Colored scalars and the neutron electric dipole moment

Svjetlana Fajfer

Department of Physics, University of Ljubljana,
Jadranska 19, 1000 Ljubljana, Slovenia and
J. Stefan Institute, Jamova 39, P. O. Box 3000, 1001 Ljubljana, Slovenia

Jan O. Eeg

Department of Physics, University of Oslo,
P.O.Box 1048 Blindern, N-0316 Oslo, Norway

(Dated: March 21, 2014)

We investigate new contributions to the neutron electric dipole moment induced by colored scalars. As an example, we evaluate contributions coming from the color octet, weak doublet scalar, accommodated within a modified Minimal Flavor Violating framework. These flavor non-diagonal couplings of the color octet scalar might account for the measured asymmetry $a_{CP}(D^0 \rightarrow K^-K^+) - a_{CP}(D^0 \rightarrow \pi^+\pi^-)$ at tree level. The same couplings constrained by this asymmetry also induce two-loop contributions to the neutron electric dipole moment. We find that the direct CP violating asymmetry in neutral $D$-meson decays is more constraining on the allowed parameter space than the current experimental bound on neutron electric dipole moment. We comment also on contributions of higher dimensional operators to the neutron electric dipole moment within this framework.

I. INTRODUCTION

The neutron electric dipole moment (NEDM) plays a special role in current searches of physics beyond the Standard Model (SM). The low energy flavor physics puts extremely tight bounds on possible non-standard model contributions to the NEDM, and in particular special attention has been paid to new sources of CP violation. The NEDM gives a great

*Electronic address: svjetlana.fajfer@ijs.si
†Electronic address: j.o.eeg@fys.uio.no
opportunity to learn about additional sources of CP violation. There are many studies on the NEDM (for a review see [1]). In addition to CP violating effects, the NEDM still offers many puzzles for the study of non-perturbative QCD effects within the SM.

Measurements of the CP asymmetry in $D \to K^+K^-/\pi^+\pi^-$ larger than SM expectations have attracted many theoretical studies. The LHCb collaboration recently updated their analysis leading to a decreased value of the world average CP asymmetry [2–4]:

$$\Delta a_{CP} = (-0.329 \pm 0.121)\%,$$

with $\Delta a_{CP} = a_{K^+K^-} - a_{\pi^+\pi^-}$ and the definition

$$a_f \equiv \frac{\Gamma(D^0 \to f) - \Gamma(D^0 \to f)}{\Gamma(D^0 \to f) + \Gamma(D^0 \to f)}.$$  

Many theoretical studies were performed in order to explain the apparent discrepancy [5–14]. Some of these approaches have explained the observed asymmetry by SM effects [13, 14], while in the rest of them, possible new physics (NP) effects were considered. In ref. [14] the authors argued that penguin contraction power corrections might significant enhance the decay amplitudes. They also found that the same mechanism might consistently explain the branching ratios for singly Cabibbo-suppressed $D \to PP$ decays. The authors of ref. [5] pointed out that most likely the effective operators explaining CP asymmetry in charm decays are color-magnetic dipole operators, although color octet scalar as well as two Higgs doublet models in specific parameter space are still allowed options [15]. The important result of these studies [5–12, 16] is that apparently one needs an additional source of CP violation and in particular in the charm sector. On the other hand, CP violation in the $B$ system has been related to the NEDM [17], while in [15] CP violation in $D^0 - \bar{D}^0$ oscillation was related to the NEDM. In the study of the NEDM usually the lowest dimensional operators were considered [1]. Recently, the authors of [18] found that higher dimensional operators might lead to rather important contributions to NEDM, as noticed some time ago in [19, 20].

The authors of [18] also noticed [21] that CP violation in charm decays can induce an increase of the NEDM, if these higher dimensional operators are included. Experimental results agree remarkably well with SM predictions in the case of the $K$ and $B_{d,s}$ systems. The Minimal Flavor Violation (MFV) framework [17, 22], in which flavor-changing transitions in the quark sector are entirely described by two quark Yukawa couplings, has been used
to parametrize NP effects. The MFV principle does not forbid appearance of new flavor
blind CP violating phases in addition to the unique phase of the CKM matrix. The Flavor
Changing Neutral Currents (FCNC) coming from down-like quark sectors are very well
known and constrained by experiment. In addition to the well understood \( s \to d \) transition
within K meson physics, rare \( B \) decays are excellent probes of NP beyond the SM. The
recent model independent study of ref. [23] has been aimed to constrain NP contributions
in rare \( B \) meson decays. Among the relevant observables used in \( B_s \to \mu^+\mu^- \), \( B \to K\mu^+\mu^- \),
\( B \to K^*\mu^+\mu^- \), and \( B \to X_s\gamma \) decays, the measured CP asymmetry in \( B \to X_s\gamma \) is specially
informative. Although the hadronic uncertainties are still large, the \( A_{CP}(B \to X_s\gamma) \) might
constrain CP violating NP contributions. Namely, the measured branching ratios are not
itself constraining enough for the possible CP violating NP effects. The strongest bound on
new CP violating parameters comes from \( b \to s\gamma \) decay [24]: \( A_{CP}^{\text{exp}}(b \to s\gamma) = (-1.2 \pm 2.8)\% \).

Using the constraint from charm decays [1], we consider contributions of the additional
heavy colored scalar meson to the NEDM. As a particular example we consider a color octet,
weak doublet scalar introduced in ref. [25]. These authors have analyzed a model with the
most general scalar structure of the standard model (SM). The main request on the model
in [25] was that it maintains the smallness of flavor changing neutral currents (FCNC), or
in another words it is a model which fulfills principles of MFV. The new scalar state might
contain a new source of CP violation. The phenomenology of SM modified by the presence
of the color octet scalar state was investigated in [25, 26]. The constraints on of the color
octet scalar coupling to up-type of quarks that arises from precision electroweak data on \( R_b \)
in \( Z \to b\bar{b} \) decays, have been considered in [26]. This study is not constraining the mass
of the color-octet. The NEDM gets its largest additional contribution within the proposed
model [25] by the color electric dipole moment (CEDM) of the \( b \)-quark. When integrated
out, the CEDM of the \( b \)-quark [27] induces Weinberg’s three-gluon CP violating operator
[28]. The electric dipole operator of the \( d \)-quark also induces a NEDM, as discussed in [29].
Namely, the existence of two different couplings of \( d \)-quark with the \( u \)-quark types lead to a
one-loop penguin-like contribution, and it can then be present in the case of \( d \)-quark EDM.

The effective Lagrangian for a dipole moment of a fermion \( f \) has the generic form

\[
\mathcal{L}_{\text{eff}} = \frac{i}{2} d_f \bar{f} \sigma_{\mu\nu} F^{\mu\nu} \gamma_5 f , \tag{3}
\]

where \( d_f \) is the EDM of the fermion, \( f \) is the fermion field, \( F^{\mu\nu} \) is the electromagnetic field
and \( \sigma_{\mu
u} = i[\gamma_\mu, \gamma_\nu]/2 \). Contributions to the NEDM come from EDMs of single quarks, and contributions due to interplay of quarks within the neutron. In the valence approximation, the contributions to NEDM from EDMs of single quarks \((d_q)\) are given by

\[
d_n = \frac{4}{3} d_d - \frac{1}{3} d_u.
\]

(4)

The existing experimental bound from ref. [30] for the NEDM is:

\[
d_n^{\exp} \leq 2.9 \times 10^{-26} \text{e cm},
\]

(5)

which bounds the appropriate imaginary parts of couplings of quarks with colored scalars. Motivated by the explanation of direct CP violation in decay of the neutral \(D\)-meson by the presence of a color octet, we investigate the impact of the imaginary couplings of quarks to the color scalars on the NEDM. The color octet scalar can be searched at hadron colliders in di-jet events [25, 31–33]. Current experimental searches at LHC based on the dijet analysis at ATLAS [34] and CMS [35], as well as former at Tevatron excludes existence of a color octet scalar with mass below 1.86 TeV, while four-jet CMS searches do not observe it at a low-energy region [36] from 250 – 740 GeV. Also, ATLAS searches exclude it in low energy region [37]. However, if the color octet scalar decays in more than two light quarks (two jets), top-quark and jet, or \(t\bar{t}\), then the existing bound on the color octet scalar mass would be different. Then even masses of order 400 GeV can not be excluded, as noticed in [38].

In Sec. II we first remind on the NEDM within the SM. Sec. III is devoted to the study of NEDM contributions obtained within a modified color octet model [39], which produces CP violation in charm at tree level. In Sec. IV we discuss obtained results.

II. NEDM IN THE SM

The NEDM has been studied for many years, both within and beyond the SM (for a review see [1]). As it turned out that EDMs in the SM were small, calculations within new physics scenarios were also performed. To obtain a CP violating amplitude within the SM, two weak interactions are needed, and at least one of these must be a penguin-like interaction. The EDMs of single quarks, which are three loop diagrams with double GIM-cancellations and at least one gluon exchange, are of order \(\alpha_s G_F^2\) and proportional to quark
masses and an imaginary CKM factor. A typical diagram is shown in Fig. 1. However, in the SM, EDMs of single quarks were found to be very small, namely of order $10^{-34} \, e \, cm$ according to studies of refs. [40, 41].

It was shown that a mechanism with interplay of weak and strong interactions would give a nonzero contribution. Typically, such amplitudes were written as baryon poles with two weak interactions, an ordinary W-exchange and a CP-violating penguin interaction, with for instance a negative parity strange baryon as intermediate baryon, and a soft photon emitted from somewhere [42–44]. There were also contributions due to amplitudes at baryonic/mesonic level where the photon was emitted from an intermediate pion or kaon within a chiral loop [45, 46]. Within such mechanisms, the $d_n/e$ was estimated to be of order $10^{-33}$ to $10^{-31} \, cm$. It was pointed out that if the pole diagrams were interpreted at quark level it would correspond to a “diquark mechanism” [19, 20]. This means a CP-violating two loop diagram for the quark process $ud \rightarrow du \gamma$, as shown in Fig. 2. Here the contribution was obtained in terms of a two loop factor proportional to $\alpha_s G_F^2$, and a CKM factor, written as $F_{CKM} = \text{Im}(V_{ub}V_{tb}V_{td},V_{ud})$. Such diagrams lead to an eight dimensional operator for the NEDM given by

$$Q_{di-quark} = F_{\alpha \beta} \epsilon^{\alpha \beta \mu \nu} \bar{u}_L \gamma_\mu d_L \, \bar{d}_L \gamma_\nu u_L ,$$

(6)
where $F_{\alpha\beta}$ is the electromagnetic tensor. The matrix element of this operator is found to be of order $(1 - 6)\times 10^{-3} \times m^3_N$, where $m_N$ is the nucleon mass. The obtained result is 

$$d_n/e \sim 10^{-33} \text{ to } 10^{-32} \text{ cm} \quad [19, 20, 47].$$

Within this mechanism the NEDM is suppressed by a small hadronic matrix element, but had logarithmic GIM instead of power-like, compensating for hadronic suppression.

Using the operator product expansion technique, one might consider operators of higher dimensions. Recently, it was found [18] that a tree-level higher dimensional operator might give a significant contribution to the NEDM comparable to the loop induced contributions. Their contribution is illustrated in Fig. 3.

The operator in (3) has dimension five and it is the lowest dimension operator in SM, if the $\theta$ problem in QCD is rotated away as described in ref. [1]. If the electromagnetic field is replaced by gluonic field $G^a_{\mu\nu}$ then one has also CEDMs of the same dimension. However, operators of higher dimension, as the well-known Weinberg operator [28, 48, 49]

$$L^W_{\text{eff}} = \frac{C_W}{6} f^{abc} G^{a}_{\mu\nu} \epsilon^{\nu\beta\rho\sigma} G^{b}_{\rho\sigma} G^{c}_{\mu\nu},$$

(7)

contribute significantly to the NEDM [1, 25]. An additional dimension six operator is

$$L^{ff'}_{\text{eff}} = C_{ff'} (\bar{\psi}_f \psi_f) (\bar{\psi}_{f'} i\gamma_5 \psi_{f'}),$$

(8)

which has been considered in ref. [1, 17]. There are also so-called Barr-Zee contributions [50] on the two loop level contributing to the NEDM within the SM, being competitive in size with the Weinberg’s three-gluon operator.
III. CONTRIBUTIONS FROM THE COLOR OCTET SCALARS WITHIN MODIFIED MFV FRAMEWORK

The extension of the SM introduced in [25, 51] allows couplings of a color octet, weak doublet state with weak hyper-charge \(1/2 (8, 2)_{1/2}\) to quarks written as:

\[
\mathcal{L} = -\sqrt{2} \eta_U \bar{u}^i_R \frac{m^i_U}{v} T^A u^i_L \phi^{A0} + \sqrt{2} \eta_U \bar{d}^i_R \frac{m^i_D}{v} T^A d^i_L \phi^{A+} - \sqrt{2} \eta_D \bar{d}^i_R \frac{m^i_D}{v} T^A V^\ast_{ij} u^i_L \phi^{A0} - \sqrt{2} \eta_D \bar{d}^i_R \frac{m^i_D}{v} T^A V^\ast_{ij} d^i_L \phi^{A+} + h.c. .
\]

The couplings \(\eta_U,D\) are universal complex numbers, \(T^A\) are color SU(3) generators, lower case roman letters denote mass eigenstate fields and \(\phi^{A+}\) and \(\phi^{A0}\) are charged and neutral component of the scalar color octet. Further, \(v\) is the vacuum expectation value (VEV) of the SM Higgs doublet and \(m^i_{U,D}\) are quark masses. The scalar octet state as a weak doublet might slightly modify precision electroweak parameter \(S\) as discussed in ref. [25]. The Higgs boson production and decay \(H \rightarrow \gamma\gamma\) were discussed [32, 52]. For masses of color octet scalar \(\sim 1\)TeV these modifications are negligible. On the other hand, low energy flavor physics gets contribution from virtual colored scalar states. The contribution proportional to the quadratic coupling \(\eta_U^2\) enters in \(K^0 - \bar{K}^0\) mixing quantities and should therefore be very small [25]. The \(\eta_D\) coupling appears in \(B\) physics, in particular a contribution proportional to a product of \(\eta_U \eta_D\) gives rise to weak radiative \(B\)-meson decays. In ref. [25, 33] it was observed that the phase of \(\eta_U \eta_D\) can contribute to the CP violating asymmetry in \(D\)-decays.

Following [5] one can write an effective Hamiltonian for non-leptonic charm decays

\[
\mathcal{H}_{|\Delta C=1|} = \frac{G_F}{\sqrt{2}} \sum_{i=1,2,5,6} \sum_{q} (C^q_i Q^q_i + C^q_{i'} Q^q_{i'}) + \frac{G_F}{\sqrt{2}} \sum_{i=7,8} (C_i Q_i + C_{i'} Q_{i'}) + h.c. .
\]

The full expression for the quark operators \(Q_i\) are given in [5]. Among all possibilities it was found in [5, 53] that most likely, a candidate to explain discrepancy between the experimental result and the SM prediction is the color-magnetic dipole operator

\[Q_s = \frac{m_c}{4\pi^2} \bar{u}_L \sigma_{\mu\nu} T^A g_s G^\mu\nu_A e_R.\]

The same authors found that NP is entering into the Wilson coefficient \(C_8\) such that \(\Delta a_{CP}^{\exp} = -1.8 \ \text{Im} \ C_8^{NP}(m_c)\). Using universal couplings \(\eta_U\) and \(\eta_D\), it was found [25, 33] that the bounds on \(\text{Im}(\eta_U^* \eta_D^*)\) determined from \(b \rightarrow s \gamma\) give a NEDM consistent with the current experimental bound, while we found that the corresponding parameters constrained by the direct CP violation in the charm sector given by Eq. (1) are already ruled out by the experimental result on NEDM. In ref. [25] the most important
 contribution to the NEDM was found to arise from Weinberg’s CP-violating three-gluon operator operator. This contribution is generated by the CEDM of the $b$-quark when it is integrated out to obtain the Weinberg operator. This means that the $b$-quark EDM inducing Weinberg’s operator gives bounds on $\text{Im}(\eta^{\ast}_b\eta^*_D)$. These bounds are strong enough to exclude an explanation of the CP violation in charm physics.

Therefore, the explanation of the direct CP violation in $D$ decays by the presence of the color dipole operator, induced by the color octet scalar, is already excluded. A viable possibility is to deviate from the generation universality for the color octet couplings to the quark fields, or to modify the original MFV set-up of [25, 33]. Thus, the authors of [39] suggested small deviation from the MFV ansatz by allowing flavor changing $u \leftrightarrow c$ quark interactions with a neutral color octet scalar. Following ref. [39] the interaction Lagrangian is given by:

$$\mathcal{L}_{\text{eff}} = G(c \rightarrow u)\bar{u}_L T^A \Phi^A c_R + X_d \bar{d}_L t^A d_R \Phi^A + \text{h.c.},$$

(11)

where the couplings $G(c \rightarrow u)$ and $X_d$ are proportional to quark masses:

$$G(c \rightarrow u) \equiv [X_u]_{12} = \zeta_u y_u X_{cu}; \quad X_{cu} \sim V_{cs}V_{us}^{\ast}; \quad X_d = \zeta_d y_d,$$

(12)

where $\zeta_u,d$ are numbers (to be determined by CP-violation in the charm sector) and $y_q = m_q/v$, where $v$ is the VEV of the Higgs, and $m_q$ is the mass of quark $q$. In this framework the $D^0 - \bar{D}^0$ mixing are not present at tree level. In the case of $D^0 \rightarrow \pi^+\pi^-$, the effect is negligible in comparison with the $D^0 \rightarrow K^+K^-$ amplitude due to the smallness of the down quark mass compared to the bigger mass of the strange quark. In the $D^0 \rightarrow K^+K^-$ decay amplitude, it was found that two operators $\hat{O}^{1}\vec{S}_1 = (\bar{u}P_R s)(\bar{s}P_R c)$ and $\hat{O}^{1}\vec{S}_2 = (\bar{u}_\alpha P_R s_\beta)(\bar{s}_\beta P_R c_\alpha)$ contribute to the effective Hamiltonian as described in [39]. Motivated by the model of [39], we use this tree level effective $(cu \Phi_S)$ coupling bounded by the charm CP asymmetry. Then we find that there is a new two-loop contribution, which we consider in the following subsection. We also find that these couplings induce higher dimensional operators present in NEDM. It is important to notice that such couplings cannot affect any other low energy observable, as already discussed in ref. [39]. The Barr-Zee mechanism induced by the color octets has already been considered in ref. [54] for the EDM of the electron.
FIG. 4: Two loop diagram for an EDM of a d-quark.

A. Two-loop contributions to EDM

Two-loop contributions are induced by the presence of the \( c \to u \) flavor changing color octet couplings described in the modified MFV framework above. In our calculation we use the effective fermion (quark) propagator in a soft gauge field\(^{[55]}\):

\[
S_1(k, F) = \left( \frac{-e_q}{4} \right) \frac{\{ (\gamma \cdot k + m_q), \sigma \cdot F \}}{(k^2 - m_q^2)^2},
\]

(13)

where \( F \) is the electromagnetic field, \( \{A, B\} \) denotes the anti-commutator, \( k \) is the four momentum and \( m_q \) the mass of the quark \( q \).

We have found that the two loop contribution in Fig. 4 induces a dimension 5 electric dipole operator for the \( d \)-quark which might be written as an effective Lagrangian in the following way:

\[
\mathcal{L}_1(d \to d \gamma)_{\Phi} = K \left( \bar{d}_L \sigma \cdot F d_R \right).
\]

(14)

The quantity \( K \) is given by

\[
K = C_3 [g_W^2 V_{ud} V_{cd}^* G(c \to u)X_d] 2m_c e_c I_{2\text{-loop}},
\]

(15)

where \( C_3 \equiv \langle T^A T^A \rangle = 4/3 \) is a color factor and the leading logarithmic approximation of the two loop integral is:

\[
I_{2\text{-loop}} \simeq \left( \frac{1}{16 \pi^2} \right)^2 \frac{1}{M^2_\Phi} \left( \left[ \frac{\ln M^2_\Phi}{m_c^2} \right]^2 - \left[ \frac{\ln M^2_W}{m_c^2} \right]^2 \right).
\]

(16)

There is also a contribution from the crossed diagram with the result:

\[
\mathcal{L}_2(d \to d \gamma)_{\Phi} = K^* \left( \bar{d}_R \sigma \cdot F d_L \right),
\]

(17)
such that there will be an EDM of the $d$-quark equal to $(d_n/e)_\Phi^{2\text{-loop}} = 2Im(K)$. Note that these diagrams appear due to the presence of appropriate chiralities in the interacting Lagrangian (11). The EDM of the $u$-quark is suppressed by extra factors of small masses. We have checked that the opposite chirality of the the scalars lead to the two-loop helicity suppressed amplitudes. Following the work of [39], we have found that one can write the asymmetry in Eq. (11), assuming maximal phase $\Phi_f$, as

$$\Delta a_{CP} = \frac{2}{9} \frac{\zeta^2}{M_\Phi^2} m_c^2 C_{RGE} C_H$$

(18)

where $C_{RGE}$ denotes the factor which includes running of the Wilson coefficient ($C_{RGE} = 0.85$, for the running from the scale $M_\Phi \sim 1$ TeV down to the scale equal $m_c$), while $C_H$ stands for the possible enhancement of the hadronic matrix elements, assumed to be as large as $C_H \simeq 3$ in comparison with the naive factorization estimate leading to $C_H = 1$ as explained in [39]. We denote $\zeta^2 = \zeta_u \zeta_d$, which appears in both expressions for NEDM and $\Delta a_{CP}$. Assuming that our result for the $d$ quark electric dipole moment gives arise to NEDM as given in (4), we bound our new contribution to NEDM by the current experimental result given in (5). Both constraints are presented on Fig. 5.

In the case that the mass of color octet is bounded in the TeV regime [34, 35], the parameter $\zeta^2$ should be scaled accordingly. We have checked that even for a mass of $M_\Phi \geq 1.86$ TeV, perturbativity is still valid for the couplings in (11). That means that one can safely use the effective Lagrangian (11) for this purpose. We point out that the loop diagram in Fig. 4 is finite due to the chiral structure of (11), and that we have only given the relevant leading-log result. Of course the leading logarithmic result in (17) will be modified when taking into account higher orders by means of renormalization group techniques, but at the present stage we use this result. For fixed asymmetry - i.e. $\frac{\zeta^2}{M_\Phi^2}$ fixed - we obtain the relation

$$(d_n/e)_{\Phi}^{2\text{-loop}} \simeq \left( \frac{\lambda^2 m_d}{8\pi^4} \right) \frac{M_\Phi^2 m_c^2}{v^4 m_K^2} \Delta a_{CP} C_{RGE} C_H \left( \left[ \ln \frac{M_\Phi^2}{m_c^2} \right]^2 - \left[ \ln \frac{M_\Phi^2}{m_c^2} \right]^2 \right).$$

(19)

Numerically, we obtain the range

$$(d_n/e)_{\Phi}^{2\text{-loop}} \simeq (1.0 - 2.3) \times 10^{-26} \text{ cm},$$

(20)

for $M_\Phi$ in the range 400 GeV to 2 TeV, assuming $C_H \simeq 3$. It is interesting that for $C_H = 1$, we would get $d_n$ three times bigger and violate the experimental bound in Eq.(5)!

Let us comment that this result should also be valid for other flavor changing $(c u \Phi)$ colored scalar (triplet, sextet) couplings, with the appropriate color factor replacement.
FIG. 5: Regions in the $\zeta - M_\Phi$ plane compatible with the data on $\Delta a_{CP}$ (dark green, $C_H = 1$ and pale brown for $C_H \simeq 3$) and on the current experimental lower bound on $NEDM$ (pale green).

FIG. 6: One loop diagram for $W^+ d \rightarrow u \gamma$.

B. Higher-dimensional operators

We now consider the di-quark mechanism and calculate one loop diagrams for $W^- u \rightarrow d \gamma$ and $W^+ d \rightarrow u \gamma$ and afterwards attach left-handed $u \rightarrow d$ and $d \rightarrow u$ currents. The contribution from the one loop diagram $W^+ d \rightarrow u \gamma$ (from Fig. 6), leading to di-quark mechanism when $W$ is connected to a left-handed current, is

$$\mathcal{L}(du \rightarrow u d \gamma) = C_3 \left[ \frac{\epsilon_c g^2_W}{4 M_W^2} V_{ud} V_{cd}^* G(c \rightarrow u) X_d \right] 2m_c \ I_{1-loop} \ Q(ud \rightarrow d u \gamma)_{\Phi} , \quad (21)$$
where $C_3$ is the color factor defined above, $g_W$ is the $W$ coupling,

$$Q(ud \to du \gamma)_{\Phi} = (\bar{u} \sigma \cdot F \gamma_\mu L \gamma^\nu d) iD^\nu (\bar{d} \gamma^\mu L u)$$ (22)

is a generated dimension 9 operator, and

$$I_{1-loop} \simeq \frac{1}{16\pi^2} \frac{1}{m_c^2 M_W^2}.$$ (23)

The four quark operator part of $Q(ud \to du \gamma)_{\Phi}$ is of the same type as obtained in eq. (6) and can be estimated by using, say an $N^*$ resonance between the two currents. The covariant derivative, corresponding to the $W$-momentum, will contribute with a momentum of order the constituent quark mass in nucleons, i.e. $\hat{m} \equiv m_{constit} \simeq 350$ MeV. There are also 3 more relevant diagrams, leading to a NEDM

$$d_n \sim C_3 \frac{(e_c - e_d) g_W^2}{4M_W^2} V_{ud} V_{cd}^* G(c \to u) X_d [2m_c I_{1-loop} F_{Hadr},$$ (24)

where $F_{Hadr}$ is a pure hadronic factor

$$F_{Hadr} = \hat{F} \hat{m} M_s I_{sp},$$ (25)

where $I_{sp}$ is the phase space integral for the intermediate $N^*$ (with mass $M_s \simeq 1.440$ GeV) found to be $I_{sp} \simeq 0.9 \cdot 10^{-2}$ GeV$^2$. The quantity $\hat{F}$ is a product of transition form factors (from $N^*$ to $N$), which takes care of the damping within these, and is expected to be between 0.3 and 0.6. We find

$$d_n/e \simeq 3 \times 10^{-31} \text{ cm},$$ (26)

which is a bit higher than the corresponding SM value based on the operator in (6). As noticed by [18] one has to be careful in neglecting higher dimension operators.

**IV. CONCLUSIONS**

We have investigated contribution to the NEDM induced by the presence of the non-MFV flavor changing $(c u \Phi_8)$ coupling. The relevant couplings can be constrained by the world average CP asymmetry found in $D \to K^+ K^-/\pi^+ \pi^-$ decays or by current experimental bound on NEDM. It is remarkable that the CP violating asymmetry in charm decays and our new NEDM two-loop contribution allow sizable parts of the parameter space in the $\zeta - M_\phi$ plane. Still the CP violating asymmetry is more constraining than the bound on
NEDM. We comment on the higher dimensional operator contribution to NEDM and found that this contribution might be $10$ to $10^2$ larger than the NEDM values within the SM. Our study is applicable for flavor non-diagonal ($c u$) couplings to a scalar, by replacing the color factor in the amplitude.

The further LHC searches might set new bounds on the masses of colored scalars. On the other hand, new measurements of the CP asymmetry in charm physics would shed more light on the possible new source of CP violation in flavor physics. At the same time improvements of the experimental value for the NEDM might help in clarifying the role of new physics induced CP violating phases.

**Acknowledgments**

JOE is supported in part by the Norwegian research council. The work of SF is partially supported by Slovenian research agency ARRS.

[1] M. Pospelov and A. Ritz, Annals Phys. 318 (2005) 119 [hep-ph/0504231].

[2] http://www.slac.stanford.edu/xorg/hfag/

[3] R. Aaij et al. [LHCb Collaboration], Phys. Rev. Lett. 108 (2012) 111602 [arXiv:1112.0938 [hep-ex]].

[4] T. Aaltonen et al. [CDF Collaboration], Phys. Rev. D 85 (2012) 012009 [arXiv:1111.5023 [hep-ex]].

[5] G. Isidori, J. F. Kamenik, Z. Ligeti and G. Perez, Phys. Lett. B 711 (2012) 46 [arXiv:1111.4987 [hep-ph]].

[6] Y. Grossman, A. L. Kagan and Y. Nir, Phys. Rev. D 75 (2007) 036008 [hep-ph/0609178].

[7] H. -Y. Cheng and C. -W. Chiang, Phys. Rev. D 85 (2012) 034036 [Erratum-ibid. D 85 (2012) 079903] [arXiv:1201.0785 [hep-ph]].

[8] H. -n. Li, C. -D. Lu and F. -S. Yu, Phys. Rev. D 86 (2012) 036012 [arXiv:1203.3120 [hep-ph]].

[9] E. Franco, S. Mishima and L. Silvestrini, JHEP 1205 (2012) 140 [arXiv:1203.3131 [hep-ph]].

[10] D. Pirtskhalava and P. Uttayarat, Phys. Lett. B 712 (2012) 81 [arXiv:1112.5451 [hep-ph]].
[11] B. Bhattacharya, M. Gronau and J. L. Rosner, Phys. Rev. D 85 (2012) 054014 [arXiv:1201.2351 [hep-ph]].
[12] T. Feldmann, S. Nandi and A. Soni, JHEP 1206 (2012) 007 [arXiv:1202.3795 [hep-ph]].
[13] H. -Y. Cheng and C. -W. Chiang, Phys. Rev. D 86 (2012) 014014 [arXiv:1205.0580 [hep-ph]].
[14] J. Brod, Y. Grossman, A. L. Kagan and J. Zupan, JHEP 1210 (2012) 161 [arXiv:1203.6659 [hep-ph]].
[15] W. Altmannshofer, A. J. Buras and P. Paradisi, Phys. Lett. B 688 (2010) 202 [arXiv:1001.3835 [hep-ph]].
[16] A. Lenz, [arXiv:1311.6447 [hep-ph]].
[17] A. J. Buras, G. Isidori and P. Paradisi, Phys. Lett. B 694 (2011) 402 [arXiv:1007.5291 [hep-ph]].
[18] T. Mannel and N. Uraltsev, Phys. Rev. D 85 (2012) 096002 [arXiv:1202.6270 [hep-ph]].
[19] J.O. Eeg, J.O. and I. Picek, Phys. Lett. B130 (1983) 308
[20] J.O. Eeg, and I. Picek, Nucl. Phys. B244 (1984) 77
[21] T. Mannel and N. Uraltsev, JHEP 1303 (2013) 064 [arXiv:1205.0233 [hep-ph]].
[22] G. D’Ambrosio, G. F. Giudice, G. Isidori and A. Strumia, Nucl. Phys. B 645 (2002) 155 [hep-ph/0207036].
[23] W. Altmannshofer and D. M. Straub, JHEP 1208 (2012) 121 [arXiv:1206.0273 [hep-ph]].
[24] “Heavy Flavor Averaging Group Collaboration” (D. Asner et al) [arXiv:1010.1589 [hep-ex]].
[25] A. V. Manohar and M. B. Wise, Phys. Rev. D 74 (2006) 035009 [hep-ph/0606172].
[26] M. I. Gresham and M. B. Wise, Phys. Rev. D 76 (2007) 075003 [arXiv:0706.0009 [hep-ph]].
[27] G. Boyd, A. K. Gupta, S. P. Trivedi and M. B. Wise, Phys. Lett. B 241, 584 (1990).
[28] S. Weinberg, Phys. Rev. Lett. 63 (1989) 2333.
[29] J. M. Arnold, B. Fornal and M. B. Wise, Phys. Rev. D 87 (2013) 075004 [arXiv:1212.4556 [hep-ph]].
[30] C. A. Baker, D. D. Doyle, P. Geltenbort, K. Green, M. G. D. van der Grinten, P. G. Harris, P. Iaydjiev and S. N. Ivanov et al., Phys. Rev. Lett. 97 (2006) 131801 [hep-ex/0602020].
[31] S. Bertolini, L. Di Luzio and M. Malinsky, Phys. Rev. D 87 (2013) 085020 [arXiv:1302.3401 [hep-ph]].
[32] J. Cao, P. Wan, J. M. Yang and J. Zhu, JHEP 1308 (2013) 009 [arXiv:1303.2426 [hep-ph]].
[33] M. Trott and M. B. Wise, JHEP 1011 (2010) 157 [arXiv:1009.2813 [hep-ph]].
[34] G. Aad et al. [ATLAS Collaboration], Eur. Phys. J. C 73 (2013) 2263 [arXiv:1210.4826 [hep-ex]].

[35] S. Chatrchyan et al. [CMS Collaboration], JHEP 1301 (2013) 013 [arXiv:1210.2387 [hep-ex]].

[36] K. Yi [CMS Collaboration], J. Phys. Conf. Ser. 455 (2013) 012034 [arXiv:1307.1400 [hep-ex]].

[37] G. Aad et al. [ATLAS Collaboration], JHEP 1301 (2013) 029 [arXiv:1210.1718 [hep-ex]].

[38] Z. Heng, L. Shang, Y. Zhang and J. Zhu, [arXiv:1312.4260] [hep-ph].

[39] W. Altmannshofer, R. Primulando, C. -T. Yu and F. Yu, JHEP 1204 (2012) 049 [arXiv:1202.2866 [hep-ph]].

[40] E.P. Shabalin, E.P., Yad.Fiz. 31 (1980)1665-1679 (Sov. J. Nucl.Phys. 31 (1980) 864).

[41] A. Czarnecki, and B. Krause, Phys.Rev.Lett. 78 (1997) 4339 [hep-ph/9704355].

[42] D.V. Nanopoulos, A. Yildiz, and P. H. Cox, Phys.Lett. B87 53

[43] B.F. Morel, Nucl.Phys. B157 (1979),23

[44] M.B. Gavela, A. Le Yaouanc, L. Oliver,O. Pene J.C. Raynal, and T.N. Pham Phys. Lett. B109 (1982) 215

[45] I.B. Khriplovich and A.R. Zhitnitsky, Moment Phys.Lett. B109 (1982) 490

[46] B.H.J. McKellar,S.R. Choudhury, X.-G. He and S. Pakvasa, Phys.Lett., B197 (1987) 556

[47] C. Hamzaoui and A. Barroso, Phys. Lett. B 154, 202 (1985).

[48] M. Jung and A. Pich, [arXiv:1308.6283] [hep-ph].

[49] D. A. Demir, M. Pospelov and A. Ritz, Phys. Rev. D 67 (2003) 015007 [hep-ph/0208257].

[50] S. M. Barr and A. Zee, Phys. Rev. Lett. 65 (1990) 21 [Erratum-ibid. 65 (1990) 2920].

[51] G. Degrassi and P. Slavich, Phys. Rev. D 81 (2010) 075001 [arXiv:1002.1071 [hep-ph]].

[52] S. Fajfer, A. Greljo, J. F. Kamenik and I. Mustac, JHEP 1307 (2013) 155 [arXiv:1304.4219 [hep-ph]].

[53] G. Isidori and J. F. Kamenik, Phys. Rev. Lett. 109 (2012) 171801 [arXiv:1205.3164 [hep-ph]].

[54] J. H. Heo and W. -Y. Keung, Phys. Lett. B 661 (2008) 259 [arXiv:0801.0231 [hep-ph]].

[55] L.J Reinders, H. Rubinstein, and S. Yazaki, Phys.Rept. 127 (1985) 1