CP VIOLATION WITHOUT ELEMENTARY SCALAR FIELDS

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ABSTRACT

Dynamically broken gauge theories of electroweak interactions provide a natural mechanism for generating CP violation. Even if all vacuum angles are unobservable, strong CP violation is not automatically avoided. In the absence of strong CP violation, the neutron electric dipole moment is expected to be of order 10^{-24} e cm.

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In this letter, we show that there is a natural mechanism for generating CP violation in dynamically broken electroweak gauge theories. Our proposal requires no new gauge interactions beyond those discussed previously.^{1,2} Spontaneous CP violation can appear when we carry out Dashen's³ procedure to identify the correct chiral vacuum. Whether CP violation occurs is determined in principal by only the gauge group and the fermion representation content of the theory.

These theories have no elementary scalar fields and no fermion bare masses. Even though there are no observable vacuum angles,⁶ strong CP violation⁷ is not automatically avoided.⁸ We state a criterion for the absence of strong CP violation; if it is satisfied, CP violating phases in the quark mass matrix are naturally suppressed by a factor of order 10^{-9} . Additional CP violation appears in the electroweak interaction and in the gauge interaction responsible for chiral symmetry breaking. We predict that the electric dipole moment of the neutron is of order 10^{-28} ecm.

In a theory of weak interactions without elementary scalar fields, a new gauge interaction^{1,2} ("hypercolor", with gauge group G_H) is required, in addition to the familiar color ($G_C = SU(3)$) and electroweak ($G_W = SU(2) \times U(1)$) interactions. Hypercolor becomes a strong interaction at a mass scale $m_H = 1$ TeV, and drives the breaking of G_W down to $U(1)_{EM}$.

A theory with gauge group $G_H \times G_C \times G_W$ alone cannot be realistic. A "side-ways" interaction^{1,2} (with group G_S) is needed to break explicitly all chiral symmetries not gauged by G_W . G_S is dynamically broken to a subgroup containing $G_H \times G_C$ at a mass scale $m_S = 100$ TeV. (We need not speculate here on the origin of the G_S breakdown.)^{9,10}

All fermions are in at most four irreducible representations of G_S .³ In the effective gauge theory which describes physics below 100 TeV, each

of these representations transforms reducibly under $G_H \times G_C$. If we neglect the broken sideways interactions and the weak interactions, the $G_H \times G_C$ -invariant effective Hamiltonian $\mathcal{H}_0$ respects a global (chiral) flavor-symmetry group G_f . G_W is a subgroup of G_f .

When hypercolor and color become strong, G_f is dynamically broken to a subgroup S_f . Many Goldstone bosons result. Three of these are eaten by the weak $W^\pm$ and Z^0 bosons.¹² The remaining Goldstone bosons acquire mass from the chiral symmetry breaking perturbation $\mathcal{K}$ generated by the weak and sideways interactions.

The ground state of $\mathcal{H}_0$ is highly degenerate; the vacua are parametrized by the coset space G_f/S_f . The perturbation $\mathcal{K}$ lifts the degeneracy and picks out the true chiral-perturbative vacuum, the limit of the ground state of $\mathcal{H}_0 + \epsilon \mathcal{K}$ as $\epsilon \rightarrow 0$. As Dashen³ explained, we can identify the correct vacuum by minimizing an effective potential which, in lowest order perturbation theory, is

$$V(g) = \langle \Omega | U^{-1}(g) \mathcal{K} U(g) | \Omega \rangle. \quad (1)$$

Here $g \in G_f$, $U(g)$ represents G_f in the Hilbert space of states, and $|\Omega\rangle$ is the S_f -invariant vacuum. In the calculations described below, it is convenient to regard the symmetry group of the vacuum as fixed. Then, from among all G_f -equivalent perturbations $\mathcal{K}(g) = U^{-1}(g) \mathcal{K} U(g)$, we choose that one which minimizes the energy of the S_f -invariant vacuum $|\Omega\rangle$.

Now we see how spontaneous CP violation can arise naturally. We assume that the fermion representation content under G_S is such that all G_S vacuum angles⁶ can be simultaneously rotated to zero, and hence are unobservable. In particular, G_S might be simple.¹¹ We further assume that the breaking of G_S does not introduce CP violation, so that the effective gauge theory below 100 TeV is CP-invariant.^{12,13} Thus the vacuum $|\Omega\rangle$ and the perturbation $\mathcal{K}$ are CP-invariant, and the effective potential is CP-symmetric. However, the

energy might be minimized by a CP-violating $\mathcal{K}(g)$; then $V(g)$ has a degenerate minimum, and CP is spontaneously broken.¹⁴ The minimum of $V(g)$, and whether spontaneous CP violation occurs, is determined once the pattern of G_S and G_f breaking is known.

The approximate flavor group G_f is determined by the $G_H \times G_C$ representation content of the fermions. We will assume that all nontrivial representations of $G_H \times G_C$ are complex, that the color and hypercolor interactions are vectorial, and that the maximal isospin is left unbroken when hypercolor and color get strong. (These assumptions ensure that G_W breaks down to $U(1)_{\Box}$, and that the relation $M_W/M_Z = \cos\theta_W$ is satisfied.²) Then the fermions may be denoted $\psi_{Lr1}^{(\rho)}$, $\psi_{Rr1}^{(\rho)}$. The index ρ identifies the irreducible representation $\mathcal{D}^{(\rho)}$ according to which $\psi^{(\rho)}$ transforms under $G_H \times G_C$. The gauge group $G_H \times G_C$ acts on the index 1. The index r ($r=1,2,\dots,n_\rho$) labels the various flavors of fermions which transform as the representation $\mathcal{D}^{(\rho)}$.

The flavor group is

$$G_f = \prod_\rho [SU(n_\rho) \times SU(n_\rho) \times U_V^{\rho}(1)] \times U_A(1)'s, \quad (2)$$

which breaks down to $S_f = \prod_\rho [SU(n_\rho) \times U_V^{\rho}(1)]$. Under G_f , the fermions transform as

$$\psi_{Lr}^{(\rho)} \rightarrow W_{r'r}^{L(\rho)} \psi_{Lr'}^{(\rho)}, \quad \psi_{Rr}^{(\rho)} \rightarrow W_{r'r}^{R(\rho)} \psi_{Rr'}^{(\rho)} \quad (3)$$

where $W^{L(\rho)}$ and $W^{R(\rho)}$ are unitary $n_\rho \times n_\rho$ matrices. If S_f is the diagonal subgroup with $W^L = W^R$, then the elements of G_f/S_f can be labeled by a set of unitary matrices $\{W^{(\rho)} = W^{L(\rho)\dagger} W^{R(\rho)}\}$.

Each representation $\mathcal{D}^{(\rho)}$ of $G_H \times G_C$ has a hypercolor and color anomaly given by

$$\partial^\mu J_{5\mu}^{(\rho)} = T_\rho^H \frac{g_H^2}{8\pi^2} \text{tr } F_H \tilde{F}_H + T_\rho^C \frac{g_C^2}{8\pi^2} \text{tr } F_C \tilde{F}_C \quad (4)$$

where T_ρ^H (T_ρ^C) is the trace of the square of the hypercolor (color) generators in the representation $\mathcal{D}^{(\rho)}$. (We have assumed that G_H is simple.) All linear

combinations of the $J_{5\mu}^{(\rho)}$'s with vanishing hypercolor and color anomalies generate $U_A(1)$ symmetries which are included in G_f . Because G_f contains only anomaly-free $U_A(1)$ factors, the $W^{(\rho)}$'s satisfy the constraint

$$\prod_\rho [\det W^{(\rho)}]^{T_\rho^H} = \prod_\rho [\det W^{(\rho)}]^{T_\rho^C} = 1. \quad (5)$$

The weak interactions play a secondary role in determining the minimum of the effective potential;¹⁵ we need consider only the broken sideways interactions. We integrate out the massive sideways gauge bosons to obtain a $G_H \times G_C \times G_W$ -invariant effective Lagrangian. The leading G_f -breaking terms in the effective Lagrangian are four-fermion operators. Higher dimension operators are suppressed by additional powers of m_g^{-1} . We simplify the discussion by assuming that $SU(2)_W$ commutes with G_S , so that there is an exact global $U(1)$ symmetry which distinguishes weak doublets from weak singlets. After Fierz rearrangements, the most general four-fermion operator invariant under $G_H \times G_C \times G_W$ and the global $U(1)$ which contributes to the effective potential has the form¹⁶

$$\mathcal{K} = \Gamma_{rr's's'}^{\rho\sigma,a} \bar{\psi}_{Lr1}^{(\rho)} \psi_{Rr'j}^{(\rho)} \bar{\psi}_{Rsk}^{(\sigma)} \psi_{Ls'l}^{(\sigma)} \epsilon_{ijkl}^{\rho\sigma,a}. \quad (6)$$

Here the $\epsilon_{ijkl}^{\rho\sigma,a}$ are $G_H \times G_C$ -invariant tensors, indexed by a , and $\Gamma_{rr's's'}^{\rho\sigma,a} = O(m_g^{-2})$ is a G_W -invariant tensor which can be calculated to arbitrary order in the sideways coupling g_g^2 . To lowest order in g_g^2 , $\mathcal{K} = g_g^2 (\mu_g^{-2})_{ab} J_{L(R)}^{a\mu} J_{R(L)}^{b\mu}$, where μ_g^2 is the massive sideways gauge boson mass matrix, and the $J_{L(R)}^{a\mu}$ are the left (right)-handed broken sideways currents. This interaction can be Fierz-transformed into the form of Eq. (6).

To compute the effective potential, we note that the S_f -invariant vacuum $|\Omega\rangle$ has the property $\langle\Omega|\bar{\psi}_{Lr1}^{(\rho)} \psi_{Rr'j}^{(\sigma)}|\Omega\rangle = \delta^{\rho\sigma} \delta_{rr'} \delta_{1j} \Delta^\rho$. Therefore S_f -invariance implies that

$$\langle\Omega|\bar{\psi}_{Lr1}^{(\rho)} \psi_{Rr'j}^{(\rho)} \bar{\psi}_{Rsk}^{(\sigma)} \psi_{Ls'l}^{(\sigma)} \epsilon_{ijkl}^{\rho\sigma,a}|\Omega\rangle = \Lambda^{\rho\sigma,a} \delta_{rr'} \delta_{ss'} + \Delta^{\rho,a} \delta^{\sigma\sigma} \delta_{rs'} \delta_{r's}. \quad (7)$$

The second term on the right-hand side of Eq. (7) is G_f -invariant. Hence, the leading term in the effective potential is

$$V(W) = \Delta^{00,a} \Gamma_{rr'ss'}^{00,a} W_{rr'}^{(p)} W_{ss'}^{(p)\dagger} + \text{constant}. \quad (8)$$

Hermiticity and CP invariance of the effective Hamiltonian require $\Gamma_{rr'ss'}^{00,a} = \Gamma_{s'sr'r}^{00,a} = \Gamma_{rr'ss'}^{00,a*}$. CP invariance of the vacuum $|\Omega\rangle$ implies $\Delta^{00,a} = \Delta^{00,a*}$. Therefore, $V(W) = V(W^*)$. In general, we expect that for some range of sideways gauge boson masses, the minimum of $V(W)$ occurs for W complex, $W \neq W^*$. Then the minimum is degenerate, and CP is spontaneously broken.

Now we minimize $V(W)$ to determine the correct chiral perturbation $J'(W)$. $V(W)$ is stationary subject to the constraint in Eq. (5) if¹⁷

$$(M^{(p)} - M^{(p)\dagger})\Delta^p = i(v_H T_H^p + v_C T_C^p)1, \quad (9)$$

where $M_{rr'}^{(p)}\Delta^{(p)} = W_{rm}^{(p)}(\partial V/\partial W_{r'm}^{(p)}) = \Lambda^{(00,a)}\Gamma_{r'mss'}^{(00,a)}W_{rm}^{(p)}W_{ss'}^{(p)\dagger}$. Here v_H and v_C are Lagrange multipliers. (Taking the trace of both sides of Eq. (9) and summing on p we obtain a relation between v_H and v_C . If G_S is simple, then $v_H = -v_C$.)

If we include contributions from higher dimension operators to $V(W)$, the condition for an extremum of $V(W)$ is still of the form given in Eq. (9), but with a modified matrix $M^{(p)}$. For quarks, this matrix $M^{(q)}$ differs from the "current-algebra" quark mass matrix by approximately one part in 10^9 . This relation holds because the mass $m_C \sim 300$ MeV at which color becomes strong is small compared to the mass $m_H \sim 1$ TeV at which hypercolor becomes strong.

Physics below $m_H \sim 1$ TeV can be described by an effective theory involving only quarks, leptons, gluons, electroweak bosons, and pseudo-Goldstone bosons. This effective Lagrangian is obtained by integrating out hypergluons and hyperfermions, as well as massive sideways bosons. We must also sum up all hard (order m_H) gluon exchanges. A series of G_C -invariant operators is generated. The operators which contribute to the effective potential are

$A_{r'r'}\bar{q}_{Lr}q_{Rr'}$, $B_{r'r'}\bar{q}_{Lr}\sigma_{\mu\nu}\lambda^a q_{Rr'}G^{\mu\nu a}$, $C_{rr'ss'}\bar{q}_r q_r \bar{q}_s q_s$, and higher dimension operators. (Some operators can be eliminated by the equations of motion.) The coefficients A , B , C depend on the hyperfermion $W^{(p)}$. A is the quark mass matrix; it contributes $2\Delta^q \text{Re Tr}(W^{(q)}A)$ to $V(W)$. Relative to this term, the contribution from four-quark operators is suppressed by $(m_C/m_H)^8 \sim 10^{-10}$. The operator $\bar{q}_L\sigma_{\mu\nu}\lambda^a q_R G^{\mu\nu a}$ is generated by graphs involving one hard-gluon exchange;¹⁵ its contribution is suppressed by $[\alpha_C(m_H)/2\pi](m_C/m_H)^2 \sim 10^{-9}$. Hence, the part of $V(W)$ involving $W^{(q)}$ is $2\Delta^q \text{Re Tr } W^{(q)}\Lambda(1+O(10^{-9}))$ and therefore $M^{(q)} = W^{(q)}\Lambda(1+O(10^{-9}))$.

Equation (9) is solved by several sets of unitary matrices $W = \{W^{(p)}\}$. From all solutions, we choose the one (or ones) which minimizes $V(W)$. The W which minimizes the effective potential will satisfy one of three conditions:

- (i) $W = W^*$. In this case, CP is not spontaneously broken, and no CP violation occurs at all.
- (ii) $W \neq W^*$, $v_C \neq 0$. In this case, strong CP violation occurs. The effective Hamiltonian contains the CP-violating term $i(v_C/\Delta^q)\bar{q}i\gamma_5 q$; a current algebra calculation¹⁸ of the neutron electric dipole moment D_n yields $D_n \sim 4 \times 10^{-16} (v_C/m_U \Delta^q) \text{ e cm}$, where $m_U \sim 5$ MeV is the up quark mass. The natural scale of v_C/Δ^q is of order m_U so we expect $D_n = 10^{-15} \text{ e cm}$, which exceeds the experimental bound^{19,20} by a factor of 10^2 .
- (iii) $W \neq W^*$, $v_C = 0$. In this case, CP is spontaneously broken, but there is no large contribution to D_n . The coefficients B and C have phases of order one, but since $M^{(q)}$ is hermitian, $W^{(q)}A$ is hermitian up to corrections of order 10^{-9} . The antihermitian part of the quark mass matrix contributes of order 10^{-28} e cm to D_n . A similar contribution comes from the operator $\bar{q}_L\sigma_{\mu\nu}q_R F^{\mu\nu}$, where $F^{\mu\nu}$ is the photon field. Its coefficient is expected to

have a phase of order one and a magnitude $[\alpha_C(m_H)/2\pi](m_U/m_H^2)e \approx 10^{-28} \text{ e cm.}^{15,21}$ The Lagrange multiplier v_C vanishes if the minimum of the effective potential remains a stationary point when we remove the constraint that $U_A(1)$ rotations with a color anomaly are not allowed.

In any given model, the dynamics determines which possibility is realized. Only models of type (iii) can have CP violating interactions consistent with experiment.

In perturbation theory about the true chiral vacuum, the $W^{(\rho)}$'s appear explicitly in the weak and sideways currents. Minimizing the effective potential determines $W^{(\rho)} = W^{L(\rho)\dagger} W^{R(\rho)}$ (up to possible discrete degeneracies), but does not determine $W^{L(\rho)}$ and $W^{R(\rho)}$ separately. For quarks, we define $W^{L(q)}$ and $W^{R(q)}$ by demanding that the mass matrix $W^{(q)}_\Lambda$ be diagonal. The quark matrices $W^{L,R(q)}$ commute with electric charge; there are unitary matrices $W^{L,R}_u, W^{L,R}_d$ for up and down quarks. The Kobayashi-Maskawa²² mixing matrix is $W^{L\dagger}_u W^L_d$. If it contains phases which cannot be removed by redefining the phases of the quark fields, then the weak interactions violate CP.

Mixing matrices appear in the broken sideways currents also. Typically, phases in these matrices cannot be absorbed by redefining fields, because both left-handed and right-handed fermions transform nontrivially under G_S . Sideways gauge bosons can couple to flavor-changing neutral currents. If the operator $C\bar{S}d\bar{S}d$ occurs in the effective interaction, $\text{Im } C$ must be suppressed. If C is comparable to the coefficient of the term which generates the down quark mass,^{10,23} and $\arg C = O(1)$, this operator induces a CP-violating part of $K^0\bar{K}^0$ mixing which is larger by 10^4 than what is observed.

The effective Lagrangian below 1 TeV contains pseudo-Goldstone bosons whose couplings to fermions can be CP-violating. Exchange of these bosons can generate milliweak CP violation.²⁴

Spontaneous CP violation typically does not occur if the sideways interaction is vectorial.¹⁵ But the sideways interaction must be nonvectorial to generate realistic quark masses.^{15,23} A toy model in which spontaneous CP violation does occur has $G_S = \text{SU}(9) \times \text{SU}(9)$ and $G_H = \text{SU}(3)$. $\text{SU}(3)$ is embedded in the diagonal $\text{SU}(9)$ such that the 9 transforms as $3+3+3$. The fermions are in the representation $(1,9)_L + (9,1)_R$. (Anomalies can be canceled without changing the results.) The flavor group is $G_f = \text{SU}(3) \times \text{SU}(3)$. If $V(W)$ is calculated to lowest order in g_S^2 , it is minimized by a complex matrix for a natural range of sideways gauge boson masses. The broken sideways interactions violate CP, but the Lagrange multiplier v_H vanishes.

We have proposed a natural mechanism for generating CP violation in a gauge theory without elementary scalar fields or fermion bare masses. Although the strong CP problem is not automatically solved, strong CP violation may be avoided in particular models. However, CP violation in the broken sideways interactions is unavoidable, if weak CP violation occurs at all. Hence, if the electric dipole moment of the neutron is not discovered soon,²⁰ the general approach to CP violation which we advocate in this paper will become untenable.

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References

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1. S. Weinberg, Phys. Rev. D 13, 974 (1976).
2. S. Dimopoulos and L. Susskind, Phys. Rev. D 20, 2619 (1979); S. Weinberg, Phys. Rev. D 19, 1277 (1978).
3. E. Eichten and K. Lane, Phys. Lett. 90B, 125 (1980).
4. S. Dimopoulos and L. Susskind, Nucl. Phys. B155, 237 (1979); S. Weinberg (unpublished).
5. R. Dashen, Phys. Rev. D 3, 1879 (1971).
6. G. 't Hooft, Phys. Rev. Lett. 37, 8 (1976); R. Jackiw and C. Rebbi, Phys. Rev. Lett. 37, 172 (1976); C. Callan, R. Dashen, and D. Gross, Phys. Lett. 63B, 334 (1976).
7. R. Peccei and H. Quinn, Phys. Rev. Lett. 38, 1440 (1977); S. Weinberg, Phys. Rev. Lett. 40, 223 (1978); F. Wilczek, Phys. Rev. Lett. 40, 279 (1978).
8. Here we disagree with previous proposals for solving the strong CP problem within these theories; H. Pagels, Phys. Lett. 87B, 222 (1979), and Ref. 11.
9. M. Peskin (unpublished).
10. S. Dimopoulos, S. Raby and L. Susskind, "Tumbling Gauge Theories," Stanford Preprint ITP-653 (1979).
11. S. Dimopoulos and L. Susskind, "A Technicolored Solution to the Strong CP Problem," Columbia Preprint 79-0196 (1979).
12. The sideways interactions are isolated from the CP-violating interactions which generate the baryon excess in the early universe. See M. Yoshimura, Phys. Rev. Lett. 41, 281 (1978), 42, 746(e) (1979).

13. A model in which CP is spontaneously broken at 100 TeV is considered by B. Holdom, "A Realistic model with Dynamically Broken Symmetries," Harvard Preprint HUTP-80/A010.
14. This possibility was first suggested by Dashen, Ref. 5. See also T. D. Lee, Phys. Rev. D 8, 1226 (1973).
15. E. Eichten, K. Lane, and J. Preskill, in preparation.
16. If $n_p = 2$ or 4, other forms occur, but our conclusions are not altered.
17. This generalizes a theorem of J. Nuyts, Phys. Rev. Lett. 26, 1604 (1971).
18. V. Baluni, Phys. Rev. D 19, 2227 (1979); R. Crewther, P. Di Vecchia, G. Veneziano, and E. Witten, Phys. Lett. 88B, 123 (1979).
19. W. B. Dress et al., Phys. Rev. D 15, 9 (1977); I. S. Altarev et al., Leningrad Nucl. Phys. Inst. Preprint 430, 1 (1978).
20. N. F. Ramsey, Phys. Rpt. 43C, 409 (1978).
21. Rare processes such as $\mu \rightarrow e\gamma$, $\mu \rightarrow e$ conversion, and $K_L \rightarrow \mu e$ are also induced by operators appearing in the effective Lagrangian at rates comparable to present experimental upper limits.
22. M. Kobayashi and K. Maskawa, Prog. Theor. Phys. 49, 652 (1973).
23. V. Baluni, "Quark Mass Spectrum and CP Nonconservation," IAS Preprint in preparation.
24. S. Weinberg, Phys. Rev. Lett. 37, 657 (1976). See also Lee, Ref. 14.