Dynamical Nonsupersymmetry Breaking

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Abstract

We emphasize the role that anomalous power-law scaling of 4-fermion operators, occurring in the presence of new strong interactions, could have in the generation of quark and lepton masses.

1 Introduction

Recent advances in the understanding of strong dynamics and symmetry breaking in supersymmetric theories, as reported at this conference, have made clear an impressive diversity of phenomena which can arise in strongly interacting theories. This has led to renewed efforts in the search for a plausible dynamical theory of quark and lepton masses, especially in theories incorporating supersymmetry. But the wide range of phenomena uncovered in strongly interacting supersymmetric theories might be considered to suggest a corresponding range of phenomena in strong gauge dynamics in general. In particular one might expect that the diversity of the phenomena would be at least as great when the constraints of supersymmetry are removed. This possibility is somewhat at odds with the typical viewpoint taken in attempts to incorporate dynamical symmetry breaking, without supersymmetry, into realistic models. Much of the previous work over the years have focused on the construction of models with dynamics resembling that of QCD as closely as possible.

In this talk we will illustrate the role that non-QCD-like dynamics, in particular dynamics which gives rise anomalous power-law scaling, could play in realistic models of fermion mass. We describe how a model with a simple symmetry structure can give rise to a wide variety of 4-fermion operators, which in the presence of anomalous

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∗Talk given at the International Workshop on Perspectives of Strong Coupling Gauge Theories (SCGT96), November 1996, Nagoya, Japan.
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scaling, can give rise to a mass spectrum with nontrivial, and potentially realistic, mixings and hierarchies.

There are other reasons why a fresh perspective on dynamical symmetry breaking may be useful. We begin by listing some statements and implicit assumptions sometimes made in connection with dynamical electroweak symmetry breaking, and suggest why these statements may not be completely correct.

- “The top mass is so large that it must be involved with the electroweak symmetry breaking dynamics”.
  — Not necessarily, since 175 is small compared to 1000. If a dynamical fermion mass is responsible for the $W$ and $Z$ masses, then we expect that fermion mass to be closer to 1000 GeV than to 175 GeV.

- “If the $t$ mass has a dynamical origin then $\delta \rho$ is too large.”
  — Not necessarily, if contributions to $\delta \rho$ are suppressed by powers of 175/1000.

- “The dynamical origin of the $t$ mass is distinct from that of other quarks.” For example in “topcolor” models it is postulated that the top mass comes from $\bar{t}tt$ operators rather than $\bar{H}Htt$ operators. Here $H$ is the TeV-mass fermion responsible for electroweak symmetry breaking.
  — Not necessarily, since $\bar{H}Htt$ operators are by themselves still okay; we will see that these operators make contributions to $\delta \rho$ which are suppressed by powers of 175/1000. The real problem is somewhat removed from $t$ mass generation and has to do with is explaining why isospin-violating operators of the form $\bar{H}H\bar{H}H$ are not too large. We will argue that this problem may be resolved in a situation where some operators enjoy more anomalous scaling than other operators.

- “The scale of new flavor physics must be well above a TeV.”
  — Not necessarily, since a TeV scale for new flavor physics coupling mainly to third family fermions is not yet ruled out.

- “A guaranteed signature of dynamical electroweak symmetry breaking is some kind of $\rho$-like resonance at a TeV.”
  — Not necessarily, if the new strong interaction at a TeV is itself broken. The effects of strong dynamics in that case may not look much like QCD dynamics.

- “Dynamical symmetry breaking destroys the concept of coupling constant unification.”
  — We will suggest below why the observed values of the couplings may not be completely accidental, even when new strong interactions are present.
We will be illustrating all the above points in an explicit model. At the very least, the model will hopefully illustrate why the following two statements are false.

- “Elementary scalar fields must be introduced to build realistic theories of flavor.”
- “Fine-tuned 4-fermion operators must be introduced to build realistic theories of flavor.”

These two statements are given together because the fine-tuning sometimes invoked in connection with 4-fermion operators is closely related to the fine-tuning associated with light elementary scalar fields. It is well known that when fine-tuned 4-fermion operators are arranged to generate small dynamical fermion masses, they also generate light composite scalar particles. We clearly want to avoid such fine-tuning.

2 Anomalous Scaling

The key point for the following discussion is the possibility that the anomalous scaling present in strongly interacting theories can turn 4-fermion operators into relevant, or nearly relevant, operators. That is, 4-fermion operators have an effective dimension close to four instead of six. Such operators can play the role of Yukawa couplings. They provide an alternative method for feeding down flavor physics at high scales into the quark and lepton mass spectrum, without the introduction of scalar fields.

Relevant 4-fermion operators were found to arise in quenched QED in ladder approximation. But the phenomenon appears more general, and it is closely related to the tendency for the scaling dimension of the mass operator $\bar{\psi}\psi$ to become two (anomalous dimension $\gamma_m = 1$) when the gauge coupling is above some critical value, $\alpha > \alpha_c$. Arguments supporting this conclusion have been made using the operator product expansion and using general properties of the Schwinger-Dyson equation. If this is accepted then a large anomalous dimension for $\bar{\psi}\psi\bar{\psi}\psi$, perhaps close to two, can also be expected. The essential requirement is a $\beta$-function which is small over some range of momentum scales. Nontrivial infrared fixed-points are clearly of interest in this context.

The important point is that fine-tuning is not required. This is quite different from the fine-tuning of 4-fermion operators in the Nambu-Jona-Lasinio model which produces an effective anomalous dimension twice as large: $\gamma_m = 2$. The latter is related to the existence of light scalar $\bar{\psi}\psi$ bound states. In our case 4-fermion operators are relevant at low energies because the fermions feel a new strong interaction, and not because the coefficients of the 4-fermion interactions are finely tuned. Instead of fine-tuning the physics at one scale to produce desired results at some much larger distance scale, the anomalous scaling in our case is arising due to strong interactions acting over a range of scales.
The operators of interest involve normal quarks and leptons, and thus the new
gauge interactions acting on quarks and leptons at high energy scales must clearly
be broken. These strong interactions must also be prevented from producing large
masses for quarks and leptons. This is reasonable in a chiral gauge theory in which all
masses break gauge symmetries. The 4-fermion operators which are being enhanced,
on the other hand, need not break any gauge symmetries.
The basic framework will share a few features with extended technicolor theories.

• A fermion condensate is responsible for electroweak symmetry breaking.

• Four-fermion operators, in the presence of this condensate, produce all other
quark and lepton masses, including the top.

The minimal version we will describe differs from basic ETC in the following ways.

• There are new strong gauge interactions above a TeV, but they do not survive
below a TeV. There is no unbroken technicolor.

• We may call the new gauge interactions flavor interactions. They break down to
a remnant flavor symmetry at some scale like 100 or 1000 TeV, and the remnant
survives down to a TeV.

• The third family, as well as a new fourth family, couple to the remnant flavor
interactions, while the light two families couple only to the flavor interactions
broken at the higher scale.

• The dynamical fourth family quark masses are responsible for electroweak sym-
metry breaking. There are no technifermions.

• The 4-fermion operators responsible for the quark masses are not of the form
$\overline{H}_L H_R \overline{q}_R q_L$, but are rather of the form $\overline{q}_L H_R \overline{H}_L q_R$ where $H$
denotes a fourth family quark. (These operators cannot arise by explicit gauge boson exchange
and must therefore originate dynamically.)

• The 4-fermion operator responsible for the $t$ mass is composed of third and
fourth family quarks, and is strongly enhanced by the remnant flavor inter-
actions. This particular operator may be close to being a relevant 4-fermion
operator, as discussed above.

• The dynamical fourth family quark masses are not in a singlet channel with
respect to the remnant flavor interactions.
3 A Model of Flavor Physics

We will illustrate these ideas by considering the most minimal flavor gauge symmetry we can imagine:

\[ U(1)_V \times SU(2)_V, \]  

(1)
such that the complete quark content of the theory is given by

\[ Q : (+, 2) \quad \bar{Q} : (-, \bar{2}). \]  

(2)

Each \( Q \) field also transforms under the standard model gauge group as a normal quark, and thus there are enough fields here to describe the quarks of four families.\(^1\) The flavor symmetry breaks at a scale \( \Lambda \approx 100 \) to 1000 TeV to \( U(1)_X \), which is an unbroken combination of \( U(1)_V \) and the diagonal generator of \( SU(2)_V \). Then the \( X \) boson only couples to the third and fourth family quarks \( Q_1 \) and \( \bar{Q}_1 \), and not to the light quarks \( Q_2 \) and \( \bar{Q}_2 \).

Flavor changing neutral currents involving the light two families are suppressed by inverse powers of \( \Lambda \). We shall argue below that most of the mass mixing between families occurs in the up-sector. This combined with the fact that flavor interactions do not induce transitions between the \( Q_2 \) and \( \bar{Q}_2 \) fields will suppress further the most problematical FCNC effect, \( K - \bar{K} \) mixing.

It is important to note that none of the fields \( Q_1, Q_1, Q_2, Q_2 \) are mass eigenstates. In fact the fourth family quark masses will correspond to

\[ \overline{Q}_{1L}Q_{1R} + \text{h.c.} \]  

(3)

As we have said, this is not a singlet under the remnant flavor symmetry \( U(1)_X \), and it can only form once the \( U(1)_X \) breaks (we discuss this breakdown below). Although the \( U(1)_X \) is strong it is unclear whether it, along with possible 4-fermion interactions, is sufficient to produce the mass in (3). We will continue our discussion of the minimal flavor interactions and leave open for now the question of whether additional interactions are required. In the appendix we describe how an enlarged color interaction can help to produce these fourth family quark masses.

We now consider a set of operators which will feed mass down from the fourth family quarks to the second and third family quarks.

\[
\begin{align*}
(\overline{Q}_L D_R) e (\overline{Q}_L U_R) & \quad B \\
(\overline{Q}_L U_R) e (\overline{Q}_L D_R) & \quad \bar{B} \\
(\overline{Q}_L D_R) e (\overline{Q}_L \hat{e} U_R) & \quad C \\
(\overline{Q}_L U_R) e (\overline{Q}_L \hat{e} D_R) & \quad \bar{C}
\end{align*}
\]

\(^1\)To actually make this a chiral gauge group a further \( U(1) \) interaction may be added, or alternatively some relevant 4-fermion operator may play an equivalent role in making the dynamics chiral.\(^5\)
The $\epsilon$ contracts $SU(2)_L$ indices while the $\hat{\epsilon}$ contracts $SU(2)_V$ indices. Note that all these operators are $SU(2)_V$ singlets, i.e. in an attractive channel with respect to $SU(2)_V$ interactions. Thus although these operators are not generated by an explicit $SU(2)_V$ gauge boson exchange, they could be expected to be produced dynamically by strong $SU(2)_V$ interactions. There may well be other operators generated, but these are the operators of most interest for the generation of mass.

These operators have differing flavor structure, as indicated by the sprinkling of the $\hat{\epsilon}$'s, and this results in nontrivial mass matrices, to which we now turn. We first consider $2 \times 2$ blocks of the full $4 \times 4$ mass matrices. The $c-t$ submatrix, in terms of the original fields, is

\[
\left( \begin{array}{cc} U_{2L} U_{2R} & \overline{U}_{2L} U_{1R} \\ U_{1L} U_{2R} & \overline{U}_{1L} U_{1R} \end{array} \right).
\]  

(4)

The operators as labeled above contribute in the following way.

\[
\left( \begin{array}{cc} \mathcal{E} & \mathcal{D} \\ \mathcal{C} & \mathcal{B} \end{array} \right) \Rightarrow c \text{ and } t \text{ masses}
\]  

(5)

Note that these elements are being fed down from the $b'$ mass. Similarly the $t'$ mass is feeding down to the $s$ and $b$ masses as follows.

\[
\left( \begin{array}{cc} \mathcal{E} & \tilde{\mathcal{D}} \\ \tilde{\mathcal{C}} & \tilde{\mathcal{B}} \end{array} \right) \Rightarrow s \text{ and } b \text{ masses}
\]  

(6)

Not only do we get a diverse set of operators with different structures emerging from a fairly simple starting point, but each operator contributes to a different element of a mass matrix. We see only one common element in the up- and down-type matrices, and thus $t-b$ mass splitting is permitted. In addition, the different off-diagonal elements imply nontrivial mass mixing between the 2nd and 3rd families. But the main point about the way masses are being generated here is the way the dynamics generates mass hierarchies. For example the generation of the $\mathcal{B}$ and $\tilde{\mathcal{B}}$ operators are favored because among the operators we have listed, only they are singlets under $U(1)_V$. More importantly, the operators experience different amounts of anomalous scaling, due to the remnant $U(1)_X$ interaction, as they are run down from scale $\Lambda$ where they are created, to a TeV. We expect that the $\mathcal{B}$ and $\tilde{\mathcal{B}}$ operators are enhanced the most while the $\mathcal{E}$ operator is enhanced the least. In particular at one-loop, the $U(1)_X$ corrections enhance the $\mathcal{B}$ and $\tilde{\mathcal{B}}$ operators, they cancel out for the $\mathcal{C}$ and $\mathcal{D}$ operators, and they resist the $\mathcal{E}$ operator. The following hierarchy follows.

\[
\mathcal{B}, \tilde{\mathcal{B}} > \mathcal{C}, \mathcal{D}, \tilde{\mathcal{C}}, \tilde{\mathcal{D}} > \mathcal{E}
\]  

(7)
What we have not explained is the origin of the $t-b$ mass splitting. That must be due to the larger size of the $(B, C, D)$ operators compared to the $(\bar{B}, \bar{C}, \bar{D})$ operators, due to physics at the high scale. This could for example be related to a dynamical breakdown of $SU(2)_R$ in an otherwise left-right symmetric theory. But whatever the origin of the isospin breaking, it is manifested at a TeV in various 4-fermion operators. We find that the most dominant of these operators do not feed directly into $\delta \rho$; in fact four insertions of the dominant $B$ operator is necessary to produce a contribution to $\delta \rho$. This produces a suppression factor of order $(m_t/m_t')^4$. It may be checked that other possible operators contributing to $\delta \rho$ are not of the form which is strongly enhanced by the $U(1)_X$ scaling effect. This applies to the $\Pi H \Pi H$ operators mentioned in the introduction, which in the present case corresponds to operators with four fourth-family quarks. In particular it is not possible to write these operators as a product of Lorentz scalars which are neutral under $U(1)_X$. We thus see how the dynamics protects $\delta \rho$ from the isospin breaking physics responsible for the large $t$ mass.

We return to the question of the breakdown of $U(1)_X$, where we note that this symmetry is broken by the $C$, $D$ and $E$ operators. Since these operators are arising dynamically, the implication is a dynamical contribution to the $X$ boson mass. What contribution does a 4-fermion condensate give to a gauge boson mass? Using naive dimensional analysis which tries to keep track of $4\pi$ factors in loops) it has been argued that it is natural to expect a contribution to the mass over coupling of the gauge boson, $M/g$, an order of magnitude or two less than the natural scale of the operator. In other words, 4-fermion condensates seem to produce a softer breaking of gauge symmetries than do 2-fermion condensates. In this way we find that $M/g$ for the $X$-boson could be as low as a TeV or so.

## 4 The Role of Leptons

We may introduce leptons in the same way as quarks, such that under $U(1)_V \times SU(2)_V$

\[
Q \text{ and } L \text{ transform as } (+, 2),
\]

\[
\bar{Q} \text{ and } \bar{L} \text{ transform as } (-, \bar{2}).
\]

Operators of interest involving leptons are as listed. (We do not consider operators involving right-handed neutrinos, since the latter may not exist in the theory at a TeV.)

\[
(\mathcal{L}_L E_R) \epsilon (\mathcal{Q}_L U_R) \quad \mathcal{F}
\]

\[
(\mathcal{L}_L E_R) \epsilon (\mathcal{Q}_L U_R)
\]

\[
(\mathcal{L}_L E_R) \epsilon (\mathcal{Q}_L U_R)
\]

\[\text{This would also imply that 4-fermion condensates are less resisted by repulsive gauge interactions than 2-fermion condensates.}\]
These operators respect all gauge symmetries, like the $B$ operators. And as we shall show, the $F$ and $G$ operators make additional contributions to the quark masses if the $\tau'$ mass corresponds to $\overline{E}_1 L_1 R$. We therefore assume that that is the case.

It turns out that the operators in (8), along with our previous operators, are sufficient to break all global symmetries. This removes the necessity of some kind of quark-lepton unification at relatively low scales, and thus allows the standard $SU(3) \times SU(2) \times U(1)$ (or left-right symmetric version) to survive to high scales. Since the fermions of our model come in standard model families, the relative running of the $SU(3) \times SU(2) \times U(1)$ couplings remains the same to lowest order in these couplings, to all orders in the new strong interactions. Thus there remains the tendency for the three couplings to become closer at high energies. In other words, if there is some form of coupling constant unification then the observed values of the three couplings are not completely accidental in this picture in spite of new strong interactions.

We may now display the full up- and down-type mass matrices.

$$M_u = \begin{pmatrix}
0 & F_2 & 0 & 0 \\
G_2 & \mathcal{E} & \mathcal{D} & 0 \\
0 & \mathcal{C} & B & F_1 \\
0 & 0 & G_1 & A
\end{pmatrix}$$

(9)

$$M_d = \begin{pmatrix}
\mathcal{H} & 0 & I & 0 \\
0 & \mathcal{E} & \tilde{D} & 0 \\
I & \tilde{C} & \tilde{B} & 0 \\
0 & 0 & 0 & \mathcal{A}
\end{pmatrix}$$

(10)

The $\mathcal{A}$ entry is the dynamical fourth-family quark mass in (3). The other new entries in the up-type matrix are due to the quark-lepton operators, which are feeding mass down from the $\tau'$. Note that the matrix is in general not symmetric. We are postulating that most of the KM mixing angles have their origin in the up-mass matrix. In particular Cabibbo mixing is due mostly to the $F_2$ entry, while the small size of the $u$ mass is due to the smaller $G_2$ entry.

In the down-type matrix, the $\mathcal{H}$ entry is due to an operator similar to the $E$ operator, but which is feeding mass down from the $t$ rather than the $t'$. (The $\mathcal{H}$ entry in the up-type matrix is negligible since it is feeding down from the $b$.) This leads to the mass relation

$$\frac{m_d}{m_t} \approx \frac{m_s}{m_{\tau'}}.$$  

(11)

It is rather novel to find that the $u$ mass is being fed down from the $\tau'$, while the $d$ is being fed down from the $t$. The $I$ operator, and perhaps the $\mathcal{E}$ and $\mathcal{H}$ operators as well, may be arising through loops involving other operators we have listed. It may be shown that mass matrices of the above form have sufficient structure to produce realistic quark masses and mixings.
In the lepton sector we are able to identify 4-fermion operators which could give mass to the charged leptons. The small mass of left-handed neutrinos is probably associated with a large mass for right-handed neutrinos. But we are unable to make predictions for the three light neutrino masses other than that they are unlikely to be much smaller than an eV.

5 Signatures

We return to the point that the $E$ entry is the same for the up- and down-type matrices, which is due to the fact that the $E$ operator is intrinsically isospin conserving. To avoid unnatural cancellation in the $b$–$s$ mass matrix, this entry should be of order the $s$ mass. This means that the $c$ mass must be due mostly to $c$–$t$ mixing induced by the $C$ and $D$ entries. These entries are thus of order $\sqrt{m_t m_c}$, and this in turn has an interesting phenomenological consequence.

If we take one of our operators and close off the $t'$ or $b'$ lines we get a contribution to the quark mass matrix. If we also attach a photon or $Z$ to the heavy quark loop then we generate an anomalous magnetic-moment-type coupling; in particular we can get the flavor-changing coupling

$$\overline{t} \sigma_{\mu \nu} q^\nu (1, \gamma_5) c (A^\mu, Z^\nu) + h.c.$$  

Since the size of this operator is closely related to the corresponding off-diagonal mass elements, we estimate that its coefficient is of order

$$\frac{em_c m_t}{\lambda^2}$$  

where $\lambda \approx$TeV. This is larger than the conventional ETC estimate, since ETC mass operators do not directly generate these couplings. It is also larger than estimates based on multi-Higgs models, since in that case the new coupling requires an extra loop compared to the tree-level mass.

This coupling yields

$$R^{ct} \equiv \frac{\sigma(e^+ e^- \rightarrow t\bar{\tau} + t\bar{c})}{\sigma(e^+ e^- \rightarrow \gamma \rightarrow \mu^+ \mu^-)}$$

$$\propto \frac{(s + 2m_t^2)(s - m_t^2)^2}{3s^2 m_t^2} (1 - \cos(\theta)^2)$$

with an magnitude which is just barely detectable at LEP2, and easily detectable at 500 GeV. Chromomagnetic moments may also be considered, and in that case the flavor-diagonal anomalous coupling of a gluon to top quarks may be of interest to the production of top quarks in hadron colliders.

Generic to this picture are the remnant flavor interactions, which implies at least one new massive gauge boson coupling to the third family with a mass as low as a
In the minimal model this is the $X$-boson, and we find that it mixes with the $Z$. As we have described it, the $X$-boson has axial couplings to the $t$ and $b$ quarks and vector couplings to the $\tau$. The mixing then results in the following wide range of fractional shifts in various electroweak parameters, which occur in the ratios,

$$A_\tau : R_b : \Gamma_{\nu_\tau} : \mathcal{A}_b : \Gamma_\tau$$

$$+20 : +2.0 : -1.5 : -0.5 : +0.2$$

The magnitude of the shift of the $Z$ coupling to a third family fermion is

$$\delta g_Z = \frac{e}{scf_t} \frac{g_X^2}{M_X^2} \cdot (16)$$

$f_t \approx 60$ GeV is determined from the $t$ loop with an effective TeV cutoff supplied by the momentum dependence of the $t$ mass. For example, an observed $2\%$ shift in $R_b$ or a $20\%$ shift in $A_\tau$ would imply that

$$\frac{M_X}{g_X} \approx 1 \text{ TeV.} \quad (17)$$

What is perhaps most surprising is that such large anomalies have not yet been completely ruled out.

We note in passing that the new flavor physics is not associated (at least in the minimal $U(1)_X$ picture presented here) with any kind of new confining interaction. This new flavor physics could be quite different from QCD-like dynamics and the implied $\rho$-like resonance typically expected in theories of dynamical symmetry breaking.

And finally, this picture leads a fourth left-handed neutrino, with a dynamical mass expected to be comparable to the other fourth-family members. Assuming that the right-handed neutrino is either absent or much heavier, we are left to consider a large Majorana mass for the left-handed neutrino. This might seem to be a disaster for electroweak corrections, and for this reason the effects of Majorana masses are normally only considered for right-handed neutrinos. We note in particular the current anomaly in $A_\tau/A_{e,\mu}$ extracted from the forward-backward lepton asymmetries at LEP.  

\[ ^3 \text{The latter is a byproduct of the choice } E_{\ell L}/E_{\ell R} \text{ for the } \tau' \text{ mass} \]

\[ ^4 \text{We note in particular the current anomaly in } A_\tau/A_{e,\mu} \text{ extracted from the forward-backward lepton asymmetries at LEP.} \]
may be useful, given the fact that other new contributions to $T$ in this context are typically positive.

In conclusion, we have described a picture in which flavor physics originates at a scale of order 100 or 1000 TeV. We have emphasized the role that anomalous scaling of 4-fermion operators can have in transforming the flavor physics at this scale into the complicated pattern of observed fermion masses.

**Appendix**

We present here an extension of the minimal model which may have more appealing dynamics. In particular we replace the $U(1)_X$ with $SU(3)_X$. Each third and fourth family fermion is now one member of an $SU(3)_X$ triplet. We also expand color into a direct product of $SU(3)$’s, such that under the gauge symmetry

$$SU(3)_X \times SU(3)_{C1} \times SU(3)_{C2},$$

the third and fourth family quarks are members of the following representations.

$$Q_L \ (3,3,1) \quad Q_L \ (\bar{3},1,3)$$

$$Q_R \ (3,1,3) \quad Q_R \ (\bar{3},3,1)$$

The symmetry in (18) is not a symmetry of the C and D operators, and as before we assume that the gauge bosons corresponding to the broken generators have masses in the TeV range. The symmetry left unbroken by the C and D operators is $SU(2)_M \times SU(3)_{C}$, where the former is a “metacolor” and the latter is standard color. We may assume that $SU(3)_{C1}$ is strong enough to help induce the $SU(3)_{C1}$ conserving and $SU(3)_{X}$ violating condensate $\langle \bar{Q}_{1L}Q_{1R} \rangle$. The condensate lies in the 6 of $SU(3)_{X}$, and it can be transformed to the given ‘11’ form. The result is the fourth family quark masses in (20), and the main point is that there is now an obviously attractive gauge interaction in this channel.

Since $SU(3)_{2C}$ is likely weaker than $SU(3)_{1C}$ in order to get the correct $\alpha_C$, the ‘broken-color’ interactions will end up coupling more strongly to the second and fourth families than to the first and third. As for metacolor, we assume it grows strong somewhat below a TeV and that it confines the two families of metafermions. If metafermions become massive then they make an unwelcome contribution to the $S$ parameter. On the other hand their chiral symmetries are already explicitly broken by the broken gauge interactions, and thus the confined metafermions could remain massless due to unbroken discrete chiral symmetries without constraints from chiral anomalies.

This whole structure can easily be incorporated into the larger gauge group present at 100 or 1000 TeV. The $SU(3)_X$ becomes embedded into $SU(4)_Y$, which acts also on the two light families. All the 4-fermion operators we have considered may be
chosen to transform as singlets under an $SU(2)$ subgroup of $SU(4)_V$, disjoint from the unbroken $SU(2)_M$ subgroup. The operators are singlets under this $SU(2)$ subgroup in the same way they were singlets under the $SU(2)_V$ of the minimal model. It may be checked that they are therefore in attractive channels with respect to $SU(4)_V$ (even though they do not respect $SU(4)_V$), just as they were in attractive channels with respect to $SU(2)_V$. Finally we note that the fermion content of the model implies that $SU(4)_V$ may be a candidate for a nontrivial infrared fixed point according to the analysis in Ref. [13].

Acknowledgments

I thank the organizers of the Strong Coupling Gauge Theories 96 meeting in Nagoya for a stimulating conference and for their support. This research was supported in part by the Natural Sciences and Engineering Research Council of Canada.

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Figure 1: Lines of constant $T$ as a function of the $N$ and $E$ masses in TeV. Thick and thin lines are for $\Lambda = 1.5m_N$ and $\Lambda = 2m_N$ respectively. In each case, from top to bottom, $T = -2, -1, 0$.

Figure 2: Thick and thin lines are lines of constant $S$ and $U$ respectively as a function of the $N$ and $E$ masses in TeV. From top to bottom in each case $S = 1/6\pi, 0, -1/6\pi$ and $U = -1/12\pi, 0, 1/6\pi$. 

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