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Lepton dipole moments in extended technicolor models

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Abstract

We analyze the diagonal and transition magnetic and electric dipole moments of charged leptons in extended technicolor (ETC) models, taking account of the multiscale nature of the ETC gauge symmetry breaking, conformal (walking) behavior of the technicolor theory, and mixing in the charged-lepton mass matrix. We show that mixing effects dominate the ETC contributions to charged lepton electric dipole moments and that these can yield a value of $|d_e|$ comparable to the current limit. The rate for $\mu \to e\gamma$ can also be close to its limit. From these and other processes we derive constraints on the charged lepton mixing angles. The constraints are such that the ETC contribution to the muon anomalous magnetic moment, which includes a significant lepton mixing term, can approach, but does not exceed, the current sensitivity level.

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We study the magnetic and electric dipole moments of charged leptons in a class of extended technicolor (ETC) models [1–4]. We also analyze the transition moments and the resultant electromagnetic decays. Charged lepton mixing plays a crucial role in determining each electric dipole moment (EDM). For the electron, the EDM can be comparable to the current experimental upper limit. Bounds on charged lepton mixing are derived from the constraint that the electron EDM be smaller than this limit and from upper limits on the decays $\mu \to e\gamma$, $\tau \to (e, \mu)\gamma$.

In technicolor theories, electroweak symmetry breaking (EWSB) arises from a new, strongly coupled gauge interaction at TeV energy scales [5]. Quark and lepton mass matrices arise from the embedding of technicolor in a larger gauge theory, extended technicolor (ETC) [6], which must break sequentially as the energy decreases from energies on the order of 10^3 TeV down to the TeV level. Precision measurements place tight constraints on these theories, suggesting a small number of new degrees of freedom at the TeV scale and non-QCD-like behavior of the technicolor theory. With this motivation, some attention has been focused on walking technicolor theories, which exhibit an approximate conformal behavior, with large anomalous dimensions, in the infrared [7].

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While this provides an attractive framework, it has been a challenge to construct explicit models along these lines, with, for example, the necessary ingredients to effect the requisite ETC symmetry breaking at each stage. In Refs. [1–4] a class of ETC models was developed which has these ingredients. The models are based on an ETC gauge group SU(5) which commutes with the Standard Model gauge group and breaks in stages corresponding to the three fermion generations, to a residual $SU(2)_{TC}$ technicolor gauge theory, naturally producing a hierarchy of charged lepton and quark masses. The models also exhibit charged-current flavor mixing, intra-family mass splittings, a dynamical origin of CP-violating phases in the quark and lepton sectors, and a see–saw mechanism for light neutrinos without the presence of a grand unified scale [2]. The choice $SU(2)_{TC}$ (i) minimizes the TC contributions to the electroweak S parameter, (ii) with a Standard Model family of technifermions in the fundamental representation of $SU(2)_{TC}$, can yield an approximate infrared fixed point and associated walking behavior, and (iii) makes possible the mechanism of [2] explaining light neutrinos.

The breaking to $SU(2)_{TC}$ is driven by the ETC interaction itself, a chiral gauge theory, along with one additional, strong SU(2) gauge interaction. Each requisite breaking is shown to be plausible, within strong-interaction uncertainties. A set of Standard Model-singlet fermions, including right-handed neutrinos, is naturally part of the models, and it is these particles that condense at the various ETC breaking scales, producing the ETC gauge boson masses and mixing, and leaving the residual unbroken $SU(2)_{TC}$ technicolor theory. At the scale Λ_{TC}, technifermion condensates break the electroweak symmetry, yielding $m_W^2 = (g^2/4) f_F^2 (N_c + 1)$ where g is the $SU(2)_L$ gauge coupling, $f_F \simeq 130$ GeV, and $\Lambda_{TC} \simeq 300$ GeV.

The class [1–4] has not yet led to a fully realistic model, yielding, for example, the measured fermion masses and mixings. Also, the models have a small number of phenomenologically unacceptable Nambu–Goldstone bosons, arising from spontaneously broken U(1) global symmetries [4]. To give them masses above current bounds, some additional interactions that explicitly break the U(1) symmetries must be invoked at energies above the highest ETC scales. However, all the models share interesting generic features, including a mechanism for CP violation, that are worth studying in their own right.

The bilinear fermion condensates at each stage of ETC breaking have, in general, non-vanishing phases, providing a natural, dynamical source of CP violation. The underlying theory consists of massless fermions and gauge fields, and is free of gauge anomalies. There are global chiral symmetries, some of which are anomalous, and hence are broken by instantons. The $F_{\mu\nu}\tilde{F}^{\mu\nu}$ terms associated with each (non-Abelian) gauge interaction may be rotated away by chiral transformations through the relevant global anomalies. The phases that develop at each stage of ETC breaking should be calculable [8] in a non-perturbative treatment of a fully realistic model. Here we take the phases to be arbitrary. Phases of order unity seem natural in our context and are consistent with the fact that in the quark sector the CP-violating quantity $\sin\delta$ in the CKM matrix is not small.

Below the electroweak breaking scale, the effective theory includes the Standard Model interactions, dimension-3 mass terms for the quarks and charged leptons, a dimension-3 mass term for the electroweak-doublet and Standard Model-singlet neutrinos, and a tower of operators of dimension-5 and higher describing the new physics of the model(s). The mass matrices are complex and have both diagonal and off-diagonal entries, inheriting these features from the ETC gauge boson mixings and phases arising at each breaking stage. Among the dimension-5 operators is one describing the electric and magnetic dipole moments of the charged leptons. It, too, involves a complex matrix with diagonal and off-diagonal entries. The QCD interactions include a $G_{\mu\nu}\tilde{G}^{\mu\nu}$ term with its associated strong CP problem, as well as a dimension-5 term describing chromo-magnetic and chromo-electric dipole operators for the quarks. Whether a resolution of the strong CP problem will emerge in the class of models [1–4] is not yet clear [8]. In this Letter, we focus on the charged lepton sector, including its CP violation.

The charged lepton mass matrix $M^{(\ell)}$ appears in the dimension-3 operator

$$\mathcal{L}_m = -\bar{\ell}_{L,j} M^{(\ell)}_{jk} \ell_{R,k} + \text{h.c.}, \tag{1}$$

where $\ell = (e, \mu, \tau)$ denotes the technisinglet ETC interaction eigenstates. We will estimate $M^{(\ell)}$ from the underlying ETC theory. It can be brought to real, diagonal form by the bi-unitary transformation[1]

$$U_L M^{(\ell)} U_R^{-1} = M_d^{(\ell)}. \qquad (2)$$

Hence, the interaction eigenstates are mapped to mass eigenstates ψ via $\ell_L = U_L^{-1}\psi_L$, $\ell_R = U_R^{-1}\psi_R$.

The magnetic and electric dipole matrices of the charged leptons are given by the dimension-5 operator

$$\mathcal{L}_{\text{dip}} = \frac{1}{2}\bar{\ell}_{j,L}\tilde{D}_{jk}\sigma_{\mu\nu}\ell_{k,R}F_{\text{em}}^{\mu\nu} + \text{h.c.} \qquad (3)$$

The matrix $\tilde{D}$ will also be estimated from the underlying ETC theory. In terms of mass eigenstates,

$$\bar{\ell}_L\tilde{D}\sigma_{\mu\nu}\ell_R F_{\text{em}}^{\mu\nu} + \text{h.c.} = \bar{\psi}_L D\sigma_{\mu\nu}\psi_R F_{\text{em}}^{\mu\nu} + \text{h.c.}, \qquad (4)$$

where

$$D = U_L\tilde{D}U_R^{-1}. \qquad (5)$$

Decomposing D into Hermitian and anti-Hermitian parts, $D = D_H + D_{AH}$, where $D_H = (D + D^\dagger)/2$ and $D_{AH} = (D - D^\dagger)/2$, the dipole operator is

$$\frac{1}{2}\left[\bar{\psi}D_H\sigma_{\mu\nu}\psi + \bar{\psi}D_{AH}\sigma_{\mu\nu}\gamma_5\psi\right]F_{\text{em}}^{\mu\nu}. \qquad (6)$$

Then $a_{\psi_j} = (g_{\psi_j} - 2)/2 = (2m_{\psi_j}/e)D_{H,jj}$ (where $e = -|e|$ is the lepton charge) and $d_{\psi_j} = -iD_{AH,jj}$ [9].

The M and D matrices, arising from physics at momentum scales $\gtrsim \Lambda_{\text{TC}}$, are defined at that scale. There are a variety of other quantum corrections to the physical mass and dipole matrices, coming from momentum scales $\lesssim \Lambda_{\text{TC}}$ and involving virtual particles with masses in this range. These arise from Standard Model interactions and iterations of the higher-dimension operators (e.g., [10]). We focus here on the contribution of physics at scales $\gtrsim \Lambda_{\text{TC}}$ incorporated in the above operators.

The mass operator (1) may be estimated from a graph in which an incoming lepton ℓ_k goes to an internal technifermion and ETC vector boson that has become massive at the stage in the breaking of the full ETC gauge group to $SU(2)_{\text{TC}}$ corresponding to the generation index k. The technifermion develops a soft dynamical mass as the technicolor interactions break the associated chiral symmetry, and it then recombines with an ETC vector boson to give the outgoing lepton ℓ_j. There is in general complex mixing among the ETC group eigenstates [1–4] arising at the various stages of ETC breaking, producing complex off-diagonal terms in the mass matrix $M_{jk}^{(\ell)}$.

The mass matrix $M_{jk}^{(\ell)}$ may be expressed as

$$M_{jk}^{(\ell)} \simeq \frac{8\pi\Lambda_{\text{TC}}^3\Pi_{jk}\eta}{3\Lambda_j^2\Lambda_k^2}, \qquad (7)$$

where Λ_{TC} is the technicolor confinement scale, Λ_j is the ETC scale associated with the family index j, and Π_{jk} is a complex function of the ETC scales with mass-squared dimensions arising from the ETC gauge boson mixing. The $8\pi/3$ factor is derived in Ref. [4]; this factor and coefficients of other expressions herein are subject to uncertainties due to the strong-coupling nature of the ETC interactions. The magnitude of Π_{jk} is no greater than $\min(\Lambda_j^2, \Lambda_k^2)$, and hence may be treated perturbatively through a single insertion. The factor η is $O(1)$ in a QCD-like technicolor theory, while in a theory with walking from Λ_{TC} to the lowest ETC scale Λ_3 and with anomalous dimension 1, $\eta = O(\Lambda_3/\Lambda_{\text{TC}})$. The off-diagonal terms in $M_{jk}^{(\ell)}$ determine the structure of U_L and U_R, the former entering along with the corresponding transformation connecting the neutrino interaction and mass eigenstates in the observed lepton mixing matrix.

To estimate $\tilde{D}$, we note that the relevant graph is obtained from the corresponding mass graph by the coupling of a photon to the internal technifermion.[2] It is more convergent than the mass graph by two powers of the loop momentum. The fact that the necessary technifermion mass is soft above Λ_{TC} then leads to the convergence of the momentum integral at momenta of this order—well below the ETC scales. Thus a potential contribution of order $1/\Lambda_{\text{ETC}}^2$ may be estimated by setting to zero the momentum flowing through the ETC gauge boson propagator including the mixing; this effectively replaces the ETC vector boson exchange by a four-fermion interaction, which

[1] In specific ETC models, the structure of $M^{(\ell)}$ may be such that U_L and U_R are related. We keep our present discussion more general.

[2] Since the ETC gauge bosons are neutral, there is no contribution to neutrino dipole moment matrices. Also, since in [2,4] the neutrinos are generically Majorana particles, the diagonal dipole moments (μ_ν and d_ν) vanish.

does not yield a contribution to the dipole moment operator.

The leading contribution to $\tilde{D}_{jk}$ therefore has additional inverse powers of the ETC mass scales and arises from integration momenta on this order. The resultant integral producing this leading term has the same power counting as the integral for $M^{(\ell)}_{jk}$ (7) in the momentum range from $\Lambda_{\rm TC}$ up to the lowest relevant ETC scale, but at this scale and above it is quite different in detail. Taking account of multiple ETC scales for different generations, the elements of the dipole matrix $\tilde{D}_{jk}$ are

$$\tilde{D}_{jk} \simeq \frac{eM^{(\ell)}_{jk}}{\Lambda^2_{jk}}, \tag{8}$$

where Λ_{jk} is of order the scale at which the relevant ETC propagator including the mixing function becomes soft. It is no greater than $\min(\Lambda_j, \Lambda_k)$, and can be less.

We note that Eqs. (7) and (8) can be obtained also using an effective ETC theory, employing, for $E < \Lambda_j$, local operators of dimension-six and higher. Dimension-six (four-fermion) operators generate $M^{(\ell)}_{jk}$ while dimension-eight operators (four-fermion with derivative couplings) similarly generate $\tilde{D}_{jk}$. However, ETC gauge theories, such as ours, operative at ETC scales and above, provide more information about the scales in Eqs. (7) and (8).

Since $\tilde{D}_{jk}$ is not, in general, $\propto M^{(\ell)}_{jk}$, it is not diagonalized by the transformation that diagonalizes $M^{(\ell)}$. The transformation yields instead the (non-diagonal and complex) dipole matrix D (5). A principal result of our analysis is that mixing generically has an important effect on charged-lepton dipole moments in ETC theories. This is true even for relatively small mixing.

For numerical estimates of the dipole matrix, we take $\Lambda_1 \simeq 10^3$ TeV, $\Lambda_2 \simeq 10^2$ TeV, and $\Lambda_3 \simeq 4$ TeV, as in Refs. [2,4]. These values can yield realistic ranges for quark and lepton masses, since for down quarks and charged leptons there can be a natural suppression so that $|\Pi_{jk}| \lesssim \min(\Lambda^2_j, \Lambda^2_k)$. A mechanism for this suppression, using relatively conjugate ETC representations for these fields, is given in [4], although none of the models yet yields exactly the requisite values for the Π_{jk}'s.

Each of the matrices U_χ, $\chi = L, R$, depends on three rotation angles θ^χ_{mn}, $mn = 12, 13, 23$, and six phases $e^{i\alpha^\chi_j}$, $e^{i\beta^\chi_j}$, $j = 1, 2, 3$. We use the conventions of [11] for the θ^χ_{mn} and write $U_\chi = P^\chi_\alpha R_{23}(\theta^\chi_{23}) \times R_{13}(\theta^\chi_{13}) R(\theta^\chi_{12}) P^\chi_\beta$, where $R_{mn}(\theta^\chi_{mn})$ is the rotation through θ^χ_{mn} in the mn subspace and $P^\chi_\alpha = \mathrm{diag}(e^{i\alpha^\chi_1}, e^{i\alpha^\chi_2}, e^{i\alpha^\chi_3})$. The D_{jk} are independent of P^χ_β, $\chi = L, R$. The rotation angles are model-dependent, typically small if the off-diagonal Π_{jk}'s are more suppressed than the diagonal ones, and large if this is not the case.

For the electron EDM, using Eqs. (5) and (8), we find

$$\frac{d_e}{e} \simeq \frac{m_\tau\, \mathrm{Im}(F_{11,3})}{\Lambda^2_3}, \tag{9}$$

where we have kept only the term with the largest (τ) lepton mass in the numerator and the smallest ETC scale in the denominator. Here, $F_{11,3}$ is a dimensionless function of the parameters in U_L and U_R, of $O(1)$ for generic values of these parameters, which vanishes if all phases are 0 mod π or if mixing angles vanish. The complex phases remain in $D = U_L \tilde{D} U_R$ because of the non-proportionality of $\tilde{D}$ and $M^{(\ell)}$. If mixing were absent, the phases in $\tilde{D}$ would be the same as the phases in $M^{(\ell)}$ (both diagonal) and would be removed by the transformation that makes the latter real. In a series expansion in small rotation angles up to quadratic order, $F_{jk,3} = e^{i[(\alpha^L_j-\alpha^L_3)-(\alpha^R_k-\alpha^R_3)]}\theta^L_{j3}\theta^R_{k3}$ for $j, k \neq 3$.

The current upper limit on the electron EDM is $|d_e| < 1.6 \times 10^{-27}$ e cm [12], and ongoing experiments project sensitivities down to 10^{-30} e cm or better [13–15]. Comparing Eq. (9) to the upper limit, with $\Lambda_3 \simeq 4$ TeV, we see that $\mathrm{Im}(F_{11,3})$ must be much less than $O(1)$; in fact, $\mathrm{Im}(F_{11,3}) \lesssim 0.7 \times 10^{-6}$. Taking the phases to be generic, of $O(1)$, we conclude that the mixing angles must be small. To bound them, we use the above expression for $F_{jk,3}$, neglecting terms beyond quadratic order in the products of the various θ's. (Higher-order terms can be included in a more detailed analysis.) We then have $|\theta^L_{13}\theta^R_{13}| \lesssim 10^{-6}$.[3] Values of the angles in this range are not unexpected, given the suppression of the off-diagonal mixings Π_{jk}

[3] These mixings also enter $a_{e,\mathrm{ETC}}$ so that it is far below the experimental accuracy $\Delta a_e = 4 \times 10^{-12}$ [11].

in some of the models of Refs. [2,4]. Even if the product $|\theta^L_{13}\theta^R_{13}|$ is somewhat below the upper bound, the ETC contribution to d_e can naturally lie in the range accessible to ongoing experiments.

The off-diagonal elements D_{jk} produce flavor-changing radiative lepton decays. From the upper bounds on these, in conjunction with the above values of the ETC scales Λ_j chosen to yield appropriate fermion masses, we can derive additional bounds on charged lepton mixing. We have $\Gamma(\psi_k \to \psi_j\gamma) = (|D_{jk}|^2 + |D_{kj}|^2)m^3_{\psi_k}/(8\pi)$. For example, for the case $k = 2$, $j = 1$, i.e., $\mu \to e\gamma$, the terms with the dominant ETC scale dependence are

$$\frac{D_{jk}}{e} \simeq \frac{m_\tau F_{jk,3}}{\Lambda_3^2}, \quad jk = 12, 21. \tag{10}$$

We infer an upper bound from the limit $B(\mu \to e\gamma) < 2.1 \times 10^{-11}$ [11,16], viz., $|\theta^L_{13}\theta^R_{23}|$, $|\theta^R_{13}\theta^L_{23}| \lesssim 5 \times 10^{-6}$. Again, this range of values is consistent with some models in [2,4].[4]

Since U_L enters along with neutrino mixing into the observed lepton mixing in neutrino oscillations, our bounds on θ^L_{jk} suggest,[5] that mixing in the neutrino sector is the primary source of the large measured lepton mixing angles θ_{23} and θ_{12}. Large neutrino-sector mixing can emerge naturally from some of the models of Refs. [2,4]. It is also the case that large leptonic CP violation, as could be observed in future neutrino oscillation experiments, is natural in these models.

From the limits $B(\tau \to \mu\gamma) < 3.1 \times 10^{-7}$ [17] and $B(\tau \to e\gamma) < 2.6 \times 10^{-7}$ [11] we obtain the respective bounds $|\theta^L_{23}|$, $|\theta^R_{23}| \lesssim 0.02$ and $|\theta^L_{13}|$, $|\theta^R_{13}| \lesssim 0.06$. The linear form of these bounds is due to the fact that one of the external particles is the τ. The dominant terms in the amplitude for the respective decays, in a small-θ expansion, are $D_{j3} \simeq e^{i(\alpha^L_j - \alpha^L_3)}\theta^L_{j3}m_\tau/\Lambda_3^2$ and $D_{3j} \simeq e^{i(\alpha^R_3 - \alpha^R_j)}\theta^R_{j3}m_\tau/\Lambda_3^2$, with $j = 2, 1$.

For the muon $g - 2$, keeping the dominant terms,

$$\frac{a_\mu}{2m_\mu} \simeq \frac{m_\mu}{\Lambda_{22}^2} + \frac{m_\tau \operatorname{Re}(F_{22,3})}{\Lambda_3^2}, \tag{11}$$

where Λ_{22} is the softness scale of the relevant ETC exchange, ranging, in the models explored, from Λ_3 ($\simeq 4$ TeV) to Λ_2 ($\simeq 10^2$ TeV). The first term in (11) would be present even without mixing [18]. For $\Lambda_{22} \simeq \Lambda_2$ and $|\theta^L_{23}|$ and $|\theta^R_{23}|$ bounded as above, the second term can dominate the first, but the ETC contribution to a_μ is $\lesssim 10^{-11}$, well below the current uncertainty $\sim 10^{-9}$ in the comparison of theory and experiment [19]. However, if $\Lambda_{22} = O(\Lambda_3)$ as in some models, the first term would dominate and, interestingly, would be of order 10^{-9}.

For the muon EDM, we find

$$\frac{d_\mu}{e} \simeq \frac{m_\tau \operatorname{Im}(F_{22,3})}{\Lambda_3^2}. \tag{12}$$

With $|\theta^L_{23}\theta^R_{23}|$ near the upper limit 4×10^{-4}, we estimate that $|d_\mu|$ could be $\simeq 10^{-24}$ e cm. This is well below the current limit $|d_\mu| < 3.7 \times 10^{-19}$ e cm [11,20] but might be observable in the proposed experiment of [21].

For the τ-lepton, the bounds $|a_\tau| < 0.06$ and $|d_\tau| < 3.1 \times 10^{-16}$ e cm [11] are not sensitive to the contributions described here, arising from physics at scales $\geq \Lambda_{TC}$.

In summary, we have analyzed the magnetic and electric dipole moments of charged leptons in a class of ETC models with lepton mixing and dynamically generated CP-violating phases. We have shown that the ETC contribution to the electron EDM is dominated by terms from charged lepton mixing and can be comparable to the current experimental limit. We have used current limits on $|d_e|$ and radiative lepton decays to set bounds on charged lepton mixing angles. We have noted that these constraints are such that the ETC contribution to the muon anomalous magnetic moment, which includes a significant lepton mixing term, can approach, but does not exceed, the current sensitivity level.

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[4] We also considered limits from bounds on $\mu \to e$ conversion; these are less stringent than those from $\mu \to e\gamma$.

[5] See Footnote 1.

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