'} \qquad \text{(III.11)}$$

Furthermore, since for the preons ϕ_α^p the global symmetry H = SU(4) x SU(2) x $U(1)_{x'}$, is a chiral symmetry, it is necessary that in the theory there should be m = o fermionic bound states to match the anomalies at the preon level of the currents in H.

A possible pattern of GB, QGF and QGB from the breakdown in (III.11) is given by

$$\begin{aligned} GB &\sim (4,2) + (\bar{4},2) + (1,1) \\ QGF &\sim (4,2) + (1,1) \\ QGB &\sim (1,1) \end{aligned} \qquad \text{(III.12)}$$

This pattern is in fact dynamically favored by complementarity /41/ and produces a set of left handed bound state fermions, in the (4,2) representation, which can be identified with the left-handed quarks and leptons of one generation (the (1,1) state is an extra state - the novino). Remarkably, one can check that the bound state fermions in

Eq. (III.12) are precisely those necessary to match the chiral anomalies of H at the preon level. Hence, in this model one has found a set of bound state massless fermions which are massless both because of supersymmetry and because of chirality (Double protection).

The novino model, although very instructive, is quite primitive. For instance only left handed quarks and leptons of one generation are bound. Furthermore some extra state, the novino, is also produced *. This model can be fairly naturally enlarged, so that also right handed quarks and leptons are bound /42/. One just replaces the underlying SU(2) gauge interaction by an SU(2) x SU(2)', and introduces another set of preons $\phi^{\prime p'}_{\alpha'}$ transforming under the SU(2)' to give the right handed states. Furthermore, by introducing a left-right binding field $\chi_{\alpha\alpha'}$, which transform as (2,2) under the gauge group, one can even contemplate vacuum expectation values

$$v_F = \langle \phi^1_\alpha \chi^{\alpha\alpha'} \phi^{1'}_{\alpha'} \rangle = \langle \phi^2_\alpha \chi^{\alpha\alpha'} \phi^{2'}_{\alpha'} \rangle \qquad \text{(III.13)}$$

which can provide a breaking of the electroweak gauge group, so that the preon model acts as technicolor. A much harder task, however, is to incorporate generations. This can be done, rather mechanically, by extending the "flavor" index from p = 1,...6 to p = 1,... $4n_f + 2$, so that in fact one gets $(4n_f,2) \supset n_f(4,2)$ bound states. Perhaps more imaginatively, one can replace the gauged SU(2) group by an SU(6) group, and then try to get generations as extra replications due to anomaly

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* The novino is quite elusive. In principle it acts very much like a neutrino, so it is difficult to imagine directly detecting it.

matching /43/. Both of these "solutions", besides being somewhat ad hoc, have dynamical difficulties (For $n_f > 2$ the model loses asymptotic freedom. The condensates required to get the breakdown wanted in the SU(6) case are very contrived and probably do not bind). Rather than belabor this point here, I prefer to examine the crucial family problem in the next section along with the problem of how to eventually pass from m = o quarks and leptons to quarks and leptons with small finite mass.

IV. Mass Generation and Family Issues

The dynamical achievement of most composite models in the literature, so far, is to get m = o fermion bound states, which have the quantum numbers of the quarks and the leptons. Even at this stage the models have some difficulties:

i) There are, in general, other states besides the quarks and leptons at zero mass

ii) The family structure in the models is either non existent or quite artificial.

Both of these problems are difficult to bypass. Even succeeding in this, the hardest problem is ahead: how does one go from a theory with a set of massless bound states, to a realistic theory, in which one has a quark and lepton spectrum with $m_{q,\ell} \ll \Lambda_c$?

There are, in my opinion, three principal difficulties to be surmounted to achieve the above mentioned goal:

1) One has to device ways to generate some small mass for the fermions, which were previously massless.

2) One has to generate, furthermore, a hierarchy of intrafamily (e.g. $m_c : m_s : m_\mu \sim 15 : 1.7 : 1$) and interfamily (e.g. $m_\tau : m_\mu : m_e \sim 3400 : 200 : 1$) splittings.

3) Unwanted m = o states, which are not quarks or leptons, have to be removed by some physical mechanism to masses much above the masses of these states.

Let me address each of these points in turn:

Generating any small mass - having imposed a mass protection mechanism - is difficult. This is particularly true if chirality is the mass protection mechanism, since it cannot be easily broken perturbatively. In particular, gauging part of the global chiral group will <u>not</u> generate mass. Gauge interactions are L-L (or R-R) operators, while fermion mass terms connect L with R. Thus unless chirality is already broken at some level, any subsequent gauging will not generate fermion masses *. In this respect there is an enormous difference between fermions and bosons. Bosons, which are at some stage massless because of a mass protection mechanism, can become massive by breaking this mass protection by gauging. A good example is provided by charged pions in massless QCD. These states are massless, since they are the Goldstone bosons of the $SU(2)_L \times SU(2)_R \rightarrow SU(2)_{L+R}$ breakdown of QCD. However, on turning on electromagnetism they pick up a mass, relative to the neutral pions.

$$m_{\pi^+}^2 - m_{\pi^0}^2 \sim \alpha \Lambda_{QCD}^2 \qquad (IV.1)$$

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* The composite Higgs way of Georgi and collaborators /12/ of generating mass for fermions by gauging, uses an already broken chiral group and vacuum misallignment to avoid the above objection.

Along with L. Mizrachi /44/ I have been investigating if it is possible to break chirality a "little" spontaneously in preon models. The idea, very simply put, is to imagine that not all condensates which form in a non Abelian gauge theory are forced to have the size of the dynamical scale in the theory. That is, even though this scale is the only relevant mass scale, there may be very large <u>numerical</u> suppression factors when one deals with condensates of objects which are already singlets under the confining group. For instance, if a lepton is a bound state of preons, one can imagine that lepton-lepton condensates may form, scaling like Λ_c^3:

$$\langle \bar{l} l \rangle = c' \Lambda_c^3 \qquad (IV.2)$$

but one would <u>not</u> expect c' to be of o(1). Most likely c' = o, but it could be that due to the non zero extension of the composite leptons c' $\neq$ o. If the above scenario obtains then one would expect /44/ that

$$m_l \sim c' \Lambda_c \ll \Lambda_c \qquad (IV.3)$$

since the condensate (IV.2) (which we have dubbed an irrelevant condensate) has a numerically very small size.

In the above scenario, as a first approximation, one neglects altogether the irrelevant condensates and one has a set of massless bound states (protection mechanism). Including the effect of the irrelevant condensates gives one a weak breaking of the protective symmetry, thereby generating small fermion masses. Of course, at this stage, everything is purely speculative since we are not able to calculate dynamically any of the irrelevant condensates - even in toy models! Nevertheless, the idea is rather appealing, although it also has some potential dangerous drawbacks. In some sense, the irrelevant condensates way of giving quarks and leptons a mass is very much like what happens in extended technicolor /3/. Hence it may well suffer from some of the drawbacks of ETC. To wit:

i) Unless the model has some very special residual interactions, the existence of irrelevant condensates will give rise to flavor changing neutral currents at an unacceptable level /44/.

ii) The existence of the irrelevant condensates (IV.2) - or analogous ones for quarks - will in general break spontaneously some of the protective symmetries. Thus Goldstone bosons will appear in the theory. Most of these states will gain some mass when SU(3) x SU(2) x U(1) is turned on (in a similar way to what happens with pions in QCD when e.m. is turned on). Typically one expects for these pseudo Goldstone bosons masses of the order of:

$$m^2_{PGB} \sim \alpha_i \Lambda_c^2 \qquad \text{(IV.4)}$$

For Λ_c in the TeV range these states could well be accessible soon. However, some of the Goldstone bosons caused by the irrelevant condensates may well remain massless (or nearly massless). These states are quite similar to the Majorons and Familons /10/ /11/ and they could be troublesome if they coupled strongly enough. Perhaps even more interestingly, in many cases, some of the pseudo Goldstone bosons have both lepton and quark quantum numbers (leptoquarks) and this should be quite distinctive. If one is an optimist, the existence of all this activity in the spin zero sector is bound to provide the experimental tell tale sign for the origin of the quark and lepton masses.

The second difficulty, connected with mass generation, concerns the intra and interfamily splittings which one must produce. This issue is at present too difficult to tackle, since we do not have a clear understanding of why families appear, in the first place. Nevertheless, it might be useful to give at least an example of the type of thinking with which one hopes to achieve some understanding of the quark and lepton mass patterns. Exceptional coset spaces - i.e. coset spaces involving exceptional groups - can provide rather naturally a setting to incorporate families, in the context of quasi Goldstone fermion models /45/. A particular interesting coset space is provided by E_7 / SU(5) x SU(3) x U(1) which has as QGF precisely three repetitions of ($\bar{5}$ + 10) of SU(5) *. These states can only get mass via some explicit breaking of supersymmetry and of the original group. One can argue /46/ that such a breaking is given by operators of the form,

$$\mathcal{L}_{Break.} \sim \frac{f^3}{f_i f_j} \bar{\psi}_i \psi_j \qquad \text{(IV.5)}$$

where the f_i describe the scales at which the E_7 sequentially breaks down to SU(5):

$$E_7 \underset{f_1}{\rightarrow} E_6 \underset{f_2}{\rightarrow} SO(10) \underset{f_3}{\rightarrow} SU(5) \qquad \text{(IV.6)}$$

Since at the first breakdown one gets two $\bar{5}$ and a 10, at the second a $\bar{5}$ and a 10, and at the third the last fermionic 10, Eq. (IV.5) implies the hierarchy

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* It has also an extra 5 state, which is problematic /45/.

$$m_d \sim m_u \sim \frac{1}{f_1^2} < m_s \sim \frac{1}{f_1 f_2} < m_c \sim \frac{1}{f_2^2} <$$
$$< m_b \sim \frac{1}{f_2 f_3} < m_t \sim \frac{1}{f_3^2} \qquad \text{(IV.7)}$$

which is certainly what one observes in nature. However, it is much harder to pronounce oneself on whether this result is enough to justify believing that this kind of model has in it some deep truth!

The last difficulty I mentioned above concerning mass generation, of getting rid of unwanted states, is particularly problematic for quasi Goldstone fermion models. After all, in these models the m = o Goldstone bosons are there because of the $G \rightarrow H$ breakdown, while the QGF are only massless because of supersymmetry. Hence, if one were to break supersymmetry without breaking G, the result would be disastrous: The unwanted Goldstone bosons would remain massless while the QGF (i.e. the quarks and leptons) would get a mass of the order of the supersymmetry breaking scale! It turns out, that even breaking supersymmetry and G at the same time one still runs into trouble since, typically, one obtains comparable masses for the fermions and bosons. Only if fermions have double protection /36/ (chirality plus supersymmetry), like in the Novino model, can one maintain the fermions lighter than the bosons. Constructing a semirealistic model, however, is rather difficult /47/. For instance, one runs into troubles with neutrinos *, obtaining typically neutrino masses of the order of the charged lepton masses.

* Unless ν_R is really missing from the spectrum

V. Concluding Remarks

It is very clear that we are far from understanding how quarks and leptons get their masses. The two most promising avenues discussed in these lectures (Superstrings and Composite Models) approach this problem entirely differently. In superstrings the compactification of the 10 dimensional string theory at scales of $O(M_{Planck})$ provides values for the Yukawa couplings h^f_{ij} which then, knowing the scale of the weak breaking v, fixes the quark and lepton masses. The difficulties here reside both in whether one will ever be able to calculate the h^f_{ij} coefficients and whether one will really be able to generate all the necessary dynamics below M_{Planck} (for example the scale v!) to get the physics out. In contrast, in composite models one tries to compute the quark and lepton masses as solutions of a bound state problem. However, to guarantee that essentially quark and leptons look rather pointlike ($\Lambda_c \gg m_{q,l}$) one focuses on models where dynamically it is possible to get states, with the quantum numbers of quarks and leptons, which are massless. From this point on, however, the dynamical task is very hard and one really does not know yet how to get small masses and an intra and interfamily splitting for the quarks and leptons.

In my opinion, the decision on whether superstrings or composite models are correct will have to come from experiments. Life could be easy and, for example, excited quarks and leptons could be found in the next generation of high energy accelerators (LEP, SLC, HERA, TEVATRON). Most likely, however, we shall only know what is the truth from the properties of the scalar sector. Superstring theories require the existence of scalar partners for all quarks and leptons. However, their Higgs sector is probably comparatively simple , being composed most likely of two Higgs doublets at low energy. In the compositeness case, in contrast, there may well exist many possible pseudo Goldstone bosons in the spectrum, including leptoquarks and various sorts of technipions, some of which may be very light.

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