Bounds on Microscopic Physics from P and T
Violation in Atoms and Molecules

Willy Fischler, Sonia Paban and Scott Thomas
Theory Group
Department of Physics
University of Texas
Austin, Texas 78712

Abstract

Atomic and molecular electric dipole moments are calculated within the minimal supersymmetric standard model. Present experiments already provide strong bounds on the combination of phases responsible for the dipole moments of the neutron and closed shell atoms. For a supersymmetry breaking scale of 100 GeV, these phases must be smaller than $\sim 10^{-2}$.

*Research supported in part by the Robert A. Welch Foundation and NSF Grant PHY 9009850.

†paban@utaphy.bitnet
‡thomas@utaphy.bitnet
The experimental bounds on electric dipole moments (edm’s) of the neutron, atoms, and molecules, are reaching a level of precision that makes them useful probes of P and T violation that may originate at scales beyond the standard model. Aside from the QCD vacuum angle, the only source of T (or CP) violation in the standard model is the phase in the quark mass matrix. This generates edm’s orders of magnitude beyond experimental reach [1]. Extensions of the standard model however generally have CP violating phases that produce edm’s of the order of, or larger than, present experimental bounds. A finite $\bar{\theta}_{QCD}$ that saturates the bound from the neutron edm [2], would also produce atomic edm’s of the same order. In this paper the main sources for atomic and molecular edm’s in the minimal supersymmetric extension of the standard model will be identified.

In order to identify the microscopic sources of P and T violation responsible for the edm’s, it is instructive to present the analysis starting from the largest relevant scale, namely the atomic scale. The next step is then to proceed down to the nuclear scale. This will enable finally a discussion of the origin of P and T violation at the subnuclear level.

At the atomic scale an atom or molecule is a composite system of electrons and nuclei. These constituents contribute to the edm via T and P odd electromagnetic moments and local non-electromagnetic interactions. Some of the electromagnetic moments are suppressed. In the nonrelativistic limit, the contributions from the electron and nuclear edm’s vanish. This is known in the literature as Schiff’s theorem [3]. The electron edm can contribute though in heavy atoms where the electrons are relativistic [4, 5]. But in closed shell atoms and molecules with paired electron spins, the electron edm contributes only through hyperfine interactions with the nucleus, and is therefore suppressed [6]. Schiff’s theorem does not apply to higher moments such as the nuclear magnetic quadrupole moment. In addition there is another T and P odd moment known as Schiff’s moment, which is proportional to the offset of electric charge and dipole distributions of the nucleus. It leads to a local electromagnetic coupling of the electron to the nucleus [3].
At the nuclear scale all of the contributions discussed above are contained in the following Hamiltonian

\[ H = \frac{1}{2} d e \bar{e} \sigma_{\mu \nu} i \gamma_5 e F^{\mu \nu} + \frac{1}{2} d_N \bar{N} \sigma_{\mu \nu} i \gamma_5 N F^{\mu \nu} + C_{en}^{FS,PS} \frac{G_F}{\sqrt{2}} (\bar{e} i \gamma_5 e)(\bar{N} N) \]

\[ + C_{en}^{T,P T} \frac{G_F}{\sqrt{2}} (\bar{e} \sigma_{\mu \nu} e)(\bar{N} \sigma_{\mu \nu} i \gamma_5 N) \]

\[ + Q_{en,V,P V}^{en} (\bar{e} \gamma_\mu e)(\bar{N} \bar{\gamma}_5 N) + Q_{PV,V}^{en} (\bar{e} \bar{\gamma}_\mu \gamma_5 e)(\bar{N} \gamma_\mu N) \]

\[ + C_{en}^{S,P S} G_F \sqrt{2} (\bar{N} i \gamma_5 N) + Q_{PV,P V}^{mn} (\bar{N} \gamma_\mu N)(\bar{N} \bar{\gamma}_5 N) \]

(1)

where \( N = (p, n) \) and isospin violation is ignored. The first and second terms are the electron and nucleon edm’s respectively. The remaining terms are local interactions among the electrons and nucleons. The nucleon edm and nucleon-nucleon couplings contribute to the Schiff and magnetic quadrupole moments [7]. The electron-nucleon couplings of course contribute to the electron-nucleus coupling [8]. Calculations of the contributions of the terms in \( H \) to the edm’s considered below may be found in Ref.s [5,7-11] and are summarized in Table 1. The contributions from \( Q_{en,V,P V}^{en} \), \( Q_{PV,V}^{en} \), and \( Q_{PV,P V}^{en} \) have not been considered in these references although \( Q_{en,V,P V}^{en} \) does contribute directly to the Schiff moment. These derivative operators have the same nonrelativistic limit as some of the nonderivative operators. This allows a rough estimate of the corresponding contributions to the edm’s.

In this paper the microscopic origin of the operators in (1) will be calculated within the context of the minimal supersymmetric extension of the standard model [12]. In addition to the KM like phases, this model has other CP violating phases in the superpotential and soft supersymmetry breaking terms. The superpotential contributes a phase from the term \(|\mu|e^{i\phi_\mu} \hat{H}_1 \hat{H}_2\). The phases in the soft breaking terms arise from the gaugino masses, \(-\frac{1}{2}|\bar{m}_a|e^{i\phi_a}\lambda_{\bar{a}}\lambda_{\bar{a}} + \text{h.c.} \), and the mass parameters coupling left and right handed squarks (sleptons), \(-f_L m_f |A_f|e^{i\phi_{A_f}} f_R + \text{h.c.} \), where \( m_f \) is the corresponding quark (lepton) mass and \( A_f \) is defined by

\[ m_f A_f = m_f \left( \frac{A_F}{h_f} + \mu \frac{v_1}{v_2} \right) \]

(2)
$A_F$ is the scalar trilinear soft breaking coupling and $h_f$ the Yukawa coupling. The importance of these new phases is that P and T odd operators can be generated irrespective of generational mixing.

The supersymmetric origin of the terms in the Hamiltonian (1) will now be considered.

1. The electron edm is a chirally violating operator of effective dimension 6. It arises at the one loop level from electroweak gaugino exchange as shown in Fig. 1. The sum of all such diagrams contains several unknown masses and phases [13]. In order to display the order of magnitude, only the photino contribution will be considered [14],

$$d_e = -e \frac{\alpha}{24\pi} \frac{m_e |A_e|}{m_\gamma^3} \sin(\phi_{A_e} - \phi_\gamma) f(x) + \cdots$$

(3)

where $x = m_e^2/m_\gamma^2$, $f(x)$ is a loop function of order unity, $f(1) = 1$ [14], and $+ \cdots$ represents the contributions from the other electroweak gauginos. Numerically, for $m_e = m_\gamma$ Eq. (3) gives [14]

$$d_e \simeq -1 \times 10^{-25} \frac{(100 \text{ GeV})^2}{m_\gamma^3 |A_e|} \sin(\phi_{A_e} - \phi_\gamma) e \text{ cm } + \cdots$$

(4)

2. The lowest dimension operators which potentially contribute to the nucleon edm are the light quark electric and chromoelectric dipole moments [15, 16], and Weinberg’s 3-gluon operator [17], [18]

$$\mathcal{O}_1 = \frac{1}{2} \bar{q} \sigma_{\mu \nu} i \gamma_5 q F^{\mu \nu}$$

$$\mathcal{O}_2 = \frac{1}{2} \bar{q} \sigma_{\mu \nu} i \gamma_5 T_a q G_a^{\mu \nu}$$

$$\mathcal{O}_3 = \frac{1}{2} f_{abc} \tilde{G}_{a \mu \nu} G_{b}^{\mu \rho} G_{c \rho}^{\nu}$$

where $\tilde{G}_{a \mu \nu} = \frac{1}{2} \epsilon_{\mu \nu \rho \sigma} G_a^{\rho \sigma}$ and $T_a$ are the SU(3)$_C$ generators. As for the electron edm the quark dipole moments are effectively dimension-6. The operators $\mathcal{O}_1$ and $\mathcal{O}_2$ are generated at the one loop level through gluino and electroweak gaugino exchange [14]. The sum of all such diagrams is again a function of several masses.

\[\text{footnote}{\text{1If the hidden sector is Polonyi and the electroweak gaugino masses vanish at tree level, the supersymmetric phases can be rotated away.}}\]
and phases. To establish the order of magnitude only the gluino and photino exchange will be considered

$$C_1(\mu) = \frac{e Q_q}{24\pi} m_q |A_q| \left( \frac{\alpha_s}{m_\lambda^2} \sin(\phi_{A_q} - \phi_\lambda) \frac{4}{3} f(y) + \frac{\alpha Q^2_q}{m_\gamma^2} \sin(\phi_{A_q} - \phi_\gamma) f(z) + \cdots \right) \zeta_1$$

$$C_2(\mu) = -\frac{g_s}{24\pi} m_q |A_q| \left( \frac{\alpha_s}{m_\lambda^2} \sin(\phi_{A_q} - \phi_\lambda) (h(y) + \frac{1}{6} f(y)) - \frac{\alpha Q^2_q}{m_\gamma^2} \sin(\phi_{A_q} - \phi_\gamma) f(z) + \cdots \right) \zeta_2$$

where $Q_q$ is the quark charge, $y = m_\gamma^2/m_\lambda^2$, $z = m_\gamma^2/m_\gamma^2$, $f$ is the same function that appears in Eq. (3), and $h(1) = 1$. $\zeta_i$ are renormalization group corrections for the evolution of the operators from the supersymmetric scale, $M$, to the hadronic scale, $\mu$. Using the known anomalous dimensions with $\alpha_s(M) = .1$ and $\alpha_s(\mu) = 4\pi/6$ gives $\zeta_1 \simeq 1.6$, and $\zeta_2 \simeq 3.6$. Note that the photino contribution is down by $\sim \alpha^2\alpha_s^{-1}$ compared with the gluino. Barring any cancellation among the phases, and if all the mass parameters are of the same order, the gluino contribution will dominate the quark dipole moments.

The 3-gluon operator is generated at the two loop level. The largest contributions come from integrating out the chromoelectric dipole moments of quarks with mass $m_Q > \Lambda_{QCD}$, $C_3(\mu) = \frac{\alpha_s(m_Q) C_2(m_Q)}{8\pi m_Q} \zeta_3$.

Keeping only the charm quark as the dominant contribution gives $\zeta_3 \simeq .38$.

An estimate of the edm requires an evaluation of the matrix elements of these operators on the nucleon. There is however no systematic approximation scheme to reliably calculate hadronic matrix elements of this type. Here, some empirical rules known as “naive dimensional analysis” that keep track of factors of $4\pi$ and mass scales will be used. This is the most uncertain aspect of the entire analysis. Even so, there is no physical reason to expect a substantial suppression or enhancement as compared to estimates given by these rules. Using these rules, the nucleon edm is

$$d_N \simeq C_1(\mu) + \frac{e}{4\pi} C_2(\mu) + \frac{e}{4\pi} \Lambda_\chi C_3(\mu)$$
where $\Lambda_\chi \simeq 4\pi f_\pi$ is the chiral symmetry breaking scale. Numerically, keeping only the gluino contributions with $m_{\tilde{q}} = m_{\tilde{\lambda}} = m_{\tilde{\gamma}} \equiv \tilde{m}$, and $A_u = A_d = A_c \equiv A_q$, Eq. (8) then gives

$$|d_N| \simeq (2.2 + 1.1 + .1) \frac{(100 \text{ GeV})^2}{\tilde{m}^3/|A_q|} \sin(\phi_{A_q} - \phi_{\tilde{\lambda}}) 10^{-23} \text{ e cm}$$

(9)

where the terms on the right side come from $O_1$, $O_2$, and $O_3$ respectively.

3. The nucleon-nucleon couplings are just the T odd component of the nuclear force. Following Weinberg’s analysis of nuclear forces [21], it is useful to think of these as arising from the T odd exchange of low lying mesons and direct contact interactions. For illustrative purposes only the exchange of the lightest meson, the pion, will be considered; and an estimate of the T odd pion-nucleon coupling, $\bar{g}_{\pi NN} \bar{N}\pi N$, arising from the microscopic physics will be made. The leading contribution will come from the light quark chromoelectric dipole moment. Using “naive dimensional analysis” the coupling is

$$\bar{g}_{\pi NN} \simeq g_{\pi NN} \frac{\Lambda_\chi}{4\pi} C_2(\mu)$$

(10)

where $g_{\pi NN}$ is the usual pseudoscalar pion-nucleon coupling constant. Integrating out the pion, the contribution to the nonderivative nucleon-nucleon coupling is

$$C_{S,PS}^{nn} \sim \bar{g}_{\pi NN} g_{\pi NN} \frac{1}{m_{\pi}^2} \sqrt{2} G_F$$

(11)

Numerically, keeping only the gluino contribution with $m_{\tilde{\lambda}} = m_{\tilde{q}}$, Eqs. (6), (10), and (11) give

$$C_{S,PS}^{nn} \simeq .7 \frac{(100 \text{ GeV})^2}{m_{\tilde{q}}^3/|A_q|} \sin(\phi_{A_q} - \phi_{\tilde{\lambda}})$$

(12)

The exchange of heavier mesons will also contribute to this coupling and the derivative nucleon-nucleon coupling. This is not expected to substantially change the estimates given below based solely on the nonderivative coupling from pion exchange.

4. The electron-nucleon couplings in (1) arise from two classes of operators. The first operators involve an electromagnetic infrared enhancement and are suppressed
by two powers of a heavy mass. The second class of operators are local at the atomic scale and suppressed by four powers of a heavy mass.

The first class of operators arise from the effective nucleon-photon couplings $\bar{N} N F_{\mu
u} \tilde{F}^{\mu\nu}$ and $\bar{N} i\gamma_5 NF_{\mu\nu} F^{\mu\nu}$. The two photons couple to electrons to produce an effective electron-nucleon interaction. Integrating out the photons gives an infrared enhancement, cutoff by the electron mass. One way in which such an effective interaction arises is through the pion pole as shown in Fig. 2, where the CP violation occurs in the pion-nucleon coupling. This contributes to the electron-nucleon pseudoscalar-scalar coupling

$$C_{PS,S}^{en} \sim g_{\pi ee} \bar{g}_{\pi NN} \frac{1}{m_e^2} \frac{\sqrt{2}}{G_F}$$

where $\bar{g}_{\pi NN}$ is given in (10), $g_{\pi ee}$ is the effective pion-electron coupling computed in Ref. [22]

$$g_{\pi ee} \sim -\frac{3\alpha^2 m_e}{2\pi^2 f_\pi \ln \left(\frac{\Lambda}{m_e}\right)}$$

and the approximation that the electron is on shell has been used. Note that the effective electron-nucleon interaction is suppressed by only two powers of a heavy mass (from the light quark chromoelectric dipole moment). A similar diagram with the CP violation in the vertex $\pi^0 F_{\mu\nu} F^{\mu\nu}$, contributes to the scalar-pseudoscalar coupling. The edms in Table 1. are however somewhat less sensitive to this coupling.

The nucleon-photon couplings can also arise from the following microscopic operators which are not suppressed by a light quark mass

$$O_4 = G_{a\mu\nu} G^\mu\nu_a F_{\rho\sigma} \tilde{F}^{\rho\sigma}$$
$$O_5 = G_{a\mu\nu} \tilde{G}^\mu\nu_a F_{\rho\sigma} F^{\rho\sigma}$$

These operators arise at the two loop level. The dominant contribution to $O_4$ comes from integrating out the electric dipole moment of a quark with mass $m_Q > \Lambda_{QCD}$, as shown in Fig. 3. This involves an infrared enhancement, cutoff of by the quark mass. The coefficient is therefore suppressed by only two powers of a
heavy mass. Similar operators have been considered in Ref. [23]. Here,

\[ C_4 \simeq \frac{e g_s^2 Q Q}{265\pi^2 m_Q^3} C_1(m_Q) \]  

(15)

Using “naive dimensional analysis” then

\[ C_{en}^{PS,S} \simeq -6 \frac{\alpha}{\pi} m_e \Lambda_\chi \ln \left( \frac{m_Q}{m_e} \right) \frac{\sqrt{2}}{G_F} C_4 \]  

(16)

A similar discussion applies to \( O_5 \) with the chromoelectric dipole moment of heavy quarks. All these operators scale like \((m_Q M)^{-2}\).

It should be noted that the infrared enhancement from the photon-electron loop comes from momenta \( \sim m_e \). Treating the resulting interactions as local is somewhat dubious for heavy atoms for which \( p_e \sim Z \alpha m_e \). This is not expected to significantly alter the estimates of the edm’s given below. Among the first class of operators, the estimate in Eq. (13) turns out to be numerically most important,

\[ C_{PS,S}^{en} \sim 2 \times 10^{-8} \frac{(100 \text{ GeV})^2}{m_\chi^3/|A_q|} \sin(\phi_{A_q} - \phi_\chi) \]  

(17)

The second class of local operators that lead to electron-nucleon couplings include

\[ O_6 = \bar{e} i \gamma_5 e \bar{q} q \quad O_7 = \bar{e} e i \gamma_5 q \]

\[ O_8 = \bar{e} \sigma_{\mu \nu} e \bar{q} \sigma^{\mu \nu} i \gamma_5 q \quad O_9 = \bar{e} i \gamma_5 e G_{a \mu \nu} G^{a \mu \nu} \]

\[ O_{10} = \bar{e} e G_{a \mu \nu} G_{a}^{\mu \nu} \quad O_{11} = \bar{e} \gamma_\mu e \bar{q} \gamma_\mu q \]

\[ O_{12} = \bar{e} \gamma_\mu \gamma_5 e \bar{q} \gamma^\mu q \]

Because of the chiral properties, all these operators are effectively dimension 8. The operators \( O_6, O_7, \) and \( O_8 \) are generated, after Fierz reordering, by box diagrams\(^\text{[f]}\) of the type shown in Fig. 4. Neglecting running corrections, the coefficients are then related by \( C_6 = C_7 = \frac{1}{2} C_8 \equiv C \). A typical contribution is of order

\[ C \sim \frac{\alpha}{4 \pi} G_F \frac{m_e m_q |\mu|}{\sin 2\beta} \frac{m_\gamma^2}{m_\gamma^2} \sin(\phi_\mu - \phi_\gamma) \]  

(18)

where \( \tan \beta = v_2/v_1 \). With all the supersymmetric mass parameters of order \( M \), this scale like \( G_F M^{-2} \).

\(^{2}\)Tree level Higgs exchange is CP conserving in the minimal supersymmetric standard model and will not contribute to the electron-nucleon couplings.
The chiral suppression of a light quark mass is avoided in the operators $O_9$ and $O_{10}$. These can arise from the box diagrams of Fig. 4 with the light quarks replaced by heavy quarks, $Q$. The heavy quarks are then integrated out as shown in Fig. 5. In the limit $M >> m_Q$,

$$C_9 = \sum_Q -\frac{2}{3} \frac{\alpha_s}{4\pi} \frac{C_6(m_Q)}{m_Q}$$

$$C_{10} = \sum_Q \frac{\alpha_s}{4\pi} \frac{C_7(m_Q)}{m_Q}$$

There are no dimension-8 operators without a light quark mass suppression that give a tensor-pseudotensor Lorentz structure. There is however an effective dimension-10 operator, $O_{13} = \frac{1}{3}\bar{e}\sigma_{\mu\nu}e_{idabc}\bar{G}_{ab}^{\sigma}G_{c}^{\rho\sigma}G_{\rho\mu}^{\nu}$, which is generated analogously to $O_9$ and $O_{10}$, but with at least three gluons on the heavy quark loop. With the rules of “naive dimensional analysis” this operator is not much suppressed compared with $O_9$ and $O_{10}$. It is therefore included in the estimate below. Using “naive dimensional analysis” the contributions of $O_6$ through $O_{10}$, and $O_{13}$, to the nonderivative electron-nucleon couplings are

$$C_{en,PS,S} = \frac{\sqrt{2}}{G_F} \left( C_6(\mu) - \sum_Q \frac{2}{3} \frac{\alpha_s}{4\pi} \frac{A_x}{m_Q} C_6(m_Q) \right)$$

$$C_{en,S,PS} = \frac{\sqrt{2}}{G_F} \left( C_7(\mu) - \sum_Q \frac{\alpha_s}{4\pi} \frac{A_x}{m_Q} C_7(m_Q) \right)$$

$$C_{en,T,PT} = \frac{\sqrt{2}}{G_F} \left( C_8(\mu) - \sum_Q \frac{1}{24} \left( \frac{g_s}{4\pi} \right)^3 \left( \frac{A_x}{m_Q} \right)^3 C_8(m_Q) \right)$$

(20)

where again running corrections have been neglected.

The operators $O_{11}$ and $O_{12}$ are generated by box diagrams of the type shown in Fig. 6a,

$$C_{11} \sim \frac{\alpha}{4\pi} e^2 Q_q^2 m_q |A_q| \frac{m_\gamma}{m_\gamma^2} \sin(\phi_\gamma - \phi_{A_q})$$

$$C_{12} \sim \frac{\alpha}{4\pi} e^2 Q_q^2 m_e |A_e| \frac{m_\gamma}{m_\gamma^2} \sin(\phi_\gamma - \phi_{A_e})$$

(21)
There are also contributions that result from a weak interaction between the electron (quark) and the weak edm of the quark (electron) as shown in Fig. 6b.

\[ C_{11} = \pm (1 - 4|Q_q|\sin^2\theta_w)(1 - 4\sin^2\theta_w) \frac{G_F C_1(\mu)}{\sqrt{2}} \frac{1}{e} \]

\[ C_{12} = \pm (1 - 4|Q_q|\sin^2\theta_w)\cot\theta_w \frac{G_F d_e^{(W)}}{\sqrt{2}} \frac{1}{e} \]

(22)

where \( \pm \) refers to up or down type quarks, \( d_e^{(W)} \) is the electron edm arising from wino exchange, and \( C_1(\mu) \) is given above. These operators scale like either \( G_F M^{-2} \) or \( M^{-4} \). Using “naive dimensional analysis” the contribution to the derivative electron-nucleon couplings in (1) are

\[ Q_{V,P,V}^{en} = C_{11} \]

\[ Q_{P,V,V}^{en} = C_{12} \]

(23)

All of the local operators turn out to be numerically somewhat less important than the estimate in Eq. (13).

The atomic and molecular edm’s can now be estimated as functions of the microscopic parameters. In order to identify the most important effects at the atomic scale, the contributions from the electron edm, nucleon edm, nonderivative nucleon-nucleon coupling, and the largest electron-nucleon coupling (i.e. from Eq. (13)) will be retained. Since an edm is proportional to spin, the leading contributions in open and closed shell atoms are different. For open shell atoms we consider \(^{133}\text{Cs} \) and \(^{205}\text{Tl} \) since good experimental bounds are available [24, 25]. Putting together the results cataloged above with the results of the atomic and nuclear calculations from Table 1.,

\[ d_{\text{Cs}} \simeq ( -1.2 \sin(\phi_{A_q} - \phi_{\gamma}) + 2 \times 10^{-3}\sin(\phi_{A_q} - \phi_{\lambda}) + 1 \times 10^{-1}\sin(\phi_{A_q} - \phi_{\bar{\lambda}}) \]

\[ + 1 \times 10^{-3}\sin(\phi_{A_q} - \phi_{\bar{\lambda}}) ) \left( \frac{100 \text{ GeV}}{M} \right)^2 10^{-23} \text{ e cm} \]  

(24)

\[ d_{\text{Tl}} \simeq ( 6 \sin(\phi_{A_q} - \phi_{\gamma}) + 2 \times 10^{-4}\sin(\phi_{A_q} - \phi_{\lambda}) + 2 \times 10^{-3}\sin(\phi_{A_q} - \phi_{\bar{\lambda}}) \]

\[ - 1 \times 10^{-2}\sin(\phi_{A_q} - \phi_{\bar{\lambda}}) ) \left( \frac{100 \text{ GeV}}{M} \right)^2 10^{-23} \text{ e cm} \]

(25)
where, for simplicity of notation, all the supersymmetric mass parameters have been assumed to be of order $M$. The quantities on the right hand side arise respectively from the electron edm, nucleon edm, nonderivative nucleon-nucleon coupling, and pseudoscalar-scalar electron-nucleon coupling. In both cases, the electron edm gives the dominant contribution since it is enhanced in heavy open shell atoms. For $^{133}$Cs, with a nuclear quadrupole moment, the nucleon-nucleon coupling is only a factor $\sim 10$ less important than the electron edm (again assuming all the masses and phases are the same order). The present experimental bounds are $^{24, 25}$

$$|d_{\text{Cs}}| < 7.2 \times 10^{-24} \text{ e cm}$$

$$|d_{\text{Tl}}| < 6.6 \times 10^{-24} \text{ e cm}$$

The $^{205}$Tl result gives a bound on the phases contributing to the electron edm of

$$\sin(\phi_{A_e} - \phi_{\tilde{\gamma}}) \left( \frac{100 \text{ GeV}}{M} \right)^2 < 0.1 \quad (100 \text{ GeV} \leq M)$$

For closed shell atoms, good experimental bounds are available for $^{129}$Xe and $^{199}$Hg. Combining our results with those from Table 1.,

$$d_{\text{Xe}} \simeq (8 \times 10^{-3} \sin(\phi_{A_e} - \phi_{\tilde{\gamma}}) - 7 \times 10^{-2} \sin(\phi_{A_q} - \phi_{\tilde{\lambda}}) + 3 \sin(\phi_{A_q} - \phi_{\tilde{\lambda}})$$

$$+ 2 \times 10^{-4} \sin(\phi_{A_q} - \phi_{\tilde{\lambda}}) ) \left( \frac{100 \text{ GeV}}{M} \right)^2 10^{-26} \text{ e cm} \quad (26)$$

$$d_{\text{Hg}} \simeq (-1.2 \times 10^{-2} \sin(\phi_{A_e} - \phi_{\tilde{\gamma}}) - 3 \times 10^{-2} \sin(\phi_{A_q} - \phi_{\tilde{\lambda}}) + 4 \sin(\phi_{A_q} - \phi_{\tilde{\lambda}})$$

$$+ 2 \times 10^{-4} \sin(\phi_{A_q} - \phi_{\tilde{\lambda}}) ) \left( \frac{100 \text{ GeV}}{M} \right)^2 10^{-25} \text{ e cm} \quad (27)$$

The electron edm is here suppressed since the electron spins are paired. The leading contribution comes from the nucleon-nucleon coupling. The present experimental bounds are $^{26-29}$,

$$|d_{\text{Xe}}| < 1.4 \times 10^{-26} \text{ e cm}$$

$$|d_{\text{Hg}}| < 3 \times 10^{-27} \text{ e cm} \quad (95\% \text{ C.L.}) \quad (28)$$

\(^{3}\text{Unless stated otherwise, the experimental bounds given here are the sum of the reported measurement and experimental error. To date, all measurements are consistent with zero.}\)
The $^{199}$Hg result gives a bound on the phases contributing to the light quark chromoelectric dipole moment,

$$\sin(\phi_{Aq} - \phi_{\tilde{\lambda}}) \left(\frac{100 \text{ GeV}}{M}\right)^2 < .008$$

It should be noted that, due to the uncertainties in the hadronic matrix elements and nuclear calculations, this bound is much more uncertain than that from the electron edm given above.

At present, the best bound on a molecular edm is for $^{205}$TlF. Again combining our results with Table 1,.

$$d_{\text{TIF}} \simeq \left(-8 \times 10^{-2}\sin(\phi_{Ae} - \phi_{\tilde{\gamma}}) - 1.5 \times 10^{-1}\sin(\phi_{Aq} - \phi_{\tilde{\lambda}}) + \sin(\phi_{Aq} - \phi_{\tilde{\lambda}}) - 1 \times 10^{-3}\sin(\phi_{Aq} - \phi_{\tilde{\lambda}}) \right) \left(\frac{100 \text{ GeV}}{M}\right)^2 10^{-22} \text{ e cm}$$

(28)

The leading contribution again comes from the nucleon-nucleon coupling since the electron spins are paired. The current experimental bound \cite{10, 30} 

$$|d_{\text{TIF}}| < 4.6 \times 10^{-23} \text{ e cm}$$

does not yet yield a substantial bound on the phases if $M \simeq 100$ GeV.

Because of the different structure of open and closed shell atoms, the experiments bound different microscopic CP violating parameters. Open shell atoms are sensitive to phases in the weak gaugino sector contributing to the electron edm. Closed shell atoms and molecules with paired electron spin get the leading contribution from nuclear effects. This provides a stringent bound on the gluino-squark phases contributing to the light quark chromoelectric dipole moment. For comparison, the neutron edm receives contributions from both the light quark electric and chromoelectric dipole moments \cite{10}. The experimental bound \cite{2} of

$$|d_n| < 8 \times 10^{-26} \text{ e cm},$$

with the estimate (9), gives

$$\sin(\phi_{Aq} - \phi_{\tilde{\lambda}}) \left(\frac{100 \text{ GeV}}{M}\right)^2 < .003$$

This is of the same order as the bound from $^{199}$Hg. Taken together, the experimental bounds provide information on different combinations of CP violating phases.
Other sources of CP violation would give rise to different patterns of edm’s \([31]\). For example, a nonzero \(\bar{\theta}_{QCD}\) would contribute to the edm’s mainly through the T odd pion-nucleon coupling \([7, 32]\). This would give a definite relation among the edm’s of open and closed shell atoms and the neutron.

The prospects for improving the present measurements are encouraging. The ultimate sensitivity of the current \(^{199}\text{Hg}\) experiment is expected to reach the level of \(3 \times 10^{-28} \text{ e cm} \) \([29]\). New techniques may allow further improvement in sensitivity for open shell atoms \([33]\). An unexplored area where the phases contributing to the light quark chromoelectric dipole moment could be measured is in light atoms. In this case the electrons are nonrelativistic and the primary contributions are nuclear. In summary, the experimental study of atomic edm’s provides a promising probe for exploring physics beyond the standard model.

We would like to thank D. Heinzen and V. Kaplunovsky for useful discussions. We would also like to thank N. Fortson for providing us with the unpublished bound on \(^{199}\text{Hg}\). This research was supported in part by the Robert A. Welch Foundation and NSF Grant PHY 9009850.
References

[1] B. McKellar, S. Chodoudhury, X-G. He, and S. Pakvasa, Phys. Lett. B 197 (1987) 556; S. Barr and W. Marciano, in CP Violation edited by C. Jarlskog (World Scientific, Singapore, 1989); F. Hoogeveen, Nucl. Phys. B 431 322 (1990).

[2] I. Altarev et al., JETP Lett. 44 (1986) 460; K. Smith et al., Phys. Lett. B 234 (1990) 191.

[3] L. Schiff, Phys. Rev. 132 (1963) 2194.

[4] P. Sandars, Phys. Lett, 14 (1965) 194; P. Sandars, Phys. Lett. 22 (1966) 290; P. Sandars and R. Sternheimer, Phys. Rev. A 11 (1975) 473; V. Flambaum, Yad. Fiz. 24 (1976) 383 [Sov. J. Nucl. Phys. 24 (1976) 199].

[5] W. Johnson, D. Guo, M. Idrees, and J. Sapirstein, Phys. Rev. A 32 (1985) 2093; W. Johnson, D. Guo, M. Idrees, and J. Sapirstein, Phys. Rev. A 34 (1986) 1043; A. Martensson-Pendrill and P. Öster, Phys. Scr. 36 (1987) 444; W. Johnson, Phys. Scr. 36 (1987) 765; B. Das, Recent Advances in Many Body Theory (Springer-Verlag, New York, 1988) p 411.

[6] E. N. Fortson, Bull. Am. Phys. Soc. 28 (1983) 1321.

[7] I. Khriplovich, Zh. Eksp. Teor. Fiz. 71 (1976) 51 [Sov. Phys. JETP 44 (1976) 25]; V. Dzuba, V. Flambaum, and P. Silvestrov, Phys. Lett. B 154 (1985) 93; V. Flambaum, I. Khriplovich, and O. Sushkov, Phys. Lett. B 162 (1985) 213; V. Flambaum, I. Khriplovich, and O. Sushkov, Nucl. Phys. A 449 (1986) 750; O. Sushkov, V. Flambaum, and I. Khriplovich, Zh. Eksp. Teor. Fiz. 87 (1984) 1521 [Sov. Phys. JETP 60 (1984) 873].

[8] C. Bouchiat, Phys. Lett B 57 (1975) 284; E. Hinds, D. Loving, and P. Sandars, Phys. Lett. B 62 (1976) 97; A. Martensson-Pendrill, Phys. Rev. Lett. 54 (1985) 1153; V. Flambaum and I. Khriplovich, Zh. Eksp. Teor. Fiz. 89 (1985) 1505 [Sov. Phys. JETP 62 (1985) 872].
[9] P. Sandars, Phys. Rev. Lett. 19 (1967) 1396; O. Sushkov and V. Flambaum, Zh. Eksp. Teor. Fiz. 75 (1978) 1208 [Sov. Phys. JETP 48 (1978) 608]; E. Hinds and P. Sandars, Phys. Rev. A 21 (1980) 471; P. Coveney and P. Sandars, J. Phys. B: At. Mol. Phys. 16 (1983) 3727.

[10] D. Cho, K. Sangster, and E. Hinds, Phys. Rev. A 44 (1991) 2783.

[11] I. B. Khriplovich, *Parity Nonconservation in Atomic Phenomena* (Gordon and Breach, Amsterdam 1991).

[12] H. Nilles, Phys. Rep. 110C (1984) 1; H. Haber and G. Kane, Phys. Rep. 117C (1985) 75.

[13] P. Nath, Phys. Rev. Lett. 66 (1991) 2565.

[14] W. Bernreuther and M. Suzuki, Rev. Mod. Phys. 63 (1991) 313.

[15] J. Ellis, S. Ferrara, and D. Nanopoulos, Phys. Lett. B 114 (1982) 231; W. Buchmuller and D. Wyler, Phys. Lett. B 121 321; J. Polchinski and M. Wise, Phys. Lett. B 125 (1983) 393.

[16] R. Arnowitt, J. Lopez, and D. Nanopoulos, Phys. Rev. D 42 (1990) 2423; R. Arnowitt, M. Duff, and K. Stelle, Phys. Rev. D 43 (1991) 3085; Y. Kizuri and N. Oshimo, Phys. Rev. D 45 (1992) 1806.

[17] S. Weinberg, Phys. Rev. Lett. 63 (1989) 2333.

[18] G. Boyd, A. K. Gupta, S. P. Trivedi and M. B. Wise, Phys. Lett. B 237 (1990) 216; M. Dine and W. Fischler, Phys. Lett. B 241 (1990) 584; J. Dai, H. Dykstra, R. G. Leigh, S. Paban, and D. Dicus, Phys. Lett. B 237 (1990) 216, B 242 (1990) 547(E).

[19] A. Morozov, Yad. Fiz. 40 (1984) 788 [Sov. J. Nucl. Phys. 40 (1984) 505]; E. Braaten, C. S. Li, and T. C. Yuan, Phys. Rev. Lett. 64 (1990) 1709; E. Braaten, C. S. Li, and T. C. Yuan, Phys. Rev. D 42 (1990) 276; D. Chang, W. Keung, C. S. Li, and T. C. Yuan, Phys. Lett. B 241 (1990) 589.
[20] A. Manohar and H. Georgi, Nucl. Phys. B 234 (1984) 189; H. Georgi and L. Randall, Nucl. Phys. B 276 (1986) 241.

[21] S. Weinberg, Phys. Lett. B 251 (1990) 288; S. Weinberg, Nucl. Phys. B 363 (1991) 3.

[22] K. Choi and J. Hong, Phys. Lett. B 259 (1991) 340.

[23] A. De Rújula, M. Gavela, O. Pène, and F. Vegas, Phys. Lett. B 245 (1990) 640; A. De Rújula, M. Gavela, O. Pène, and F. Vegas, Nucl. Phys. B 357 (1991) 311.

[24] S. Murthy, D. Krause, Z. Li, and L. Hunter, Phys. Rev. Lett. 63 (1989) 965.

[25] K. Abdullah, C. Carlberg, E. Commins, H. Gould, and S. Ross, Phys. Rev. Lett. 65 (1990) 2347.

[26] T. Vold, F. Raab, B. Heckel, and E. N. Fortson, Phys. Rev. Lett. 52 (1984) 2229.

[27] E. Oteiza, R. Hoare, T. Chupp, Bull. Am. Phys. Soc. 37 (1992) 947.

[28] S. Lamoreaux, J. Jacobs, B. Heckel, F. Raab, and N. Fortson, Phys. Rev. Lett. 59 (1987) 2275.

[29] N. Fortson, private communication.

[30] D. Cho, K. Sangster, and E. Hinds, Phys. Rev. Lett. 63 (1989) 2559.

[31] S. Barr, Phys. Rev. Lett. 68 (1992) 1822; S. Barr, Bartol preprint BA-92-09; S. Thomas, University of Texas Theory Group Preprint UTTG-04-92.

[32] W. Haxton and E. Henley, Phys. Rev. Lett. 51 (1983) 1937.

[33] D. Heinzen, private communication.
|                  | $^{133}$Cs | $^{205}$Tl | $^{129}$Xe | $^{199}$Hg | $^{205}$TlF |
|------------------|------------|------------|------------|------------|------------|
| $d/d_e$          | 120        | -600       | -.8 $\times 10^{-3}$ | 1.2 $\times 10^{-2}$ | 80         |
| $d/d_N$          | $7 \times 10^{-4}$ | $7 \times 10^{-5}$ | $2 \times 10^{-5}$ | $\sim 10^{-4}$ | .5         |
| $d(e\text{ cm})/C_{S,PS}^{en}$ | $2 \times 10^{-24}$ | $3 \times 10^{-26}$ | $5 \times 10^{-26}$ | $6 \times 10^{-25}$ | $2 \times 10^{-22}$ |
| $d(e\text{ cm})/C_{PS,S}^{en}$ | $7 \times 10^{-19}$ | $-5 \times 10^{-18}$ | $9 \times 10^{-23}$ | $1 \times 10^{-21}$ | $-6 \times 10^{-18}$ |
| $d(e\text{ cm})/C_{T,PT}^{en}$ | $9 \times 10^{-21}$ | $5 \times 10^{-21}$ | $5 \times 10^{-21}$ | $-6 \times 10^{-20}$ | $1 \times 10^{-16}$ |
| $d(e\text{ cm})/C_{S,PS}^{en}$ | -         | -1 $\times 10^{-23}$ | $3 \times 10^{-23}$ | $-4 \times 10^{-19}$ | -         |

Table 1. Contributions to electric dipole moments from the Hamiltonian (1).
Figure Captions

**Fig. 1** A typical contribution to $d_e$ from photino exchange.

**Fig. 2** The contribution to an effective electron-nucleon coupling from pion exchange with the T odd pion-nucleon coupling.

**Fig. 3** The contribution to the operator $G_{a\mu\nu}G^{\mu\nu}_a F_{\rho\sigma} \tilde{F}^{\rho\sigma}$ from the electric dipole moment of a heavy quark. Other diagrams related by gauge invariance are not shown.

**Fig. 4** A typical contribution to the electron-quark operators.

**Fig. 5** The contribution to the electron-gluon operators from the electron-heavy quark operators. Other diagrams related by gauge invariance are not shown.

**Fig. 6** Typical contributions to the operator $\bar{e} \gamma^\mu \gamma_5 \gamma^\rho \bar{q} q$; (6a) from a box diagram, (6b) from Z-exchange and the electron electroweak dipole moment.