Implications of R-parity violating supersymmetry for atomic and hadronic electric dipole moments

Amand Faessler,[1] Thomas Gutsche,[1] Sergey Kovalenko,[2] and Valery E. Lyubovitskij[1,*]

[1] *Institut für Theoretische Physik, Universität Tübingen, Auf der Morgenstelle 14, D-72076 Tübingen, Germany*
[2] *Departamento de Física, Universidad Técnica Federico Santa María, Casilla 110-V, Valparaíso, Chile*

(Received 26 July 2006; published 19 October 2006)

We calculate the electric dipole moments (EDM) of the neutral ^{199}Hg atom, deuteron, nucleons and neutral hyperons Λ, Σ^0 and Ξ^0 in the framework of a generic SUSY model without R-parity conservation ($\not{R}_p$SUSY) on the basis of the SU(3) version of chiral perturbation theory (ChPT). We consider CP-violation in the hadronic sector induced by the chromoelectric quark dipole moments and CP-violating 4-quark effective interactions. From the null experimental results on the neutron and ^{199}Hg atom EDMs we derive limits on the imaginary parts of certain products $\mathrm{Im}(\lambda' \lambda'^*)$ of the trilinear $\not{R}_p$-couplings and demonstrate that they are more stringent than those existing in the literature. Using these limits we give predictions for the EDMs of neutral hyperons. We also estimate the prospects of future storage ring experiments on the deuteron EDM and show that the expected improvement of the above limits in these experiments may reach several orders of magnitude.

DOI: 10.1103/PhysRevD.74.074013 PACS numbers: 12.39.Fe, 11.30.Er, 13.40.Em, 14.20.Dh

I. INTRODUCTION

Electric dipole moments (EDMs) of neutral atoms, hadrons and leptons are flavor blind CP-odd observables which are recognized to be sensitive probes of physics beyond the standard model (SM). As is known, the SM predictions for these observables are at least 6–7 orders of magnitude below the current experimental limits. Thus an observation of EDMs at larger values would witness the presence of a non-SM source of CP-violation (CPV). Physics beyond the SM brings in new complex parameters and, therefore, new sources of CPV which may contribute to EDMs. In supersymmetric (SUSY) models, these parameters come from the soft SUSY breaking sector and the superpotential μ-term and, if R-parity is not imposed, additional CPV phases appear from the R-parity violating trilinear and bilinear parameters.

During the last few years significant progress has been achieved in experimental studies of various EDMs [1–4]. Presently there exist stringent upper bounds on the neutron EDM, d_n, [2] and the EDM, d_{Hg}, of the neutral ^{199}Hg atom [3]:

$$|d_n| \leq 3.0 \times 10^{-26}\ e \cdot \mathrm{cm}, \tag{1}$$

$$|d_{\mathrm{Hg}}| \leq 2.1 \times 10^{-28}\ e \cdot \mathrm{cm}. \tag{2}$$

Recently it was also proposed to measure the deuteron EDM, d_D, in storage ring experiments [4] with deuteron ions instead of neutral atoms. The advantage of these experiments is the absence of Schiff screening, which introduces significant uncertainties in the case of neutral atoms. This allows a direct probe of the value for d_D. In the near future it is hoped to obtain the experimental upper bound of

$$|d_D| \leq (1 \div 3) \times 10^{-27}\ e \cdot \mathrm{cm}. \tag{3}$$

The upper limits for the EDMs, derived from the above null experimental results, stringently constrain or even reject various models of CPV [5]. For the case of SUSY models with the superpartner masses around the electroweak scale ~100 GeV–1 TeV, these limits imply that the CPV SUSY phases are very small. Various aspects of the calculation of the EDMs within the popular versions of SUSY models with [6,7] and without [8–12] R-parity conservation have been studied in the literature.

In the present paper we are studying the EDMs of the ^{199}Hg atom and the deuteron as well as of the light baryons (nucleon and neutral hyperons Λ, Σ^0 and Ξ^0) in the framework of a generic SUSY model without R-parity conservation ($\not{R}_p$SUSY) on the basis of the $SU(3)$ version of chiral perturbation theory (ChPT) [7,13–15]. We consider CP-violation in the hadronic sector originating from the quark chromoelectric dipole moments (CEDMs) and CPV 4-quark effective interactions which are induced by the complex phases of the trilinear $\not{R}_p$-couplings λ'. From the experimental bounds of Eqs. (1) and (2) we derive upper limits on the imaginary parts of the products of the trilinear $\not{R}_p$-couplings and compare these limits to the existing ones. On this basis we predict the values for the EDMs of the light neutral hyperons. We also discuss the prospects of the deuteron EDM experiments (3) from the view point of their ability to improve these limits.

II. HADRONIC EDMS IN CHIRAL PERTURBATION THEORY

Here we briefly outline the formalism we use for the calculation of the EDMs of the neutral ^{199}Hg atom, the deuteron and light baryons n, p, Λ, Σ^0 and Ξ^0.

*On leave of absence from Department of Physics, Tomsk State University, 634050 Tomsk, Russia.

The ^{199}Hg is a diamagnetic atom with a closed electron shell. Its EDM is dominated by the nuclear CP-violating effects characterized by the Schiff moment $S_{\rm Hg}$, generating a T-odd electrostatic potential for atomic electrons. The ^{199}Hg atomic EDM is given by [16]

$$d_{\rm Hg} = -2.8 \times 10^{-4} S_{\rm Hg} \cdot {\rm fm}^{-2}. \qquad (4)$$

The deuteron EDM is a theoretically rather clean problem [17] since the deuteron represents the simplest nucleus with a well understood dynamics. The corresponding EDM can be written as the sum of the three terms

$$d_D = d_p + d_n + d_D^{NN}, \qquad (5)$$

where d_n, d_p are the neutron and proton EDMs, respectively, and d_D^{NN} is due to the CP-violating nuclear forces.

For the evaluation of the proton-neutron CP-odd nuclear term, d_D^{NN} and the Schiff moment $S_{\rm Hg}$ we are using a one-meson (π or η) exchange model with CP-odd meson-nucleon interactions [7,17–19].

The baryon EDM d_B is given by the value of the corresponding form factor at zero recoil, i.e. $d_B = D_B(0)$. The EDM form factor $D_B(Q^2)$ is defined in the standard way through the baryon matrix element of the electromagnetic current:

$$\begin{aligned}\langle B(p')|J_\mu(0)|B(p)\rangle = \bar{u}_n(p')\bigg[& \gamma_\mu F_B^1(Q^2) \\ & + \frac{i}{2m_B}\sigma_{\mu\nu} q^\nu F_B^2(Q^2) \\ & - \sigma_{\mu\nu}\gamma_5 q^\nu D_B(Q^2) \\ & + (\gamma_\mu q^2 - 2m_N q_\mu)\gamma_5 A_B(Q^2)\bigg] u_n(p),\end{aligned} \qquad (6)$$

where, in addition, $F_B^1(Q^2)$ and $F_B^2(Q^2)$ are the well-known CP-even electromagnetic form factors and $A_B(Q^2)$ is the baryon anapole moment form factor.

In what follows we evaluate the EDMs of light baryons, d_n, d_p, d_Λ, d_{Σ^0} and d_{Ξ^0} on the basis of the $SU(3)$ version of Chiral Perturbation Theory (ChPT) [7,14,15]. We use the Lagrangian of $SU(3)$ ChPT of order $O(p)$ and restrict ourselves to the meson-loop diagrams given in Fig. 1. As it is known and was discussed before in the literature (see e.g. Ref. [7]), an accurate calculation of the baryon EDMs should also include the contribution of the unknown low-energy constants (LECs) which parameterize the short-distance effects and remove the ultraviolet divergences. However, we assume that at the level of accuracy necessary for the analysis of the $\not{R}_p$SUSY in hadronic EDMs the one-loop meson-cloud approximation [7,14,15] is adequate.

The CP-conserving vertices of the diagrams in Fig. 1 correspond to the terms of the ChPT Lagrangian which is given by the sum of leading meson and meson-baryon pieces:

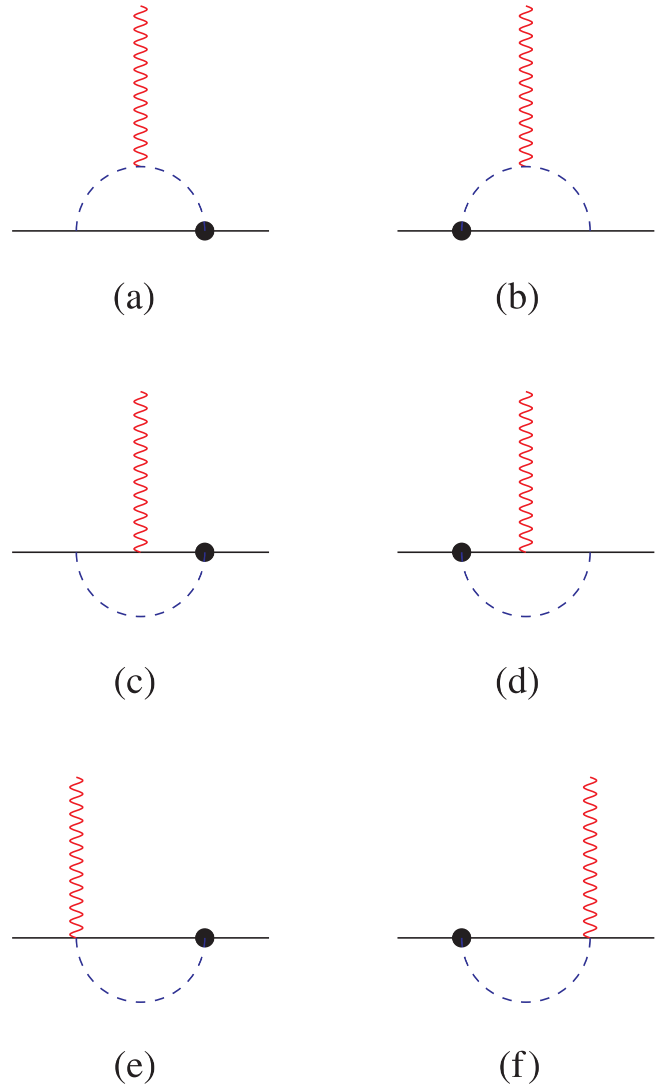

FIG. 1 (color online). Meson-loop diagrams contributing to the EDMs of baryons. Solid, dashed and wiggly lines refer to baryons, pseudoscalar mesons and electromagnetic field, respectively. A CP-violating vertex is denoted by a black filled circle.

$$\begin{aligned}\mathcal{L} = & \frac{F_\pi^2}{4}\langle u_\mu u^\mu + \chi_+\rangle + \langle \bar{B}(i\not{D} - m_B)B\rangle \\ & + \frac{D}{2}\langle \bar{B}\gamma^\mu\gamma^5\{u_\mu B\}\rangle + \frac{F}{2}\langle \bar{B}\gamma^\mu\gamma^5[u_\mu B]\rangle\end{aligned}$$

where $D = 0.80$ and $F = 0.46$ are the baryon axial coupling constants, m_B is the baryon mass in the chiral limit, the symbols $\langle\ldots\rangle$, $\{\ldots\}$ and $[\ldots]$ denote the trace over flavor matrices, anticommutator and commutator, respectively.

We use the standard notation for the basic blocks of the ChPT Lagrangian [13] where

$$U = u^2 = \exp(iP\sqrt{2}/F_\pi) \qquad (7)$$

is the chiral field collecting pseudoscalar fields P in the exponential parametrization with $F_\pi = 92.4$ MeV being the octet leptonic decay constant, D_μ denotes the chiral and gauge-invariant derivative, $u_\mu = iu^\dagger D_\mu U u^\dagger$, $\chi_\pm = u^\dagger \chi u^\dagger \pm u \chi^\dagger u$, $\chi = 2B(s+ip)$, $s = \mathcal{M} + \ldots$ and $\mathcal{M} = \mathrm{diag}\{m_u, m_d, m_s\}$ are the charge and the mass matrix of current quarks, respectively; B is the quark vacuum condensate parameter. The explicit form of the octet matrices of pseudoscalar mesons P and baryons B can be found e.g. in Refs. [7,14,15]. In our analysis we take into account $\pi^0 - \eta$ meson mixing [13] with the corresponding mixing angle ε given by

$$\tan 2\varepsilon = \frac{\sqrt{3}}{2}\frac{m_d - m_u}{m_s - \hat{m}} \tag{8}$$

where $\hat{m} = (m_u + m_d)/2$. In the numerical calculations we use the standard set of current quark masses: $m_u = 5$ MeV, $m_d = 9$ MeV and $m_s = 175$ MeV. For the pion and kaon masses we use the values of the charged mesons: $M_\pi = M_{\pi^\pm} = 139.57$ MeV and $M_K = M_{K^\pm} = 493.677$ MeV. For the baryon masses we use the universal parameter, the value of the octet baryon mass in the chiral limit, which for convenience is identified with the proton mass: $m_B = m_p = 938.27$ MeV. Note that the Lagrangian (7) generates the CP-even meson-baryon, photon-meson and photon-baryon coupling.

The CP-odd meson-baryon Lagrangian has been derived in Ref. [7] where one can find its complete form. Here we explicitly only show the CPV pion-nucleon terms:

$$\begin{aligned}\mathcal{L}^{\rm CPV}_{\rm MBB} &= \bar{g}_{\rm MBB}\bar{B}MB \\ &= \bar{N}\{\bar{g}^{(0)}_{\pi NN}\vec{\pi}\,\vec{\tau} + \bar{g}^{(1)}_{\pi NN}\pi^0 + \bar{g}^{(2)}_{\pi NN}(\vec{\pi}\,\vec{\tau} - 3\pi^0\tau^3)\}N \\ &\quad + \ldots\end{aligned} \tag{9}$$

where $\bar{g}^{(0)}_{\pi NN}$, $\bar{g}^{(1)}_{\pi NN}$ and $\bar{g}^{(2)}_{\pi NN}$ are the corresponding isoscalar, isovector and isotensor coupling constants.

III. EDMS IN $\not{R}_p$ SUSY: QUARK LEVEL

The effective CP-odd Lagrangian in terms of quark, gluon and photon fields up to operators of dimension six, normalized at the hadronic scale ~1 GeV, has the following standard form:

$$\begin{aligned}\mathcal{L}^{\rm CPV} &= \frac{\bar{\theta}}{16\pi^2}\,\mathrm{tr}(\tilde{G}_{\mu\nu}G^{\mu\nu}) - \frac{i}{2}\sum_{i=u,d,s} d_i\bar{q}_i\sigma^{\mu\nu}\gamma_5 eF_{\mu\nu}q_i \\ &\quad - \frac{i}{2}\sum_{i=u,d,s}\tilde{d}_i\bar{q}_i\sigma^{\mu\nu}\gamma_5 g_s G^a_{\mu\nu}T^a q_i \\ &\quad - \frac{1}{6}C_W f^{abc}G^a_{\mu\alpha}G^{b\alpha}_{\nu}G^c_{\rho\sigma}\varepsilon^{\mu\nu\rho\sigma},\end{aligned} \tag{10}$$

where $G^a_{\mu\nu}$ is the gluon stress tensor, $\tilde{G}_{\mu\nu} = \frac{1}{2}\epsilon_{\mu\nu\sigma\rho}G^{\sigma\rho}$ is its dual tensor, and T^a and f^{abc} are the $SU(3)$ generators and structure constants, respectively. In this equation the first term represents the SM QCD θ-term, while the last three terms are the nonrenormalizable effective operators induced by physics beyond the SM. The second and third terms are the dimension-five electric and chromoelectric dipole quark operators, respectively, and the last term is the dimension-six Weinberg operator. The light quark EDMs and CEDMs are denoted by d_i and $\tilde{d}_i$, respectively. In what follows we adopt the Peccei-Quinn mechanism, eliminating the $\bar{\theta}$-term as an independent source of CPV.

We also consider the 4-quark CPV interactions of the form [20,21]

$$\begin{aligned}\mathcal{L}^{\rm CPV}_{4q} &= \sum_{i,j}\{C^P_{ij}(\bar{q}_i q_i)(\bar{q}_j i\gamma_5 q_j) + C^T_{ij}(\bar{q}_i\sigma_{\mu\nu}q_i) \\ &\quad \times(\bar{q}_j i\sigma^{\mu\nu}\gamma_5 q_j)\},\end{aligned} \tag{11}$$

where the sum runs over all the quark flavors $i, j = u, d, s, c, b, t$. The operators of the above Lagrangians in Eqs. (10) and (11) can be induced by physics beyond the SM at loop or tree level after integrating out the heavy degrees of freedom.

Here we are studying the CPV effects in the hadronic sector induced by the trilinear interactions of $\not{R}_p$SUSY models. The corresponding part of the R_p-violating superpotential reads:

$$W_{\not{R}_p} = \lambda'_{ijk}L_i Q_j D^c_k, \tag{12}$$

where the summation over the generation indices i, j, k is understood, L, Q and D^c are the superfields of lepton-sleptons, quarks-squarks and CP-conjugated quarks-squarks, respectively, and λ'_{ijk} are the complex coupling constants violating lepton number conservation. Equation (12) results in the interactions

$$\begin{aligned}\mathcal{L}_{\lambda'} &= -\lambda'_{ijk}(\tilde{\nu}_{iL}\bar{d}_k P_L d_j + \tilde{d}_{jL}\bar{d}_k P_L \nu_i + \tilde{d}_{kR}\bar{d}_j P_R \nu^c_i \\ &\quad - \tilde{l}_{iL}\bar{d}_k P_L u_j - \tilde{u}_{jL}\bar{d}_k P_L l_i - \tilde{d}_{kR}\bar{u}_j P_R l^c_i) + \mathrm{H.c.}\end{aligned} \tag{13}$$

with $P_{L,R} = (1 \mp \gamma_5)/2$.

The interactions of the Lagrangian (13) generate the terms in the effective CPV Lagrangians (10) and (11) at certain orders in the λ'-couplings. It is straightforward to derive the corresponding contribution to the 4-quark contact terms (11). It arises from a tree level contribution induced by sneutrino ($\tilde{\nu}$) exchange given by

$$\mathcal{L}^{\rm CPV}_{4q} = [C^P_{sd}(\bar{s}s) + C^P_{bd}(\bar{b}b)](\bar{d}i\gamma_5 d) \tag{14}$$

with

$$C^P_{sd} = \sum_i \frac{\mathrm{Im}(\lambda'_{i22}\lambda'^*_{i11})}{2m^2_{\tilde{\nu}(i)}}, \qquad C^P_{bd} = \sum_i \frac{\mathrm{Im}(\lambda'_{i33}\lambda'^*_{i11})}{2m^2_{\tilde{\nu}(i)}}, \tag{15}$$

where $m_{\tilde{\nu}}$ is the sneutrino mass. Note, that the four-quark term involving only d-quarks is absent in (14) due to $\mathrm{Im}(\lambda'_{i11}\lambda'^*_{i11}) \equiv 0$.

The interactions of the Lagrangian (13) generate the quark EDMs, d_q, and CEDMs, $\tilde{d}_q$, starting from 2-loops

[9,10] and the dominant $\not{R}_p$-contributions are of second order in the λ'-couplings. It was shown in Ref. [10] that the up-quark EDM and CEDM are suppressed by the light quark mass and mixing angles, which, therefore, can be neglected. The quark EDMs are irrelevant for our study based on the pion-exchange model with the interaction Lagrangian (9). We also do not consider the Weinberg term, which does not appear at the order of $O(\lambda'^2)$ unlike the quark CEDMs and 4-quark contact terms. In our analysis we use for the d-quark CEDMs the 2-loop result of Ref. [10]. The dominant contribution coming from the virtual b-quark takes the form:

$$\tilde{d}_k = 6.2 \times 10^{-7}\ (\mathrm{GeV}^{-1}) \sum_i \mathrm{Im}(\lambda'_{i33}\lambda'^{*}_{i11}) \mathcal{F}\left(\frac{m_b^2}{m_{\tilde{\nu}(i)}^2}\right), \tag{16}$$

where $k = 1, 2$ and $d_1 \equiv d_d$, $d_2 \equiv d_s$. The scaling factor $\mathcal{F}$ originates from the loop integration and can be written as

$$\mathcal{F}(\tau) = \frac{F(\tau)}{F(\tau_{300})}, \qquad F(\tau) = \tau\left[\frac{\pi^2}{3} + 2 + \ln\tau + (\ln\tau)^2\right], \tag{17}$$

where $\tau_{300} = (m_b/300\ \mathrm{GeV})^2$ and $m_b = 4.5$ GeV. For convenience the scaling factor $\mathcal{F}$ is normalized as $\mathcal{F}(\tau_{300}) = 1$ which corresponds to $m_{\tilde{\nu}} = 300$ GeV.

IV. EDMS IN $\not{R}_p$SUSY: HADRONIC LEVEL

In order to link the CP-violation at the quark and hadronic levels we have to relate the parameters of the Lagrangian in Eq. (9) to the quark CEDMs, $\tilde{d}_q$, and the CPV 4q-couplings C_{ij}^P. Towards this end we apply the standard matching of the quark-level (10) and (11) and hadronic-level (9) Lagrangians. This allows us to express the CP-odd meson-baryon couplings in terms of the quark CEDMs [7] and the CPV 4q-couplings C_{ij}^P as

$$\bar{g}^{(0)}_{\pi NN} = \frac{\langle \bar{u}u - \bar{d}d\rangle}{2F_\pi}\left\{A_u + A_d - \frac{\varepsilon}{3\sqrt{3}}(A_u - A_d)\right\}, \tag{18}$$

$$\bar{g}^{(1)}_{\pi NN} = \frac{\langle \bar{u}u + \bar{d}d\rangle}{2F_\pi}\left\{A_u - A_d - \frac{\varepsilon}{\sqrt{3}}(A_u + A_d)\right\} + \frac{\langle \bar{s}s\rangle}{2F_\pi}\frac{4\varepsilon}{\sqrt{3}}A_s - F_\pi\frac{M_\pi^2}{2m_d}(C_{sd}^P\langle\bar{s}s\rangle + C_{bd}^P\langle\bar{b}b\rangle), \tag{19}$$

$$\bar{g}^{(2)}_{\pi NN} = \frac{\langle \bar{u}u - \bar{d}d\rangle}{2F_\pi}\frac{\varepsilon}{3\sqrt{3}}(A_u - A_d), \qquad \bar{g}_{K^+n\Sigma^-} = \frac{\langle \bar{s}s - \bar{d}d\rangle}{2F_\pi}(A_u + A_s), \cdots \tag{20}$$

If the Peccei-Quinn (PQ) symmetry is imposed the parameters A_q are expressed through the quark CEDMs $\tilde{d}_q$ as $A_q = -0.27\tilde{d}_q$ GeV. Here $\langle\bar{q}q\rangle \equiv \langle p|\bar{q}q|p\rangle$ are the scalar quark condensates in the proton [7,22,23]:

$$\langle\bar{u}u\rangle = 3.5, \qquad \langle\bar{d}d\rangle = 2.8, \qquad \langle\bar{s}s\rangle = (0.64 \div 3.9), \qquad \langle\bar{b}b\rangle = 9\times 10^{-3}. \tag{21}$$

Note that the values of strange and bottom condensates in the nucleon are subject to significant uncertainties. In our analysis we use the estimates from Refs. [22,23]. For the value of $\langle\bar{s}s\rangle$ we indicate the interval according to Ref. [22]. For $\langle\bar{b}b\rangle$ we only need an order of magnitude estimate since it is associated with the subdominant term not essential for our analysis.

Now we are in the position to calculate the diagrams which contribute to the baryon EDMs (see Fig. 1). The calculation of the Feynman diagrams in Fig. 1 is straightforward and discussed before e.g. in Refs. [14,15]. The diagrams in Fig. 1(a) and 1(b) contribute to the chiral logarithms [24], the constant terms plus higher-order terms which can be neglected in the chiral expansion: $\sim[\log(m_B/M_P) - 1 + O(M_P)]$. The diagrams in Fig. 1(c) and 1(d) are dominated by the constant terms in the chiral expansion: $\sim[1 + O(M_P)]$. Finally, the diagrams in Fig. 1(e) and 1(f) cancel each other. For the neutral baryons (n, Λ, Σ^0 and Ξ^0) both sets of diagrams in Fig. 1(a)–1(d) contribute in such a way that the constant terms cancel each other. This is not the case for the EDMs of charged baryons where both chiral logarithms and constant terms give a nontrivial contribution. In the case of the neutron EDM the leading-order contributions of the diagrams in Figs. 1(a)–1(d) in terms of CP-even and CP-odd meson-baryon coupling constants are [7]:

$$d_n^{1(a+b)} = \frac{eg_{\pi NN}\bar{g}_{\pi NN}}{4\pi^2 m_B}\left[\log\frac{m_B}{M_\pi} - 1\right] - \frac{eg_{KN\Sigma}\bar{g}_{K^+n\Sigma^-}}{4\pi^2 m_B}\left[\log\frac{m_B}{M_K} - 1\right] \tag{22}$$

and

$$d_n^{1(c+d)} = \frac{eg_{\pi NN}\bar{g}_{\pi NN}}{4\pi^2 m_B} - \frac{eg_{KN\Sigma}\bar{g}_{K^+n\Sigma^-}}{4\pi^2 m_B}, \tag{23}$$

where

$$g_{\pi NN} = \frac{m_B}{F_\pi}(D + F), \qquad g_{KN\Sigma} = \frac{m_B}{F_\pi}(D - F), \tag{24}$$

$$\bar{g}_{\pi NN} = \bar{g}^{(0)}_{\pi NN} + \bar{g}^{(2)}_{\pi NN}. \tag{25}$$

The complete result for the neutron EDM is:

$$d_n = \frac{eg_{\pi NN}\bar{g}_{\pi NN}}{4\pi^2 m_B}\log\frac{m_B}{M_\pi} + \frac{eg_{KN\Sigma}\bar{g}_{K^+n\Sigma^-}}{4\pi^2 m_B}\log\frac{m_B}{M_K}. \tag{26}$$

Substituting the expressions of the baryon-meson couplings Eqs. (18)–(20) in terms of the parameters of the ChPT Lagrangian and quark CEDMs [7] we have

$$d_n = \beta\tilde{d}^+_{ud}c^-_{ud}(D+F)\log\frac{m_B}{M_\pi} + \beta\tilde{d}^+_{us}c^-_{ds}(D-F)\log\frac{m_B}{M_K} \tag{27}$$

where $\tilde{d}^{\pm}_{q_1q_2} = \tilde{d}_{q_1} \pm \tilde{d}_{q_2}$, $c^{\pm}_{q_1q_2} = \langle\bar{q}_1q_1 \pm \bar{q}_2q_2\rangle$ and $\beta = 0.27e/(8\pi^2F^2_\pi)$. In Eq. (26) we neglected the pion-eta meson mixing ($\varepsilon = 0$).

In a similar way we calculate the EDMs of other baryons. Here we indicate the final results of these calculations:

$$\begin{aligned} d_p = & -\beta\tilde{d}^+_{ud}c^-_{ud}(D+F)\log\frac{m_B}{M_\pi} - \beta\tilde{d}^+_{us}\Big[Fc^-_{us} + \frac{D}{3}(c^-_{ud} - c^-_{ds})\Big]\log\frac{m_B}{M_K} \\ & + \frac{\beta}{3}\tilde{d}_u[D(5c^-_{ud} + c^+_{us}) + 3F(2c^-_{ud} + 2c^-_{us} + c^+_{ds})] + \frac{\beta}{3}\tilde{d}_d[D(5c^-_{ud} - 5c^-_{ds} - 2c^+_{us}) + 3Fc^-_{us}] \\ & + \frac{\beta}{3}\tilde{d}_s[D(c^+_{us} - 5c^-_{ds}) + 3F(2c^-_{us} + 2c^-_{ds} - c^+_{ud})], \end{aligned} \tag{28}$$

$$d_\Lambda = -d_{\Sigma^0} = \frac{\beta}{2}\tilde{d}^+_{us}[Dc^-_{us} + F(c^-_{ud} - c^-_{ds})]\log\frac{m_B}{M_K}, \tag{29}$$

$$d_{\Xi^0} = \beta\tilde{d}^+_{ud}c^-_{ds}(D-F)\log\frac{m_B}{M_\pi} + \beta\tilde{d}^+_{us}c^-_{ud}(D+F)\log\frac{m_B}{M_K}. \tag{30}$$

As it is known [7] the proton and neutron contributions to the deuteron EDM cancel out in leading order of the chiral expansion in the $SU(2)$ version of ChPT [25]. However it does not hold in the $SU(3)$ extension [7]. Therefore, the contribution of the strangeness sector to the deuteron EDM becomes important. Note, in the final expressions for the neutron and proton EDMs, Eqs. (27), we disagree with the results of Ref. [7] by a factor 2. We discuss this issue in the Appendix.

The deuteron EDM also receives a contribution from P- and T-odd proton-neutron forces generated by π- and η-meson exchange between two nucleons with one CP-even and one CP-odd vertex. With the corresponding potential one can calculate the d^{NN}_D term in Eq. (5) as the expectation value of $e\mathbf{r}/2$, where $\mathbf{r}$ is the relative proton-neutron coordinate. In this way one obtains the following result [17,18]:

$$d^{NN}_D = -\frac{eg_{\pi NN}\bar{g}^{(1)}_{\pi NN}}{12\pi m_\pi}\frac{1+\xi}{(1+2\xi)^2}, \tag{31}$$

where $\xi = \sqrt{M_NE_B}/m_\pi$ and $E_B = 2.23$ MeV is the deuteron binding energy. Expressing the CP-even and CP-odd pion-nucleon couplings in terms of the CEDMs and parameters of the chiral Lagrangian we get:

$$\begin{aligned} d^{NN}_D \simeq & \gamma^{(1)}_D\Big[\tilde{d}^-_{ud}c^+_{ud} + \frac{4\varepsilon}{\sqrt{3}}\tilde{d}_s\langle\bar{s}s\rangle\Big] - \gamma^{(2)}_D[C^P_{sd}\langle\bar{s}s\rangle \\ & + C^P_{bd}\langle\bar{b}b\rangle] \end{aligned} \tag{32}$$

where

$$\begin{aligned} \gamma^{(1)}_D &= 0.13(F+D)\frac{em_B}{24\pi M_\pi F^2_\pi}, \\ \gamma^{(2)}_D &= 0.48(F+D)\frac{em_BM_\pi}{24\pi m_d}. \end{aligned} \tag{33}$$

Recently, the Schiff moment $S_{\rm Hg}$ [19] has been calculated within a reliable nuclear structure model which takes full account of core polarization on the basis of the P- and T-odd one-pion-exchange potential. Note that in Ref. [7] it was shown that the contribution of the eta-meson exchange to the Schiff moment is suppressed by a factor 10^{-3} with respect to the pion exchange. The result for the Schiff moment, taking into account a finite range interaction and the core polarization effect is

$$\begin{aligned} S_{\rm Hg} = & -0.055g_{\pi NN}\{0.007\bar{g}^{(0)}_{\pi NN} + \bar{g}^{(1)}_{\pi NN} \\ & - 0.16\bar{g}^{(2)}_{\pi NN}\}\ e\cdot{\rm fm}^3. \end{aligned} \tag{34}$$

Therefore, only the isovector channel is sufficient for the evaluation of the Schiff moment [7,19]. In terms of the quark EDMs and ChPT parameters the latter is given by

$$\begin{aligned} S_{\rm Hg} \simeq & \Big\{\gamma^{(1)}_{Hg}\Big[\tilde{d}^-_{ud}c^+_{ud} + \frac{4\varepsilon}{\sqrt{3}}\tilde{d}_s\langle\bar{s}s\rangle\Big] \\ & - \gamma^{(2)}_{Hg}[C^P_{sd}\langle\bar{s}s\rangle + C^P_{bd}\langle\bar{b}b\rangle]\Big\}\ e\cdot{\rm fm}^3, \end{aligned} \tag{35}$$

where

$$\begin{aligned} \gamma^{(1)}_{\rm Hg} &= 0.015(F+D)\frac{m_B}{2F^2_\pi}, \\ \gamma^{(2)}_{\rm Hg} &= 0.055(F+D)\frac{m_BM^2_\pi}{2m_d}. \end{aligned} \tag{36}$$

V. CONSTRAINTS ON $\not{R}_p$ SUSY FROM HADRONIC EDMS

Let us summarize the formulas for the considered hadronic EDMs in terms of the trilinear $\not{R}_p$-couplings. Using Eqs. (4), (5), (27)–(30), (32), and (35) we get the following expressions with numerical coefficients:

$$d_p = -10^{-20} \times \mathcal{F}\left(\frac{m_b^2}{m_{\tilde{\nu}(i)}^2}\right)[(1.67 \div 2.21)\,\mathrm{Im}(\lambda'_{i33}\lambda'^{*}_{i11}) + (0.23 \div 0.48)\,\mathrm{Im}(\lambda'_{i33}\lambda'^{*}_{i22})]\; e\cdot\mathrm{cm}, \quad (37)$$

$$d_n = 10^{-20} \times \mathcal{F}\left(\frac{m_b^2}{m_{\tilde{\nu}(i)}^2}\right)[0.82\,\mathrm{Im}(\lambda'_{i33}\lambda'^{*}_{i11}) + (-0.12 \div 0.23)\,\mathrm{Im}(\lambda'_{i33}\lambda'^{*}_{i22})]\; e\cdot\mathrm{cm}, \quad (38)$$

$$d_{\mathrm{Hg}} = 10^{-23} \times \mathcal{F}\left(\frac{m_b^2}{m_{\tilde{\nu}(i)}^2}\right)[11.4\,\mathrm{Im}(\lambda'_{i33}\lambda'^{*}_{i11}) + (0.28 \div 1.62)\,\mathrm{Im}(\lambda'_{i33}\lambda'^{*}_{i22})]\; e\cdot\mathrm{cm} - (0.90 \div 5.49) \times 10^{-23}\left(\frac{300\ \mathrm{GeV}}{m_{\tilde{\nu}(i)}}\right)^2 \times \mathrm{Im}(\lambda'_{i22}\lambda'^{*}_{i11})\; e\cdot\mathrm{cm}, \quad (39)$$

$$d_D = 10^{-20} \times \mathcal{F}\left(\frac{m_b^2}{m_{\tilde{\nu}(i)}^2}\right)[(11.79 \div 12.34)\,\mathrm{Im}(\lambda'_{i33}\lambda'^{*}_{i11}) + (-0.41 \div 0.03)\,\mathrm{Im}(\lambda'_{i33}\lambda'^{*}_{i22})]\; e\cdot\mathrm{cm} - (0.4 \div 2.5) \times 10^{-20}\left(\frac{300\ \mathrm{GeV}}{m_{\tilde{\nu}(i)}}\right)^2 \times \mathrm{Im}(\lambda'_{i22}\lambda'^{*}_{i11})\; e\cdot\mathrm{cm}, \quad (40)$$

$$d_\Lambda = -d_{\Sigma^0} = (0.08 \div 0.25) \times 10^{-20}\mathcal{F}\left(\frac{m_b^2}{m_{\tilde{\nu}(i)}^2}\right)\,\mathrm{Im}(\lambda'_{i33}\lambda'^{*}_{i22})\; e\cdot\mathrm{cm}, \quad (41)$$

$$d_{\Xi^0} = 10^{-20}\mathcal{F}\left(\frac{m_b^2}{m_{\tilde{\nu}(i)}^2}\right)[(-0.35 \div 0.69)\,\mathrm{Im}(\lambda'_{i33}\lambda'^{*}_{i11}) + 0.28\,\mathrm{Im}(\lambda'_{i33}\lambda'^{*}_{i22})]\; e\cdot\mathrm{cm}. \quad (42)$$

In the above equations the summation over $i = 1, 2, 3$ is implied. The uncertainties in the coefficients are due to the variation of the strange quark sea in the proton. Note that the contribution from $\mathrm{Im}(\lambda'_{i22}\lambda'^{*}_{i11})$ appears solely via the 4-quark CPV interactions (14).

Comparing Eqs. (38) and (39) with the corresponding experimental bounds Eqs. (1) and (2) we derive constraints on the imaginary parts of the products of $\not{R}_p$-couplings given in Table I. In the last column of Table I we also show for comparison the existing limits on $|\lambda'_{i33}\lambda'^{*}_{i11}|$, $|\lambda'_{i22}\lambda'^{*}_{i11}|$ and $|\lambda'_{i33}\lambda'^{*}_{i22}|$ [26]. It is seen that the presently most stringent limits on $|\mathrm{Im}(\lambda'_{ikk}\lambda'^{*}_{i11})|$, $|\mathrm{Im}(\lambda'_{i33}\lambda'^{*}_{i22})|$ come from the ^{199}Hg atom EDM (2). The forthcoming experiments on the deuteron EDM (3) are going to improve these limits by about one to 3 orders of magnitude. Note, that we obtained about 1-order of magnitude improvement for the limit $|\mathrm{Im}(\lambda'_{i33}\lambda'^{*}_{i11})| \leq 1.2 \times 10^{-5}$ previously derived in Ref. [10] from the neutron EDM constraint (1) on the basis of the $SU(6)$ quark model. The existing limits on the absolute values of the corresponding products do not exclude the values of $|\mathrm{Im}(\lambda'_{ikk}\lambda'^{*}_{i11})|$, $|\mathrm{Im}(\lambda'_{i33}\lambda'^{*}_{i22})|$ within the limits derived from EDMs. Using the limits from Table I we can predict on the basis of Eqs. (41) and (42) for the EDMs of neutral light hyperons the following upper limits:

$$|d_\Lambda| = |d_{\Sigma^0}| \leq 1.9 \times 10^{-25}\; e\cdot\mathrm{cm}, \qquad |d_{\Xi^0}| \leq 2.4 \times 10^{-25}\; e\cdot\mathrm{cm}$$

which might have some future phenomenological implications.

TABLE I. Upper limits on the imaginary parts of the products of the trilinear $\not{R}_p$-couplings derived from the experimental bounds on the EDMs of the neutron [1], the neutral ^{199}Hg atom [3] and the deuteron [4]. The existing constraints from other experiments on the absolute values of the corresponding products of $\not{R}_p$-coupling are taken from Ref. [26]. The scaling factor $\mathcal{F}$ is defined in Eq. (17) and takes the values $\mathcal{F} = 1$, 0.34 and 0.15 for $m_{\tilde{\nu}} = 300$ GeV, 600 GeV and 1 TeV, respectively.

Couplings	d_n [1]	d_{Hg} [3]	d_D [4]	Existing limits [26]
$\lvert\mathrm{Im}(\lambda'_{i33}\lambda'^{*}_{i11})\rvert \cdot \mathcal{F}(\frac{m_b^2}{m_{\tilde{\nu}(i)}^2})$	$\leq 3.6 \times 10^{-6}$	$\leq 1.8 \times 10^{-6}$	$\leq (0.8 \div 2.5) \times 10^{-8}$	$\lvert\lambda'_{133}\lambda'_{111}\rvert \leq 4.5 \times 10^{-5}$ $\lvert\lambda'_{233}\lambda'_{211}\rvert \leq 5.4 \times 10^{-3}$ $\lvert\lambda'_{333}\lambda'_{311}\rvert \leq 1.3 \times 10^{-3}$
$\lvert\mathrm{Im}(\lambda'_{i22}\lambda'^{*}_{i11})\rvert \cdot (\frac{300\ \mathrm{GeV}}{m_{\tilde{\nu}(i)}})^2$		$\leq (0.4 \div 2.3) \times 10^{-5}$	$\leq (0.4 \div 7.5) \times 10^{-7}$	$\lvert\lambda'_{122}\lambda'_{111}\rvert \leq 4.5 \times 10^{-5}$ $\lvert\lambda'_{222}\lambda'_{211}\rvert \leq 1.3 \times 10^{-3}$ $\lvert\lambda'_{322}\lambda'_{311}\rvert \leq 1.3 \times 10^{-3}$
$\lvert\mathrm{Im}(\lambda'_{i33}\lambda'^{*}_{i22})\rvert \cdot \mathcal{F}(\frac{m_b^2}{m_{\tilde{\nu}(i)}^2})$	$\leq (1.3 \div 2.5) \times 10^{-5}$	$\leq (1.3 \div 7.5) \times 10^{-5}$	$\leq 2.4 \times 10^{-7}$	$\lvert\lambda'_{133}\lambda'_{122}\rvert \leq 4.0 \times 10^{-5}$ $\lvert\lambda'_{233}\lambda'_{222}\rvert \leq 2.5 \times 10^{-3}$ $\lvert\lambda'_{333}\lambda'_{322}\rvert \leq 3.0 \times 10^{-3}$

VI. SUMMARY

We have studied the contributions of the trilinear $\not{R}_p$-couplings to the EDMs of ^{199}Hg atom, deuteron, nucleon and neutral hyperons within the $SU(3)$ ChPT, applying the meson-exchange model of CPV nuclear forces. We have analyzed the $\not{R}_p$-contributions via the d-quark CEDM and CPV 4-quark interactions. We have shown that the latter contribute only to the nuclear EDMs via the CPV nuclear forces and are irrelevant for the EDMs of the nucleon and neutral hyperons. We have also found that these two mechanisms give rise to a dependence of the hadronic EDMs proportional to different λ'-couplings. Therefore, taking into account both mechanism allows one to obtain a complimentary information on the imaginary parts of the products of the λ'-couplings. The corresponding upper limits from the null experimental results on measurements of the above mentioned hadronic EDMs are given in Table I. On the basis of the derived constraints on the trilinear $\not{R}_p$-couplings we have given predictions for the EDMs of neutral hyperons which might have some phenomenological implications in future.

We have also demonstrated that the present limits from the ^{199}Hg EDM experiments are by a factor ~ 6 more stringent than those from the experiments on the neutron EDM and that the planned storage ring experiments with the deuterium ions would be able to significantly improve these limits.

ACKNOWLEDGMENTS

This work was supported in part by the FONDECYT project 1030244, by the DFG under the contracts No. FA67/31-1 and GRK683. This research is also part of the EU Integrated Infrastructure Initiative Hadronphysics project under the contract number RII3-CT-2004-506078 and President grant of Russia "Scientific Schools" No. 5103.2006.2.

APPENDIX: ON NEUTRON AND PROTON EDMS

As we mentioned in Sec. IV our final expressions, Eqs. (27) and (28), for the neutron and proton EDMs disagree with the results of Ref. [7] by a factor of 2. In order to check these results we compare them with the well-known result of chiral perturbation theory in the two-flavor scheme, involving only pion loops. In this case all chiral approaches (see Refs. [14,15,24]) give the same model-independent expression for the leading-order neutron and proton EDMs in the chiral expansion, the so-called "chiral logarithm". Neglecting kaon loops in our formulas Eqs. (27) and (28) we reproduce the result of the chiral approaches:

$$d_n = -d_p = \frac{e g_{\pi NN} \bar{g}_{\pi NN}}{4\pi^2 m_p} \log\frac{m_p}{M_\pi}. \tag{A1}$$

On the contrary, the result of Ref. [7] given in their Eq. (50)

$$d_n = -d_p = \frac{e}{4\pi^2 F_\pi^2}(D+F)(A_u + A_d)(\langle \bar{u}u\rangle - \langle \bar{d}d\rangle)\log\frac{m_p}{M_\pi} \tag{A2}$$

differs from this formula by a factor 2. Indeed, using the expressions for $g_{\pi NN}$ and $\bar{g}_{\pi NN}$:

$$g_{\pi NN} = (D+F)\frac{m_p}{F_\pi}, \qquad \bar{g}_{\pi NN} = (A_u + A_d)\frac{\langle \bar{u}u\rangle - \langle \bar{d}d\rangle}{2F_\pi} \tag{A3}$$

one can rewrite Eq. (A2) in the form

$$d_n = -d_p = \frac{e g_{\pi NN} \bar{g}_{\pi NN}}{2\pi^2 m_p} \log\frac{m_p}{M_\pi}, \tag{A4}$$

which disagrees with Eq. (A1) by the factor 2 in the denominator.

[1] P. G. Harris *et al.*, Phys. Rev. Lett. **82**, 904 (1999).
[2] C. A. Baker *et al.*, Phys. Rev. Lett. **97**, 131801 (2006).
[3] M. V. Romalis, W. C. Griffith, J. P. Jacobs, and E. N. Fortson, Phys. Rev. Lett. **86**, 2505 (2001).
[4] Y. K. Semertzidis *et al.* (EDM Collaboration), AIP Conf. Proc. **698**, 200 (2004).
[5] I. B. Khriplovich and S. K. Lamoreaux, *CP Violation Without Strangeness. Electric Dipole Moments of Particles, Atoms, and Molecules* (Springer-Verlag, Berlin, 1997), p. 230.
[6] T. Falk and K. A. Olive, Phys. Lett. B **375**, 196 (1996); **439**, 71 (1998); T. Ibrahim and P. Nath, Phys. Lett. B **418**, 98 (1998); Phys. Rev. D **57**, 478 (1998); **58**, 019901(E) (1998); **60**, 079903(E) (1999); **60**, 119901(E) (1999); **58**, 111301 (1998); **60**, 099902(E) (1999); M. Brhlik, G. J. Good, and G. L. Kane, Phys. Rev. D **59**, 115004 (1999); M. Brhlik, L. Everett, G. L. Kane, and J. Lykken, Phys. Rev. Lett. **83**, 2124 (1999); A. Bartl, T. Gajdosik, W. Porod, P. Stockinger, and H. Stremnitzer, Phys. Rev. D **60**, 073003 (1999); S. Pokorski, J. Rosiek, and C. A. Savoy, Nucl. Phys. **B570**, 81 (2000); T. Falk, K. A. Olive, M. Pospelov, and R. Roiban, Nucl. Phys. **B560**, 3 (1999); V. Barger, T. Falk, T. Han, J. Jiang, T. Li, and T. Plehn, Phys. Rev. D **64**, 056007 (2001); S. Abel, S. Khalil, and O. Lebedev, Phys. Rev. Lett. **86**, 5850 (2001); Nucl. Phys. **B606**, 151 (2001).
[7] J. Hisano and Y. Shimizu, Phys. Rev. D **70**, 093001 (2004).
[8] R. Barbieri and A. Masiero, Nucl. Phys. **B267**, 679 (1986).

[9] R. M. Godbole, S. Pakvasa, S. D. Rindani, and X. Tata, Phys. Rev. D **61**, 113003 (2000); S. A. Abel, A. Dedes, and H. K. Dreiner, J. High Energy Phys. 05 (2000) 013.
[10] D. Chang, W. F. Chang, M. Frank, and W. Y. Keung, Phys. Rev. D **62**, 095002 (2000).
[11] P. Herczeg, Phys. Rev. D **61**, 095010 (2000).
[12] A. Faessler, T. Gutsche, S. Kovalenko, and V. E. Lyubovitskij, Phys. Rev. D **73**, 114023 (2006).
[13] S. Weinberg, Physica A **96**, 327 (1979); J. Gasser and H. Leutwyler, Ann. Phys. (N.Y.) **158**, 142 (1984); Nucl. Phys. **B250**, 465 (1985).
[14] A. Pich and E. de Rafael, Nucl. Phys. **B367**, 313 (1991).
[15] B. Borasoy, Phys. Rev. D **61**, 114017 (2000).
[16] V. A. Dzuba, V. V. Flambaum, J. S. M. Ginges, and M. G. Kozlov, Phys. Rev. A **66**, 012111 (2002).
[17] I. B. Khriplovich and R. A. Korkin, Nucl. Phys. A **665**, 365 (2000).
[18] V. F. Dmitriev and I. B. Khriplovich, Phys. Rep. **391**, 243 (2004).
[19] V. F. Dmitriev, R. A. Sen'kov, and N. Auerbach, Phys. Rev. C **71**, 035501 (2005).
[20] V. M. Khatsimovsky, I. B. Khriplovich, and A. S. Yelkhovsky, Ann. Phys. (N.Y.) **186**, 1 (1988); X. G. He and B. McKellar, Phys. Lett. B **390**, 318 (1997).
[21] D. A. Demir, O. Lebedev, K. A. Olive, M. Pospelov, and A. Ritz, Nucl. Phys. **B680**, 339 (2004).
[22] A. Corsetti and P. Nath, Phys. Rev. D **64**, 125010 (2001).
[23] A. Faessler, T. Gutsche, S. Kovalenko, V. E. Lyubovitskij, and I. Schmidt, Phys. Rev. D **72**, 075006 (2005).
[24] R. J. Crewther, P. Di Vecchia, G. Veneziano, and E. Witten, Phys. Lett. B **88**, 123 (1979); **91**, 487(E) (1980).
[25] O. Lebedev, K. A. Olive, M. Pospelov, and A. Ritz, Phys. Rev. D **70**, 016003 (2004).
[26] M. Chemtob, Prog. Part. Nucl. Phys. **54**, 71 (2005); R. Barbier, C. Berat, M. Besancon, M. Chemtob, A. Deandrea, E. Dudas, P. Fayet, S. Lavignac, G. Moreau, E. Perez, and Y. Sirois, Phys. Rep. **420**, 1 (2005).