Baryogenesis and EDMs: Constraining CP Violation Beyond the Standard Model

To cite this article: Christopher Lee 2007 J. Phys.: Conf. Ser. 69 012036

View the article online for updates and enhancements.

Related content
- A Pedagogical Introduction to Electroweak Baryogenesis: The baryon asymmetry
  G A White
- A Pedagogical Introduction to Electroweak Baryogenesis: Introduction
  G A White
- $n$-point functions at finite temperature
  Defu Hou, Enke Wang and Ulrich Heinz

Recent citations
- Jason Tsz Shing Yue
Baryogenesis and EDMs: Constraining CP Violation Beyond the Standard Model

Christopher Lee
Institute for Nuclear Theory, University of Washington, Seattle, WA 98195-1550 USA
E-mail: clee@phys.washington.edu

Abstract. New sources of CP violation beyond the Standard Model may help the scenario of electroweak baryogenesis (EWB) to account for the baryon asymmetry of the universe, and may be detectable in searches for permanent electric dipole moments of fundamental particles. In this talk, I focus on the calculation of the sources and relaxation rates in the transport equations for particle densities in the closed time path formalism of quantum field theory, suitable for the finite temperature and out-of-equilibrium conditions present at the electroweak phase transition. Applying the methods of closed time-path quantum field theory to derive transport equations relevant for EWB in the MSSM, we find regions of the parameter space where generation of the baryon asymmetry is resonantly enhanced, allowing successful baryogenesis with CP-violating phases small enough to be consistent with experimental limits on electron and neutron EDMs.

1. Introduction
Experimental searches for permanent electric dipole moments (EDMs) of elementary particles are now beginning to probe CP violation at such a sensitivity as to impose tight constraints on some models of new physics beyond the Standard Model. At the same time, these new sources of CP violation may allow the scenario of electroweak baryogenesis to account for the baryon asymmetry of the universe (BAU). Therefore, it is imperative to calculate reliably both the EDMs and the BAU that can be generated by these sources. Here, I describe some recent approaches to calculating the BAU using the tools of nonequilibrium quantum field theory [1, 2, 3], and in the context of the MSSM, compare the size of CP-violating phases required for baryogenesis to the limits on those phases implied by EDM searches.

2. Electroweak Baryogenesis
In the scenario of electroweak baryogenesis (see [4] for a review), the baryon asymmetry is generated during a first-order electroweak phase transition, during which bubble nucleation of regions of broken electroweak phase occurs. Outside of these bubbles, baryon-number violating sphaleron processes are active, generating a baryon asymmetry that is subsumed by the bubbles as they grow, inside of which the sphaleron processes turn off, so that the BAU is frozen in. These sphaleron processes, however, require a nonzero chiral fermion density to be already present, and this is generated by CP-violating interactions inside the bubble wall, which then diffuses outside faster than the bubble wall velocity in order to seed the baryon number violating sphalerons. In these ways, electroweak baryogenesis satisfies the three Sakharov criteria for generation of a nonzero baryon asymmetry: baryon-number violation, C and CP violation, and departure
from thermal equilibrium (assuming CPT invariance). The baryon density $\rho_B$ is governed by a diffusion equation:

$$D_q \rho_B''(\bar{z}) - v_w \rho_B'(\bar{z}) - \theta(-\bar{z}) \mathcal{R} \rho_B(\bar{z}) = \theta(-\bar{z}) \frac{n_F}{2} \Gamma_{ws} n_L(\bar{z}),$$  \hspace{1cm} (1)$$

where $D_q$ is the quark diffusion constant, $\mathcal{R}$ is a relaxation coefficient, $n_F$ is the number of families of left-handed fermion doublets present in the bubble wall, and $\Gamma_{ws}$ is the rate of weak sphaleron transitions. The theta function $\theta(-\bar{z})$ accounts for the fact that these sphalerons are active only outside the wall ($\bar{z} < 0$). We have assumed a small curvature for the wall so that densities are a function only of $\bar{z} = z + v_w t$, the distance from the wall boundary. For simplicity we have chosen a step function profile for the bubble wall dividing the regions where weak sphalerons are active or not.

Generation of a nonzero baryon density requires a pre-existing nonzero density of left-handed fermions $n_L$ in the vicinity of the walls of bubbles of nucleating broken electroweak phase, seeding baryon-violating sphaleron processes outside the bubbles. Thus we must also find equations governing the generation of $n_L$ due to CP-violating interactions inside the bubble wall.

3. Transport Equations from Closed-Time-Path Quantum Field Theory

Let us derive the quantum transport equations from the most basic possible level in quantum field theory. For bosonic or fermionic particles, we seek to calculate the expectation values: $\langle \partial_\mu j^\mu(x) \rangle$, where the current densities are:

$$j^\mu_\phi = i(\phi^* \partial^\mu \phi - \phi \partial^\mu \phi^*), \quad j^\mu_\psi = i\bar{\psi} \gamma^\mu \psi,$$  \hspace{1cm} (2)$$

for bosons and fermions, respectively, and the expectation values are taken with respect to a density matrix $\rho$:

$$\langle O \rangle = \frac{1}{Z} \text{Tr}(\rho O),$$  \hspace{1cm} (3)$$

where $Z = \text{Tr}(\rho)$.

In the Minimal Supersymmetric Extension of the Standard Model (MSSM), which we will use shortly as our illustrative example, the relevant densities we must consider are the linear combinations of particles and superparticle densities: $Q = n_{L} + n_{L^c} + n_{L^R} + n_{L^R^c}$, $T = n_{R} + n_{R^c}$, and $H = n_{H} - n_{H^c} - n_{H^R} - n_{H^R^c}$. We choose these combinations making the assumption that gauge and supergauge interactions keep gauge partner and superpartner densities in thermal equilibrium [5]. A more rigorous analysis relaxing these assumptions can be pursued. The density $H$ is associated with the vector Higgsino current. A density $h$ associated with the axial current may also be considered.

These densities evolve out of thermal equilibrium, but assuming the departure from equilibrium is small, we approximate the distribution of states by the equilibrium form but with a spacetime-varying chemical potential, $n_{B,F}(\omega, \mu(x))$. We can relate the particle densities to their associated chemical potentials, through the relation:

$$n_i = g_i \int \frac{d^3 k}{(2\pi)^3} \left[ n_{B,F}(\omega, \mu_i) - n_{B,F}(\omega, -\mu_i) \right],$$  \hspace{1cm} (4)$$

where $g_i$ is the number of internal degrees of freedom of particle type $n_i$ (e.g. color), leading to the relation, $n_i = k_i T^2 \mu_i / 6$, where $k_i$ is a statistical factor dependent on the mass and type of particle.

The evolution of the particle densities can be derived beginning from Schwinger-Dyson equations in the Closed Time Path formalism of quantum field theory [6]. This formalism
and anti-time-ordered Green’s functions. Green’s functions for fermionic fields are similarly defined, with an additional minus sign in the path-ordering operator $\mathcal{P}$ acts on the operators labeled by $+$ and $-$ by placing all $-$ fields to the left of $+$ fields, time-ordering $+$ fields amongst themselves, and anti-time-ordering $-$ fields amongst themselves. The time integrals in (5) go, as usual, from $-\infty$ to $\infty$. By defining a “closed time path” contour $C$ going from $-\infty$ to $\infty$ and back again to $-\infty$, we can combine the terms in (5) into:

$$\langle n|\mathcal{P}\left\{ \mathcal{O}(x)\exp\left(i\int_{C} d^{4}x \mathcal{L}_{+} - i\int_{C} d^{4}x \mathcal{L}_{-}\right)\right\} |n\rangle,$$

(6)

where the closed time-path contour $C$ goes from $-\infty$ to $\infty$ ($C_{+}$) and back again to $-\infty$ ($C_{-}$), as shown in Fig. 1. The operator $\mathcal{P}$ orders fields from right to left in the order they appear on the path $C$. Green’s functions of fields on this contour, $G(x,y) = \langle \mathcal{P}\phi(x)\phi^{\ast}(y) \rangle$, are really four possible functions, depending on the relative location of $x,y$ on the two time branches:

$$G^{>}(x,y) = \langle \phi_{-}(x)\phi_{+}^{\ast}(y) \rangle \quad G^{<}(x,y) = \langle \phi_{+}^{\ast}(y)\phi_{+}(x) \rangle$$

$$G^{t}(x,y) = \langle T\phi_{+}(x)\phi_{+}^{\ast}(y) \rangle \quad G^{t}(x,y) = \langle \overline{T}\phi_{-}(x)\phi_{-}^{\ast}(y) \rangle,$$

(7)

which are conveniently expressed in the matrix form,

$$\bar{G}(x,y) = \begin{pmatrix} G^{t}(x,y) & -G^{<}(x,y) \\ G^{>}(x,y) & -G^{t}(x,y) \end{pmatrix}. $$

(8)

Green’s functions for fermionic fields are similarly defined, with an additional minus sign in the $<$ and anti-time-ordered Green’s functions.

The matrix of Green’s functions $\bar{G}(x,y)$ satisfies a Schwinger-Dyson equation,

$$\bar{G}(x,y) = \bar{G}^{0}(x,y) + \int d^{4}z \int d^{4}w \bar{G}^{0}(x,z)\Sigma(z,w)\bar{G}(w,y)$$

(9)

$$\bar{G}(x,y) = \bar{G}^{0}(x,y) + \int d^{4}z \int d^{4}w \bar{G}(x,z)\bar{\Sigma}(z,w)\bar{G}^{0}(w,y),$$

(10)
where $\tilde{G}_0$ are the free Green’s functions. The time integrals here are over the normal path, $-\infty$ to $\infty$. By applying the Klein-Gordon operator $(\Box_x + m^2)$ to the first equation and $(\Box_y + m^2)$ to the second, subtracting the two resulting equations, and taking the limit $x = y$, we find that the equations are turned into a quantum Boltzmann equation:

$$\frac{\partial n}{\partial X_0} + \nabla \cdot j(X) = \int d^3z \int_{-\infty}^{X_0} dz^0 \left[ \Sigma^> (X, z) G^< (z, X) - G^> (X, z) \Sigma^< (z, X) \right] \tag{11}$$

with the quantity on the right-hand side giving the source for the current $j^\mu (X)$. A nonzero source is induced by $CP$-violating interactions in the self-energies $\Sigma$, or by a nonzero chemical potential in spectral representation of the Green’s functions $G^\Sigma$:

$$G^\Sigma (x, y) = \int \frac{d^4k}{(2\pi)^4} e^{-ik \cdot (x-y)} f^\Sigma_B (k^0, \mu_\phi) \rho_\phi (k), \tag{12}$$

where $f_B^\Sigma (k^0, \mu_\phi) = 1 + n_B (k^0 - \mu_\phi)$ and $f_B^\Sigma (k^0, \mu_\phi) = n_B (k^0 - \mu_\phi)$. The spectral function $\rho_\phi (k)$ contains the pole structure of the Green’s functions, which are determined by the thermal masses and widths of the excitations of $\phi$ in the spectral function. Similar results hold for fermionic fields as well. We will make the approximation of small densities, so that $\mu/T \ll 1$. In an expansion in powers of $\mu/T$, we find that the $CP$-violating source enters at zeroth-order, while at first-order, we find terms describing relaxation of the densities.

### 4. Baryogenesis in the Minimal Supersymmetric Standard Model

In the minimal supergravity supersymmetry-breaking scenario of the MSSM (see [7]), there are two independent $CP$-violating phases in the Lagrangian that could contribute to a nonzero source through the self-energies in Eq. (11). These are in the $\mu$ parameter in the $\mu H_u H_d$ term in the MSSM superpotential, and the $A_t$ parameter in the soft supersymmetry-breaking Lagrangian. We look only at the contribution of third-generation squarks due to their large Yukawa coupling $y_t \sim O(1)$. The interactions of squarks involving the $\mu$ and $A_t$ parameters are given by:

$$\mathcal{L}_t = y_t \tilde{t}_L (A_t H_u^0 - \mu^* H_d^{0*}) \tilde{t}_R + \text{h.c.}, \tag{13}$$

where $\mu = |\mu| e^{i\phi_\mu}$ and $A_t = |A_t| e^{i\phi_A}$. There are also $CP$-violating phases in the Lagrangians for Higgsinos and gauginos.

After electroweak symmetry breaking, the Higgs fields obtain a vacuum expectation value (vev). We treat vev insertions as interaction vertices, assuming $v/T$ is small, which give rise to the $CP$-violating source and relaxation terms in the transport equations [2], through the

![Figure 2](image.png)

**Figure 2.** Self-energies from scattering from Higgs vevs. The self-energies for (a) fermions or (b) bosons inserted into the quantum Boltzmann equations generated both $CP$-violating sources for squark and Higgsino densities in the MSSM as well as relaxation rates $\Gamma_{M,H}$. 


self-energies illustrated in Fig. 2. Interactions with the physical Higgs particles gives rise to terms in the quantum transport equations that change (s)quark and Higgs(ino) densities. This effect is vital to create a net quark density in scenarios where the CP-violating Higgsino sources is much larger than the squark source.

Figure 3. One-loop self-energies. The contribution of these self-energies inserted into the quantum Boltzmann equations generated the $\Gamma$ flavor-transfer terms in the transport equations mixing (s)quark and Higgs(ino) densities. This effect is vital to create a net quark density in scenarios where the CP-violating Higgsino sources is much larger than the squark source.

Plugging the leading-order self-energies induced by interactions in the MSSM Lagrangian into the quantum Boltzmann equation (11) for each species of particle, expanding to first-order in $v/T$ and $\mu/T$, we obtain transport equations of the form:

$$v_u T'(\bar{z}) - D_u T''(\bar{z}) = S^\text{OP}_{tR} + S_{s,s} - \Gamma_M(\mu_T - \mu_Q) + \Gamma_Y(\mu_Q - \mu_T + \mu_H) \tag{14}$$

$$v_u Q'(\bar{z}) - D_u Q''(\bar{z}) = -S^\text{OP}_{tR} - 2S_{s,s} - \Gamma_M(\mu_Q - \mu_T) - \Gamma_Y(\mu_T - \mu_Q + \mu_H) \tag{15}$$

$$v_u H'(\bar{z}) - D_u H''(\bar{z}) = S^\text{OP}_{\tilde{H}} - \Gamma_H \mu_H - \Gamma_Y(\mu_Q - \mu_T + \mu_H), \tag{16}$$

where $S^\text{OP}_{tR}$ are the CP-violating sources for squarks and Higgsinos in the MSSM, $S_{s,s}$ is a source induced by strong sphalerons, the $\Gamma_M, H$ terms are relaxation terms induced by scattering with the background Higgs vev, and the $\Gamma_Y$ terms are induced by Yukawa or tri-scalar $y_t$-proportional interactions. On the left-hand side we have used that the densities depend on time $t$ only through $\bar{z} = z + v_u t$, and made the diffusion approximation $j_{T,Q} = D_q \nabla n_{T,Q}$, $j_H = D_h \nabla n_H$ for the currents.

At zeroth order in $\mu/T$, we obtain from the right-hand side of Eq. (11) the CP-violating source:

$$S^\text{OP}_{tR} = \frac{N_C y_t^2}{4\pi^2 T} \left| A_t v_u(x) - \mu v_d(x) \right|^2 \int_0^{\infty} \frac{dk^2}{\omega^2_{i_R} \omega^2_{i_L}} \text{Im} \left\{ \frac{n_B(\mathcal{E}_R) - n_B(\mathcal{E}_L)}{(\mathcal{E}_R - \mathcal{E}_L)^2} + \frac{1 + n_B(\mathcal{E}_L) + n_B(\mathcal{E}_R)}{(\mathcal{E}_L + \mathcal{E}_R)^2} \right\}, \tag{17}$$

and at first-order, the relaxation coefficient:

$$\Gamma^\pm_M = \frac{N_C y_t^2}{4\pi^2 T} |A_t v_u(x) - \mu v_d(x)|^2 \int_0^{\infty} \frac{dk^2}{\omega^2_{i_R} \omega^2_{i_L}} \text{Im} \left\{ \frac{h_B(\mathcal{E}_L) \mp h_B(\mathcal{E}_R)}{\mathcal{E}_L - \mathcal{E}_R} + \frac{h_B(\mathcal{E}_L) \mp h_B(\mathcal{E}_R)}{\mathcal{E}_L + \mathcal{E}_R} \right\}, \tag{18}$$

where $N_C$ is the number of colors of squarks, $v(x)^2 = v_u(x)^2 + v_d(x)^2$, $\tan \beta(x) = v_u(x)/v_d(x)$, $\omega_i = (k^2 + m_i(T)^2)^{1/2}$ contains the thermal squark mass, $\mathcal{E}_i = \omega_i - i\Gamma_i$ contains the thermal squark width, $n_B(x) = 1/(e^{x/T} - 1)$ and $h_B(x) = T d n_B(x)/dx$. Similar but more involved expressions are found for the Higgsino source $S^\text{OP}_{\tilde{B}}$ [1] and relaxation coefficient $\Gamma_B$ [2], which are shown just in graphical form in Fig. 4. Both the source and relaxation coefficient are resonantly enhanced when $m_{i_L} = m_{i_R}$. However, precision electroweak constraints force $m_{i_L}$...
to be heavy while a first-order electroweak phase transition in the MSSM requires \( m_{\tilde{t}_R} \) to be light. A resonance in the Higgsino source and relaxation terms when Higgsino and weak gaugino masses \(|\mu|\) and \(M_2\) are equal may still be realized. The resonance of the \(CP\)-violating source is stronger than that of the relaxation term, so as the two effects compete in the final resulting baryon asymmetry, the resonance in the source wins out. When the Higgsino source is dominant, then the \(\Gamma_Y\) terms in the transport equations (14) have the effect of transferring the \(H\) density to \(Q, T\) densities, which are necessary for generation of the baryon asymmetry. The effects of varying the value of \(\Gamma_Y\) are investigated further in [3].

In extensions of the MSSM with an extra gauge singlet scalar field, the phase transition is strengthened by the extra field. (For a study of electroweak baryogenesis in such an extension, see [10].) Then the requirement that the right-handed squark mass \(m_{\tilde{t}_R}\) be small while \(m_{\tilde{t}_L}\) is large may be relaxed, and a resonance in the squark source becomes possible.

The coupled transport equations (14) may be solved for the densities \(Q, T, H\), assuming \(\Gamma^\pm_M\) is constant in the broken phase, \(\bar{z} > 0\), and zero outside, while the sources \(S^{OP}_H\) are constant inside the bubble wall \(0 < \bar{z} < L_w\), and zero elsewhere. This allows for an analytical solution to the equations, although they can be solved for a more general wall profile numerically. The total left-handed fermion density in the end is given by \(n_{\tilde{L}}(\bar{z}) = 5Q(\bar{z}) + 4T(\bar{z})\) [5], which is plugged into Eq. (1), which is then solved for the baryon density. The dependence of \(\rho_B\) on the phases \(\phi_{\mu, A}\) takes the form:

\[
\rho_B(\bar{z} > 0) = F_1 \sin \phi_\mu + F_2 \sin(\phi_\mu + \phi_A),
\]

where \(F_1\) and \(F_2\) are the \(CP\)-violating phases.

5. Electric Dipole Moments

The \(CP\)-violating phases \(\phi_{\mu, A}\) in the MSSM may also induce EDMs of elementary particles through loop effects, such as those shown in Fig. 5. The \(CP\)-violating phases enter the mixing matrices for mass eigenstates of the various superpartners. These graphs generate terms in a low-energy effective Lagrangian of the form:

\[
\mathcal{L}_E = -\frac{i}{2} d_f \bar{\psi}\sigma_{\mu\nu}\gamma_5 \psi F^{\mu\nu},
\]

where \(d_f\) is the electric dipole moment of fermion \(\psi\), and

\[
\mathcal{L}_C = -\frac{i}{2} \bar{q}_C \sigma_{\mu\nu}\gamma_5 T^A q^A \Gamma_{\mu\nu},
\]

(19)...

(20)
Figure 5. Sample one- and two-loop graphs in the MSSM generating electric dipole moments of fermions. $CP$ violation in the chargino and neutralino mixing matrices contribute to EDMs through these graphs. The contribution of the one-loop graphs decreases with increasing sfermion ($\tilde{f}$) masses, so that for large enough sfermion masses, the two-loop graphs contribute dominantly to the EDM.

where $d_{C}$ is the chromoelectric dipole moment of quark $q$. The EDM of the neutron is related to the electric and chromoelectric dipole moments of quarks [11]:

$$d_n = \frac{4}{3}d_d - \frac{1}{3}d_u + 1.1eg_s(d_\tilde{d} + 0.5\tilde{d}_u).$$

(21)

The present best experimental limits on the electron [12] and neutron [13] EDMs are $|d_e| < 1.9 \times 10^{-27} e \cdot cm$, and $|d_n| < 3.6 \times 10^{-26} e \cdot cm$, at 95% confidence level. The constraints on the phases $\phi_{\mu,A}$ implied by these EDM limits taking the EDMs to be given by the one-loop contributions in [14], one of which is illustrated in Fig. 5, are shown in Fig. 6, together with the phases required to account successfully for the observed ratio of baryon-to-entropy density, $\rho_B/s = (7.3 \pm 2.5) \times 10^{-11}$ [15]. The two cases shown are for parameters on the Higgsino-gaugino resonance, $|\mu| = M_2 = 200$ GeV, and away from the resonance, moving $|\mu|$ to 250 GeV. The values of the other MSSM parameters used to make Fig. 6 are given in [2].

Often the constraints on $CP$-violation from EDM searches are evaded by increasing the sfermion masses $m_{\tilde{f}}$ by reducing the size of the one-loop graph like that shown in Fig. 5 by increasing the postulated sfermion mass $m_{\tilde{f}}$. However, Refs. [16, 17, 18] showed that there are two-loop graphs contributing to the EDMs containing no sfermions, and thus independent of the sfermions masses. Above some value of sfermion masses (for our parameter choices, above 2–3 TeV), then, the two-loop graphs, such as that in Fig. 5, take over as the dominant contribution to the EDMs, giving an unavoidable constraint on the $CP$-violating phases.

The results illustrated by Fig. 6 imply that electroweak baryogenesis in the minimal supergravity SUSY-breaking scenario of the MSSM is still consistent with current EDM limits. However, improvements in the sensitivity of EDM searches by two-to-three orders of magnitude, for example with neutron EDM search planned at SNS [19], together with accelerator searches for the supersymmetric particles, should be able to decisively probe this scenario of baryogenesis.

6. Conclusions
Experiments today and in the near future at accelerators and precision low-energy searches can point us to the correct scenario of baryogenesis, constraining the space of $CP$-violating and other parameters that could account for the observed baryon asymmetry and remain consistent with limits from permanent electric dipole moment searches. Theoretical developments applying nonequilibrium quantum field theory to the quantum transport of $CP$-violating densities during the electroweak phase transition have revealed parameter regions where baryon production
Figure 6. Regions in the $\phi_\mu$-$\phi_A$ plane implied by consistency with the 95% C.L. limits on electron (solid) and neutron (dashed) EDMs, and the baryon asymmetry (shaded) implied by BBN [15] at 95% C.L. On the left baryon production is resonantly enhanced, with $|\mu| = M^2 = 200$ GeV, while on the right, $M^2 = 200$ GeV and $|\mu| = 250$ GeV, away from the resonance.

becomes resonantly enhanced, requiring smaller $CP$-violating phases to generate the correct baryon asymmetry. For example, in the minimal supergravity scenario of the MSSM the phases required for successful electroweak baryogenesis lie just within the region still allowed by the constraints from EDM searches, but will be decisively probed by a new generation of experiments. These same methods can be applied to other models, so that we may be prepared to follow the direction to which new experiments and searches for new physics will point.

Acknowledgments
I would like to thank Vincenzo Cirigliano, Michael Ramsey-Musolf, and Sean Tulin for their collaboration on the work presented here. This work was supported in part by the U.S. Department of Energy under grant number DE-FG02-00ER41132.

References
[1] Riotto A 1998 *Phys. Rev.* D 58 095009
[2] Lee C, Cirigliano V and Ramsey-Musolf M J 2005 *Phys. Rev.* D 71 075010
[3] Cirigliano V, Lee C, Ramsey-Musolf M J, Tulin S 2006 *Phys. Rev.* D 73 115009
[4] Trodden M 1999 *Rev. Mod. Phys.* 71 1463
[5] Huet P and Nelson A 1996 *Phys. Rev.* D 53 4578
[6] Chou K c, Su Z b, Hao B l and Yu L 1985 *Phys. Rept.* 118 1
[7] Martin S P 1997 A supersymmetry primer Preprint hep-ph/9709356
[8] Carena M, Moreno J M, Quiros M, Seco M and Wagner C E M 2001 *Nucl. Phys.* B 599 158
[9] Konstandin T, Prokopec T, Schmidt M G and Seco M 2005 *Nucl. Phys.* B 738 1
[10] Huber S J, Konstandin T, Prokopec T and Schmidt M G 2006 *Nucl. Phys.* B 757 172
[11] Pospelov M and Ritz A 2001 *Phys. Rev.* D 63 073015
[12] Regan B C, Commins E D, Schmidt C J and DeMille D 2002 *Phys. Rev. Lett.* 88 071805
[13] Baker C A *et al.* 2006 *Phys. Rev. Lett.* 97 131801
[14] Ibrahim T and Nath P 1998 *Phys. Rev.* D 57 478 (Erratum-ibid. 58 019901, 60 079903, 60 119901)
[15] Eidelman S *et al.* [Particle Data Group Collaboration] 2004 *Phys. Lett.* B 592 1
[16] Chang D, Keung W Y and Pilaftsis A 1999 *Phys. Rev. Lett.* 82 900
[17] Chang D, Chang W F and Keung W Y 2002 *Phys. Rev.* D 66 116008
[18] Pilaftsis A 2002 *Nucl. Phys.* B 644 263
[19] Ito T 2006 in these Proceedings