Probing non-standard top-quark couplings via optimal-observable analyses at LHC

Zenrō Hioki¹,² and Kazumasa Ohkuma³,⁴

¹Institute of Theoretical Physics, University of Tokushima, Tokushima 770-8502, Japan
²Department of Information Science, Fukui University of Technology, Fukui 910-8505, Japan

Abstract. Focusing attention on top-quark pair production and its decay processes at the LHC, non-standard top-quark couplings are studied based on the effective Lagrangian constructed with \( SU(3) \times SU(2) \times U(1) \) invariant dimension-6 operators. The optimal-observable analysis is carried out for the charged-lepton distributions in \( pp \to t\bar{t}X \to \ell^+X' (\ell = e \text{ or } \mu) \) in order to estimate the expected statistical uncertainties in measurements of those non-standard \( t\bar{t}g \) and \( tbW \) couplings that contribute to this process in the leading order.

1 Introduction

Almost twenty years have passed since the top quark was discovered at the Tevatron [1, 2]. Even today, however, this quark is still the heaviest in the experimentally-confirmed elementary particles. Therefore, precise measurements of the top quark are expected to be one of the most promising ways for finding possible signals from new physics beyond the standard model in the current situation where there has been no discovery of non-standard particles [3]. The Large Hadron Collider (LHC) is now the unique facility for those top-quark studies in stead of the Tevatron.

In any processes, non-standard signals mean non-negligible differences between actual experimental data and the corresponding standard-model prediction. Such deviations might be originated from quantum effects of new particles which are not belonging to the standard-model framework. Those quantum effects are usually characterized as new form factors corresponding to the non-standard couplings. The size of the form factors are calculable if you take a specific model which is an extension of the standard model. We are also able to treat those form factors as parameters in the manner of effective-Lagrangian approach, in which its low-energy form reproduces the standard model, is one of the most general methods to describe new-physics phenomena when the energy of our experimental facility is not high enough to produce new particles. Assuming any non-standard-model particles too heavy to appear as real ones, the effective Lagrangian is given as

\[
\mathcal{L}_{\text{eff}} = \mathcal{L}_{\text{SM}} + \frac{1}{A^2} \sum_i \left( C_i O_i + C_i^\ast O_i^\ast \right),
\]

where \( \mathcal{L}_{\text{SM}} \) is the standard-model Lagrangian, \( O_i \) mean \( SU(3) \times SU(2) \times U(1) \) gauge-invariant operators of mass-dimension 6 involving only the standard-model fields [7, 8] and their coefficients \( C_i \) parameterize quantum effects of new particles at an energy less than the assumed new-physics scale \( A \).

2 Strategy

In this section, we explain briefly the effective Lagrangian and the optimal-observable analysis, both of which play important roles in our analysis.

2.1 Effective Lagrangian

The effective-Lagrangian approach, in which its low-energy form reproduces the standard model, is one of the most general methods to describe new-physics phenomena when the energy of our experimental facility is not high enough to produce new particles. Assuming any non-standard-model particles too heavy to appear as real ones, the effective Lagrangian is given as

\[
\mathcal{L}_{\text{eff}} = \mathcal{L}_{\text{SM}} + \frac{1}{A^2} \sum_i \left( C_i O_i + C_i^\ast O_i^\ast \right),
\]

where \( \mathcal{L}_{\text{SM}} \) is the standard-model Lagrangian, \( O_i \) mean \( SU(3) \times SU(2) \times U(1) \) gauge-invariant operators of mass-dimension 6 involving only the standard-model fields [7, 8] and their coefficients \( C_i \) parameterize quantum effects of new particles at an energy less than the assumed new-physics scale \( A \). Since dimension-5 operators violate the lepton-number conservation, they are not treated hereafter. Therefore, we deal with dimension-6 operators, which give the largest contributions in relevant processes.

\[\text{Article available at }\text{http://www.epj-conferences.org/epjconf/20134917010}\]
A full list of the dimension-6 operators were initially presented by Buchmüller and Wyler [7], and also by Arzt et al. [8]. However, it was pointed out that some operators there are related with others through equations of motion, which means that they are not independent of each other [9]. Then, the whole related operators were rearranged to get rid of this redundancy in Refs. [10, 11].

Following the notation of Ref. [10], the relevant effective Lagrangian which describes non-standard interactions of the third generation for the parton-level process \(qq/gg \rightarrow t\bar{t} \rightarrow \ell^+\ell^-X\) is given in [5] as

\[
\mathcal{L}_{\text{eff}} = \mathcal{L}_{\text{dbi}} + \mathcal{L}_{\text{dbW}} : \\
\mathcal{L}_{\text{dbi}} = -\frac{1}{4} g_s \sum_a \left[ \bar{\psi}_i(x) \gamma^\mu \partial_\mu \psi_i(x) \phi(x) G^a_\mu(x) \\
- \bar{\psi}_i(x) \gamma^\mu \partial_\mu \phi(x) G^a_\mu(x) \right], \\
\mathcal{L}_{\text{dbW}} = -\frac{1}{\sqrt{2}} \left[ \bar{\psi}_i(x) \gamma^\mu (f_{1L}^W P_L + f_{1R}^W P_R) \gamma_\mu \phi(x) W^-_\mu(x) + \bar{\psi}_i(x) \gamma^\mu \phi(x) \partial_\mu W^-_\mu(x) \right],
\]

(2)

where \(g_s\) and \(g\) are the \(SU(3)\) and \(SU(2)\) coupling constants, \(P_{L,R} \equiv (1 \mp \gamma_5)/2\), \(d_\ell, \lambda\), and \(f_{1,2}L\) are form factors.

2.2 Optimal-observable analysis

The optimal-observable analysis is a method that could systematically estimate the expected statistical uncertainties of measurable parameters [12–15]. Here we give a brief review of this procedure.

First, assume that we have a cross section

\[
\frac{d\sigma}{d\phi} (\equiv \Sigma_\phi) = \sum_i c_i f_i(\phi),
\]

(6)

where \(f_i(\phi)\) are known functions of the final-state variables \(\phi\) and \(c_i\)'s are model-dependent coefficients. The goal is to determine the \(c_i\)'s. This can be done by using appropriate weighting functions \(w_i(\phi)\) such that \(\int w_i(\phi) \Sigma(\phi) d\phi = c_i\). Then, we determine \(w_i(\phi)\) as the next step. In general, different choices for \(w_i(\phi)\) are possible, but there is a unique choice for which the resultant statistical error is minimized. Such functions are given by

\[
w_i(\phi) = \sum_j X_{ij} f_j(\phi)/\Sigma(\phi),
\]

(7)

where \(X_{ij}\) is the inverse matrix of \(M_{ij}\) which is defined as

\[
M_{ij} \equiv \int d\phi f_i(\phi) f_j(\phi)/\Sigma(\phi).
\]

Finally, using Eqs. (7) and (8), the statistical uncertainty of \(c_i\) is obtained as

\[
|\delta c_i| = \sqrt{X_{ii}} \sigma_T/N,
\]

(9)

where \(\sigma_T \equiv \int (d\sigma/d\phi) d\phi\) and \(N\) is the total number of events.

Let us apply this technique to the process \(pp \rightarrow t\bar{t}X\), expressing its differential cross section (the angular and energy distribution of the charged-lepton \(\ell^+\)) which we derived in [5] as follows:

\[
\frac{d^2\sigma}{dE_d\cos \theta_\ell} = f_{SM}(E_\ell, \cos \theta_\ell) + d_L f_{dL}(E_\ell, \cos \theta_\ell) + d_R f_{dR}(E_\ell, \cos \theta_\ell) + \cdots,
\]

(10)

The explicit forms of \(f_1(E_\ell, \cos \theta_\ell)\) at the parton level are easily found in the relevant formulas in [5].

Since the magnitude of \(d_L\) and \(d_R\) has been shown small [5, 6], we neglect any contribution from terms quadratic (or higher) in those non-standard model parameters hereafter.\(^2\) Note that all \(\Delta\) contributions disappear under this linear approximation, because there is no term proportional to \(\Delta\). Thus, the angular and energy distributions of a decay lepton are written as

\[
\frac{d\sigma_\ell}{dE_\ell d\cos \theta_\ell} = g_1(\cos \theta_\ell) + d_\ell g_2(\cos \theta_\ell),
\]

(12)

\[
\frac{d\sigma_\ell}{dE_\ell} = h_1(E_\ell) + d_\ell h_2(E_\ell) + d_R h_3(E_\ell),
\]

(13)

where \(g_1(\cos \theta_\ell)\) and \(h_1(E_\ell)\) are given by

\[
g_1(\cos \theta_\ell) = \int dE_\ell f_1(E_\ell, \cos \theta_\ell),
\]

\[
h_1(E_\ell) = \int dE_\ell h_1(E_\ell, \cos \theta_\ell)
\]

with \(i = 1, 2\) and 3 corresponding to \(I = SM, d_\ell\) and \(d_R\), respectively. Here should be one comment about the

\(^2\)The quadratic terms of \(f_{1,2}L\) have already been neglected in our previous papers. Indeed, studies at the Tevatron suggest that those contributions are tiny: See [16] for the latest data.
angular distribution: We can thereby probe exclusively the $d_\ell$ term, since any contribution from $d_\ell$ disappears within our approximation as a result of the decoupling theorem found in [17–19].

We are now ready to calculate the following matrices:

$$M^E_{ij} = \sqrt{X^E_{ij}} \frac{\sigma_i(N_i)}{N_t} = \sqrt{X^E_{ij}/L}$$

and their inverse matrices $X^{-1}_{ij}$, all of which are apparently symmetric. Then the statistical uncertainties for the measurements of couplings $d_\ell$ and $d_\ell$ could be estimated by

$$|\delta d_\ell| = \sqrt{X^E_{ij}/N_t} = \sqrt{X^E_{ij}/L}$$

through the angular distribution, and

$$|\delta d_\ell| = \frac{\sqrt{X^E_{ij}/N_t}}{N_t} = \frac{\sqrt{X^E_{ij}/L}}{N_t}$$

via the energy distribution, where $\sigma_i$, $N_t$, and $L$ denote the total cross section, the number of events and the integrated luminosity for the process, respectively.

3 Numerical results and discussion

Below we show the elements of $M^E_{ij}$ computed for $\sqrt{s} = 7, 8, 10$, and $14$ TeV, assuming $m_t = 173$ GeV [20].

1 The angular distribution

(1-1) $\sqrt{s} = 7$ TeV

$$M^E_{11} = 23.102, \quad M^E_{12} = -245.412, \quad M^E_{22} = 2607.340.$$  

(1-2) $\sqrt{s} = 8$ TeV

$$M^E_{11} = 33.234, \quad M^E_{12} = -353.598, \quad M^E_{22} = 3762.753.$$  

(1-3) $\sqrt{s} = 10$ TeV

$$M^E_{11} = 59.333, \quad M^E_{12} = -632.179, \quad M^E_{22} = 6736.735.$$  

(1-4) $\sqrt{s} = 14$ TeV

$$M^E_{11} = 134.052, \quad M^E_{12} = -1428.300, \quad M^E_{22} = 15220.286.$$  

2 The energy distribution

(2-1) $\sqrt{s} = 7$ TeV

$$M^E_{11} = 23.102, \quad M^E_{12} = -245.412, \quad M^E_{13} = 0.000, \quad M^E_{22} = 2607.658, \quad M^E_{23} = -1.974, \quad M^E_{33} = 15.323.$$  

(2-2) $\sqrt{s} = 8$ TeV

$$M^E_{11} = 33.234, \quad M^E_{12} = -353.598, \quad M^E_{13} = 0.000, \quad M^E_{22} = 3763.252, \quad M^E_{23} = -2.880, \quad M^E_{33} = 21.124.$$  

(2-3) $\sqrt{s} = 10$ TeV

$$M^E_{11} = 59.333, \quad M^E_{12} = -632.179, \quad M^E_{13} = 0.000, \quad M^E_{22} = 6737.696, \quad M^E_{23} = -5.120, \quad M^E_{33} = 35.233.$$  

(2-4) $\sqrt{s} = 14$ TeV

$$M^E_{11} = 134.052, \quad M^E_{12} = -1428.300, \quad M^E_{13} = 0.000, \quad M^E_{22} = 15222.443, \quad M^E_{23} = -10.953, \quad M^E_{33} = 72.624.$$  

Here, all $M^E_{13}$ became zero for the same reason as vanishing $d_\ell$ terms in the angular distribution, i.e., the decoupling theorem [17–19]. Using the inverse matrices calculated from the above elements and Eq.(16) – Eq.(18), we can estimate the statistical uncertainties of the relevant couplings.

Before giving the results, however, we should comment about instabilities of inverse-matrix computations: We noticed that the results fluctuate to a certain extent (beyond our expectation) depending on to which decimal we multiply the tree cross sections by the different inputs, we taking (1-1) as an example: For those different inputs, we can estimate the statistical uncertainties of the relevant couplings.

Let us explain these treatments more specifically by taking (1-1) as an example: For those different inputs, we get

$$\sqrt{X^E_{22}} = 1.73 \text{ for } M^E_{11} = 23.102, \cdots, \quad 2.57 \text{ for } M^E_{11} = 23.10, \cdots,$$

which lead to the mean value 2.15, the maximum 2.57, and the minimum 1.73, and resultant errors ±0.42, which are derived by 2.57 – 2.15 and 1.73 – 2.15.

We also have to explain how we can take into account QCD higher-order corrections: All the numerical computations of $M^E_{ij}$ in this section were done with the tree-level formulas. In order to include QCD corrections there, we multiply the tree cross sections by the $K$-factor. This factor disappears in the combination $X^{E,F}_{i,j}$ and remains only in $N_i$ ($=L_{\ell}r_i$) when we estimate $\delta d_\ell$. Therefore the luminosity $L$ discussed in the following should be understood as an effective one including $K$ (and also the lepton detection efficiency $\epsilon$).

We now show the whole results in Table 1 and Table 2 for the angular and energy distributions, respectively. The uncertainties we encountered here are hard to eliminate, but still the results will tell us the necessary luminosity for reaching the precision which we aim to realize. For example, we need at least $L \approx 1500 \text{ pb}^{-1} (= 1.5 \text{ fb}^{-1})$ in order to achieve $|\delta d_\ell| = O(10^{-2})$ in case of measuring the angular distribution at the LHC whose colliding energy is 8 TeV. As mentioned, those $L$ should be divided by $K \approx 1.5$ and $\epsilon$. If we assume $\epsilon = 0.5$, the resultant $L$ increases slightly, which however hardly affects our conclusion.
the necessary luminosity and new-physics scale thereby to a certain extent. However, our conclusion about $ff$-effects do not change the order of the results. Since the angular distribution is a function accuracy of $d\ell$, we got the following knowledge about the determination for performing the optimal-observable analysis. Although the $d_A$ terms disappeared under this approximation, we got the following knowledge about the determination accuracy of $d\ell$ and $d_R$:

- Measuring the energy distribution could study both $d\ell$ and $d_R$ at the same time with a higher precision than the case of the angular distribution.
- Since the angular distribution is affected by only $d\ell$ contributions, one is suitable for exploring the mechanism of top-pair productions exclusively through the precision of $d\ell$ is slightly lower than that from the energy distribution.

Furthermore, it was pointed out that there were some ambiguities in inverse-matrix computations depending on where to round the input data (i.e., $M_{ij}^{\ell\ell}$-elements) off, and consequently estimated statistical errors were affected thereby to a certain extent. However, our conclusion about the necessary luminosity and new-physics scale $\Lambda$ do not receive that serious influence, because those ambiguities do not change the order of the results.

Finally, if the LHC is going to be steadily upgraded and $L = 500$ fb$^{-1}$ is achieved at $\sqrt{s} = 14$ TeV, $d\ell$ and $d_R$ could be determined with $|\hat{d}\ell| \sim O(10^{-3})$ and $|\hat{d}_R| \sim O(10^{-4})$ respectively from the energy distribution: Here, $|\hat{d}\ell| \sim 0.001$ ($|\hat{d}_R| \sim 0.0001$) means that the contribution from effective operator $O_{\ell\ell}^{33}$ ($O_{sw}^{33}$) is suppressed by $\Lambda \gg 7$ (20) TeV if the center value of measured $d\ell$ ($d_R$) is close to zero and $C_{\ell\ell}^{\ell\ell}$ ($C_{sw}^{\ell\ell}$) $\sim 1$, which we estimated from Eqs. (5) and (11).

On the other hand, concerning $d_A$ terms, which were neglected in this work, they induce CP-violating interactions, therefore those effects would be better off being probed via some kind of asymmetric observables [22].

**Acknowledgments**

This work was partly supported by the Grant-in-Aid for Scientific Research No. 22540284 from the Japan Society for the Promotion of Science. Part of the algebraic and numerical calculations were carried out on the computer system at Yukawa Institute for Theoretical Physics (YITP), Kyoto University.

**References**

[1] F. Abe et al. (CDF Collaboration), Phys. Rev. Lett. 74 (1995) 2626 (hep-ex/9503002).
[2] S. Abachi et al. (D0 Collaboration), Phys. Rev. Lett. 74 (1995) 2632 (hep-ex/9503003).
[3] For recent review J.F. Kamenik, J. Shu and J. Zupan, arXiv:1107.5257 [hep-ph].
[4] Z. Hioki and K. Ohkuma, Eur. Phys. J. C 65 (2010) 127 (arXiv:0910.3049 [hep-ph]); Eur. Phys. J. C 71 (2011) 1535 (arXiv:1011.2655 [hep-ph]).
[5] Z. Hioki and K. Ohkuma, Phys. Rev. D 83 (2011) 114045 (arXiv:1104.1221 [hep-ph]).
[6] Z. Hioki and K. Ohkuma, Phys. Lett. B 716 (2012) 310 (arXiv:1206.2413 [hep-ph]).
[7] W. Buchmuller and D. Wyler, Nucl. Phys. B 268 (1986) 621.
[8] C. Arzt, M.B. Einhorn and J. Wudka, Nucl. Phys. B 433 (1995) 41 (hep-ph/9405214).
[9] B. Grzadkowski, Z. Hioki, K. Ohkuma and J. Wudka, Nucl. Phys. B 689 (2004) 108 (hep-ph/0310159).
[10] J.A. Aguilar-Saavedra, Nucl. Phys. B 812 (2009) 181 (arXiv:0811.3842 [hep-ph]); Nucl. Phys. B 821 (2009) 215 (arXiv:0904.2387 [hep-ph]).
[11] B. Grzadkowski, M. Iskrzynski, M. Misiak and J. Rosiek, JHEP 1010 (2010) 085 (arXiv:1008.4884 [hep-ph]).
[12] D. Atwood and A. Soni, Phys. Rev. D 45 (1992) 2405.
[13] M. Davier, L. Duflot, F. Le Diberder and A. Rouge, Phys. Lett. B 306 (1993) 411.
[14] M. Diehl and O. Nachtmann, Z. Phys. C 62 (1994) 397.
[15] J.F. Gunion, B. Grzadkowski and X.-G. He, Phys. Rev. Lett. 77 (1996) 5172 (hep-ph/9605326).
[16] T. Aaltonen et al. [CDF and D0 Collaborations], Phys. Rev. D 85 (2012) 071106 (arXiv:1202.5272 [hep-ex]).
[17] B. Grzadkowski and Z. Hioki, Phys. Lett. B 476 (2000) 87 (hep-ph/9911505); Phys. Lett. B 529 (2002)
82 (hep-ph/0112361); Phys. Lett. B 557 (2003) 55 (hep-ph/0208079).
[18] S.D. Rindani, Pramana 54 (2000) 791 (hep-ph/0002006).
[19] R.M. Godbole, S.D. Rindani and R.K. Singh, JHEP 0612 (2006) 021 (hep-ph/0605100).
[20] O. Brandt [CDF and D0 Collaborations], arXiv:1204.0919 [hep-ex].
[21] Z. Hioki and K. Ohkuma, in preparation.
[22] G. Valencia, PoS HQL 2012 (2012) 050 (arXiv:1301.0962 [hep-ph]).