

# What the neutron knows about the top at the LHC

Jernej F. Kamenik,<sup>1,2,\*</sup> Michele Papucci,<sup>3,4,†</sup> and Andreas Weiler<sup>4,5,‡</sup>

<sup>1</sup>*J. Stefan Institute, Jamova 39, P. O. Box 3000, 1001 Ljubljana, Slovenia*

<sup>2</sup>*Department of Physics, University of Ljubljana, Jadranska 19, 1000 Ljubljana, Slovenia*

<sup>3</sup>*Ernest Orlando Lawrence Berkeley National Laboratory, University of California, Berkeley, CA 94720*

<sup>4</sup>*CERN, Theory Division, CH-1211, Geneva 23, Switzerland*

<sup>5</sup>*DESY, Notkestrasse 85, D-22603 Hamburg, Germany*

(Dated: October 5, 2017)

We investigate the direct and indirect bounds on dipole operators involving the top quark. A careful analysis shows that the experimental upper limit on the neutron electric dipole moment strongly constrains the chromo-electric dipole of the top. We improve previous bounds by two orders of magnitude. This has significant implications for new physics models and it also means that CP violation in top pair production mediated by dipole operators will not be accessible at the LHC. The CP conserving chromo-magnetic dipole moments are constrained by recent measurements of the  $t\bar{t}$  spectrum by the ATLAS collaboration. We also update the indirect constraints on electric and magnetic dipole moments from radiative  $b \rightarrow s$  transitions, finding that they can be considerably larger than their colored counterparts.

## Introduction

In many scenarios the top quark provides a preferred window on physics beyond the Standard Model (SM), given its large coupling to the physics responsible for the Electro-Weak Symmetry Breaking (EWSB). Top quarks will be copiously produced at the LHC - the current rate is of the order of five tops per minute. While the top quark has been studied in some detail at the Tevatron, many of its properties besides the mass, spin and the color and electric charges are still poorly constrained. Significant new insights on top quark properties will therefore be one of the tasks of the LHC.

A fascinating possibility is that the top quark shows deviations from its behavior predicted by the SM. The leading contributions are encoded in the (chromo)-electric and (chromo)-magnetic dipole moments, (C)EDM and (C)MDM in the following.

A particularly interesting scenario is realized if some of the quarks are partially or fully composite [1]. The top, being the most massive quark, is the most natural candidate, as the mass and the amount of compositeness are often related, see e.g. [2]. If the compositeness scale is in its natural range, large CMDM and CEDM are expected [3].

Supersymmetric models on the other hand, can also lead to enhanced dipole moments of the top. This happens if the supersymmetric partners of the top are not too heavy, which is a requirement of naturalness and possibly electro-weak baryogenesis, see e.g. [4]. The presence of sizable flavor-blind phases accessible to the third-generation quarks can also explain the recent hints of CP violation in  $B_s$  mixing in the context of Minimal Flavor Violation [5].

We leave the implications of our results for specific scenarios to future work. In this Letter we investigate the direct and indirect constraints on dipole operators

involving the top quark in a model independent way. In particular we will critically reanalyze the phenomenologically relevant question whether CP violation in top pair production can be mediated by dipole operators at a level accessible for the LHC.

As it is well known, dipole moments before EWSB are encoded in dimension-6 operators

$$\begin{aligned} & (c_{LR,c} g_s \bar{Q} H \sigma^{\mu\nu} T^a U + \text{h.c.}) G_{\mu\nu}^a, \\ & (c_{LR,w} g \bar{Q} \tau^a H \sigma^{\mu\nu} U + \text{h.c.}) W_{\mu\nu}^a, \\ & (c_{LR,y} g' \bar{Q} H \sigma^{\mu\nu} U + \text{h.c.}) B_{\mu\nu}, \end{aligned} \quad (1)$$

with dimensional couplings  $c_i \sim 1/\Lambda_i^2$ . In particular,  $\text{Im}(c_{LR,i}) \neq 0$  would signal CP violation.

In the following, we will employ a phenomenological Hamiltonian which can be easily translated to the more physical  $SU(2) \times U(1)$  gauge invariant basis in Eq. (1),

$$\begin{aligned} \mathcal{H}_{\text{eff}} = & -\frac{1}{2} \bar{\psi}_q \left[ (F_{\mu\nu} \sigma^{\mu\nu})(\mu_q + i\gamma_5 d_q) \right. \\ & \left. + g_s (G_{\mu\nu}^a t^a \sigma^{\mu\nu})(\tilde{\mu}_q + i\gamma_5 \tilde{d}_q) \right] \psi_q \\ & - \frac{1}{6} w f^{abc} \varepsilon^{\mu\nu\lambda\rho} G_{\mu\sigma}^a G_{\nu}^{b\sigma} G_{\lambda\rho}^c, \end{aligned} \quad (2)$$

where  $q = u, d, s, c, b, t$  and  $\varepsilon^{0123} = 1$ . We denote  $d_q$  and  $\tilde{d}_q$  as the EDM and the CEDM of the quark  $q$ , while  $\mu_q$  and  $\tilde{\mu}_q$  are the corresponding MDM and CMDM. We have included the CP violating Weinberg operator [6], which will be crucial later. Correspondingly, we have omitted the terms involving the charged gauge bosons, since they will not play an important role in the following and have been already investigated elsewhere [7].

## Indirect constraints

We first consider the present indirect constraints on the top CEDM  $\tilde{d}_t$ . The operators in Eq. (2) run

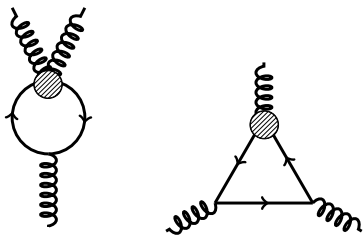


FIG. 1: Diagrams generating the contribution to the Weinberg operator at the top threshold [9]. The grey blob denotes the insertion of the chromo-electric dipole operator.

and mix under QCD renormalization group (RG) evolution. At present, these effects are known to NLL accuracy [8]. In particular the Weinberg operator mixes into the (C)EDMs of quarks, but not vice versa. Nevertheless, it has been known for some time [9] that the CEDMs induce a *finite* threshold correction to the Weinberg operator when a heavy quark is integrated out, as shown in Fig. 1,

$$\delta w^{(q)} = \frac{g_s^3}{32\pi^2} \frac{\tilde{d}_q}{m_q}. \quad (3)$$

In the case of the top quark, the combined effects of the finite shift in  $w$  and the subsequent RG evolution to the hadronic scale will induce non-zero contributions also for the (C)EDMs of the light quarks.

At the hadronic scale  $\mu_H \sim 1$  GeV we thus obtain  $d_{u,d}(\mu_H)$ ,  $\tilde{d}_{u,d}(\mu_H)$  and  $w(\mu_H)$  in terms of  $\tilde{d}_t(m_t)$

$$\begin{aligned} d_u &= -3.1 \cdot 10^{-9} e \tilde{d}_t, & d_d &= 3.5 \cdot 10^{-9} e \tilde{d}_t, \\ \tilde{d}_u &= 8.9 \cdot 10^{-9} \tilde{d}_t, & \tilde{d}_d &= 2.0 \cdot 10^{-8} \tilde{d}_t, \\ w &= 1.0 \cdot 10^{-5} \text{ GeV}^{-1} \tilde{d}_t, \end{aligned} \quad (4)$$

where we have used  $m_t = 173.3$  GeV [10],  $m_d^{\overline{MS}}(2 \text{ GeV}) = (4.7 \pm 0.1) \text{ MeV}$ ,  $m_u^{\overline{MS}}(2 \text{ GeV}) = (2.1 \pm 0.1) \text{ MeV}$  [11] and  $\alpha_s^{\overline{MS}}(m_Z) = 0.1184$  [12].

Presently, the most sensitive observables are the atomic EDMs of mercury ( $d_{Hg} < 3.1 \cdot 10^{-29} e \text{ cm}$  at 95% C.L. [13]) and of the neutron ( $d_n < 2.9 \cdot 10^{-26} e \text{ cm}$  at 95% C.L. [14]). Following [15], we evaluate the relevant contributions as

$$d_{Hg} = -1.8 \cdot 10^{-4} \text{ GeV}^{-1} e \tilde{g}_{\pi NN}^{(1)}, \quad (5a)$$

$$\begin{aligned} d_n &= (1 \pm 0.5) [1.1 e (\tilde{d}_d + 0.5 \tilde{d}_u) + 1.4 (d_d - 0.25 d_u)] \\ &\quad + (22 \pm 10) \cdot 10^{-3} \text{ GeV} e w, \end{aligned} \quad (5b)$$

where  $\tilde{g}_{\pi NN}^{(1)} = 4_{-2}^{+8} (\tilde{d}_u - \tilde{d}_d) \text{ GeV}$ . All quantities are evaluated at the scale  $\mu_H \sim 1$  GeV. The values and uncertainty estimates for the relevant matrix elements, particularly for the Weinberg operator contribution to  $d_n$ , have been evaluated using QCD sum rule techniques [16].

Inserting (4) into the above expressions and treating all the relevant theoretical uncertainties as flat distributions

within the stated errors, we find that the neutron EDM constrains the top CEDM to be

$$|\tilde{d}_t| < 2.3 \cdot 10^{-19} \text{ cm} \quad (95\% \text{ C.L.}), \quad (6)$$

i.e.  $|\tilde{d}_t m_t| < 2.0 \cdot 10^{-3}$ . The constraint from the neutron EDM is dominated by the contribution to the Weinberg operator, which amounts to roughly 85% of the total effect of  $\tilde{d}_t$  in  $d_n$ , even though the light quark (C)EDM contributions are not totally negligible. Furthermore, the constraint from  $d_{Hg}$  provides a (two orders of magnitude) weaker bound on  $\tilde{d}_t$ , since it is not sensitive to the Weinberg operator and also comparatively weaker than  $d_n$  for the light quark CEDMs.

The indirect constraints on the other top dipole moments in Eq. (2) are considerably weaker. The EDM of the top,  $d_t$ , induces light quark EDMs only through weak interactions and is suppressed by flavor mixing factors [17] resulting in  $d_d = 2.4 \times 10^{-12} d_t$ , and consequently we find a weak bound of

$$|d_t| < 1.7 \times 10^{-14} e \text{ cm} \quad (95\% \text{ C.L.}). \quad (7)$$

A stronger limit comes from  $b \rightarrow s\gamma$  and  $b \rightarrow s\ell^+\ell^-$  processes, since the leading SM contribution carries the same loop and flavor suppressions. Following [18], we obtain

$$\Delta C_{7\gamma}(m_W) = 6.5 \cdot 10^{-2} (\mu_t - 2.65 d_t) m_t, \quad (8)$$

where we have included the effects of the top MDM which is also constrained.  $\Delta C_{7\gamma}$  is the new contribution to the Wilson coefficient of the magnetic operator mediating the  $b \rightarrow s$  transition. Using the semi-analytic formulae of [19] for the SM NNLO prediction of  $Br(B \rightarrow X_s \gamma)$  [20] and the NNLL result for  $Br(B \rightarrow X_s \ell^+\ell^-)$  for the low  $q^2$  region [21], we obtain

$$-1.83 < (\mu_t - 2.65 d_t) m_t < 0.53 \quad (95\% \text{ C.L.}). \quad (9)$$

These processes are also sensitive to the top CMDM, through the operator  $\mathcal{O}_{8g}$ . However, due to the additional  $\alpha_s$  suppression, the indirect bound is significantly weaker and will be superseded by the direct one derived in the next section.

### Collider constraints

The gluonic dipole couplings of the top directly affect  $t\bar{t}$  production at hadron colliders [22]. The complementarity of the total production cross-section measurements at the Tevatron and the LHC in constraining such contributions has recently been pointed out in [23]. In [24], the CMDM is constrained using Tevatron data in combination with four fermi operators and spin correlations are shown to be promising. A number of  $T$ -odd correlation observables can be constructed which are sensitive

to the top CEDM (c.f. [25]). At present however, the most sensitive observable available at the Tevatron and the LHC are still the total production cross-sections and their differential spectrum as a function of the invariant mass of the top pair ( $m_{t\bar{t}}$ ).

At the Tevatron the recent combination of the CDF analyses yields [26],

$$\sigma_{\text{exp}}^{\text{Tevatron}} = (7.50 \pm 0.48) \text{ pb} \quad (10)$$

for an assumed top mass of  $m_t = 172.5$  GeV. We have combined the estimated statistical and systematic errors in quadrature. This is consistent with the most recent theoretical SM prediction for this observable [27],

$$\sigma_{\text{SM}}^{\text{Tevatron}} = (6.75_{-0.42}^{+0.08}) \text{ pb} \quad (11)$$

based on approximate NNLO QCD calculation using the same top mass and the MSTW2008 PDFs [28]. As pointed out in [29], the most significant information in the high  $m_{t\bar{t}}$  region at the Tevatron is the one derived from the next-to-highest measured bin [30],

$$\sigma(700 \text{ GeV} < m_{t\bar{t}} < 800 \text{ GeV})_{\text{exp}} = (80 \pm 37) \text{ fb}, \quad (12)$$

to be compared with the SM theory prediction of [31],

$$\sigma(700 \text{ GeV} < m_{t\bar{t}} < 800 \text{ GeV})_{\text{SM}} = (80 \pm 8) \text{ fb}. \quad (13)$$

At the LHC at 7 TeV, presently the most precise measurement [32] for the total production cross-section yields

$$\sigma_{\text{exp}}^{\text{LHC}} = (180 \pm 18) \text{ pb}, \quad (14)$$

in agreement with the SM prediction of [33],

$$\sigma_{\text{SM}}^{\text{LHC}} = (165_{-16}^{+11}) \text{ pb}. \quad (15)$$

Recently the ATLAS collaboration published a study [34] of the  $m_{t\bar{t}}$  spectrum using  $200 \text{ pb}^{-1}$  collected luminosity. In the high  $m_{t\bar{t}} > 1$  TeV region they report  $N_{\text{exp}}^{1\text{TeV}} = 77 \pm 9$  events. The results were not unfolded. Among the systematic uncertainties, the  $b$ -tagging efficiency (11%) dominates the total inclusive cross-section measurement and stays almost constant with  $m_{t\bar{t}}$ . The other important source of error is the jet energy scale uncertainty. It is subleading in the total inclusive measurement (9%), with a mild  $m_{t\bar{t}}$  dependence. Since the ATLAS result is not unfolded, one needs to take into account the invariant mass resolution, and the reconstruction efficiency and acceptance, ( $A \cdot \epsilon$ ). We model the former by smearing our partonic  $m_{t\bar{t}}$  distributions with a Gaussian kernel. We estimate a  $\mathcal{O}(0.25 \cdot m_{t\bar{t}})$  width for this smearing by comparing to the reconstructed invariant mass resolutions of a sample of narrow  $Z'$  models in the same ATLAS study. Finally, to estimate  $A \cdot \epsilon$  we

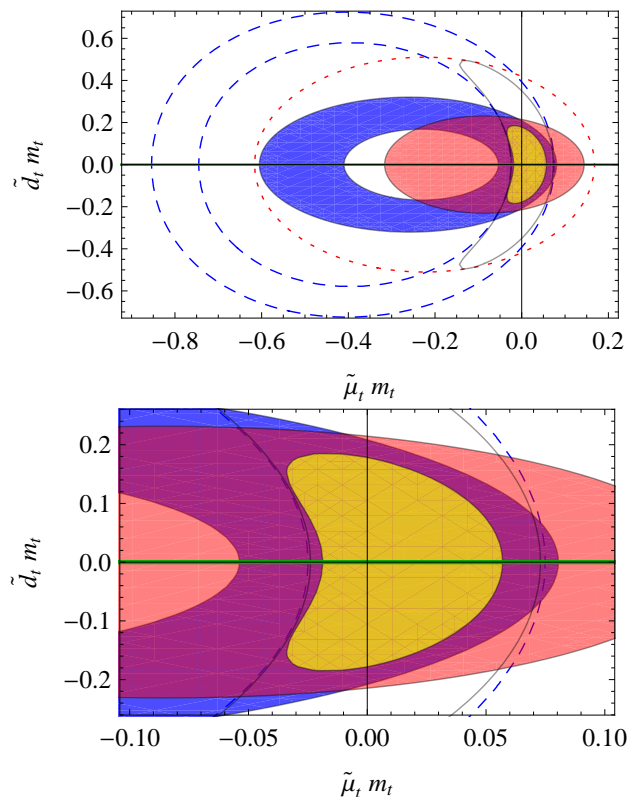


FIG. 2: Combined LHC and Tevatron 95% C.L. constraints on the top CDM ( $\tilde{\mu}_t$ ) and the CEDM ( $\tilde{d}_t$ ) (shaded in yellow). Individual constraints come from the total cross-section and  $m_{t\bar{t}}$  spectrum measurements at the Tevatron (dashed blue and dotted red), as well as the LHC (shaded blue and red). The combination of only Tevatron constraints is drawn in black. Finally the CEDM indirect constraint is presented in green.

compare the smeared SM  $m_{t\bar{t}}$  distribution, computed using known approximate NNLO QCD results [31], with the reconstructed SM background distribution presented in the ATLAS Note. Using  $Br(t\bar{t} \rightarrow 4j + \ell) = 0.3$  (where  $\ell = e, \mu$ ) we extract a constant  $A \cdot \epsilon \simeq 0.3$  for the  $m_{t\bar{t}}$  bins between  $1 \text{ TeV} < m_{t\bar{t}} < 1.6 \text{ TeV}$ , which can be now used to compare the signal with the data.

We find for the measured  $t\bar{t}$  cross-sections including statistical and our estimates for the systematic uncertainties

$$\sigma(1 \text{ TeV} < m_{t\bar{t}} < 1.2 \text{ TeV}) = (2.9 \pm 0.6) \text{ pb}, \quad (16a)$$

$$\sigma(1.2 \text{ TeV} < m_{t\bar{t}} < 1.4 \text{ TeV}) = (1.0 \pm 0.3) \text{ pb}, \quad (16b)$$

$$\sigma(1.4 \text{ TeV} < m_{t\bar{t}} < 1.6 \text{ TeV}) = (0.45 \pm 0.19) \text{ pb}, \quad (16c)$$

which corresponds to

$$\sigma(m_{t\bar{t}} > 1 \text{ TeV}) = (4.5 \pm 0.79) \text{ pb}. \quad (17)$$

These results may be directly compared to partonic  $m_{t\bar{t}}$  distributions smeared with  $0.25 \cdot m_{t\bar{t}}$ -wide Gaussians.

We evaluate the effect of the top CDM and CEDM on the relevant Tevatron and LHC observables at LO

in QCD, using the known partonic cross-section formulae [22] convolved with MSTW2008 PDFs [28]. We normalize our SM values to approximate NNLO results [27, 31, 33] including theoretical uncertainties. Additionally we have checked the residual theoretical uncertainty in the relative NP contributions by varying the factorization and renormalization scales and finding negligible differences. We compare these estimates of the inclusive and high  $m_{t\bar{t}}$  cross-sections with the corresponding measurements both at the Tevatron and the LHC in Fig. 2.

We observe that the new ATLAS result on the high  $m_{t\bar{t}}$  region at the LHC sizably shrinks the allowed region in the  $(\tilde{\mu}_t, \tilde{d}_t)$  plane relative to previous results [23] or compared to using only Tevatron data. Marginalizing over the CEDM values, we obtain a new best bound on the top CMDM of

$$|\tilde{\mu}_t| m_t < 0.05 \quad (95\% \text{ C.L.}). \quad (18)$$

The CEDM of the top is constrained to

$$|\tilde{d}_t| m_t < 0.16 \quad (95\% \text{ C.L.}), \quad (19)$$

or  $|\tilde{d}_t| < 1.9 \cdot 10^{-17}$  cm, which is almost two orders of magnitude weaker than our new indirect bound (6). A remark on the consistency of our EFT expansion is in order here. The CP violating CEDM does not interfere with the SM and its contribution to the cross section starts at  $\sim 1/\Lambda^4$ , and not at  $\sim 1/\Lambda^2$  as the CP conserving CMDM. It is therefore of the same order as interfering dimension eight operators, which in principle should have been included. Fortunately the far dominant constraint arises from indirect observables and we can safely ignore this issue.

The bounds can be expressed in terms of the gauge-invariant basis of Eq. (1). The minimum scale for new physics contributing to the gluonic dipole moments has to be

$$\begin{aligned} \text{Re } \Lambda_{LR,c}^{\text{direct}} &> 1.1 \text{ TeV} \\ \text{Im } \Lambda_{LR,c}^{\text{direct}} &> 0.62 \text{ TeV} \\ \text{Im } \Lambda_{LR,c}^{\text{neutron}} &> 5.5 \text{ TeV} \end{aligned}$$

where we have separated the minimum scale for real and imaginary contributions.

Let us now consider the prospects of probing a CEDM at the LHC as small as required by the indirect bound (6). In [35], it was estimated that a  $5\sigma$  detection for a value of  $\tilde{d}_t m_t = 0.05$  requires  $10 \text{ fb}^{-1}$ . We see that we would therefore need at least  $\mathcal{O}(1 \text{ ab}^{-1})$  luminosity to detect CP violation at  $5\sigma$  from a top CEDM. This is likely beyond current LHC capabilities, unless further collider studies can improve the reach.

## Conclusions

In this Letter we have investigated the constraints on dipole operators involving the top quark. The neutron EDM imposes a stringent bound on the top CEDM. We have also derived a direct bound on the CMDM of the top from the invariant mass distribution in  $t\bar{t}$  events at the LHC. These two bounds have significant implications for models of top compositeness and supersymmetric models with light stops which we will address in an upcoming paper. Finally, we have updated the constraints on the EDM/MDM of the top coming from  $b \rightarrow s$  transitions with the most recent data.

## Acknowledgements

This work was initiated at the tmini workshop at the Weizmann Institute of Science. We thank the organizers for the inspiring atmosphere and great hospitality. We would like to thank Georgios Choudalakis, Christophe Grojean, Elin Bergeaas Kuutmann, Michele Redi, and Pekka Sinervo for useful discussions. M.P. would like to thank the Aspen Center for Physics where part of this work was completed. The work of M.P. was supported in part by the Director, Office of Science, Office of High Energy and Nuclear Physics, of the US Department of Energy under Contract DE-AC02-05CH11231. The work of J.F.K. was supported in part by the Slovenian Research Agency. The work of A.W. was supported in part by the German Science Foundation (DFG) under the Collaborative Research Center (SFB) 676.

---

\* Electronic address: [jernej.kamenik@ijs.si](mailto:jernej.kamenik@ijs.si)

† Electronic address: [mpapucci@lbl.gov](mailto:mpapucci@lbl.gov)

‡ Electronic address: [andreas.weiler@desy.de](mailto:andreas.weiler@desy.de)

- [1] D. B. Kaplan, Nucl. Phys. B **365** (1991) 259.
- [2] Y. Grossman and M. Neubert, Phys. Lett. B **474** (2000) 361 [[arXiv:hep-ph/9912408](https://arxiv.org/abs/hep-ph/9912408)]; T. Gherghetta and A. Pomarol, Nucl. Phys. B **586** (2000) 141 [[arXiv:hep-ph/0003129](https://arxiv.org/abs/hep-ph/0003129)]; R. Contino and A. Pomarol, JHEP **0411** (2004) 058 [[arXiv:hep-th/0406257](https://arxiv.org/abs/hep-th/0406257)]; B. Lillie, J. Shu and T. M. P. Tait, JHEP **0804** (2008) 087 [[arXiv:0712.3057](https://arxiv.org/abs/hep-ph/0712.3057) [hep-ph]]; K. Kumar, T. M. P. Tait and R. Vega-Morales, JHEP **0905** (2009) 022 [[arXiv:0901.3808](https://arxiv.org/abs/hep-ph/0901.3808) [hep-ph]]; A. Pomarol and J. Serra, Phys. Rev. D **78** (2008) 074026 [[arXiv:0806.3247](https://arxiv.org/abs/hep-ph/0806.3247) [hep-ph]].
- [3] K. Agashe, G. Perez and A. Soni, Phys. Rev. D **75** (2007) 015002 [[arXiv:hep-ph/0606293](https://arxiv.org/abs/hep-ph/0606293)].
- [4] M. Carena, G. Nardini, M. Quiros and C. E. M. Wagner, Nucl. Phys. B **812** (2009) 243 [[arXiv:0809.3760](https://arxiv.org/abs/hep-ph/0809.3760) [hep-ph]].
- [5] A. L. Kagan, G. Perez, T. Volansky and J. Zupan, Phys. Rev. D **80** (2009) 076002 [[arXiv:0903.1794](https://arxiv.org/abs/hep-ph/0903.1794) [hep-ph]]; Z. Ligeti, M. Papucci, G. Perez and J. Zupan, Phys. Rev. Lett. **105** (2010) 131601 [[arXiv:1006.0432](https://arxiv.org/abs/hep-ph/1006.0432) [hep-ph]].
- [6] S. Weinberg, Phys. Rev. Lett. **63**, 2333 (1989).

- [7] B. Grzadkowski, M. Misiak, Phys. Rev. **D78**, 077501 (2008); [[arXiv:0802.1413](#) [hep-ph]]; J. Drobnak, S. Fajfer, J. F. Kamenik, Phys. Rev. **D82**, 114008 (2010). [[arXiv:1010.2402](#) [hep-ph]]; Phys. Lett. **B701**, 234-239 (2011). [[arXiv:1102.4347](#) [hep-ph]]; J. A. Aguilar-Saavedra, N. F. Castro, A. Onofre, Phys. Rev. **D83**, 117301 (2011). [[arXiv:1105.0117](#) [hep-ph]].
- [8] G. Degrassi, E. Franco, S. Marchetti, L. Silvestrini, JHEP **0511**, 044 (2005). [[hep-ph/0510137](#)].
- [9] E. Braaten, C. S. Li and T. C. Yuan, Phys. Rev. Lett. **64** (1990) 1709; D. Chang, T. W. Kephart, W. Y. Keung and T. C. Yuan, Phys. Rev. Lett. **68** (1992) 439.
- [10] [CDF and D0 Collaboration], [[arXiv:1007.3178](#) [hep-ex]].
- [11] J. Laiho, [[arXiv:1106.0457](#) [hep-lat]].
- [12] S. Bethke, Eur. Phys. J. **C64**, 689-703 (2009). [[arXiv:0908.1135](#) [hep-ph]].
- [13] W. C. Griffith, M. D. Swallows, T. H. Loftus, M. V. Romalis, B. R. Heckel, E. N. Fortson, Phys. Rev. Lett. **102**, 101601 (2009).
- [14] C. A. Baker, D. D. Doyle, P. Geltenbort, K. Green, M. G. D. van der Grinten, P. G. Harris, P. Iaydjiev, S. N. Ivanov *et al.*, Phys. Rev. Lett. **97**, 131801 (2006). [[hep-ex/0602020](#)].
- [15] M. Pospelov, A. Ritz, Annals Phys. **318**, 119-169 (2005). [[hep-ph/0504231](#)].
- [16] D. A. Demir, M. Pospelov and A. Ritz, Phys. Rev. D **67** (2003) 015007 [[arXiv:hep-ph/0208257](#)].
- [17] A. Cordero-Cid, J. M. Hernandez, G. Tavares-Velasco and J. J. Toscano, J. Phys. G **35** (2008) 025004 [[arXiv:0712.0154](#) [hep-ph]].
- [18] J. L. Hewett and T. G. Rizzo, Phys. Rev. D **49** (1994) 319 [[arXiv:hep-ph/9305223](#)].
- [19] S. Descotes-Genon, D. Ghosh, J. Matias and M. Ramon, JHEP **1106** (2011) 099 [[arXiv:1104.3342](#) [hep-ph]].
- [20] M. Misiak *et al.*, Phys. Rev. Lett. **98** (2007) 022002 [[arXiv:hep-ph/0609232](#)].
- [21] C. Bobeth, P. Gambino, M. Gorbahn and U. Haisch, JHEP **0404** (2004) 071 [[arXiv:hep-ph/0312090](#)]. T. Huber, E. Lunghi, M. Misiak and D. Wyler, Nucl. Phys. B **740** (2006) 105 [[arXiv:hep-ph/0512066](#)].
- [22] D. Atwood, A. Kagan, T. G. Rizzo, Phys. Rev. **D52**, 6264-6270 (1995). [[hep-ph/9407408](#)]; P. Haberl, O. Nachtmann and A. Wilch, Phys. Rev. D **53** (1996) 4875 [[arXiv:hep-ph/9505409](#)].
- [23] Z. Hioki, K. Ohkuma, Eur. Phys. J. **C65**, 127-135 (2010). [[arXiv:0910.3049](#) [hep-ph]]; Eur. Phys. J. **C71**, 1535 (2011). [[arXiv:1011.2655](#) [hep-ph]].
- [24] C. Degrande, J. M. Gerard, C. Grojean, F. Maltoni and G. Servant, JHEP **1103** (2011) 125 [[arXiv:1010.6304](#) [hep-ph]].
- [25] D. Atwood, S. Bar-Shalom, G. Eilam and A. Soni, Phys. Rept. **347** (2001) 1 [[arXiv:hep-ph/0006032](#)].
- [26] T. Aaltonen *et al.* [CDF Collaboration], public conf. note no. 9913, October 2009.
- [27] V. Ahrens, A. Ferroglia, B. D. Pecjak and L. L. Yang, [[arXiv:1105.5824](#) [hep-ph]].
- [28] A. D. Martin, W. J. Stirling, R. S. Thorne, G. Watt, Eur. Phys. J. **C63**, 189-285 (2009). [[arXiv:0901.0002](#) [hep-ph]].
- [29] K. Blum *et al.*, [[arXiv:1102.3133](#) [hep-ph]].
- [30] T. Aaltonen *et al.*, Phys. Rev. Lett. **102**, 222003 (2009) [[arXiv:0903.2850](#) [hep-ex]].
- [31] V. Ahrens, A. Ferroglia, M. Neubert *et al.*, JHEP **1009**, 097 (2010). [[arXiv:1003.5827](#) [hep-ph]].
- [32] [ATLAS Collaboration], ATLAS-CONF-2011-040, March 22, 2011.
- [33] M. Aliev, H. Lacker, U. Langenfeld, S. Moch, P. Uwer, M. Wiedermann, Comput. Phys. Commun. **182**, 1034-1046 (2011). [[arXiv:1007.1327](#) [hep-ph]].
- [34] [ATLAS Collaboration], ATLAS-CONF-2011-087, June 5, 2011.
- [35] S. K. Gupta, A. S. Mete and G. Valencia, Phys. Rev. D **80** (2009) 034013 [[arXiv:0905.1074](#) [hep-ph]].