Fermion Masses and *CP* Violation in a Model with Scale Unification

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It is observed that the scale-unifying model based on supersymmetry and compositeness provides a natural reason for the family mass hierarchy $\bar{m}_e \ll \bar{m}_\mu \ll \bar{m}_\tau$ and links spontaneous *CP* violation to the nonvanishing mass of the electron family. Some of its predictions include (i) K^0-$\bar{K}^0$ and $K_L \to \bar{\mu} e$ are normal, but (ii) $Z \to t\bar{c}$, $c\bar{u}$, and $\mu\bar{e}$ have observably large strengths *allowing for single top production at the CERN e^+e^- collider LEP II*, and (iii) $d_n \approx (0.5-10)\times 10^{-26}\, e\,$cm.

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It has recently been shown [1,2] that the idea that quarks, leptons, and Higgs bosons are composites and that their constituents possess local supersymmetry can be realized in the context of a *viable* and most economical preon model which has many attractive features. These include (i) a common origin of all the diverse scales from $M_{\rm Pl}$ to m_ν [1]; (ii) a simple and compelling reason, based on supersymmetry, for replication of chiral families [2]; and (iii) an explanation based on the index theorem for the protection of quark-lepton masses [3].

The purpose of this Letter is to probe into the origins of interfamily mass hierarchy, family mixing, and *CP* violation, within this scale-unifying model. In the process, we observe that the model not only provides a natural reason for the progressive hierarchy $\bar{m}_e \ll \bar{m}_\mu \ll \bar{m}_\tau$, but links *CP* violation that arises spontaneously within the model to the small but nonvanishing mass of the electron family, and leads to a host of testable predictions.

To discuss the fermion masses, we first need to recall a few salient features of the model [1,2]. The model assumes $N=1$ local supersymmetry (SUSY) at the Planck scale. It introduces six positive and six negative massless chiral preonic superfields $\Phi^{a,\sigma}_{\pm} = (\varphi,\psi,F)^{a,\sigma}_{L,R}$, each belonging to the representation **N** of a metacolor gauge symmetry SU(N). Here σ denotes the metacolor index running from 1 to N; a denotes flavor-color quantum numbers having six values (x,y,r,y,b,l), where (x,y) provide up and down flavors and (r,y,b,l) the four colors including lepton color [4]. The symmetry $SU(N)\times SU(2)_L \times SU(2)_R \times SU(4)^C_{L+R}$ is gauged. Corresponding to an input value for the metacolor coupling $\bar{\alpha}_M \approx 0.07$ to 0.05 at $M_{\rm Pl}/10$, the asymptotically free metacolor force becomes strong and confining at a scale $\Lambda_M \simeq 10^{11}$ GeV, for $N=5$–6. At that point, it serves many purposes.

(i) It makes three light chiral families of composite quarks and leptons $(q^i_{L,R})_{i=1,2,3}$ *and two relatively heavy* (mass ~ 200 GeV–2 TeV) *vectorlike families $Q_{L,R}$ and $Q'_{L,R}$ that couple vectorially to W_L's* and *W_R's, respectively* [2,5]. There are thus altogether five $SU(2)_L$-doublet and five $SU(2)_R$-doublet families, each having the transformation properties under $SU(2)_L\times SU(2)_R \times SU(4)^C_{L+R}$ as noted below:

$$(q_L^{1,2,3},Q_L,Q_R) \sim (2_L,1,4^{*C}),\quad (q_R^{1,2,3},Q'_R,Q'_L) \sim (1,2_R,4^{*C})\,. \tag{1}$$

The members of these families are denoted by $q^e_{L,R} = (u,d,\nu_e,e)_{L,R}$, $q^\mu_{L,R} = (c,s,\nu_\mu,\mu)_{L,R}$, $q^\tau_{L,R} = (t,b,\nu_\tau,\tau)_{L,R}$, $Q_{L,R} = (U,D,N,E)_{L,R}$, and $Q'_{L,R} = (U',D',N',E')_{L,R}$.

(ii) It is assumed that the metacolor force makes a *SUSY-preserving* condensate Δ_R of the scale of Λ_M which transforms as $(1,3_R,10^{*C})$ under $SU(2)_L \times SU(2)_R \times SU(4)^C_{L+R}$. This gives superheavy Majorana masses of order $\Lambda_M \sim 10^{11}$ GeV to the three right-handed neutrinos ν^i_R's belonging to the chiral families q^i_R's and breaks $SU(2)_L\times SU(2)_R\times SU(4)^C_{L+R}$ to $SU(2)_L\times U(1)_Y \times SU(3)^C_{L+R}$ [6].

(iii) It is furthermore assumed that the metacolor force makes a few SUSY-breaking condensates as well. These include the metagaugino condensate $\langle \boldsymbol{\lambda}\cdot\boldsymbol{\lambda}\rangle$ and the matter fermion condensates $\langle \bar{\psi}^a\psi^a\rangle$, each of which breaks SUSY [3]. Noting that, within the class of models under consideration, the index theorem prohibits a dynamical breaking of supersymmetry in the absence of gravity [3,7], however, the formation of these condensates must need the collaboration between the metacolor force and gravity. As a result, each of these condensates is expected to be damped by one power of $\Lambda_M/M_{\rm Pl}$ relative to Λ_M [3,8]:

$$\langle \boldsymbol{\lambda}\cdot\boldsymbol{\lambda}\rangle = a_\lambda \Lambda_M^3(\Lambda_M/M_{\rm Pl})\,;\quad \langle \bar{\psi}^a\psi^a\rangle = a_{\psi_a}\Lambda_M^3(\Lambda_M/M_{\rm Pl})\,. \tag{2}$$

There are four $\langle\bar{\psi}\psi\rangle$ condensates corresponding to a having the values x, y, (r,y,b), or l [8]. The coefficients a_λ and a_{ψ_a}, *a priori*, are expected to be of order unity within a factor of 10 (say), although a_λ is expected to be larger than the a_ψ's, typically by a factor of 3–10, because the ψ's are in the fundamental and the λ's are in the adjoint representation of the metacolor group [1].

The condensates $\langle \boldsymbol{\lambda}\cdot\boldsymbol{\lambda}\rangle$ and $\langle \bar{\psi}^a\psi^a\rangle$ induce SUSY-breaking mass splittings $\delta m_S \sim a_\lambda\Lambda_M(\Lambda_M/M_{\rm Pl}) \sim 1$ TeV. The condensates $\langle \bar{\psi}^a\psi^a\rangle$, for $a=x$ and y, break not only SUSY but also the electroweak symmetry $SU(2)_L \times U(1)_Y$. The resulting masses of W and Z bosons are $m_W, m_Z \sim g_2 a_\psi \Lambda_M(\Lambda_M/M_{\rm Pl}) \sim 100$ GeV, where g_2 is the $SU(2)_L$ gauge coupling constant.

Masses of the vectorlike families $Q_{L,R}$ and $Q'_{L,R}$ are protected by $U(1)_X$. They acquire flavor-color-independent masses of order $a_\lambda\Lambda_M(\Lambda_M/M_{\rm Pl}) \sim 1$ TeV only through the condensate $\langle \boldsymbol{\lambda}\cdot\boldsymbol{\lambda}\rangle$, which breaks $U(1)_X$ just as needed [1,3]. But the chiral families $q^i_{L,R}$ acquire masses primarily through their mixings with the vector-like families $Q_{L,R}$ and $Q'_{L,R}$ which are induced by $\langle\bar{\psi}^a\psi^a\rangle$. This is because the direct mass terms $m^{(0)}_{\rm dir}(q^i_L \to q^i_R)$ cannot be induced through either $\langle\boldsymbol{\lambda}\cdot\boldsymbol{\lambda}\rangle$ or $\langle\bar{\psi}\psi\rangle$. These receive small contributions $\lesssim 1$ MeV at Λ_M from effective four-body condensates like $\langle\bar{\psi}\psi\varphi^*\varphi\rangle$, which are, however, damped by $(\Lambda_M/M_{\rm Pl})^2$. Thus, ignoring QCD corrections and $m^{(0)}_{\rm dir}$ for now, the Dirac mass matrices of all four sectors—i.e., q_u, q_d, l, and ν—have the form

$$M^{(0)}_{f,c} = \begin{array}{c} \\ \bar{q}^i_R \\ \bar{Q}_R \\ \bar{Q}'_R \end{array}\begin{array}{c} \begin{array}{ccc} q^i_L & Q_L & Q'_L \end{array} \\ \left[\begin{array}{ccc} 0 & X\kappa_f & Y\kappa_c \\ Y'^{\dagger}\kappa_c & \kappa_\lambda & 0 \\ X'^{\dagger}\kappa_f & 0 & \kappa_\lambda \end{array}\right]\end{array} . \tag{3}$$

Here, $f=x$ or y and $c=(r,y,b)$ or l. The index i runs over three families. The entities X, Y, X', and Y' are column matrices in the family space having entries of ~ 1 to $\frac{1}{10}$. In the above, $\kappa_f \equiv O(a_{\psi_f})\Lambda_M(\Lambda_M/M_{\rm Pl})$, $\kappa_c \equiv O(a_{\psi_c})\Lambda_M(\Lambda_M/M_{\rm Pl})$, and $\kappa_\lambda \equiv O(a_\lambda)\Lambda_M(\Lambda_M/M_{\rm Pl})$. Following the remarks made above, we expect $\kappa_\lambda \simeq (3\text{–}10)\kappa_{f,c}$. *Thus the Dirac mass matrices of all four sectors have at least an approximate seesaw structure.*

In the absence of electroweak corrections [$\sim$(5–10)%], left-right symmetry and flavor-color independence of the metacolor force guarantee (a) $X=X'$ and $Y=Y'$, and (b) the same X, Y, and κ_λ apply to q_u, q_d, l, and ν [see Eq. (3)]. This results in an enormous reduction of parameters.

We first observe that by ignoring electroweak corrections one can always rotate the chiral fermions q^i_R and q^i_L to bring the row matrices $Y^T = Y'^T$ to the simple form (0,0,1) and simultaneously $X^T = X'^T$ to the form $(0,p,1)$, with redefined κ_f and κ_c. As a result, *the* 5×5 *mass matrices of the four sectors—i.e.,* q_u, q_d, l, *and* ν*—which in general could involve a hundred parameters, are essentially determined (barring electroweak corrections and contributions from* $m^{(0)}_{\rm dir}$*) by just six effective parameters*—i.e., p, κ_u, κ_d, κ_r, κ_l, *and* κ_λ. Furthermore, we know their approximate values (within a factor of 10, say). Examining the relevant preon diagrams, one can argue that p is less than but not very much smaller than unity; $p\approx\frac{1}{2}$ to $\frac{1}{4}$ is quite natural [9].

Since we expect $\kappa_f, \kappa_c \le \kappa_\lambda/3$ (see above), we obtain the following eigenvalues in the leading seesaw limit (neglecting electroweak corrections and $m^{(0)}_{\rm dir}$):

$$\begin{aligned} &m_u^{(0)} = m_d^{(0)} = m_e^{(0)} = (\tilde{m}_{\nu_e}^{(0)}) = 0\,, \\ &(m_c^{(0)}, m_s^{(0)}) \approx (\kappa_u,\kappa_d)(\kappa_r/\kappa_\lambda)(p^2/2)\eta_{\rm QCD}\,, \\ &(\tilde{m}_{\nu_\mu}^{(0)}, m_\mu^{(0)}) \simeq (\kappa_u,\kappa_d)(\kappa_l/\kappa_\lambda)(p^2/2)\,, \\ &(m_t^{(0)}, m_b^{(0)}) \simeq (\kappa_u,\kappa_d)(\kappa_r/\kappa_\lambda)(2)\eta_{\rm QCD}\,, \\ &(\tilde{m}_{\nu_\tau}^{(0)}, m_\tau^{(0)}) \simeq (\kappa_u,\kappa_d)(\kappa_l/\kappa_\lambda)(2)\,, \\ &m(U,D,U',D') \simeq \kappa_\lambda\eta_{\rm QCD}\,, \\ &m(E,E') \approx \tilde{m}(N,N') \approx \kappa_\lambda\,. \end{aligned} \tag{4}$$

The tildes on neutrino masses denote that they are Dirac masses. Combined with the superheavy Majorana masses of ν_R's, they yield light ν_L's [10]. The QCD renormalization factors for quarks are momentum dependent. With five families and their superpartners (masses ~ 1 TeV), we obtain $\eta_{\rm QCD}(\mu)\approx 2.9$, 3.3, 4.1, and 5.2 for $\mu = 1$ TeV, 100 GeV, 5 GeV, and 1 GeV, respectively.

We see that despite the fact that the electron family is made of the same stuff as the μ *and the* τ *families, it is guaranteed to remain massless [barring contributions from* $m^{(0)}_{\rm dir} \sim (1\ {\rm MeV})\eta_{\rm QCD}$*]—a fact which is not far from the truth.* The reason is simply the rank of the matrix $M^{(0)}$. We also see that the μ-τ mass ratios (evaluating $\eta_{\rm QCD}$ at a fixed momentum for all quarks) are

$$\frac{m_c^{(0)}}{m_t^{(0)}} \approx \frac{m_s^{(0)}}{m_b^{(0)}} \approx \frac{m_\mu^{(0)}}{m_\tau^{(0)}} \approx \frac{p^2}{4}\,. \tag{5}$$

Thus, for $p\approx\frac{1}{3}$ to $\frac{1}{4}$, which is natural (see remarks above), *we obtain a rather large* μ-τ *hierarchy of about* $\frac{1}{40}$ to $\frac{1}{64}$. In this way, the model provides a natural reason for the *interfamily hierarchy* $\bar{m}_e \ll \bar{m}_\mu \ll \bar{m}_\tau$.

To accommodate the observed features of quark-lepton mass splittings within a family (e.g., m_b/m_τ) including the $\eta_{\rm QCD}$ factor and the up-down ratios [11] (e.g., m_t/m_b), one needs to assume $\kappa_r/\kappa_l \approx 0.6\pm0.2$ and $\kappa_d/\kappa_u \approx 1/(30\pm5)$. The first ratio is in a natural range but the second is outside. It is conceivable that κ_d is so small because it is generated only radiatively through κ_u [12]. To see the kind of masses which could be obtained at the tree level, consider the following choice of parameters which turns out to be near optimum: $p\approx0.31$, $\kappa_u\approx 80$ GeV, $\kappa_l/\kappa_\lambda\approx\frac{1}{3}$, $\kappa_r/\kappa_l\approx0.6$, $\kappa_d/\kappa_u\approx1/30$, and $\kappa_\lambda\approx(3\text{–}5)\kappa_u\approx 200\text{–}400$ GeV. These yield (including QCD corrections) $m_u^{(0)}=m_d^{(0)}=m_e^{(0)}=0$, $m_t^{(0)}\approx110$ GeV, $m_b^{(0)}\approx4.7$ GeV, $m_c^{(0)}\approx3.9$ GeV, $m_s^{(0)}\approx130$ MeV, $m_\tau^{(0)}\approx1.7$ GeV, and $m_\mu^{(0)}\approx40$ MeV, while $m(U,D,U',D')\approx1.5\text{–}3$ TeV and $m(E,E')\approx\tilde{m}(N,N')\approx200\text{–}400$ GeV.

While these results possess at least the desired gross pattern—i.e., $\bar{m}_e \ll \bar{m}_\mu \ll \bar{m}_\tau$, with $\bar{m}_e\approx0$—they are off in details by a factor of 2–3. In particular, m_c is too high and m_μ too low; all the other masses are reasonable. The tree-level mass matrix $M^{(0)}$ has an additional shortcoming: The Cabibbo-Kobayashi-Maskawa (CKM) matrix in the 3×3 light-family sector is found to be essentially unity, rendering $\theta_{e\mu}\approx\theta_{\mu\tau}\approx\theta_{e\tau}\approx0$. The results improve dramatically, however, with regard to both the masses and the CKM matrix by including the electroweak corrections at Λ_M and $|m^{(0)}_{\rm dir}|\sim1$ MeV.

The $SU(2)_L\times U(1)_Y$ interactions distinguish between left and right, up and down, and quarks and leptons. These corrections, evaluated at Λ_M, through preon diagrams [9], alter X^T, Y^T, $(X')^T$, and $(Y')^T$ to the general forms

$$X^T=(0,p+\delta_2,1),\quad Y^T=(0,0,1+\delta_3),\quad (X')^T=(\tilde\delta_1,p+\tilde\delta_2,1+\tilde\delta_3),\quad (Y')^T=(0,0,1)\,. \tag{6}$$

The parameters δ_i and $\tilde\delta_i$ are in principle calculable. There are eight δ's—i.e., $\delta_{2,3}^{u,d,l,\nu}$—and twelve $\tilde\delta$'s—i.e., $\tilde\delta_{1,2,3}^{u,d,l,\nu}$. Each of these δ's is a sum of several δ's (evaluated in the preon basis), and is expected to be nearly a few to 10%. (Note that $(\alpha_2/2\pi)\ln[\Lambda_M/(100\text{ GeV})]\sim\frac{1}{200}\,20\sim10\%$.) Including the δ's, the mass eigenvalues are altered as follows ($\eta_{\rm QCD}$ is suppressed):

$$\hat m_e^j=0,\quad \hat m_\mu^j\approx\kappa_f\left(\frac{\kappa_c}{\kappa_\lambda}\right)\left(\frac{p^2}{2}\right)\left(1+\frac{\delta_2^j+\tilde\delta_2^j}{p}\right),$$
$$\hat m_\tau^j\approx\kappa_f\left(\frac{\kappa_c}{\kappa_\lambda}\right)2\left(1+\frac{p^2}{4}+\frac{\delta_3^j+\tilde\delta_3^j}{2}+\frac{p(\delta_2^j+\tilde\delta_2^j)}{4}-\frac{\kappa_f^2+\kappa_c^2}{\kappa_\lambda^2}\right). \tag{7}$$

Here, $j=u,d,\nu,l$. The carets indicate that $m_{\rm dir}^{(0)}\lesssim1$ MeV is not included. We see that while the corrections to the τ-family masses are only of order (5–10)%, those to the muon family can be substantial because they are proportional to $1+(\delta_2^j+\tilde\delta_2^j)/p\approx1\pm(5\text{–}20)\%/0.3\approx1\pm(16\text{–}66)\%$. For example, if $\delta_2^u+\tilde\delta_2^u\approx-20\%$ and $\delta_2^l+\tilde\delta_2^l\approx+22\%$ (say), which are within a reasonable range, m_c could be reduced by a factor of 3 and m_μ enhanced by about 1.7, compared to the tree-level solutions, just as desired. Assuming, conservatively, that each individual $|\delta_i^j|\lesssim(5\text{–}12)\%$, and using observed values of $m_\tau/m_\mu\simeq17$, $m_c(1\text{ GeV})\simeq1.4\pm0.1$ GeV, and $m_t(\text{phys})\geq89$ GeV, we find from Eq. (7) that $0.29\lesssim p\lesssim0.33$ and $m_t(\text{phys})\lesssim180$ GeV. The model, however, typically prefers much lower values of $m_t\lesssim130$ GeV.

The CKM elements, ignoring $m_{\rm dir}^{(0)}$ still, are given by

$$\hat V_{us}\approx[(\tilde\delta_1^d-\tilde\delta_1^u)/p][1+O(\delta/p)]\,,$$
$$\hat V_{ub}\approx[(\tilde\delta_1^d-\tilde\delta_1^u)/2][1+O(\delta/p)]\,, \tag{8}$$
$$\hat V_{cb}\approx[(\tilde\delta_2^d-\tilde\delta_2^u)/2][1+O(p)]+(p/2)(\kappa_u^2-\kappa_d^2)/\kappa_\lambda^2\,.$$

These can yield a reasonable set of mixing angles for the δ's (a few to 10%). Note especially that $\hat V_{us}$, enhanced by $1/p$, is expected to be larger than both $\hat V_{ub}$ and $\hat V_{cb}$.

Let us now include the contributions from $m^{(0)}(q_L^i\to q_R^i)$, which are induced only by effective four-body condensates $\langle\bar\psi_L\psi_R\varphi_L^*\varphi_R\rangle$ and are thus of order (1 MeV) $\times\eta_{\rm QCD}$. These lead to $m_e,m_u,m_d\neq0$. Most importantly, they also permit spontaneous CP violation through the fermion mass matrix which would vanish as $m_{\rm dir}^{(0)}\to0$.

To see this, first set $m_{\rm dir}^{(0)}=0$ and introduce phases into the condensates $\langle\boldsymbol\lambda\cdot\boldsymbol\lambda\rangle$ and $\langle\bar\psi\psi\rangle$ or equivalently into the κ's—i.e., $\kappa_f=|\kappa_f|e^{i\xi_f}$, $\kappa_c=|\kappa_c|e^{i\xi_c}$, and $\kappa_\lambda=|\kappa_\lambda|e^{i\xi_\lambda}$. Simultaneously, impose the following transformations: $q_L^{f,c}\to q_L^{f,c}$, $q_R^{f,c}\to e^{i(\xi_f+\xi_c-\xi_\lambda)}q_R^{f,c}$, $Q_L^{f,c}\to e^{i(\xi_c-\xi_\lambda)}Q_L^{f,c}$, $Q_R^{f,c}\to e^{i\xi_c}Q_R^{f,c}$, $Q_L^{\prime f,c}\to e^{i(\xi_f-\xi_\lambda)}Q_L^{\prime f,c}$, and $Q_R^{\prime f,c}\to e^{i\xi_f}Q_R^{\prime f,c}$. These do not introduce any phase into V_{KM}^L because the left chiral fields are unchanged while Q_L and Q_R transform the same way for up and down. It is easy to verify that the mass matrix $M_{f,c}$ including electroweak corrections, subject to the transformations mentioned above, is rendered real if $m_{\rm dir}^{(0)}=0$. This says that neither the mass matrix nor the gauge interactions (ignoring $W_R^\pm$, which are superheavy) can generate observable CP violation if $m_{\rm dir}^{(0)}=0$. However, with $m_{\rm dir}^{(0)}$ being nonvanishing and complex, the reality of the mass matrix is in general lost and, thereby, CP conservation as well. *We thus see an interesting connection between the nonvanishing masses of the electron family and the spontaneously generated CP violation in the model.*

To explore the consequences of $m_{\rm dir}^{(0)}$, we write the mass matrix for the 3×3 light d-quark sector in the form $M^{(d)}\equiv\hat M^{(d)}+m_{\rm dir}^{(0)d}$ and choose the basis such that $\hat M^{(d)}$ (which includes electroweak corrections) is diagonal: $\hat M^{(d)}=(0,\hat m_s,\hat m_b)$. In the same basis, we denote $(m_{\rm dir}^{(0)d})_{ij}\equiv\Delta_{ij}^{(d)}$, where the Δ_{ij}'s are complex. For quarks, we expect $|\Delta_{ij}^{(q)}|\sim(1\text{ MeV})\eta_{\rm QCD}(1\text{ GeV})\approx$ a few to 15 MeV. The CKM elements for $W_L^\pm$ are now altered to

$$V_{ud}=V_{cs}^*\approx1-\frac{\tilde\delta_1^{u\,2}+\tilde\delta_1^{d\,2}}{2p^2}-\frac{\tilde\delta_1^d-\tilde\delta_1^u}{p}\left(\frac{\Delta_{12}^{*d}}{m_s}-\frac{\Delta_{12}^u}{m_c}\right),$$
$$V_{us}\approx\hat V_{us}+\left(\frac{\Delta_{12}^d}{m_s}-\frac{\Delta_{12}^u}{m_c}\right)-\frac{\Delta_{23}^{d*}}{m_b}\frac{\tilde\delta_1^d-\tilde\delta_1^u}{2},$$
$$V_{cd}\approx-\hat V_{us}+\left(\frac{\Delta_{12}^{*u}}{m_c}-\frac{\Delta_{12}^{*d}}{m_s}\right)-\frac{\Delta_{13}^{*d}}{m_b}\frac{\tilde\delta_2^d-\tilde\delta_2^u}{2},$$
$$V_{cb}\approx\hat V_{cb}+\Delta_{23}^d/m_b,\quad V_{ub}\approx\hat V_{ub}+\Delta_{13}^d/m_b\,.$$

The phase-invariant parameter $J\equiv{\rm Im}(V_{ud}V_{cs}V_{us}^*V_{cd}^*)$, relevant for CP violation in $K\to2\pi$ decay, is given by

$$J\approx\tfrac12(\tilde\delta_2^d-\tilde\delta_2^u){\rm Im}\{[(\tilde\delta_1^d-\tilde\delta_1^u)/p+\Delta_{12}^{d*}/m_s]\Delta_{13}^d/m_b\}\,. \tag{9}$$

This leads to $J\approx[0.05\text{–}0.07](1\text{–}2)(0.1\text{–}0.15)(2\times10^{-3})\xi\approx(1\text{–}4)\times10^{-5}\xi$, where ξ is the phase of Δ_{13}/m_b. This gives $|\epsilon|\sim\frac{1}{600}$ with $\xi\sim1$ [13]. Thus the suppression of ϵ is naturally explained because, essentially, $|\epsilon|\sim|\Delta_{13}^d/m_b|\sim m_d/m_b\approx2\times10^{-3}$, with a maximal ξ. As regards ϵ', it is found to receive contributions primarily from the penguin graph as in the CKM model.

Turning attention to the electric dipole moment of the neutron, d_n, it is a special property of this model that although $W_R^{\pm}$ are superheavy, right chiral currents couple to W_L's because q_R's mix with Q_R's [see Eq. (3)] belonging to the vectorlike family Q which couple to W_L's. The dominant contribution comes from $d_L \to d_R + \gamma$ with charm quark and W_L^- in the loop. This involves the vertex $d_R \to c_R + W_L^-$, for which the CKM element is given by $(\kappa_u \kappa_d/\kappa_\lambda^2) p^2 \Delta_{12}^{*d}/m_s$. Thus, we obtain

$$d_n \simeq \left[\left(\frac{e a_2}{4\pi} \right) \left(\frac{m_c}{m_W^2} \ln \frac{m_c^2}{m_W^2} \right) \sin\theta_C \right] \left(\frac{\kappa_u \kappa_d}{\kappa_\lambda^2} \right) p^2 \left| \frac{\Delta_{12}^{d*}}{m_s} \right| \sin\eta \,, \tag{10}$$

where η is the phase of Δ_{12}^d. Allowing for $\kappa_d/\kappa_u \approx 1/30$, $\kappa_u/\kappa_\lambda \approx \frac{1}{3} - \frac{1}{5}$, $|\Delta_{12}^{d*}/m_s| \approx (\frac{1}{2} - 1.5) \times 10^{-1}$ and $\eta \approx 1 - \frac{1}{10}$, we expect $d_n \approx 10^{-25}$ to $\frac{1}{2} \times 10^{-26}$ e cm [14]. This is a relatively large d_n which should be observable.

Finally, as regards flavor-changing processes, arising from the mixing of q^i's with Q and Q', we find [9] that the new contributions to processes such as K^0-$\bar{K}^0$, $K_L \to \mu^+\mu^-$, and $K_L \to \bar{\mu} e$ (through box and tree graphs) are smaller typically by 1 to 2 orders of magnitude than that of the standard model, while those for B^0-$\bar{B}^0$ are comparable to that of the standard model [15]. However, the model predicts intriguing new processes and effects such as the following: (i) $Z \to t\bar{c}$ with a coupling $\approx (g_2/\cos\theta_W)(\kappa_\mu/\kappa_\lambda)^2 p/2 \approx (g_2/\cos\theta_W)(2 - \frac{1}{2})\%$, which provides the genuine scope for observing a $t\bar{c}$ "resonance" in e^+e^- annihilation. This is, of course, the only way the top can be observed at the CERN e^+e^- collider LEP II if $m_t \gtrsim 100$ GeV. (ii) $Z \to c\bar{u}$ with a coupling $\approx (g_2/\cos\theta_W)(\kappa_u/\kappa_\lambda)^2 (p/2) \tilde{\delta}_1^u$ which gives $\Delta m(D\text{-}\bar{D}) \simeq (10-3) \times 10^{-14}$ GeV. This is at least 10 times larger than the standard model prediction and is in range for experimental detection [16]. (iii) $Z \to \bar{\mu} e$ with a coupling $\approx (g_2/\cos\theta_W)(\kappa_d/\kappa_\lambda)^2 (p/2) \tilde{\delta}_1^l$ leading to $B(\mu \to 3e) \approx (1-5) \times 10^{-13}$. (iv) Significant departures from unitarity in certain combinations occurring within the 3×3 part of the full CKM matrix which would imply a (4–10)% increase in top and τ lifetimes compared to standard model predictions.

Our dramatic prediction and hallmark of the model is, of course, the existence of vectorlike families Q and Q' [1,2] whose charged lepton and quark members have masses $\simeq$200–500 GeV and 0.6–1.5 TeV, respectively. This should provide rich new physics to be probed at the Superconducting Super Collider, the CERN Large Hadron Collider, and TeV-range e^+e^- colliders. All these show that the model not only provides a natural reason for the interfamily mass hierarchy and an attractive framework for CP violation [17], but (a) it is safe at present (unlike standard technicolor) and (b) it can be falsified in many ways, even at low energies.

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[5] The q_L^i's are made of preonic combinations such as $\psi_L^f \varphi_R^{c*} v$, $\varphi_L^f \psi_R^{c*} v$, $\psi_L^f \psi_R^{c*} \lambda$, and $\varphi_L^f \varphi_R^{c*} (\sigma_\mu \bar{\lambda})$; q_R^i's are obtained by switching $L \leftrightarrow R$ and $\lambda \leftrightarrow \bar{\lambda}$; while $Q_L \sim \psi_L^f \varphi_L^{c*} v$, $Q_R \sim \varphi_L^f \psi_L^{c*} v$, $Q_R' \sim \psi_R^f \varphi_R^{c*} v$, and $Q_L' \sim \varphi_R^f \psi_R^{c*} v$ [2]. Here f stands for flavor indices (x,y) and c for color indices (r,y,b,l) and $(v_\mu, \lambda, \bar{\lambda})$ denote metacolor gauge fields.

[6] Hereby, we are assuming a breakdown of global vectorial symmetries such as $SU(2)_{L+R}$ in SUSY QCD, which would be forbidden in ordinary QCD by the Vafa-Witten theorem. Whether such a breaking is permitted in SUSY QCD for which the proof of Vafa-Witten theorem does not apply is still an open problem.

[7] E. Witten, Nucl. Phys. **B185**, 513 (1981); **B202**, 253 (1983); E. Cohen and L. Gomez, Phys. Rev. Lett. **52**, 237 (1984).

[8] This is based on contributions from a single graviton exchange to the (mass)^2 of the corresponding composite Higgs boson [1,3]. It is worth noting that if the leading contributions to $\langle \boldsymbol{\lambda} \cdot \boldsymbol{\lambda} \rangle$ and $\langle \bar{\psi}\psi \rangle$ were damped by $(\Lambda_M/M_{\rm Pl})^2$, involving two-graviton exchange, the values of δm_S, m_W, m_Z, and the masses of quarks and charged leptons would still be unaltered if one chooses $\bar{a}_M(M_{\rm Pl}/10)$ such that $\Lambda_M^3/M_{\rm Pl}^2 \sim 1$ TeV, i.e., $\Lambda_M \sim 10^{13.7}$ GeV.

[9] Details of these will be given in a forthcoming paper by K.S. Babu, J. C. Pati, and H. Stremnitzer.

[10] K. S. Babu, J. C. Pati, and H. Stremnitzer, Phys. Lett. B (to be published).

[11] We allow $89 \lesssim m_t \lesssim 150$ GeV, where the lower limit is experimental and the upper is theoretical, in the model.

[12] A possible mechanism leading to $\kappa_d \approx \kappa_u (a/2\pi) \ln[\Lambda_M/(100 \text{ GeV})]$ will be discussed elsewhere.

[13] See, e.g., C. Albright, C. Jarlskog, and B. A. Lindholm, Phys. Rev. D **38**, 872 (1988).

[14] We have estimated that additional contributions to d_n through either an induced three-gluon or an induced θ term do not exceed the estimate of Eq. (10).

[15] As in all SUSY models, box graphs involving squarks could introduce additional contributions to K^0-$\bar{K}^0$. These would be suppressed, however, either if the squarks of the first two families with masses of order few TeV are highly degenerate (to within 10%) or if they are superheavy ($\gg$TeV)—a possibility that arises for a new allowed scenario for SUSY breaking [9].

[16] For a phenomenological discussion see P. Langacker and D. London, Phys. Rev. D **38**, 886 (1988).

[17] In addition, as noted in Refs. [1] and [2], it also provides a good reason for the origin of diverse mass scales from $M_{\rm Pl}$ to m_ν and of families.