Natural Strong *CP* Conservation in Flipped Physics

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A natural axion-free solution of the strong *CP* problem *at tree level* is noted within an E(6) grand unified theory. Using this as a springboard, it is shown that several flipped SU(5) theories which occur in superstring phenomenology contain within them a mechanism which enforces $\bar{\theta}=0$ at high accuracy.

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The most popular solution of the strong *CP* problem depends on having an axion as a physical particle. Most of the allowed mass range for the axion has been already excluded, leaving only two small windows, 10^{-5}–10^{-3} eV for the generic invisible axion and between about 2 and 5 eV for the hadronic axion which couples only to quarks. As these small windows become walled up, it is timely to seek a fallback axionless position. Searches are in progress to brick up both windows. Of all the alternative axionless scenarios, the most attractive may be the idea of arranging a real determinant of the *tree-level* mass matrix in a natural way.

Some examples of unified theories which fulfill the requisite requirements have been published long ago.[1] All of these examples were somewhat contrived. Here we shall present what we consider the most elegant class of such axion-free solutions. It is based on the choice of grand-unified-theory (GUT) group as E(6), and relies on some group-theoretic niceties peculiar to this exceptional group. Incidentally, it is the only such solution that fits closely in the philosophy of the heterotic superstring.

According to the model-building rules expounded in Ref. 1, the theory must contain a real representation of exotic fermions additional to that of the standard model, and the Yukawa couplings to the Weinberg-Salam doublet must involve only the usual quarks and leptons. Soft (spontaneous) *CP* violation occurs only by vacuum expectation values (VEV's) coupling the usual to the exotic fermions.

These requirements are simple to satisfy in E(6). For the three families of quarks and leptons we shall use three **27**'s of E(6). In the Higgs sector we shall consider only **27**'s and the adjoint **78**. More than one **27** is necessary to break unwanted symmetries and the **78** is needed to break SU(5). Within our class of models, some will avoid even the **78** altogether.

The Higgs potential for the **78** of Higgs fields ϕ^a is $(a,b=1\text{–}78)$

$$V(\Phi^a)=M^2\delta_{ab}\Phi^a\Phi^b+\Lambda\delta_{ab}\delta_{cd}\Phi^a\Phi^b\Phi^c\Phi^d\,. \tag{1}$$

For the **27** of Higgs fields ϕ^α it is $(\alpha=1\text{–}27)$

$$V(\phi^\alpha)=\mu^2\phi^\alpha\bar{\phi}_\alpha+\lambda(\phi^\alpha\bar{\phi}_\alpha)^2+\lambda'(\phi^\alpha\phi^\beta\phi^\gamma d_{\alpha\beta\gamma}+\text{H.c.})+\lambda''\phi^\alpha\phi^\beta\bar{\phi}_\gamma\bar{\phi}_\delta d^{\gamma\delta\varepsilon}d_{\alpha\beta\varepsilon}\,. \tag{2}$$

Cross terms are also possible:

$$V(27,78)=\chi\Phi_a\Phi_a\phi^\alpha\bar{\phi}_\alpha+\gamma(T_a)^\alpha{}_\beta\Phi_a\phi^\beta\bar{\phi}_\alpha+\xi[(T_a)^\alpha{}_\delta d^{\beta\gamma\delta}\Phi_a\bar{\phi}_\alpha\bar{\phi}_\beta\bar{\phi}_\gamma+\text{H.c.}]\,. \tag{3}$$

The decomposition of the **27** fermions under E(6) $\to$ SU(5) is (F denotes complex family, R denotes real exotics)

$$\mathbf{27}=(\mathbf{10}+\bar{\mathbf{5}})_F+(\mathbf{5}+\bar{\mathbf{5}}+\mathbf{1}+\mathbf{1})_R \tag{4}$$

and for each family we may take the final twelve states as the additional real representation of exotic fermions. Under E(6) $\to$ SU(6)$\times$SU(2) $\to$ SU(5)$\times$SU(2) [SU(5) is the Georgi-Glashow SU(5)] one has

$$\mathbf{27}=(\bar{\mathbf{6}},\mathbf{2})+(\mathbf{15},\mathbf{1}) \tag{5}$$

$$=(\bar{\mathbf{5}}+\mathbf{1},\mathbf{2})+(\mathbf{10}+\mathbf{5},\mathbf{1}) \tag{6}$$

$$=\psi_{Ki}+\psi_i+\psi^{KL}+\psi_K\,, \tag{7}$$

where $K,L=1\text{–}5$ of SU(5) and $i=1,2$ of SU(2). Let us identify a family as $(\psi_{K1}+\psi^{KL})_F$ and the exotics as $(\psi_{K2}+\psi^K)_R$. With the same notation for the scalars ϕ^α,

$$\phi^\alpha=\phi_{Ki}+\phi_i+\phi^{KL}+\phi^K\,, \tag{8}$$

we shall break the E(6) to the standard-model group by giving VEV's to the ϕ_i; we then break $SU(2)_L\times U(1)_Y$ by the doublet contained in ϕ_{K2}. By noticing that the Yukawa couplings have the form

$$h\psi^{KL}\psi_{Ki}\phi_{Lj}\varepsilon^{ij}\,, \tag{9}$$

it is clear that the VEV $\langle\phi_{K2}\rangle$ couples F-F but *not* F-R or R-R.

Soft spontaneous *CP* violation is accomplished by giving a *complex* VEV to ϕ_2 which couples

$$h'\psi^K\psi_{Ki}\phi_j\varepsilon^{ij}\,. \qquad (10)$$

With this arrangement, the determinant of the quark mass matrix is real. At the tree level, this means that $\bar\theta=\theta_{\rm QED}+\theta_{\rm QFD}=0$ (where QFD denotes quantum flavor dynamics). In loop diagrams, there are nonzero contributions to θ but these can easily be consistent with the bound $\bar\theta<10^{-10}$ imposed by the neutron electric dipole moment. Supersymmetry can make the corrections even smaller, as discussed later.

The above example uses the canonical embedding of the SM in E(6). It is also possible to give a natural solution of strong *CP* with the "flipped" embeddings.[2] Unlike the canonical embedding, these will no longer need an adjoint Φ_a of Higgs fields. We consider five such flipped examples.

(i) We may flip the SU(5) so that each family is contained in a **16** of SO(10) but now with the e_L^+ and v_R exchanged relative to the usual SU(5) embedding. Specifically in terms of $SU(3)_C\times SU(2)_L\times U(1)_Y$ with $Y'=\frac{1}{5}Y-\frac{6}{5}B$, the **16** becomes

$$\mathbf{16}=\mathbf{1}+\bar{\mathbf{5}}+\mathbf{10} \qquad (11)$$

$$=(1,1)+[(\bar 3,1)_{-2/3}+(1,2)_{-1/2}]$$
$$+[(1,1)_0+(\bar 3,1)_{1/3}+(3,2)_{1/6}]\,, \qquad (12)$$

while

$$\mathbf{10}=\mathbf{5}+\bar{\mathbf{5}}\,. \qquad (13)$$

We give a real VEV to the doublet within the $\bar{\mathbf{5}}$ of Eq. (13) and a complex spontaneous *CP*-violating VEV to the $(1,1)_0$ in Eq. (12). At this point one needs *either* an adjoint **78** of Higgs in E(6) *or* to use the four-dimensional superstring constriction[3] to break E(6) to the appropriate subgroup. In either case, the theory with flipped SU(5) has natural strong *CP* conservation.

(ii) There is a doubly flipped model[4] in which the embedding of $SU(3)_C\times SU(2)_L\times U(1)_{Y''}$, with Y'' a combination of the U(1) in E(6) that is not in SO(10) with the two U(1)'s inside SO(10) [unlike case (i) where the $U(1)_{Y'}$ was entirely inside SO(10)]. In this case the charges are

$$\mathbf{27}=(1,1)^F_{+1}+[(3,1)^R_{-1/3}+(1,2)^R_{-1/2}+(\bar 3,1)^F_{-2/3}+(1,2)^F_{-1/2}]$$
$$+\{(1,1)^R_0+[(\bar 3,1)^R_{1/3}+(1,2)^R_{1/2}]+[(3,2)^F_{1/6}+(\bar 3,1)^F_{1/3}+(1,1)^R_0]\}\,. \qquad (14)$$

In this, we identify F and R fermions as indicated. The VEV's are in the components of the scalar corresponding to the two $(1,1)^R_0$ pieces and to the doublet $(1,2)^R_{-1/2}$ which breaks $SU(2)_L\times U(1)_Y$. The $(1,1)^R_0$ in the **10** of the usual SU(5) may have a *CP*-violating VEV. As in the case (i) above, the breaking can be achieved only if E(6) is broken to an appropriate rank-6 subgroup as in the four-dimensional superstring,[3] or we must alternatively add an adjoint **78** of Higgs fields, which does not spoil the solution.

(iii)–(v) There are three types of "isoflipped" E(6) models[5] and we have checked that all of these also satisfy the requirements of natural strong *CP* conservation. We shall give only one example; the others are similar.

Here the idea is that the embedding is $E(6)\supset SU(5)\times SU(2)_L\times U(1)'''$ where $U(1)'''$ is contained in SU(6) of $E(6)\supset SU(6)\times SU(2)_L$. There are three independent choices of $U(1)'''$ identification. In one case the **27** contains

$$(5,2)=(3,2)^F_{1/6}+(1,2)^R_{1/2}+(1,2)^R_{-1/2}\,,$$

$$(1,2)=(1,2)^F_{1/2}\,,$$

$$(\overline{10},1)=(\bar 3,1)^F_{-2/3}+(\bar 3,1)^F_{1/3}+(3,1)^R_{-1/3}+(1,1)^R_0\,,$$

$$(\bar 5,1)=(\bar 3,1)^R_{1/3}+(1,1)^R_0+(1,1)^F_{+1}\,.$$

We identify the F and R fermions as indicated. The VEV's are in the components of the scalar corresponding to the $(1,1)_0$ in (10,1) (real), the $(1,1)_0$ in (5,1) (*CP* violating), and the $(1,2)_{1/2}$ in $(\bar 5,2)$ (real). Again this type of symmetry breaking is consistent with the four-dimensional superstring.[3]

The above analysis shows how to find natural strong *CP* conservation *at tree level* in E(6) GUT's and in superstring-inspired flipped SU(5) models.

At loop level the situation is more subtle and related to the usual gauge-hierarchy discussion. We have assumed that certain Higgs scalar components have nonzero VEV's and that all others have zero VEV's. In general, if we arrange the former to be true in an otherwise generic Higgs potential, then the latter may not be true for all of the parameter space. This would have the effect of ruining the naturalness of our scheme. But this is not the appropriate argument since in the usual discussions of superheavy and light VEV's from the same Higgs potential (the gauge-hierarchy problem) it is usually postulated that there exists *some* region of parameter space (perhaps "almost nowhere") where the required gauge hierarchy can be obtained, even at tree level. At loop level, it is similarly postulated that retuning of parameters is possible which can maintain the hierarchy.

For the natural solution of strong *CP* that we have discussed, one should be allowed equal license. If we can arrange the tree-level Higgs potential, for some limited domain of parameter space such that only those scalar components have a VEV that has been postulated (and all others are zero), then this is allowable. In the loop expansion we may postulate the possibility of retuning such that the arrangement of the VEV's remains stable under perturbation theory. This is exactly analogous to the situation with regard to the hierarchy problem.

In general, one expects loop diagrams to be small in any case, if the Yukawa and Higgs couplings are sufficiently small.[1] In supersymmetrized generalizations, loop corrections to $\bar{\theta}_{\rm QCD}$ are even more suppressed.[6] In fact, the mechanism we have discussed becomes more natural in a supersymmetrized GUT. To avoid giving unwanted VEV's even at tree level in E(6) we have tuned certain Higgs-boson self-couplings to zero; in the supersymmetric case this is more natural because radiative corrections are strongly suppressed (the coupling remaining zero in the absence of supersymmetry breaking).

This last remark about supersymmetry leads us naturally to the subject of superstring phenomenology. Such schemes, in general, involve a subgroup of E(6) in four spacetime dimensions. The most appealing such schemes are based on flipped models.[3,4] Despite the number of articles in Ref. 3 on single hyperflipped SU(5), we feel that not enough work has been done on the very exciting topic of flipped physics from the superstring because, among other things, the all important topic of strong *CP* has not been fully explored. Our present discussion has brought a new class of strong *CP*-conserving hyperflipped and isoflipped models within our grasp, all of which are *a priori* as interesting as flipped $SU(5)\times U(1)$.

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[1]A. Nelson, Phys. Lett. **136B**, 387 (1984); **143B**, 165 (1984); S. M. Barr, Phys. Rev. Lett. **53**, 329 (1984); Phys. Rev. D **30**, 1805 (1984); Phys. Rev. Lett. **55**, 2253 (1985).

[2]A. De Rujula, H. Georgi, and S. L. Glashow, Phys. Rev. Lett. **45**, 413 (1980); S. M. Barr, Phys. Lett. **112B**, 219 (1982); H. Georgi, in *Proceedings of the Twenty-First International Conference on High Energy Physics* [J. Phys. (Paris), Colloq. **43**, C3-1 (1982)]; T. W. Kephart and N. Nakagawa, Phys. Lett. **141B**, 329 (1984).

[3]I. Antoniadis, J. Ellis, J. S. Hagelin, and D. V. Nanopoulos, Phys. Lett. B **194**, 231 (1987); B. Campbell, J. Ellis, J. S. Hagelin, D. V. Nanopoulos, and K. Olive, *ibid.* **197**, 355 (1987); B. Campbell, J. Ellis, J. S. Hagelin, D. V. Nanopoulos, and R. Ticciati, *ibid.* **198**, 200 (1987); K. T. Tamvakis, *ibid.* **201**, 95 (1988); M. Drees, *ibid.* **206**, 259 (1988); **206**, 265 (1988); B. Campbell, J. Ellis, J. S. Hagelin, D. V. Nanopoulos, and K. Olive, *ibid.* **200**, 483 (1988); J. Ellis, J. S. Hagelin, D. V. Nanopoulos, and K. Olive, *ibid.* **207**, 451 (1988); J. Ellis, J. S. Hagelin, S. Kelly, D. V. Nanopoulos, and K. Olive, *ibid.* **209**, 283 (1988). For a recent review with a more extensive bibliography, see D. V. Nanopoulos, in *Proceedings of the Last Workshop on Grand Unification,* edited by P. H. Frampton (World Scientific, Singapore, 1989), p. 202.

[4]T. W. Kephart and N. Nakagawa, Phys. Rev. D **30**, 1978 (1984); J. Rizos and K. Tamyakis, Phys. Lett. B **212**, 176 (1988).

[5]T. W. Kephart and T. C. Yuan, Phys. Lett. B **231**, 275 (1989).

[6]L. J. Hall and L. J. Randall, Nucl. Phys. **B271**, 57 (1986); A. Dannenberg, L. J. Hall, and L. J. Randall, Nucl. Phys. **B274**, 157 (1986); S. M. Barr and A. Masiero, Phys. Rev. D **38**, 366 (1988); S. M. Barr, Phys. Rev. D **41**, 2129 (1990).