PHYSICAL REVIEW D **88,** 071301(R) (2013)

Top anomalous magnetic moment and the two-photon decay of the Higgs boson

Lance Labun and Johann Rafelski

Department of Physics, University of Arizona, Tucson, Arizona 85721, USA
and TH Division, Physics Department, CERN, CH-1211 Geneva 23, Switzerland
(Received 29 August 2013; published 16 October 2013)

We compute the dependence of the Higgs to two-photon decay rate $\Gamma_{h\to\gamma\gamma}$ on the top quark gyromagnetic factor g_t in the heavy top limit and evaluate the expected change for one-loop SM correction to g_t. Our results are general and allow consideration of further modifications of g_t, and we predict the resultant $\Gamma_{h\to\gamma\gamma}$.

DOI: 10.1103/PhysRevD.88.071301 PACS numbers: 14.80.Bn, 11.15.Tk, 13.40.Em, 14.65.Ha

I. INTRODUCTION

The recent discovery of a Higgs-like boson [1–4] provides a new opportunity to study the top quark coupling to photons. The reason is that the Higgs, being itself uncharged, couples to photons only through loops [5–7]. The dominant contributions due to the W boson and top quark are illustrated in Fig. 1. In determination of the relevant contributions from fermions, there is a competition between the fermion mass, which reduces loop strength, and the fermion-Higgs coupling, which is proportional to mass. Therefore, the top quark contributes a significant fraction of the Higgs–two-photon amplitude.

The possible deviation of the decay rate $\Gamma_{h\to\gamma\gamma}$ from the standard model (SM) prediction has generated a lot of interest in finding a mechanism for a deviation in $h\to\gamma\gamma$ that minimally affects other SM results. We contribute to this objective by evaluating the Higgs decay rate $\Gamma_{h\to\gamma\gamma}$ as a function of the top quark gyromagnetic ratio g_t, related in the conventional way to the particle magnetic moment

$$\mu_t = \frac{g_t}{2}\frac{Qe}{2m_t}, \qquad Qe = +(2/3)e. \tag{1}$$

The magnetic moment is neither fixed nor protected by a conservation law: the Dirac gyromagnetic ratio $g_{\rm D} = 2$ for bare pointlike fermions is well known to be modified by quantum corrections. The coupling of the top to the Higgs, $y_{\rm th} = m_t\sqrt{2}/v \simeq 0.99$ [in which the Higgs vacuum expectation value $v = 246.2$ GeV $\simeq (G_F\sqrt{2})^{-1/2}$, with G_F the Fermi constant and $m_t = 173.4$ GeV] could herald top structure capable of greatly altering the value of g_t.

The lowest-order SM calculation of g_t yields at one loop the following three contributions: (a) the standard QED result, (b) a similar QCD result [8,9], and (c) an electroweak contribution,

$$(g_t-2)^{(1L)}_{\rm QED} = \alpha_{\rm QED}/\pi = 2.5\times 10^{-3} \tag{2a}$$

$$(g_t-2)^{(1L)}_{\rm QCD} = \frac{N_c^2-1}{2N_c}\frac{\alpha_s}{\pi} = 25\times 10^{-3} \tag{2b}$$

$$(g_t-2)^{(1L)}_{\rm EW} = 7.5\times 10^{-3}. \tag{2c}$$

To obtain the electroweak (EW) numerical value, we note that all vertex corrections have the same loop integral. Therefore, we can use Eq. (3) of [10] in combination with their Eq. (1), which identifies their $\Delta\kappa \to g_t - 2$. Both the virtual gluon and Higgs exchange corrections to g_t can be relatively large considering the coupling strengths $\alpha_s(m_t) \simeq 0.11$ and $y_{\rm th}^2/4\pi = 0.08$.

Furthermore, composite particles can have gyromagnetic ratios unrelated to $g_{\rm D}$. The sensitivity of the magnetic moment to compositeness is well known [11]. Determining μ_t will provide important information about the structure of the top, whether it is pointlike or composite and whether or not it couples with particles beyond the standard model (BSM) [12–14].

Many efforts have focused on relating the anomalous moments to top quark production and decay [15–19], and the radiative decay of the bottom to strange quark $b\to s\gamma$ has also long been studied in search of indirect constraints [20,21] on g_t. The experimental data continue to be analyzed to place constraints on the dipole moments [13,14] and compositeness [22]. Recently, new opportunities provided by Higgs production and decay [23–26] have been recognized.

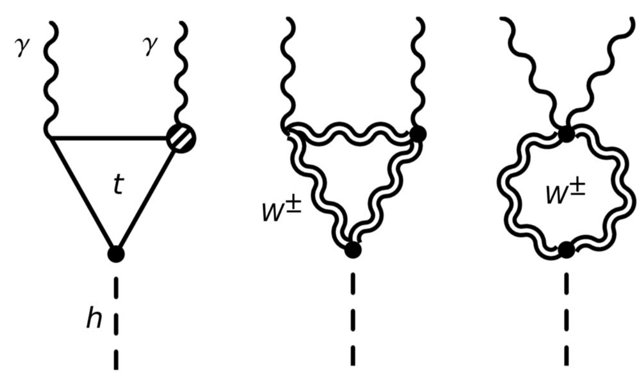

FIG. 1. Diagrams giving dominant contribution to the Higgs h decay to two photons γ. In the top loop at left, we denote the top-photon vertex with a shaded circle to signify that we consider a general value of the top gyromagnetic factor.

The Higgs–two-photon decay amplitude receives a contribution from the top given by the low-energy theorems of [5,6], evaluated employing the effective action

$$\mathcal{L}_{h\gamma\gamma} = -b_0 \frac{\alpha}{2\pi} \frac{h}{v} \frac{1}{4} F^{\mu\nu} F_{\mu\nu} \tag{3}$$

to leading order in the electromagnetic coupling constant α and in the ratio $m_h/2m_t \simeq 0.36$ for $m_h \simeq 125.5$ GeV. Here $F^{\mu\nu}$ is the electromagnetic field strength tensor, and b_0 is the one-loop beta-function coefficient in the expansion of the renormalization group, $-d\alpha/d\ln\lambda \equiv -\beta(\alpha) = (b_0/2\pi)\alpha^2 + (b_1/4\pi^2)\alpha^3 \ldots$, where the scale λ is fixed by measuring the coupling $\alpha(\lambda)$. In the form Eq. (3) the combination $\beta(\alpha) F^{\mu\nu} F_{\mu\nu}/4\alpha$ is renormalization group invariant to one loop, and hence no logarithmic corrections arise [26].

In fact, the effective Lagrangian Eq. (3) gives the amplitude to within 3% of the loop calculation [7]. Therefore, by obtaining $b_0(g_t)$, which further displays the effect of the top quark magnetic moment on the polarization of the photon, we also obtain the $h \to \gamma\gamma$ amplitude to required precision. Corrections in powers of $1/m_t$ [27] can be computed separately, as they are for the Higgs-gluon effective coupling, see e.g. [28].

II. DESCRIBING TOP ANOMALOUS MAGNETIC MOMENT

The difficulty of handling an effective non-Dirac magnetic moment is well known. We are aware of the following three possible approaches:

(1) The perturbative evaluation of g_t can be continued, summing contributions in the theory based at $g_t = 2$, adding higher-order diagrams to the top loop seen in Fig. 1. It is known that perturbative expansion is slowly convergent [8]. Moreover, this approach excludes the possibility of testing for composite structure by considering arbitrary g_t.

(2) One can complement the Dirac action for the top quark with an effective operator of mass dimension 5 (or 6, if before electroweak symmetry breaking). This procedure introduces by dimensional counting a logarithmic divergence. The logarithm can be connected with the running of the dimensionful coupling constant from the BSM scale to the Higgs scale, and the effect of this running is being discussed by several groups [24–26].

(3) Here we consider the top quark as effectively pointlike at the scale of study, m_h. Therefore we employ a description that is independent of additional scales. This is achieved by introducing the second-order theory of fermions, in essence 'squaring' the Dirac operator following a method first proposed by Schwinger [29]. The resulting doubling of Fermi degrees of freedom can be dealt with by including where appropriate a factor $1/2$. However, further refinement of the method is necessary for $g_t > 2$ [30].

The action we consider for the top coupled to the electromagnetic field is

$$\mathcal{L}_t = \bar{t}\left(\Pi^2 - m_t^2 - \frac{g_t}{2}\frac{e\sigma_{\mu\nu}F^{\mu\nu}}{2}\right)t \tag{4}$$

with $\Pi^\mu = i\partial^\mu - eA^\mu$ the Hermitian momentum operator. The fermion fields have mass dimension 1. Hence the magnetic moment interaction seen in Eq. (4) is a dimension-4 operator. Because the arbitrary magnetic moment is renormalizable in this framework, we obtain a finite well-determined result for the low-energy effective coupling b_0 for arbitrary g_t, and hence for the $h \to \gamma\gamma$ amplitude.

Equation (4) has many early mentions in literature discussed in [31]. At the point $g_t = 2$, the second-order theory produces exactly equivalent results to the first-order Dirac action for quantum electrodynamics (QED) (results agree up to the factor $1/2$ mentioned above needed to correct the number of degrees of freedom e.g. in fermion loops), and to one loop in perturbation theory it has been shown to be renormalizable for all g_t [32]. A second-order formulation of the entire standard model has recently been discussed [33].

Using a second-order theory of the top, there is a further diagram in Fig. 1 corresponding to two-top–two-photon coupling analogous to the W loop in (c). This contribution is included when deriving the low-energy theorem and evaluating the top loop amplitude but is not shown explicitly in Fig. 1.

III. g_t-DEPENDENT EFFECTIVE COUPLING

The one-loop β-function coefficient b_0 has been obtained in the second-order formulation for a fermion with arbitrary magnetic moment using the following two independent methods: (a) perturbative computation [31] and (b) the external field method [30]. The perturbative result for the vacuum polarization top-quark loop is [Eq. (3.24) of Ref. [32]],

$$b_0(g_t) = -\frac{4}{3} N_c Q^2 \left(\frac{3}{8} g_t^2 - \frac{1}{2}\right). \tag{5}$$

The $Q^2 = 4/9$ arises from the charge of the top, and $N_c = 3$ for the color trace. The $-4/3$ factor separated in front is the well-known value of b_0 for a fermion with unit charge at $g_t = 2$.

The low-energy theorem Eq. (3) was originally derived from the form of the effective potential obtained in the external field method [5]. When using the external-field method, an in-depth study of $|g_t| > 2$ is necessary following the observation that the Schwinger proper-time evaluation of the effective action [29] when adapted to arbitrary value of g_t [34] diverges for $|g_t| > 2$. This suggests that the convergence radius of the perturbative $b_0(g_t)$, Eq. (5), is $|g_t| \leq 2$.

A convergent result for the effective potential is achieved when the doubled degrees of freedom are properly separated into two half-Hilbert spaces for $|g_t| > 2$ [30]. The separation accomplishes stability and Lorentz invariance of the top quark vacuum, and the procedure is repeated at $g_t^{(N)} = 2 + 4N$ for each integer N. At these values $g_t^{(N)}$, there is a countably infinite number of level crossings occurring between states belonging to the two half-Hilbert spaces. As a result of this procedure, $b_0(g_t)$ is found to be a periodic function, repeating the fundamental domain $-2 \leq g_t \leq 2$ for $|g_t| > 2$, see in Eq. (9) of Ref. [30], and for further discussion see [35],

$$b_0(g_t) = -\frac{4}{3}N_c Q^2\left(\frac{3}{8}(g_t - 4N)^2 - \frac{1}{2}\right) - 2 \leq g_t - 4N \leq 2,\quad N \in \mathbb{Z}. \tag{6}$$

The perturbative result for $b_0(g_t)$ Eq. (5) corresponds to fixing $N = 0$ and ignoring the periodicity.

An important and unexpected observation about $b_0(g_t)$, first made in Ref. [30], is that it changes sign at $g_t - 4N = \pm 2/\sqrt{3}$ and hence is positive for $|g_t - 4N| < 2/\sqrt{3}$. For the perturbative evaluation of $b_0(g_t)$, this sign change occurs twice, when $N = 0$. The consideration of diamagnetic (originating in Π^2) and paramagnetic (originating in $\sigma_{\mu\nu}F^{\mu\nu}$) contributions in Eq. (6) reveals that this effect is due to the decrease in strength of the paramagnetic terms [36]. As g_t^2 decreases, the paramagnetic component in $b_0(g_t)$ diminishes, and the 'asymptotic safety' disappears, see Eq. (10) of [37].

IV. HIGGS TO TWO-PHOTON RATE

The top quark loop contribution to the $h \to \gamma\gamma$ decay is derived from Eq. (3) by inserting photon momentum and polarization vectors for each $F^{\mu\nu}$ in Eq. (3) [5] (see also [6]),

$$A_t(h \to \gamma\gamma) \simeq \frac{1}{v}\frac{\alpha b_0}{4\pi}(k_1^\kappa \epsilon_1^\lambda - k_1^\lambda \epsilon_1^\kappa)(k_2^\kappa \epsilon_2^\lambda - k_2^\lambda \epsilon_2^\kappa). \tag{7}$$

$k_{1,2}^\kappa$ and $\epsilon_{1,2}^\lambda$ are the 4-momenta and polarization vector of the photons. Evaluating the amplitude from Eq. (7) means an error of a few percent relative to the result from the loop amplitude [5,7] for the value $(m_h/2m_t)^2 \simeq 0.13$.

The dependence of the total $h \to \gamma\gamma$ decay amplitude on the top magnetic moment is found by combining the g-dependent top-loop contribution with the contribution from the W-boson loop. Inserting the g-dependent β function [Eq. (6)] in Eq. (7), the total amplitude for Higgs decay into two photons is

$$A_{\rm tot}(h \to \gamma\gamma) \simeq A_t(h \to \gamma\gamma) + A_W(h \to \gamma\gamma), \tag{8}$$

with the W loop contributing [5,7]

$$A_W(h \to \gamma\gamma) = f_W \frac{1}{v}\frac{\alpha}{4\pi}(k_1^\kappa \epsilon_1^\lambda - k_1^\lambda \epsilon_1^\kappa)(k_2^\kappa \epsilon_2^\lambda - k_2^\lambda \epsilon_2^\kappa),$$
$$f_W(x) = 3x(2 - x)(\arcsin(x^{-1/2}))^2 + 3x + 2,$$
$$x = \frac{4m_W^2}{m_h^2} \tag{9}$$

for which we have stated only the relevant case $4m_W^2/m_h^2 \equiv x > 1$ from [5], and we use $x = 1.641$ corresponding to $m_h = 125.5$ GeV. A similar function $f_q(x)$ (see Eq. (21) in [5]) derived from the fermion loop generalizes the low- energy Eq. (7) and with $x_t = 4m_t^2/m_h^2 \to 7.64$ gives a numerical value differing by less than 4% from the low-energy $(x \to \infty)$ result, as mentioned above.

To compare with the measured decay rate, we evaluate the decay width (Eq. (3) of [5] and Eq. (7) of [7]),

$$\Gamma_{h\to\gamma\gamma} \simeq \left| f_W\left(\frac{4m_W^2}{m_h^2}\right) + b_0(g_t)\right|^2 \left(\frac{\alpha}{4\pi}\right)^2 \frac{m_h^3}{16\pi v^2}, \tag{10}$$

with $f_W(x)$ given in Eq. (9) and $b_0(g)$ from Eq. (6). Equation (10) provides for $g_t \to 2$ the Higgs to two-photon decay rate within a few percent of the width stated in the 2011 updated Higgs Cross Section Working Group tables (which include next-to-leading-order QCD and electroweak corrections), after accounting for partial decay widths $h \to ZZ$, $h \to WW$ and to 4-fermions, see Eq. (1) of [38].

Figure 2 shows the g_t dependence of the total $h \to \gamma\gamma$ decay rate, normalized to its value at $g_t = g_D = 2$ for both nonperturbative and perturbative evaluation. Considering the periodicity of $b_0(g_t)$, the rate $\Gamma_{h\to\gamma\gamma}$ is always enhanced in the nonperturbative approach, as compared to $g = g_{\rm D}$ (and its periodic recurrences i.e. $g_t = \pm 2, \pm 6, \pm 10$, etc).

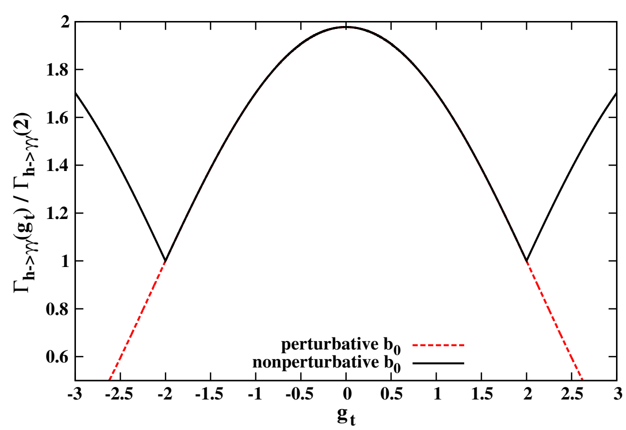

FIG. 2 (color online). Higgs to two-photon $h \to \gamma\gamma$ decay rate Eq. (10) normalized to its value at $g_t = 2$ for $x = 4m_W^2/m_h^2 = 1.64$, $m_h = 125.5$ GeV. The dashed (red) line corresponds to the perturbative evaluation [31] of $b_0(g_t)$, Eq. (5), and the solid (black) line corresponds to the external field method [30] result, Eq. (6).

However, the prediction using the perturbative computation of $b_0(g_t)$ implies for $|g_t| > 2$ a suppression of $h \to \gamma\gamma$.

The line of reasoning here applied to $h \to \gamma\gamma$ also applies to Higgs decays into two gluons, $h \to GG$, because the hGG interaction is similarly an effective interaction involving a top quark loop and, to a smaller degree, a bottom quark loop. In contrast to the $h \to \gamma\gamma$ case, in the GG decay channel the top quark loop gives the dominant contribution. The $h \to GG$ dependence on chromomagnetic moment g_t^c and related experimental opportunities are discussed in [35]. Although similar SM and BSM physics can contribute to both g_t^c and g_t, it is important to remember that there is no immediate relation between the specific values of g_t^c and g_t.

V. COMPARISON OF g_t WITH EXPERIMENT

Experimental input on the magnitude of g_t is from here considered Higgs decay, the radiative $b \to s\gamma$ decay, and top production.

Higgs $\to \gamma\gamma$ *decay*: This study was prompted by reports of possible enhancement of the Higgs candidate $\to \gamma\gamma$ decay: ATLAS recently released a report of a $\sim$55% enhancement over the SM prediction in the $h \to \gamma\gamma$ rate [2]. The combined Tevatron results [39] also find rate enhancement at 1σ level. CMS initially reported an enhancement [3], but in a later update [4] concludes there is a possible slight suppression, agreeing at 1σ with the SM prediction.

Based on the one-loop result, Eq. (2), for g_t and evaluating the amplitude, Eq. (8), we find not more than a 1.0% top decay amplitude modification and hence a 2.0% modification of the decay rate, Eq. (10). This suggests that the SM perturbative value of $g_t - 2$ predicts too small a modification (a few percent) of decay rate to explain a potential enhancement as reported by ATLAS. However, there are further contributions of higher-order QCD and virtual Higgs which have not been considered in Eq. (2), as well as BSM effects that may be visible in g_t.

The radiative $b \to s\gamma$ *decay:* The earlier work [20,21] was recently updated, see Eq. (9) in [13],

$$-1.83 < \Delta\mu_t m_t < 0.53, \qquad \frac{\Delta\mu_t}{2} = \frac{g_t - 2}{2}\,\frac{\frac{2}{3}e}{2m_t}. \tag{11}$$

In the above we have not considered the electric dipole moment term, which is constrained to be $\sim 10^6$ times smaller than μ_t. We consider the numbers in Eq. (9) of [13] dimensionless and translating into a bound on g_t,

$$-3.49 < g_t < 3.59. \tag{12}$$

This is therefore not constraining in our consideration, as it is consistent with the full range of g_t factors suggested by the enhancement of the rate shown in Fig. 2.

Top production: Limits on g_t should arise from the study of top production. In hadron colliders, top production is predominantly via strong interactions, so the LHC is more sensitive to the top chromomagnetic moment, i.e. the corresponding g_t^c (c for chromodynamic factor) [40], and g_t^c could perhaps be connected using Schwinger-Dyson equations to g_t. Only a future direct study of the top production in e^+e^- collisions appears to offer another sensitive measure of the QED top anomalous magnetic dipole moment [15,41].

VI. SUMMARY AND CONCLUSIONS

We have demonstrated that the decay rate $h \to \gamma\gamma$ depends significantly on the top quark gyromagnetic ratio g_t, presenting the following two results:

(1) The perturbative result follows from (a) the separation of b_0 into its paramagnetic and diamagnetic components (see Sec. 12.3.3 of [42]) and (b) can be obtained in explicit Feynman diagram evaluation [32]. Using this perturbative form, we presented the resultant perturbative $h \to \gamma\gamma$ decay rate.

(2) However, one must expect that the perturbative evaluation has a finite radius of convergence because the strength of the paramagnetic interaction increases as g_t^2. In the external-field method we identified a divergence of effective action [34] for $|g_t| > 2$. We believe that the appearance of this divergence sets the radius of convergence of the perturbative method at $|g_t| \le 2$. The proposed solution, yielding a finite effective action [30], leads to periodic behavior of $h \to \gamma\gamma$ decay rate as a function of g_t valid for all g_t.

This nonperturbative result, Eq. (6), implies an enhancement of $h \to \gamma\gamma$ for any $g \neq 2$, see Fig. 2. The perturbative evaluation in the domain $|g_t| > 2$ produces a suppression of the $\Gamma_{h\to\gamma\gamma}$ in the SM expected $g_t - 2 > 0$ domain. Both results are finite since we are using the "squared" Dirac operator in computations. In this second-order theory of the top quark, the magnetic moment interaction is renormalizable for any g_t, and no discussion of the scale dependence or running enters [24–26].

Irrespective of whether the perturbative or the external-field result applies, the rate $\Gamma_{h\to\gamma\gamma}$ offers the greatest sensitivity to top quark magnetic properties. When complemented with further independent measurements of g_t, a study of $\Gamma_{h\to\gamma\gamma}$ could experimentally settle the theoretical question about the behavior of $\beta(g_t)$ for large g_t. In view of the g_t dependence of the rate $\Gamma_{h\to\gamma\gamma}$ presented in Fig. 2, the Higgs-two-photon decay $h \to \gamma\gamma$ may be offering the most precise presently available information on g_t, which will reflect composite or other top quark structure and thus could be a sensitive probe of BSM physics.

A currently visible value $g_t \neq 2$ implies appearance of BSM top quark structure and the $h \to \gamma\gamma$ decay turns out to be an excellent probe of this. The question arises how $g_t \neq 2$ affects all the other observables; however, it is beyond the scope of this work to study the slight tension that exists between the directly measured Higgs mass value, $m_h = 125.5$ GeV, and the fit to precision electroweak experimental data that predicted a considerably smaller Higgs mass [43], $m_h \simeq 100$ GeV.

ACKNOWLEDGMENTS

We thank the TH division of the CERN Physics Department for hospitality while much of this work was done in 2012. This work was supported by the U.S. Department of Energy, Grant No. DE-FG02-04ER41318.

[1] G. Aad *et al.* (ATLAS Collaboration), Phys. Lett. B **716**, 1 (2012).
[2] G. Aad *et al.* (ATLAS Collaboration), Phys. Lett. B **726**, 88 (2013).
[3] S. Chatrchyan *et al.* (CMS Collaboration), Phys. Lett. B **716**, 30 (2012).
[4] CMS Collaboration, Note CMS-PAS-HIG-13-001.
[5] M. A. Shifman, A. I. Vainshtein, M. B. Voloshin, and V. I. Zakharov, Yad. Fiz. **30**, 1368 (1979) [Sov. J. Nucl. Phys. **30**, 711 (1979)].
[6] J. R. Ellis, M. K. Gaillard and D. V. Nanopoulos, Nucl. Phys. **B106**, 292 (1976).
[7] W. J. Marciano, C. Zhang, and S. Willenbrock, Phys. Rev. D **85**, 013002 (2012).
[8] W. Bernreuther, R. Bonciani, T. Gehrmann, R. Heinesch, T. Leineweber, and E. Remiddi, Nucl. Phys. **B706**, 245 (2005).
[9] W. Bernreuther, R. Bonciani, T. Gehrmann, R. Heinesch, T. Leineweber, P. Mastrolia, and E. Remiddi, Phys. Rev. Lett. **95**, 261802 (2005).
[10] R. Martinez, M. A. Perez, and N. Poveda, Eur. Phys. J. C **53**, 221 (2008).
[11] S. J. Brodsky and S. D. Drell, Phys. Rev. D **22**, 2236 (1980).
[12] W. Bernreuther, J. Phys. G **35**, 083001 (2008).
[13] J. F. Kamenik, M. Papucci, and A. Weiler, Phys. Rev. D **85**, 071501 (2012); **88**, 039903(E) (2013).
[14] A. O. Bouzas and F. Larios, Phys. Rev. D **87**, 074015 (2013).
[15] D. Atwood and A. Soni, Phys. Rev. D **45**, 2405 (1992).
[16] D. Atwood, A. Kagan, and T. G. Rizzo, Phys. Rev. D **52**, 6264 (1995).
[17] P. Haberl, O. Nachtmann, and A. Wilch, Phys. Rev. D **53**, 4875 (1996).
[18] C. Zhang and S. Willenbrock, Phys. Rev. D **83**, 034006 (2011).
[19] A. J. Larkoski and M. E. Peskin, Phys. Rev. D **83**, 034012 (2011).
[20] J. L. Hewett and T. G. Rizzo, Phys. Rev. D **49**, 319 (1994).
[21] R. Martinez and J. A. Rodriguez, Phys. Rev. D **55**, 3212 (1997).
[22] M. Fabbrichesi, M. Pinamonti, and A. Tonero, arXiv:1307.5750.
[23] D. Choudhury and P. Saha, J. High Energy Phys. 08 (2012) 144.
[24] C. Degrande, J. M. Gerard, C. Grojean, F. Maltoni, and G. Servant, J. High Energy Phys. 07 (2012) 036; 03 (2013) 032.
[25] C. Grojean, E. E. Jenkins, A. V. Manohar, and M. Trott, J. High Energy Phys. 04 (2013) 016.
[26] J. Elias-Miro, J. R. Espinosa, E. Masso, and A. Pomarol, J. High Energy Phys. 08 (2013) 033.
[27] W. Bernreuther, Ann. Phys. (N.Y.) **151**, 127 (1983).
[28] R. V. Harlander and K. J. Ozeren, J. High Energy Phys. 11 (2009) 088.
[29] J. S. Schwinger, Phys. Rev. **82**, 664 (1951).
[30] J. Rafelski and L. Labun, arXiv:1205.1835.
[31] R. Angeles-Martinez and M. Napsuciale, Phys. Rev. D **85**, 076004 (2012).
[32] C. A. Vaquera-Araujo, M. Napsuciale, and R. Angeles-Martinez, J. High Energy Phys. 01 (2013) 011.
[33] J. Espin and K. Krasnov, arXiv:1308.1278.
[34] L. Labun and J. Rafelski, Phys. Rev. D **86**, 041701 (2012).
[35] L. Labun and J. Rafelski, arXiv:1210.3150.
[36] A. Nink and M. Reuter, J. High Energy Phys. 01 (2013) 062.
[37] A. Nink and M. Reuter, Int. J. Mod. Phys. D **22**, 1330008 (2013).
[38] A. Denner, S. Heinemeyer, I. Puljak, D. Rebuzzi, and M. Spira, Eur. Phys. J. C **71**, 1753 (2011).
[39] T. Aaltonen *et al.* (CDF and D0 Collaborations), Phys. Rev. D **88**, 052014 (2013).
[40] Z. Hioki and K. Ohkuma, Eur. Phys. J. C **65**, 127 (2010); **71**, 1535 (2011).
[41] E. Devetak, A. Nomerotski, and M. Peskin, Phys. Rev. D **84**, 034029 (2011).
[42] K. Huang, *Quarks, Leptons & Gauge Fields* (World Scientific, Singapore, 1992).
[43] A. Djouadi, Phys. Rep. **457**, 1 (2008).