CORRECTIONS TO (PSEUDO)SCALARS DECAY INTO A FERMION PAIR FROM GRAVITATIONAL TORSION*

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We study the contribution of the torsion-descendent four-fermion contact interaction to the decay width of a neutral (pseudo)scalar field into a fermion pair. This new interaction comes from the existence of gravitational torsion in models with extra dimensions. Additionally, we exemplify the formalism by studying two cases: first, the variation of the considered branching ratio of the Higgs in the context of the Standard Model, and second the proper variations of the scalar and pseudoscalar fields of the type II-1 two-Higgs doublets model.

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1. Introduction

The discovery of the Higgs boson in 2012 at the Large Hadron Collider (LHC) [1–3] not only has completed the picture of the Standard Model (SM), but also has opened the possibility of real existence of fundamental scalar fields in nature. At the same time, some of the puzzles in the SM, such as neutrino mass generation and dark matter, have stimulated the scientific community to consider models with a larger scalar sector [4–16]. Extended scalar sector is also *predicted* by models such as supersymmetry [17–19], some versions of strong electroweak symmetry breaking models [20] and non-minimal composite Higgs models [21–23]. Although no deviation from the SM has yet been observed, the LHCb Collaboration has reported

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anomalies in the Lepton Flavour Universality violating ratios, R_K and R_{K^*}. These anomalies can be explained via models that include new heavy vector and scalar bosons [24–26].

At the same time, there have been other extensions of the SM motivated by the possible existence of more than three spatial dimensions [27–33]. In these scenarios, it is tempting to consider (in the bulk) an extended gravitational sector. Indeed, Einstein's theory of gravity, known as General Relativity (GR), is now viewed as a low-energy effective theory of a (yet unknown) fundamental model, in particular due to the lack of a consistent quantum version of the theory[1]. In an effort to obtain a more fundamental theory of gravity, several generalizations of GR have been proposed, from the *minimal* generalization of considering a metric compatible affine connection [41–44], models which keep the precepts of GR but in higher dimensions [45, 46], metric-affine theories [47], affine theories [48–54], models with higher order in curvature and/or torsion [55–61], *etc.*

In this letter, we shall only consider Cartan's generalization to GR, usually called Einstein–Cartan theory of gravity (ECT), in which the torsion turns out to be a non-dynamical field, and it can be integrated out of the system. When the ECT of gravity is coupled with fermionic matter, the integration of the torsion induces an effective four-fermion contact interaction [62–65], whose phenomenological effects can be observed at accelerators. It is well-known that such an induced effective interaction is strongly suppressed because it is diminished by Newton's constant, or in other words, by the inverse of the Planck mass squared. However, there are scenarios with extra dimensions which achieve naturalness between the electroweak, M_W, and the (fundamental) gravitational scales, M_*, while the known Planck's mass, $M_{\rm Pl}$, is an enhanced effective gravitational scale [66–70]. Therefore, the suppression of the torsion-descendent four-fermion interaction is not as dramatic.

Among the phenomenological aspects which can be observed from the induced four-fermion interaction, one can name the following: several cosmological problems [71–76], the origin of fermion masses [77], neutrino oscillation [78–80], impose limits on extra dimensional model [81–83], and changing one-loop observable [84, 85].

A possible effect of this four-fermion interaction is to modify, through one-loop effects, the decay width of generic (pseudo)scalar bosons into a pair of SM fermions. The aforementioned is applicable, for example, to the Higgs decay. This deviation from the standard predictions could be observed in principle, by means of precision measurements performed in future lepton colliders as the International Linear Collider (ILC) or the Compact Linear Collider (CLIC).

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[1] There are several attempts of quantize the gravitational interactions, see, for example, Refs. [34–39]. For a historical review, see Ref. [40].

The letter is organized as follows. A brief review of the theoretical setup is presented in Sec. 2. In Sec. 3, we show the one-loop corrections due to the effective interaction to the decay width for a (pseudo)scalar boson decaying into a fermion pair. In Sec. 4 and Sec. 5, we apply our result to the SM Higgs, and to the (pseudo)scalars in the framework of two-Higgs doublets model (2HDM), decaying into $\tau^+\tau^-$ and $b\bar{b}$ final states. Finally, in Sec. 6, we present a discussion of the results and the concluding remarks.

2. Effective interaction through gravitational torsion

The standard GR is interpreted as a field theory for the metric. Since the field equations for the metric are of second order, the approach is known as *second order formalism*. However, even in standard GR — where the connection is a metric potential — one can treat the metric and the connection as independent fields, and their field equations are then first order differential equations. This latter approach is called *first order formalism* or sometimes *Palatini's formalism* [86]. Although Palatini's approach can be used with the metric and connection fields, it is useful to consider an equivalent set of fields known as the *vielbein* (e^a_μ) and the *spin connection* (${\omega_\mu}^{ab}$) [2] which encode the information of how to *translate* from the curved spacetime to the tangent space, and how these tangent spaces are *connected* with those of the neighbourhood points[3]. The equivalence between the metric and the vielbein is given by

$$g_{\mu\nu} = \eta_{ab} e^a_\mu e^b_\nu \,. \tag{1}$$

Despite one can write down an explicit relation between the Christoffel connection and the spin connection, we omit it. Instead, we present the equations that define the torsion and curvature two-forms[4] *i.e.* the Cartan structure equations

$$\mathbf{d}\boldsymbol{e}^a + {\boldsymbol{\omega}^a}_b \wedge \boldsymbol{e}^b = \boldsymbol{\mathcal{T}}^a \qquad \text{and} \qquad \mathbf{d}\boldsymbol{\omega}^{ab} + {\boldsymbol{\omega}^a}_c \wedge {\boldsymbol{\omega}^b}_c = \boldsymbol{\mathcal{R}}^{ab} \,. \tag{2}$$

The vielbein and spin connection one-forms are defined as

$$\boldsymbol{e}^a = e^a_\mu \mathrm{d}x^\mu \qquad \text{and} \qquad \boldsymbol{\omega}^{ab} = {\omega_\mu}^{ab} \mathrm{d}x^\mu \,, \tag{3}$$

while the two-forms are written explicitly in components as

$$\boldsymbol{\mathcal{T}}^a = \tfrac{1}{2} {\mathcal{T}_\mu}^a{}_\nu \, \mathrm{d}x^\mu \wedge \mathrm{d}x^\nu = \tfrac{1}{2} {\mathcal{T}_m}^a{}_n \, \boldsymbol{e}^m \wedge \boldsymbol{e}^n$$

and

$$\boldsymbol{\mathcal{R}}^a{}_b = \tfrac{1}{2} {\mathcal{R}_{\mu\nu}}^a{}_b \, \mathrm{d}x^\mu \wedge \mathrm{d}x^\nu = \tfrac{1}{2} {\mathcal{R}_{mn}}^a{}_b \, \boldsymbol{e}^m \wedge \boldsymbol{e}^n \,. \tag{4}$$

[2] The name "spin connection" is historical, and it is not necessarily related with the spin of the fields. For this reason, some authors prefer to call it *Lorentz connection.*

[3] The vielbein field ensures the validity of the equivalence principle.

[4] We make extensive use of the formalism of differential forms [87–90].

In the following, in order to distinguish among quantities in higher or four-dimensional spacetimes, we shall use the notation defined in Refs. [77, 83, 85], where hatted quantities refer to objects (or indices) lying in the former, while unhatted quantities refer to objects (or indices) lying in the latter. It is worth to mention that we denote by γ^* the four-dimensional chiral matrix and multi-index gamma matrices represent the antisymmetrized product of gamma matrices, *i.e.*, $\gamma_{\mu_1\cdots\mu_n} = \gamma_{[\mu_1}\cdots\gamma_{\mu_n]}$.

As a starting point, we consider the D-dimensional action which includes ECT of gravity coupled minimally with Dirac fields[5]

$$\begin{aligned} S &= \frac{1}{2\kappa^2}\int \frac{\epsilon_{\hat{a}_1\ldots\hat{a}_D}}{(D-2)!}\hat{\boldsymbol{\mathcal{R}}}^{\hat{a}_1\hat{a}_2}\wedge\hat{\boldsymbol{e}}^{a_3}\wedge\ldots\wedge\hat{\boldsymbol{e}}^{a_D} \\ &\quad -\frac{1}{2}\sum_f\int\left(\bar{\Psi}_f\hat{\boldsymbol{\gamma}}\wedge\star\hat{\boldsymbol{\mathcal{D}}}\Psi_f - \hat{\boldsymbol{\mathcal{D}}}\bar{\Psi}_f\wedge\star\hat{\boldsymbol{\gamma}}\Psi_f\right), \end{aligned} \tag{5}$$

where $\hat{\boldsymbol{\mathcal{D}}}$ is the spinorial covariant derivative in a curved spacetime, defined by[6]

$$\begin{aligned} \hat{\boldsymbol{\mathcal{D}}}\Psi &= \mathbf{d}\Psi + \tfrac{1}{4}\hat{\boldsymbol{\omega}}^{\hat{a}\hat{b}}\gamma_{\hat{a}\hat{b}}\Psi, \\ \hat{\boldsymbol{\mathcal{D}}}\bar{\Psi} &= \mathbf{d}\bar{\Psi} - \tfrac{1}{4}\bar{\Psi}\hat{\boldsymbol{\omega}}^{\hat{a}\hat{b}}\gamma_{\hat{a}\hat{b}}, \end{aligned} \tag{6}$$

where the symbol $\hat{\boldsymbol{\gamma}}$ denotes the contraction $\gamma_{\hat{a}}\hat{\boldsymbol{e}}^{\hat{a}}$, the $\star$ stands for the Hodge *star* map, and the subscript f stands for the fermion's flavour.

The field equation for the spin connection in Eq. (5) yields an algebraic equation for the components of the torsion

$$\frac{1}{2}\left[\hat{\mathcal{T}}_{\hat{b}\hat{c}\hat{a}} + \hat{\mathcal{T}}_{\hat{b}\hat{a}\hat{c}} + \hat{\mathcal{T}}_{\hat{a}\hat{b}\hat{c}}\right] \equiv \hat{\mathcal{K}}_{\hat{a}\hat{b}\hat{c}} = -\frac{\kappa^2}{4}\sum_f\bar{\Psi}_f\gamma_{\hat{a}\hat{b}\hat{c}}\Psi_f\,. \tag{7}$$

Notice that the expression in the LHS is the contorsion, whose only nontrivial component, from Eq. (7), is its totally antisymmetric part.

The contorsion is the tensor which relates the *affine* spin connection with the torsion-less spin connection, $\hat{\bar{\boldsymbol{\omega}}}^{\hat{a}}{}_{\hat{b}}$, through the equation

$$\hat{\boldsymbol{\omega}}^{\hat{a}}{}_{\hat{b}} = \hat{\bar{\boldsymbol{\omega}}}^{\hat{a}}{}_{\hat{b}} + \hat{\boldsymbol{\mathcal{K}}}^{\hat{a}}{}_{\hat{b}}, \tag{8}$$

where the contorsion one-form is defined by $\hat{\boldsymbol{\mathcal{K}}}^{\hat{a}}{}_{\hat{b}} = \hat{\mathcal{K}}_{\hat{m}}{}^{\hat{a}}{}_{\hat{b}}\,\hat{\boldsymbol{e}}^{\hat{m}}$.

5 We assume that fermion masses are developed through the Higgs mechanism, so there is no need for considering nontrivial fundamental mass terms.

6 Hereon, multi-index gamma matrices represent the totally antisymmetric product of gammas.

The advantage of Eq. (7) is algebraic, which means it can be substituted back into the action, allowing us to obtain an effective, torsion-free action. The effective action includes GR coupled minimally with the Dirac fields, plus an induced four-fermion contact interaction, namely

$$\mathscr{L}_{4\mathrm{FI}} = \frac{\kappa^2}{32} \sum_{f_1,f_2} \left(\bar{\Psi}_{f_1} \gamma_{\hat{a}\hat{b}\hat{c}} \Psi_{f_1}\right) \left(\bar{\Psi}_{f_2} \gamma^{\hat{a}\hat{b}\hat{c}} \Psi_{f_2}\right) . \tag{9}$$

In four dimensions — where $\kappa^2 = \frac{1}{M_{\mathrm{Pl}}^2}$ — the extra contact interaction is strongly suppressed by the Planck mass, as anticipated. Therefore, this effective interaction is negligible for any phenomenological effect.

Lately, the phenomenological insight of scenarios with extra dimensions has increased, boosted by works which solve the *hierarchy problem*[7] *i.e.* the huge difference between the electroweak and gravitational scales, through the introduction of a fundamental scale of gravity, $\kappa_*^{-1} = M_* \sim$ TeV, which gets enhanced in the four-dimensional effective theory, up to the Planck scale [66–70].

Within the framework of model with extra dimensions, the coupling accompanying the effective four-fermion interaction in Eq. (9), should be replaced from κ to κ_*, which permits — in principle — to obtain some particle physics phenomenology from the gravitational induced term.

In the rest of the paper, we restrict ourselves to considering a single extra dimension. As a first step, we decompose the induced four-fermion interaction in terms of four-dimensional quantities, assuming the five-dimensional Clifford algebra admits the same representation as the one in four dimensions. Therefore,

$$\left(\gamma_{\hat{a}\hat{b}\hat{c}}\right) \left(\gamma^{\hat{a}\hat{b}\hat{c}}\right) = \left(\gamma_{abc}\right) \left(\gamma^{abc}\right) + 3 \left(\gamma_{ab*}\right) \left(\gamma^{ab*}\right) . \tag{10}$$

Hence, the interaction in Eq. (9) rises an axial–axial and a tensor-axial–tensor-axial interactions [77]

$$\begin{aligned} \mathscr{L}_{\mathrm{eff}} &= \frac{3\kappa_*^2}{16} \sum_{f_1,f_2} \left(\bar{\Psi}_{f_1} \gamma_a \gamma^* \Psi_{f_1}\right) \left(\bar{\Psi}_{f_2} \gamma^a \gamma^* \Psi_{f_2}\right) \\ &\quad + \frac{3\kappa_*^2}{32} \sum_{f_1,f_2} \left(\bar{\Psi}_{f_1} \gamma_{ab} \gamma^* \Psi_{f_1}\right) \left(\bar{\Psi}_{f_2} \gamma^{ab} \gamma^* \Psi_{f_2}\right) , \end{aligned} \tag{11}$$

where γ^* is the chiral matrix in four dimensions.

[7] One of the most outstanding proposals in the context of extra dimensions is the AdS/CFT correspondence (see, for example, Refs. [91–93]), which related different physical theories living in different dimensions, the reason why it is sometimes called *holographic theory*. Nevertheless, we do not use the correspondence in this work.

3. One-loop correction of decay width for a (pseudo)scalar into a pair of fermions

The splitting of the effective interaction, Eq. (11), can be written in terms of current–current interactions, as shown in Ref. [94]

$$\mathscr{L}_{\mathrm{eff}} = \frac{3\kappa_*^2}{16} \sum_{f_1,f_2} \left(J^*_{a\,f_1}\right)\left(J^{a*}_{f_2}\right) + \frac{3\kappa_*^2}{32} \sum_{f_1,f_2} \left(J^*_{ab\,f_1}\right)\left(J^{ab*}_{f_2}\right) . \tag{12}$$

There are two different contributions to the $\varphi \to f\bar{f}$ process which will be called *s-channel* (see Fig. 1 (a)) and *t-channel* (see Fig. 1 (b)), respectively. It is worth to mention that in order to obtain chiral fermions in the effective four-dimensional theory, an orbifold condition must be imposed in the extra dimension [84], and such a condition avoid the presence of tensor-axial–tensor-axial currents in Eq. (12). Therefore, in the below analysis, only the induced axial–axial currents will be considered.

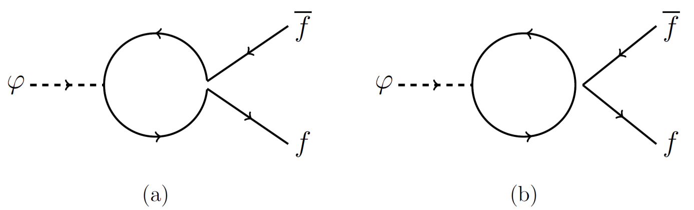


Fig. 1. Scalar to fermion pair through the four-fermion interaction in s-channel (a) and t-channel (b).

We assume that the (pseudo)scalar fields couple to fermions through generic Yukawa interactions, whose couplings are not necessarily proportional to the final-state fermion mass. Further, we assume that the scalar field φ_{s} is CP-even, and the pseudo-scalar field φ_{p} is CP-odd. Then, our Lagrangian contains the terms

$$\mathscr{L} = \sum_f y^f_{\mathrm{s}} \varphi_{\mathrm{s}} \bar{\psi}_f \psi_f + \imath \sum_f y^f_{\mathrm{p}} \varphi_{\mathrm{p}} \left(\bar{\psi}_f \gamma^* \psi_f\right) , \tag{13}$$

where $y^f_{\mathrm{s,p}}$ are real and arbitrary constants, and the f index runs for each SM fermion, without considering neutrinos. On the other hand, a (pseudo)scalar field decays into a fermion pair through a current of the form of

$$J = \bar{u}_f(\vec{p})\left(S + \imath P \gamma^*\right) v_f\left(\vec{p}\,'\right) , \tag{14}$$

where S and P are the scalar and pseudo-scalar form factors. According to the current in Eq. (14), the decay width of a (pseudo)scalar particle into a fermion pair at tree level is given by

$$\Gamma\left(\varphi \to f\bar{f}\,\right) = N_{\rm c}\frac{M_\varphi}{8\pi}\sqrt{1-\frac{4m_f^2}{M_\varphi^2}}\left(\left(y_{\rm s}^f\right)^2 S^2\left(1-\frac{4m_f^2}{M_\varphi^2}\right)+\left(y_{\rm p}^f\right)^2 P^2\right)\,, \tag{15}$$

where M_φ is the mass of the (pseudo)scalar, m_f is the fermion mass in the final state of the process and $N_{\rm c}$ is the colour factor, which in the case of decay into quarks, will take the value of $N_{\rm c}=3$ [8].

It is worth noticing that the structure of the induced four-fermion interaction in Eq. (12), the t-channel Feynman diagram — see Fig. 1 (b) — does not contribute to the decay width of the scalar nor pseudoscalar field. In the former, the trace of the product of Dirac matrices

$$\mathrm{Tr}\left[\left(\not{p}_1+\not{p}_2-\not{q}-m_i\right)\left(\not{q}-m_i\right)\gamma^\mu\gamma^*\right]$$

vanishes identically due to the presence of the chiral matrix (γ^*), while in the latter, although the trace of the amplitude is nontrivial, the Feynman integral results to be zero.

Next, we want to estimate the order of the correction to the decay width induced by the four-fermion interaction described above. For that end, we assume that the fundamental scale of gravity M_* is of the order of the *new physics* scale Λ $(M_*\sim\Lambda)$. Therefore, although our result comes from generic models with an extra dimension, we hide the details of the model, such as the size of the extra dimension and the embedding of the four-dimensional spinors into the five-dimensional ones, within this new scale of physics.

The one-loop corrections to the current in Eq. (14), δJ, through the scalar field decay into two fermions, considering the effective four-fermion interaction is

$$\delta J = \bar{u}_f(\vec{p})\,(\delta S)\,v_f\left(\vec{p}\,'\right) \quad \text{with} \quad \delta S = -\frac{3}{32}\frac{1}{\Lambda^2}\left(M_\varphi^2-2m_f^2\right)\log\left(\frac{\Lambda^2}{M_\varphi^2}\right)\,, \tag{16}$$

while for the pseudoscalar it is

$$\delta J = \bar{u}_f(\vec{p})\,(\imath P\gamma^*)\,v_f(\vec{p}\,') \quad \text{with} \quad \delta P = -\frac{3}{32}\frac{1}{\Lambda^2}\left(M_\varphi^2-6m_f^2\right)\log\left(\frac{\Lambda^2}{M_\varphi^2}\right)\,. \tag{17}$$

Keeping the original coupling (tree level) and accounting for CP invariance, these results generate corrections to the variation of the decay width

[8] We have cross-checked our calculations using the Mathematica package FeynCalc [95].

of the form of

$$\delta\Gamma^{\mathrm{S}}_{4\mathrm{FI}} = -\frac{3}{128}\frac{N_{\mathrm{c}}\left(y^f_{\mathrm{s}}\right)^2 M_\varphi}{\pi\Lambda^2}\left(M^2_\varphi - 2m^2_f\right)\left(1-\frac{4m^2_f}{M^2_\varphi}\right)^{3/2}\log\left(\frac{\Lambda^2}{M^2_\varphi}\right), \tag{18}$$

and

$$\delta\Gamma^{\mathrm{P}}_{4\mathrm{FI}} = -\frac{3}{128}\frac{N_{\mathrm{c}}\left(y^f_{\mathrm{p}}\right)^2 M_\varphi}{\pi\Lambda^2}\left(M^2_\varphi - 6m^2_f\right)\left(1-\frac{4m^2_f}{M^2_\varphi}\right)^{1/2}\log\left(\frac{\Lambda^2}{M^2_\varphi}\right). \tag{19}$$

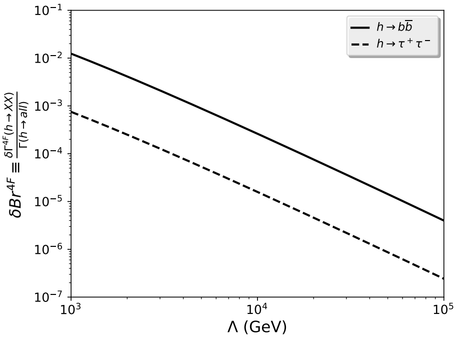
In these two cases, the original result is a function of the Passarino–Veltman integrals, however, we have written the expressions with the explicit logarithmic dependence on the scale Λ.

4. Standard Model example: correction to Higgs decay into a pair of fermions

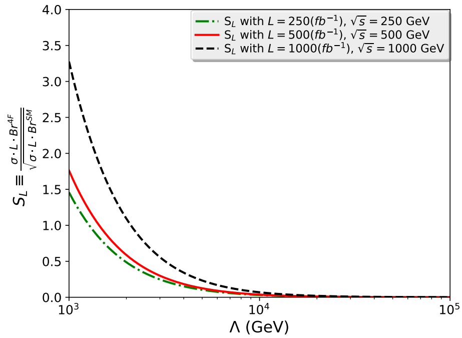
Now, we focus on a special case of the Higgs boson decay. As mentioned above, only the s-channel diagrams contribute to the variation of the Higgs decay width. Furthermore, due to the fact that SM Higgs is a scalar particle, the quantities S and P in Eq. (15) are one and zero, respectively. Since the torsion-induced four-fermion interaction comes from the kinetic term, although the dimensional reduction induces a Kaluza–Klein tower in the effective particle spectrum, indisputably the fermion around the loop has the same flavour as the outgoing particles. Therefore, none of the particles in the Kaluza–Klein tower enter in the analysis. Then, the correction to the variation of the Higgs decay width is

$$\delta\Gamma_{4\mathrm{FI}}\left(h\to f\bar{f}\right) = -\frac{3}{512}\frac{g^2 m^2_f M_h}{\pi M^2_W\Lambda^2}\left(M^2_h - 2m^2_f\right)\left(1-\frac{4m^2_f}{M^2_h}\right)^{3/2}\log\left(\frac{\Lambda^2}{M^2_h}\right). \tag{20}$$

We will focus on Higgs decays into both $\tau^+\tau^-$ and $b\bar{b}$, which are the main fermionic decay modes, in order to estimate the size of the effects and compare these corrections with the total decay width predicted by the SM. In Fig. 2 (a), we show the correction on the Higgs branching ratio into fermion pairs as a function of the gravitational scale. For fundamental gravitational scales as low as 1 TeV, the correction induced by the torsion interaction is about 1.024% for the decay channel $h\to b\bar{b}$, while for the process $h\to\tau^+\tau^-$, it decreases to 0.075%. As it is expected, for higher gravitational scales, the correction decreases due to the quadratic suppression (Λ^{-2}) in Eq. (20).



(a)



(b)

Fig. 2. (a) Variation of the Higgs boson branching ratio $\delta\mathrm{Br}^{4\mathrm{F}}$ due to the 4-fermion interaction as a function of the new physics scale Λ. The dashed line denotes $h \to \tau^+\tau^-$ decay and the solid one $h \to b\bar{b}$ decay channel. (b) Expected significance level (S_L) at ILC.

Although the dominant Higgs branching fractions come from $h \to b\bar{b}$ and $h \to \tau^+\tau^-$, the LHC coupling precision capabilities are not good enough to resolve what we are interested in, due to the presence of QCD backgrounds [96, 97]. However, the $b\bar{b}$ signal channels may be more visible at future Higgs factories, such as the ILC [97–99] or CLIC [100–103], where the QCD background is reduced. Then it is expected higher precision measurements in the Higgs sector, allowing to explore deeply the couplings and decay width, and therefore, being able to measure deviations in the Higgs decay width, eventually as low as our results.

Keeping in mind the aforementioned conditions, we estimate the expected significance level (S_L) at the ILC, coming from the four-fermion interaction. Such estimation reads

$$S_\mathrm{L} = \frac{\sigma \times L \times \mathrm{Br}^{4\mathrm{F}}(\Lambda)}{\sqrt{\sigma \times L \times \mathrm{Br}^{\mathrm{SM}}}}\epsilon_f\,, \tag{21}$$

where σ is the production cross section of the Higgs boson via Higgsstrahlung $\sigma = \sigma(e^+e^- \to hZ)$ and vector boson fusion $\sigma(e^+e^- \to \nu\bar{\nu}h)$, L is the expected luminosity for each run, and ϵ_f is the signal selection efficiency for the f-channel, which is approximatelly $\epsilon_f \simeq 0.3$ in both of the considered channels [104, 105]. As shown in Fig. 2 (b), as the gravitational scale increases, the expected significance in the number of events — due to the torsion — decreases. It tells us that the effect is observable at the ILC only if $\Lambda \sim 1$ TeV. Therefore, if ILC does not see a significant excess of events in both $b\bar{b}$ and $\tau^+\tau^-$ channels at these energies scales, either the scale of gravity is much bigger than these energy scales, or ETC gravity is not coupled minimally to fermions.

However, recent analysis on the constraints imposed by the torsion induced four-fermion interaction on the Z boson decay (see Refs. [84, 85]), the strongest limit is $\Lambda \simeq 30$ TeV. Given this stringent limit, the correction to the decay width of the Higgs drops to approximately $3.3 \times 10^{-3}\%$ and $2.2 \times 10^{-4}\%$ for bottom and tau pairs respectively. Such limits are unlike to be measured in current experiments, but could be reached at future Higgs factories.

5. Beyond Standard Model example: 2HDM

The 2HDM has in its physical spectrum two neutral scalars (h^0, H^0), one pseudo-scalar (A^0), and two charged bosons $(H^{\pm})$, see, for example, Ref. [106]. We focus on the coupling between neutral bosons and SM fermions. The parametrization of the Yukawa interactions in this context is

$$\mathcal{L}_{\text{Yuk}} = -\sum_f \frac{m_f}{v} \left(\hat{y}_f^h \bar{f} f h^0 + \hat{y}_f^H \bar{f} f H^0 - \imath \hat{y}_f^A \bar{f} \gamma^5 f A^0 \right) , \tag{22}$$

where the constants $\hat{y}_f^{h,H,A}$ are real numbers which depend on the specific model, and v is the vacuum expectation value of the Higgs field. There exist a diversity of forms of the 2HDM (Type I, II, X and Y), but we shall consider the type II in its first scenario, called Type II-1, which has the best fits to the observed data. In this scenario, the h^0 state matches with the 126 GeV resonance observed h at the LHC, then $h^0 = h$, and the $\hat{y}_f^h$ measures the deviation at tree level in the coupling between the Higgs and the SM fermions. The other neutral scalars are heavier than the corresponding Higgs boson and the coupling constants $\hat{y}_f^{H,A}$ are determined by the Type II-1 model [106].

5.1. Corrections to the Higgs decay width in Type II-1 2HDM

We compare the corrections to the Higgs decay width, induced by the torsion-descendent four-fermion interaction, in two possibles submodels: the constrained by flavour-physics and the unconstrained [106]. We summarise the values of the Yukawa couplings in both submodels in Table I.

Considering the matching between the coupling constants in Eqs. (13) and (22), *i.e.* $y_{\text{s,p}}^f = m_f \hat{y}_f^h / v$, we put these values in our master formula for the scalar decays into both $b\bar{b}$ and $\tau\bar{\tau}$. The variation of the Higgs partial width decay due to the four-fermion effective interaction in the context of the 2HDM are shown in Fig. 3. The differences in the Higgs decay width variation between the SM and the 2HDM frameworks are small for all Λ (less than 1%), because the deviations in the Yukawa coupling between the two cases are negligible.

TABLE I

Yukawa couplings to up-type quark, to down-type quark, and to charged leptons for both submodels.

Yukawa coupling	Constrained	Unconstrained
$\hat{y}_u^h$	1.28	1.05
$\hat{y}_d^h$	-0.91	-0.99
$\hat{y}_l^h$	-0.91	-0.99

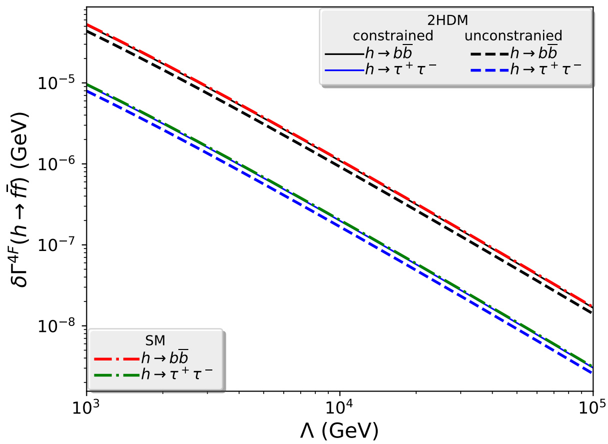


Fig. 3. Variation of the Higgs decay witdh into $b\bar{b}$ and $\tau\bar{\tau}$ at one-loop due to the 4-fermion interaction as a function of the new physics scale Λ.

5.2. Decay width corrections to the heavy neutral (pseudo)scalars in the 2HDM

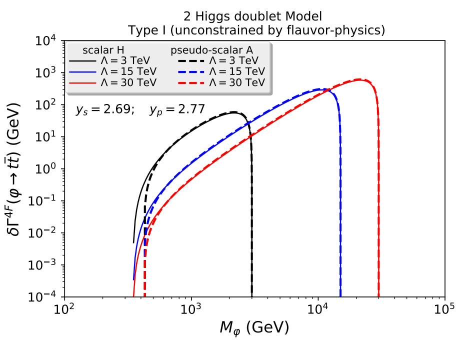
Next, we estimate the corrections to the decay width of the heavy neutral scalars (H, A). We exemplify in the unconstrained Type II-1 model, whose Yukawa couplings are summarised in Table II.

TABLE II

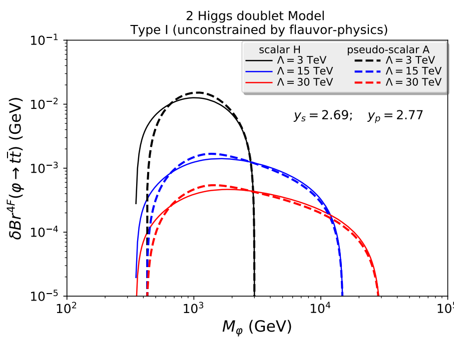
Effective Yukawa couplings for Type II-1 unconstrained model for the massive scalar H^0 and pseudoscalar A^0 to fermions: up-type quarks, down-type quarks and charged leptons.

Yukawa coupling	Scalar (H^0)	Pseudoscalar (A^0)
y_u	2.69	2.77
y_d	0.37	0.36
y_l	0.37	0.36

The first important consequence to mention is that there is no important difference in $\delta\Gamma^{4\mathrm{F}}$ between the scalar and pseudo-scalar case at any value of Λ, except near the threshold. This is because at lower scalar masses there is a bigger suppresion in the pseudoscalar variation (see (18) and (19)), making thus the scalar variation visible at lower scalar masses. Note that always the corrections are valid below Λ, which is the effective cut-off theory, and at $M_\varphi = \Lambda$ the curves in the plot fall steeply due to the logarithm behaviour in the correction. Complementary, Fig. 4 (b) shows the change in the branching fraction of the (pseudo)scalar to $t\bar{t}$ as a function of its mass, for the same cut-off values as before. As it is expected, the Λ cuadratic suppresion in the (pseudo)scalar variation makes the corrections bigger for low gravitational scales ($\sim$ 0.1 to 1%), and suppressed for $\Lambda \gtrsim$ 15 TeV, making a correction less than 0.1%. Note that the only way to distinguish the branching fraction corrections between scalar and pseudoscalar, for any Λ, is near the low-masses threshold.



(a)



(b)

Fig. 4. Variation of the decay widths (a) and branching ratios (b) into $t\bar{t}$ at one-loop for the heavy scalar H and the pseudoscalar A as a function of their mass.

It is worth pointing out that although our results for SM Higgs boson are less sensitive than the ones presented in Ref. [85] for Z boson decay, we cannot assure the same for the 2HDM model.

6. Discussion and conclusions

We have reviewed how gravitational torsion induces an effective interaction between SM fermions. This new interactions affect directly particle observables, such as their decay width. We analysed the variation induced by the torsion-descendent four-fermion interaction, in scalar and pseudoscalar particles in the SM and the Type II-1 2HDM.

Concerning SM Higgs decays, we have focused on $h \to b\bar{b}$ and $h \to \tau^+\tau^-$ decays, which are the dominant decay modes having branching ratios of $\approx 57\%$ and $\approx 6\%$, respectively. We have considered the correction to the branching ratio for these processes mediated by the effective four-fermion interaction at one-loop level. It can be seen in Fig. 2 that the contribution to both fermionic channels become smaller as the gravitational scale grows up. On the other hand, $\delta \mathrm{Br}^{4\mathrm{F}}(h \to b\bar{b})$ is roughly speaking an order of magnitude bigger than $\delta \mathrm{Br}^{4\mathrm{F}}(h \to \tau^+\tau^-)$ for any scale energy Λ, doing this channel more relevant from a phenomenological viewpoint. For gravitational scales as low as $\Lambda = 1$ TeV, the corrections to the branching ratio for $h \to b\bar{b}$ is $\sim 1\%$, meanwhile $h \to \tau^+\tau^-$ is $\sim 0.1\%$. Moreover, from Fig. 2 (a), one can see that when $\Lambda = 30$ TeV, the corrections are $3.3 \times 10^{-3}\%$ for $h \to b\bar{b}$ and $2.2 \times 10^{-4}\%$ for $h \to \tau^+\tau^-$.

However, the $b\bar{b}$ signal channel may be more visible at future Higgs factories, such as the International Linear Collider (ILC) or the Compact Linear Collider (CLIC), where the QCD background is reduced having therefore more precision in some observables. Additionally, at both ILC and CLIC, there are expected higher precision measurements in the Higgs sector than at the LHC, allowing to explore more deeply into the quantitative information of the couplings and Higgs decay width, being therefore able to measure deviations in the Higgs decay width, eventually as low as our results.

At this point, we want to remark one more time that our results have shown that the Higgs decay width is less sensitive than, for instance, the Z boson decay width [85] to the kind of corrections we are studying. This is mainly due to the higher number of degrees of freedom present in the vector case and to the fact that the properties of the Z boson have been measured with a high accuracy.

On the other hand, our results turn out to be more auspicious in the case of the 2HDM, particularly if the non-standard scalars are heavy, as shown in Fig. 4. It is even possible to distinguish between scalars and pseudo-scalars near the threshold of the decay channel if there are additional heavy fermions. The corrections $\delta\Gamma^{\mathrm{P}}_{4\mathrm{FI}}$ and $\delta\Gamma^{\mathrm{S}}_{4\mathrm{FI}}$ can be distinguished in the lower mass threshold, when we have provided $y_{\mathrm{s}} = 2.69$ and $y_{\mathrm{p}} = 2.77$. However, it is important to note that in general (arbitrary values of y_{s} and y_{p}) the condition of distinguishability is

$$y_p^2 \neq y_s^2 \, \frac{M_\phi^2 - 2m_f^2}{M_\phi^2 - 6m_f^2} \left(1 - \frac{4m_f^2}{M_\phi^2}\right) . \tag{23}$$

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