NEW PHYSICS AT THE TOP*

Bohdan Grzadkowski

Institute of Theoretical Physics, Warsaw University
Hoża 69, 00-681 Warsaw, Poland

Zenrō Hioki

Institute of Theoretical Physics, University of Tokushima
770-8502 Tokushima, Japan

Kazumasa Ohkuma

Department of Information Science, Fukui University of Technology
910-8505 Fukui, Japan

José Wudka

Department of Physics, University of California
Riverside, CA 92521-0413, USA

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We consider the possibility of using $t\bar{t}$ production at photon colliders as a probe for physics beyond the Standard Model. The angular and energy distributions of top-quark decay products are employed in the analysis that determines the accuracy with which the new physics parameters can be measured.

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1. Introduction

The Standard Model (SM) is widely believed to be the low-energy limit of a yet unspecified more fundamental theory. This conviction has led to a strenuous search for any deviation from the SM predictions without yielding any unambiguous indication of new physics effects [1]. This failure opens the possibility that future colliders will be unable to directly produce the heavy excitations. In this case new physics effects might be noticeable in

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precision measurements through virtual effects which are best parameterized in terms of the coefficients α_i of an effective Lagrangian [2]

$$\mathcal{L}_{\rm eff} = \mathcal{L}_{\rm SM} + \sum_i \frac{\alpha_i}{\Lambda^{n_i}} \mathcal{O}_i \,, \tag{1}$$

where $\mathcal{O}_i$ are local gauge-invariant operators involving the SM fields and Λ denotes the scale at which the new physics becomes apparent. The parameters α_i, though unknown, are constrained by naturality.

This approach is model independent and very general, in particular it describes all possible flavor effects that can be produced by heavy new physics. In addition, though the expansion in terms of effective operators is infinite, the theory does have predictability: due to the Λ^{-n} suppression factors only a finite number of operators can produce effects larger than a given experimental accuracy. In the following we will only require that the underlying theory be weakly coupled and decoupling[1].

Careful studies of processes involving light quarks have produced no evidence of deviations from the SM predictions [1]. The top-quark system has also shown no evidence of deviations, but with a much reduced accuracy. For this reason we will study the possibility of detecting new physics effects in processes involving the top quark. We will do so in photon–photon colliders in order to use the flexibility of these machines in specifying the polarization of the incident beams [4].

2. Process under consideration

The specific reaction we will consider is $\gamma\gamma \to t\bar{t} \to (bW^+)(\bar{b}W^-)$ with the $W^\pm$ decaying into light leptons or quarks (see Fig. 1). The vertices are derived in a standard fashion. For example, the $t\bar{t}\gamma$ effective vertex is generated by the operators [5]

$$\mathcal{O}'_{uB} = i\left(\bar{q}\sigma_{\mu\nu}u\right)\tilde{\phi}B^{\mu\nu}\,, \qquad \mathcal{O}_{uW} = i\left(\bar{q}\sigma_{\mu\nu}\tau^I u\right)\tilde{\phi}W_I^{\mu\nu}\,, \tag{2}$$

and their Hermitian conjugates[2]. Explicitly [6],

$$\gamma t\bar{t}:\ \frac{\sqrt{2}}{\Lambda^2}v\ \not{k}\gamma_\mu\left(\alpha_{\gamma 1} + i\alpha_{\gamma 2}\gamma_5\right)\,, \qquad \alpha_{\gamma 1} + i\alpha_{\gamma 2} = s_{\rm W}\alpha_{uW} + c_{\rm W}\alpha'_{uB}\,, \tag{3}$$

[1] The most popular alternative scenario considers the possibility that there is no light Higgs and that the symmetry-breaking mechanism in the SM is produced by a unitary field; such models become strongly-coupled at scales $\sim 3\,$TeV [3] and will not be considered here.

[2] q denotes a up-down left-handed quark iso-doublet, u and d the corresponding right-handed iso-singlets, B and W the U(1) and SU(2) gauge fields and ϕ the scalar iso-doublet.

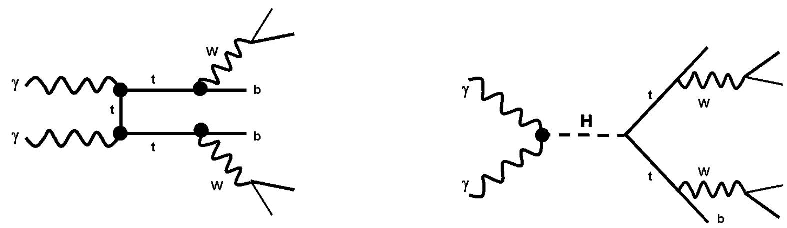

Fig. 1. Reactions being studied (heavy blobs denote SM and effective vertices).

where $s_{\rm W}$ and $c_{\rm W}$ denote sine and cosine of the Weinberg angle, respectively. The coefficient $\alpha_{\gamma 2}$ summarizes all CP-violating effects in this vertex (to order $1/\Lambda^2$ accuracy)[3]. The additional contributing vertices are [6]

$$\begin{aligned} \gamma\gamma H: &\quad -\tfrac{4}{\Lambda^2} v \Big\{ \alpha_{h1} \left[(k_1 \cdot k_2)\, \eta^{\mu\nu} - k_1^\nu k_2^\mu \right] - 2\alpha_{h2} k_1^\rho k_2^\sigma \epsilon_{\rho\sigma}{}^{\mu\nu} \Big\}, \\ Wtb: &\quad -\tfrac{g}{\sqrt{2}} \Big[\gamma^\mu f_1^{\rm L} P_L - \tfrac{1}{m_{\rm W}} i\sigma^{\mu\nu} k_\nu f_2^{\rm R} P_R \,, \Big] \end{aligned} \tag{4}$$

(a total of 8 unknown parameters) with $f_1^{\rm L} - 1,\; f_2^{\rm R} \sim O(v^2/\Lambda^2)$. We will also use the definition $\alpha_d = \Re\left(f_2^{\rm R}\right)$.

3. Cross-section calculation

In obtaining the cross section we adopt the narrow-width approximation for the top-quark propagator. Using then the Kawasaki–Shirafuji–Tsai formalism [7] we find

$$d\sigma(\gamma\gamma \to t\bar{t} \to \ell X,\; bX) \simeq d\sigma(\gamma\gamma \to t(n)\bar{t}) \cdot d\Gamma_{t\to \ell X,\; bX} \,, \tag{5}$$

where $d\sigma(\gamma\gamma \to t(n)\bar{t})$ denotes the $t\bar{t}$-production cross section for t quarks with an appropriately-chosen effective polarization vector n and $d\Gamma_{t\to\ell X,\; bX}$ denotes the unpolarized-decay width.

We also use the following collider parameters [8]:
(i) Incident photons: energy ω_0, average linear-polarization $P_{\rm lin}$ and average helicity $P_{\rm hel}$ (with $0 \le P_{\rm lin}^2 + P_{\rm hel}^2 \le 1$); The azimuthal angle associated with $P_{\rm lin}$ is denoted by φ for one beam and $\tilde{\varphi}$ for the other with $\chi = \varphi - \tilde{\varphi}$.
(ii) Incident electrons: energy E and average longitudinal-polarization P_e.
(iii) Scattered photons: energy $\omega = Ex/(1+x)$ with $0 \le x \le 4.828$ (the upper bound corresponds to the threshold for the reaction $\gamma\gamma \to e^+e^-$).

[3] This list does not include the SM $\gamma t\bar{t}$ term.

4. Results

4.1. Asymmetries

We first considered the asymmetries [9] (see Fig. 2)

$$A_{\rm cir} = \frac{\sigma(++) - \sigma(--)}{\sigma(++) + \sigma(--)}\,, \quad A_{\rm lin} = \frac{\sigma(\chi = +\pi/4) - \sigma(\chi = -\pi/4)}{\sigma(\chi = +\pi/4) + \sigma(\chi = -\pi/4)}\,, \tag{6}$$

where $\sigma(\pm\pm)$ denotes the total cross section with $P_e = P_{\rm hel} = \pm 1$ and $\sigma(\chi)$ that for $P_e = 1,\ P_{\rm lin} = P_{\rm hel} = 1/\sqrt{2}$. The resonance effect is produced by the Higgs contribution: $A_{\rm cir,lin}$ are suppressed for $m_H < 2m_t$ and $m_H > E_{\rm CM\ max}$ (where $E_{\rm CM\ max}$ denotes the maximum CM energy for the hard process) since in these regions the Higgs propagator cannot resonate.

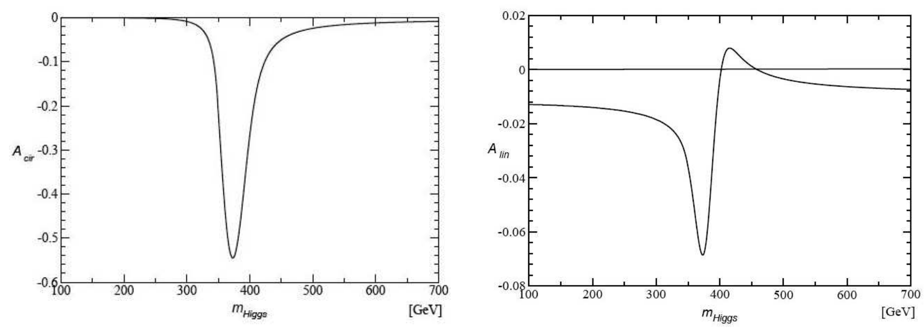

Fig. 2. Circular (left) and linear (right) asymmetries for $\Lambda = 1\,\text{TeV}, \alpha_{\gamma 2}, \alpha_{h2} = 0.1, \sqrt{s} = 0.5\,\text{TeV},\ x = \tilde{x} = 4.828$.

4.2. Optimal-observable analysis

The result of the above calculations is of the form

$$\frac{d\sigma}{d\vartheta} = f_{\rm SM}(\vartheta) + \sum_i \alpha_i f_i(\vartheta) + O(\alpha^2)\,, \tag{7}$$

(ϑ denotes a set of convenient parameters such as the scattering angles and energies of the final products). The idea [10] is to find some observables $\{w_i\}$ that allow us to extract the α_i (by satisfying $\int d\vartheta\, w_i f_j = \delta_{ij}$ and $\int d\vartheta\, w_i f_{\rm SM} = 0$) and minimizing the correlation-matrix $V_{ij} \propto \int d\vartheta\ f_{\rm SM} w_i w_j$. The observables satisfying these conditions are

$$w_i = I_{\rm SM} \sum_j M_{ij}^{-1} \left[\frac{f_j}{f_{\rm SM}} - \frac{I_j}{I_{\rm SM}}\right], \quad V = \frac{I_{\rm SM}}{N} M^{-1} + O(\alpha),$$

$$M_{ij} = \int d\vartheta \, \frac{f_i f_j}{f_{\rm SM}} - \frac{I_i I_j}{I_{\rm SM}} \,, \qquad I_{\rm SM} = \int d\vartheta \, f_{\rm SM}, I_j = \int d\vartheta \, f_j \,, \quad (8)$$

where N denotes the total number of events. For this choice of w_i the statistical uncertainty is given by $\Delta\alpha_i = \sqrt{V_{ii}}$.

We apply this method for the case where one W undergoes leptonic decays while the other decays into quarks. We take $\Lambda = 1\,\text{TeV}$, a luminosity of $500\,\text{fb}^{-1}$ for the e^+e^- collider and the CM energy that is far from the Higgs resonance[4]. We took a sample of choices for the polarization parameters to determine those which provide the greatest sensitivity. Due to numerical instabilities we restricted our calculations to the cases where only 2 or 3 of the α_i are varied (assuming that the remaining parameters are measured using other processes).

3 parameter results ($m_H = 500\,\text{GeV},\ P_e = 0,\ P_{\rm lin} = P_{\rm hel} = 1/\sqrt{2}$)

$$\text{Lepton detection: } \Delta\alpha_{\gamma 2} = \frac{73.4}{\sqrt{N}}\,, \quad \Delta\alpha_{h2} = \frac{8.6}{\sqrt{N}}\,, \quad \Delta\alpha_d = \frac{3.3}{\sqrt{N}}\,.$$

$$\text{Bottom detection: } \Delta\alpha_{\gamma 2} = \frac{125.0}{\sqrt{N}}\,, \quad \Delta\alpha_{h2} = \frac{11.1}{\sqrt{N}}\,, \quad \Delta\alpha_d = \frac{10.7}{\sqrt{N}}\,. \quad (9)$$

2 parameters results

Lepton detection ($P_{\rm lin} = 0\,,\ P_{\rm hel} = -P_e = 1$)

$$\Delta\alpha_{h1} = \frac{3.6}{\sqrt{N}}\,, \quad \Delta\alpha_d = \frac{2.3}{\sqrt{N}}\,, \quad \Delta\alpha_{h2} = \frac{46.5}{\sqrt{N}}\,, \quad \Delta\alpha_d = \frac{2.3}{\sqrt{N}}\,. \quad (10)$$

Bottom detection ($P_{\rm lin} = -P_e = 1\,,\ P_{\rm hel} = 0$)

$$\Delta\alpha_{h1} = \frac{6.4}{\sqrt{N}}\,, \quad \Delta\alpha_d = \frac{3.9}{\sqrt{N}}\,, \quad \Delta\alpha_{h2} = \frac{6.0}{\sqrt{N}}\,, \quad \Delta\alpha_d = \frac{3.0}{\sqrt{N}}\,. \quad (11)$$

5. Discussions

To gauge the meaning of these results, observe that all coefficients except $f_1^{\rm L}$ (the V–A coupling in the Wtb vertex) are necessarily generated by loops in the underlying theory; naturality [11] then demands that $\alpha_i \lesssim 1/(4\pi)^2$. Writing $\Delta\alpha_i = C_i/\sqrt{N}$ the condition $\alpha_i > \Delta\alpha_i$ then implies

$$N > 25\,000\, C_i^2\,, \quad (12)$$

which shows that these measurements will require several year's worth of data (prohibitive for the case of $\alpha_{\gamma 2}$). Even then the extraction of the

[4] The region $s_{\gamma\gamma} \sim m_H^2$ has been studied elsewhere [6,9].

signal will require for the SM parameters to be very well-measured (so that the uncertainty in $f_{\rm SM}$ is significantly smaller than the contribution from each $\alpha_i f_i$). On the theoretical side, the SM contribution should be evaluated to at least the one-loop level.

This gloomy picture is not generic. For example the on-resonance process $\gamma\gamma \to H$ will allow measurements of $\Delta\alpha_{h2}$ to a precision of 10^{-3}–10^{-4} [6]. In addition, e^+e^- colliders [12], having ~ 10 more events will be more sensitive for measuring α_d but less competitive in measuring $\alpha_{\gamma 1}$:

collider	$N_{1\text{-}\sigma}(\alpha_{\gamma_1})$	$N_{1\text{-}\sigma}(\alpha_d)$	$N_{\rm 1year}$
e^+e^-	1.5×10^6	3.6×10^4	1.5×10^5
$\gamma\gamma$	2.3×10^5	5.6×10^4	10^4

where $N_{1\text{-}\sigma}(\alpha_i)$ denotes the number of evens needed for a 1-σ signal in measuring α_i, and $N_{\rm 1year}$ is the expected number of events in one year [5].

This calculation could be improved by imposing appropriate cuts and by considering a different set of observables. Nonetheless the smallness of the α_i suggests that even then the signal will be small and difficult to extract. Note that increasing the energy of the collider will help, though one must not forget that the whole formalism is based on the assumption $\Lambda^2 \gg s$.

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