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Electric Dipole Moment of the Neutron in a Left-Right-Symmetric Theory of *CP* Nonconservation

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The electric dipole moment of the neutron (μ_n^e) is calculated in the left-right-symmetric theory of Mohapatra and Pati and related to CP-nonconserving parameters of the neutral kaon system, yielding the bound $\mu_n^e \gtrsim 10^{-22}|\eta_{+-}-\eta_{00}|$ $e\cdot$cm. Recent measurements of $K_0 \to 2\pi$ amplitudes indicate that μ_n^e may well be of order 10^{-25} $e\cdot$cm which is several orders of magnitude larger than the prediction of the Kobayashi-Maskawa model and only a few times smaller than the current experimental limit μ_n^e(expt) $< 6\times 10^{-25} e\cdot$cm.

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The observed asymmetry between left- and right-handed currents in weak interactions has often been speculated to be a low-energy phenomenon.[1] Models of electroweak interactions have been proposed in which the unbroken Lagrangian possesses a left-right symmetry (LRS) based on the group $SU(2)_L \otimes SU(2)_R \otimes U(1)$ and the parity nonconservation seen at ordinary energies is implemented by introducing different scales of symmetry breaking for the left- and right-handed gauges.[1] The observed magnitude of CP nonconservation has been readily accommodated in such a model by linking it to the suppressed right-handed currents.[2] Unfortunately, the confirmation of these ideas may have to await the construction of accelerators many times more powerful than the ones now available unless one can think of effects that for some reason are suppressed if the underlying interactions are purely left-handed, i.e., $V-A$. One such effect is the electric dipole moment (EDM) of the neutron.

In a pure $V-A$ theory with CP nonconservation arising from phases of the quark mixing matrix, e.g., the Kobayashi-Maskawa model,[3] there exists a phase cancellation in the one-loop calculation of the EDM of an elementary fermion, leading to zero contribution.[4] In fact, Shabalin has shown that in the Kobayashi-Maskawa model the sum of all two-loop graphs contributing to the EDM of a quark vanishes as well.[5] Consequently, in the Kobayashi-Maskawa model the EDM of the neutron (μ_n^e) is expected to be $\lesssim 10^{-30}\,e\cdot\text{cm}$; that is, at least five orders of magnitude smaller than the present experimental limit[6]:

$$\mu_n^e(\text{expt}) < 6\times 10^{-25}\,e\cdot\text{cm}. \qquad (1)$$

In contrast, the EDM resulting from one loop need not vanish in a theory with LRS and might therefore be expected to be larger, by several orders of magnitude, than in a pure $V-A$ theory. In recent years significant technological advances have been made in the storage of ultracold neutrons. Several experiments now in progress hope to reduce the existing limit for μ_n^e by two to three orders of magnitude in the next year or two.[7] It therefore appears useful to undertake a calculation of μ_n^e in a left-right-symmetric gauge theory.

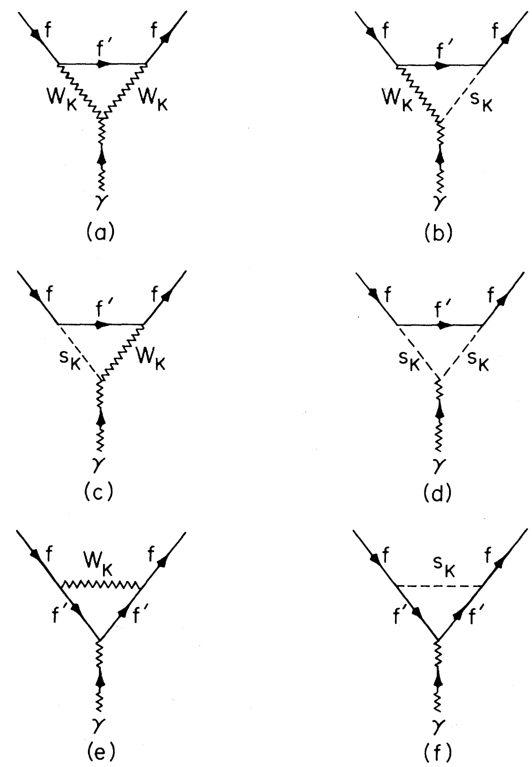

FIG. 1. Diagrams which contribute to the electric dipole moment of an elementary fermion f in the 't Hooft-Feynman gauge. W_k ($k=1,2$) are charged vector bosons; S_k are unphysical scalars.

For definiteness, we consider the four-quark left-right-symmetric model[2,8] of Mohapatra and Pati with appropriate modifications.[9] CP nonconservation can arise in such a theory through both complex fermion-scalar coupling and complex scalar vacuum expectation values. After diagonalization of the fermion mass matrices the charge-current eigenstates (d^0, s^0) are related to the mass eigenstates (d, s) by

$$\begin{pmatrix} d^0 \\ s^0 \end{pmatrix}_{L,R} = \begin{pmatrix} \cos\theta_{L,R} & \sin\theta_{L,R}e^{i\delta_{L,R}} \\ -\sin\theta_{L,R}e^{-i\delta_{L,R}} & \cos\theta_{L,R} \end{pmatrix} \begin{pmatrix} d \\ s \end{pmatrix}_{L,R}, \qquad (2)$$

where $\theta_{L,R}$ and $\delta_{L,R}$ are real parameters. The complex vacuum expectation values also give rise to a complex gauge field mass matrix. The left- and right-coupling gauge fields, $W_{L,R}$, are then related to their mass eigenstates, $W_{1,2}$, by

$$W_1 = W_L\cos\zeta - W_R\sin\zeta\, e^{i\lambda}, \quad W_2 = W_L\sin\zeta\, e^{-i\lambda} + W_R\cos\zeta, \qquad (3)$$

where ζ and λ are real angles. It is this L-R mixing through ζ that avoids the phase cancellation mentioned above, allowing a nonvanishing one-loop contribution to the EDM. The phase λ in (3) represents a second source of CP nonconservation independent of the quark sector.

In the 't Hooft–Feynman gauge, the leading contribution to the EDM of an elementary fermion (f_j) arises from the six Feynman graphs shown in Fig. 1.[10] For generality we write the flavor-changing part of the interaction as

$$L_I = -\sum_{ijk}\bar{f}_i\gamma_\mu(a_{ij}^k + b_{ij}^k\gamma_5)f_j W_k^\mu - \sum_{ijk}\bar{f}_i(c_{ij}^k + d_{ij}^k\gamma_5)f_j S_k + \text{H.c.}, \qquad (4)$$

where gauge invariance constrains the couplings of the unphysical scalars S_k to be

$$c_{ij}^k = (m_i - m_j)a_{ij}^k/M_k; \quad d_{ij}^k = (m_i + m_j)b_{ij}^k/M_k. \qquad (5)$$

Here m_i, m_j are the masses of the fermions f_i, f_j, and M_k are the masses of the charged spin-one fields W_k, $k = 1,2$.

If the external fermion masses can be neglected relative to the internal ones we obtain for the EDM of f_j the result[11]

$$\mu_j^e = \sum_{kl} \frac{e m_l}{8\pi^2 M_k^2} \operatorname{Im}(a_{jl}^k b_{jl}^{k*}) \frac{1}{(1-r)^2}\left[Q_j\left(2 - \tfrac{11}{2}r + \tfrac{1}{2}r^2 - 3r^2 \frac{\ln r}{(1-r)}\right) - Q_l(4 - 5r + r^2 + 3\ln r)\right], \tag{6}$$

where $r = m_l^2/M_k^2$, m_l is the mass of the internal fermion line, and Q_i is the electric charge of f_i in units of $|e|$.

For the neutron EDM we retain only the lowest-order terms in r. Expressing the coefficients a_{jl}^k and b_{jl}^k in terms of the parameters of the gauge model we obtain the quark EDM's:

$$\mu_{u,d}^e = \frac{e g_L g_R}{96\pi^2} \sin 2\zeta \left(\frac{1}{M_2^2} - \frac{1}{M_1^2}\right) A (m_{d,u} \cos\theta_L \cos\theta_R \sin\lambda + m_{s,c} \sin\theta_L \sin\theta_R \sin B). \tag{7}$$

where $\delta \equiv \delta_L - \delta_R$, and $A = 4$, $B = \lambda + \delta$ for the u quark and $A = 5$, $B = \lambda - \delta$ for the d quark. Adopting the nonrelativistic bound-state model for the neutron we have[12]

$$\mu_n = (4\mu_d - \mu_u)/3. \tag{8}$$

Thus

$$\mu_n^e = \frac{e g_L g_R}{72\pi^2} \sin 2\zeta \left(\frac{1}{M_1^2} - \frac{1}{M_2^2}\right) \times \{(5m_u - m_d) \cos\theta_L \cos\theta_R \sin\lambda + [5m_c \sin(\lambda - \delta) - m_s \sin(\lambda + \delta)] \sin\theta_L \sin\theta_R\}. \tag{9}$$

With a proper choice of the scalar potential we may impose a LRS on the Lagrangian such that $g_L = g_R$ and $\theta_L = \theta_R$.[13] If we also assume that $M_1^2 \ll M_2^2$ we find

$$\mu_n^e \simeq (3.6\times 10^{-21}\, e\cdot\text{cm GeV}^{-1}) \tan\zeta \{(5m_u - m_d)\cos^2\theta_C \sin\lambda + [5m_c \sin(\lambda-\delta) - m_s \sin(\lambda+\delta)]\sin^2\theta_C\}, \tag{10}$$

$$\simeq (10^{-21}\, e\cdot\text{cm}) \tan\zeta [4 \sin\lambda + 1.4 \sin(\lambda - \delta) - 0.1 \sin(\lambda + \delta)], \tag{11}$$

where θ_C is the Cabibbo angle and we have used $m_u = m_d = 300$ MeV, $m_s = 500$ MeV, and $m_c = 1.5$ GeV.

The magnitude of the EDM is crucially dependent on the left-right mixing angle ζ whose value is unknown. This uncertainty is eliminated by utilizing the CP parameters of the neutral kaon system. If $\zeta = 0$ the above model is "isoconjugate" and $\eta_{+-} = \eta_{00}$.[2,14] We thus calculate that

$$|\eta_{+-} - \eta_{00}| \simeq \tan\zeta [\sin\lambda + \sin(\lambda + \delta)] |\gamma|, \tag{12}$$

where γ is a parameter determined by strong interactions and is presumably of order 1. Thus

$$\mu_n^e \simeq (10^{-21}\, e\cdot\text{cm}) |\eta_{+-} - \eta_{00}| \frac{[4\sin\lambda + 1.4\sin(\lambda-\delta) - 0.1\sin(\lambda+\delta)]}{[\sin\lambda + \sin(\lambda+\delta)]|\gamma|}. \tag{13}$$

There are two cases in which this expression becomes particularly simple: (i) "manifest" LRS,[15] wherein the quark mass matrix is Hermitian and $\delta = 0$, and (ii) real left-right mixing, wherein the W mass matrix is real and $\lambda = 0$. We then have (i)

$$|\mu_n^e| \simeq (2.7/|\gamma|)\times(10^{-21}\, e\cdot\text{cm})|\eta_{+-} - \eta_{00}|; \quad \delta = 0, \tag{14}$$

and (ii)

$$|\mu_n^e| \simeq (1.5/|\gamma|)\times(10^{-21}\, e\cdot\text{cm})|\eta_{+-} - \eta_{00}|; \quad \lambda = 0. \tag{15}$$

For a reasonable upper bound for $|\gamma|$ we take $|\gamma| < 10$, giving the relation[16,17]

$$|\mu_n^e| \gtrsim 10^{-22} |\eta_{+-} - \eta_{00}|\; e\cdot\text{cm}. \tag{16}$$

Recent measurements by Christenson *et al.*[18,19] give the result

$$\left|\frac{\eta_{+-} - \eta_{00}}{\eta_{+-}}\right| \simeq 0.22 \pm 0.11, \tag{17}$$

indicating that the neutron EDM may be of order $10^{-25}\ e\cdot$cm, although the large uncertainties in the experimental result (17) do not allow one to make definitive predictions. We note that this measurement along with Eq. (12) leads us to place a useful lower bound on the important left-right mixing angle ζ:

$$|\zeta| \gtrsim 10^{-4}/|\gamma| \gtrsim 10^{-5}. \qquad (18)$$

This is complementary to the upper bound $|\zeta| < 0.06$ deduced in Ref. 15.

In conclusion, we find that in the left-right-symmetric theory of Mohapatra and Pati the EDM of the neutron may well be of order $10^{-25}\ e\cdot$cm. This is barely an order of magnitude less than the existing experimental limit ($<6\times 10^{-25}\ e\cdot$cm) and is several ($>5$) orders of magnitude larger than that expected in the Kobayashi-Maskawa model. In addition, the LRS theory yields a relatively unambiguous link between the EDM of the neutron and the CP-nonconserving parameters of the neutral kaon system; in particular we find $\mu_n^e \gtrsim 10^{-22}|\eta_{+-}-\eta_{00}|\ e\cdot$cm. Projected experimental measurements[7,20] of μ_n^e and η's in the foreseeable future could thus provide decisive tests of the left-right-symmetric model of electroweak interactions.

This work is an offshoot of an unpublished work[21] of Rabi Mohapatra and Jogesh Pati with one of the authors (A.S.). Helpful discussions with Myron Bander are greatly appreciated. This work was supported in part by the National Science Foundation under Grants No. PHY 78-21502 and No. PHY 79-10262.

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[9]In particular, CP nonconservation through the phase λ in W_L-W_R mixing [see Eq. (3) below] has not been included before. Note that irrespective of the choice of the Higgs sector no such CP nonconserving phase can arise through mixing among the neutral gauge fields (Z's). There will be a real mixing among the Z's but that does not contribute to the EDM or to the neutral kaon decays that are relevant to this work.

[10]We ignore the possible contribution to the EDM from physical Higgs fields. Since our estimate for the EDM comes out to be so close to the current experimental limit it seems reasonable to assume that that contribution is likely to be small to leave our estimate intact unless, of course, there are some highly unnatural cancellations.

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Evidence for Hard-Gluon Bremsstrahlung in a Deep-Inelastic Neutrino Scattering Experiment

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Data from νN interactions in the Fermilab 15-ft bubble chamber show that the mean transverse momentum (p_T) of forward hadrons in the hadron c.m. system exceeds that of backward hadrons for $W^2 > 100$ GeV^2. Events with high π_F, a measure of forward p_T, tend to have planar hadron systems and show a three-jet structure in their angular energy flow; their number exceeds that expected in two-jet models. These observations are consistent with quantum chromodynamic predictions of hard-gluon bremsstrahlung.

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Gluon bremsstrahlung can account for the results of several recent experiments. Groups at PETRA[1] observe evidence for three-jet events whose origin may be $e^+e^- \to q\bar{q}g$ (q denotes quark; g denotes gluon). The analogous process in weak (electromagnetic) deep-inelastic scattering is $W^{\pm}q(\gamma^*q) \to qg$ ($W^{\pm}$ are the intermediate vector bosons, γ^* is the exchanged virtual photon), where the remaining diquark in the target nucleon is a spectator to the interaction. Evidence that this process contributes substantially to the total deep-inelastic cross section at high energy comes from both neutrino[2,3] and muon[4] experiments. We report here the observation, in a neutrino experiment, of three-jet events whose properties are consistent with those expected for hard-gluon bremsstrahlung ($W^{\pm}q \to qg$) events.

Neutrinos (ν_μ) from the Fermilab quadrupole-triplet beam interacted in the 15-ft bubble chamber filled with a 47-at.% Ne-H mixture. Muons from charged-current interactions were identified by the two-plane external muon identifier.[5] A

