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 pseudo scalar 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)\textsuperscript{[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\textsuperscript{[4–16]}. Extended scalar sector are also predicted by models like supersymmetry\textsuperscript{[17–19]}, some versions of strong electroweak symmetry breaking models\textsuperscript{[20]} and non-minimal composite Higgs models\textsuperscript{[21–23]}. Although no deviation from the SM has yet been observed, the LHCb collaboration has reported 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\textsuperscript{[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\textsuperscript{[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 view 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\textsuperscript{[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\textsuperscript{[41–44]}, models which keep the precepts of GR but in higher dimensions\textsuperscript{[45,46]}, metric-affine theories\textsuperscript{[47]}, affine theories\textsuperscript{[48,54]}, models with higher order in curvature and/or torsion\textsuperscript{[55,61]}, et cetera.

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\textsuperscript{1} There are several attempts of quantize the gravitational interactions, see for example Refs.\textsuperscript{34–39}. For a historical review, see Ref.\textsuperscript{40}.
In this letter, we shall only consider the 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 induced effective interaction is strongly suppressed because it is diminished by the of 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, $M_{pl}$, 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 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).

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 the 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 ($e_a^\mu$) and the spin connection ($\omega_{\mu}^{ab}$) 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 neighbouhood points.

The equivalence between the metric and the vielbein is given by

$$g_{\mu\nu} = \eta_{ab}e^a_\mu e^b_\nu. \quad (1)$$

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, i.e. the Cartan structure equations,

$$de^a + \omega_{b\mu}^a e^b = \mathcal{T}^a \quad \text{and} \quad d\omega^{ab} + \omega^a_{c\mu}\omega_{b\mu} = \mathcal{R}^{ab}.$$

The vielbein and spin connection one-forms are defined as

$$e^a = e^a_\mu dx^\mu \quad \text{and} \quad \omega^{ab} = \omega^{ab}_\mu dx^\mu, \quad (3)$$

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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 [87,89].
while the two-forms are written explicitly in components as

\[
\mathcal{T}^a = \frac{1}{2} T_{\mu \nu}^a dx^\mu \wedge dx^\nu = \frac{1}{2} T_{m n}^a e^m \wedge e^n
\]

and

\[
\mathcal{R}^a_{\ b} = \frac{1}{2} R_{\mu \nu}^a b dx^\mu \wedge dx^\nu = \frac{1}{2} R_{mn}^a b e^m \wedge e^n.
\]

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 objects (or indices) lying in the latter. Worth to mention, we denote by \( \gamma^\star \) 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 starting point, we consider the \( D \)-dimensional action which includes ECT of gravity coupled minimally with Dirac fields,

\[
S = \frac{1}{2\kappa^2} \int \frac{\epsilon_{\hat{a}_1 \cdots \hat{a}_D}}{(D-2)!} \mathcal{R}^{\hat{a}_1 \hat{a}_2} e^{\hat{a}_3} \wedge \cdots \wedge e^{\hat{a}_D} - \frac{1}{2} \sum_f \int \left( \bar{\Psi}_f \gamma^\star \hat{D} \Psi_f - \hat{D} \bar{\Psi}_f \gamma^\star \Psi_f \right),
\]

where \( \hat{D} \) is the spinorial covariant derivative in a curved spacetime, defined by

\[
\hat{D} \Psi = d\Psi + \frac{1}{4} \hat{\omega}^{\hat{a}\hat{b}} \gamma_{\hat{a}\hat{b}} \Psi,
\]

\[
\hat{D} \bar{\Psi} = d\bar{\Psi} - \frac{1}{4} \bar{\Psi} \hat{\omega}^{\hat{a}\hat{b}} \gamma_{\hat{a}\hat{b}},
\]

the symbol \( \gamma^\star \) denotes the contraction \( \gamma_{\hat{a}} \hat{e}^{\hat{\alpha}} \), 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,

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

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, \( \hat{\omega}_{\hat{a} \hat{b}} \), through the equation

\[
\hat{\omega}_{\hat{a} \hat{b}} = \hat{\omega}_{\hat{a} \hat{b}} + \hat{K}_{\hat{a} \hat{b} \hat{c}}.
\]

where the contorsion one-form is defined by \( \hat{K}_{\hat{a} \hat{b} \hat{c}} = \hat{K}_{\hat{m} \hat{n} \hat{a}} \hat{e}^{\hat{m} \hat{n}} \).

The advantage of Eq. (7) been algebraic, is that 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

\[
\mathcal{L}_{4FI} = \frac{\kappa^2}{32} \sum_{f_1, f_2} (\bar{\Psi}_{f_1} \gamma_{\hat{a}\hat{b}\hat{c}} \Psi_{f_1}) (\bar{\Psi}_{f_2} \gamma^{\hat{a}\hat{b}\hat{c}} \Psi_{f_2}).
\]

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

6 Hereon, multi-index gamma matrices represent the totally anti-symmetric product of gammas.
In four dimensions—where \( \kappa^2 = \frac{1}{M_{\text{pl}}^2} \)—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, \( \kappa_\ast^{-1} = M_\ast \sim \text{TeV} \), which gets enhanced in the four-dimensional effective theory, up to the Planck scale \( \text{Pl} \)\(^6\)–\(^7\).

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

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

\[
(\gamma^a)(\gamma^b) = (\gamma^{abc})(\gamma^{abc}) + 3(\gamma^{ab})(\gamma^{ab}).
\]

Hence, the interaction in Eq. (9) rises an axial–axial and a tensor-axial–tensor-axial interactions\(^7\).

\[
\mathcal{L}_{\text{eff}} = \frac{3\kappa^2}{16} \sum_{f_1,f_2} (\bar{\Psi}_{f_1} \gamma_a \gamma^* \Psi_{f_1})(\bar{\Psi}_{f_2} \gamma^a \gamma^* \Psi_{f_2})
+ \frac{3\kappa^2}{32} \sum_{f_1,f_2} (\bar{\Psi}_{f_1} \gamma_{ab} \gamma^* \Psi_{f_1})(\bar{\Psi}_{f_2} \gamma^{ab} \gamma^* \Psi_{f_2})
\]

where \( \gamma^* \) is the chiral matrix in four dimensions.

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. \(^9\),

\[
\mathcal{L}_{\text{eff}} = \frac{3\kappa^2}{16} \sum_{f_1,f_2} (J^*_{a,f_1})(J^a_{f_2})
+ \frac{3\kappa^2}{32} \sum_{f_1,f_2} (J^*_{ab,f_1})(J^{ab}_{f_2}).
\]

There are two different contributions to the \( \varphi \rightarrow f \bar{f} \) 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\(^8\), and such 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.

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 \( \phi_s \) is CP-even, and the pseudo-scalar field \( \phi_p \) is CP-odd. Then, our Lagrangian contains the terms

\[
\mathcal{L} = \sum_f y^f_s \phi_s \bar{\psi}_f \gamma \psi_f + i \sum_f y^f_p \phi_p (\bar{\psi}_f \gamma^* \psi_f)
\]

where \( y^f_{s,p} \) 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

\(^7\) One of the most outstanding proposals in the context of extra dimensions is the AdS/CFT correspondence (see for example Ref. \(^9\))\(^1\)\(^3\)), which related different physical theories which live in different dimensions, reason why it is sometimes called holographic theory. Nevertheless, we do not use the correspondence in this work.
Fig. 1. Scalar to fermion pair through the four-fermion interaction in s-channel (a) and t-channel (b).

current of the form

\[ J = \bar{u}_f(p)(S + iP\gamma^*)v_f(p') \]

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

\[ \Gamma(\varphi \to f\bar{f}) = N_c \frac{M_\varphi}{8\pi} \sqrt{1 - \frac{4m_f^2}{M_\varphi^2}} \left( (y_f^s)^2 S^2 \left( 1 - \frac{4m_f^2}{M_\varphi^2} \right) + (y_f^p)^2 P^2 \right), \]

where \( M_\varphi \) is the mass of the (pseudo)scalar, \( m_f \) is the fermion mass in the final state of the process and \( N_c \) is the colour factor, which in the case of decay into quarks it will take the value \( N_c = 3 \).

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

\[ \text{Tr} \left[ (\gamma_1 + \gamma_2 - \gamma - m_\text{i})(\gamma^* - m_\text{i}) \gamma^* \right] \]

vanishes identically due to the presence of the chiral matrix \((\gamma^*)\), 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 \( \Lambda \) (\( M_* \sim \Lambda \)). 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), \( \delta J \), through the scalar field decay into two fermions, considering the effective four-fermion interaction is

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

while for the pseudoscalar is

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

Keeping the original coupling (tree level) and accounting for CP-invariance. These results generate corrections to the variation of the decay width of the form

\[ \delta \Gamma_{4FI}^S = -\frac{3}{128} \frac{N_c(y_f^s)^2 M_\varphi}{\pi \Lambda^2} \left( M_\varphi^2 - 2m_f^2 \right) \left( 1 - \frac{4m_f^2}{M_\varphi^2} \right)^{3/2} \log \left( \frac{\Lambda^2}{M_\varphi^2} \right), \]

\[ \delta \Gamma_{4FI}^P = -\frac{3}{128} \frac{N_c(y_f^p)^2 M_\varphi}{\pi \Lambda^2} \left( M_\varphi^2 - 6m_f^2 \right) \left( 1 - \frac{4m_f^2}{M_\varphi^2} \right)^{3/2} \log \left( \frac{\Lambda^2}{M_\varphi^2} \right). \]

\[ ^8 \text{We have cross-checked our calculations using the Mathematica package "FeynCalc" [95].} \]
and
\[ \delta \Gamma_{4\text{FI}} = -\frac{3}{128} \frac{N_c(y_f^2 M_\phi^2)}{\pi \Lambda^2} \left( M_\phi^2 - 6m_f^2 \right) \left( 1 - \frac{4m_f^2}{M_\phi^2} \right)^{1/2} \log \left( \frac{\Lambda^2}{M_\phi^2} \right). \] (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 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 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
\[ \delta \Gamma_{4\text{FI}}(h \to f\bar{f}) = -\frac{3}{512} \frac{g^2 m_f^2 M_h}{\pi M_W^2 \Lambda^2} \left( M_h^2 - 2m_f^2 \right) \left( 1 - \frac{4m_f^2}{M_h^2} \right)^{3/2} \log \left( \frac{\Lambda^2}{M_h^2} \right). \] (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 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 (\( \Lambda^{-2} \)) in Eq. [20].

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
\[ S_L = \frac{\sigma \times L \times Br^{4\text{FI}}(\Lambda)}{\sqrt{\sigma \times L \times Br^{SM}}} \epsilon_f, \] (21)

where \( \sigma \) 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 \( \epsilon_f \) is the signal selection efficiency for the \( f \) channel, which is approximately \( \epsilon_f \simeq 0.3 \) 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 tell us that the effect is observable at ILC only if \( \Lambda \sim 1 \text{ TeV} \). Therefore, if ILC does not see a significant excess of events in both \( b\bar{b} \) and \( \tau^+\tau^- \) channels this 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. \cite{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 unlikely to be measure 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 scalar ($h^0, H^0$), one pseudo-scalar ($A^0$), and two charged bosons ($H^{\pm}$), see for example Ref. \cite{106}. We focus on the coupling between neutral bosons and SM fermions. The parametrization of the Yukawa interactions in this context is

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

(22)

where the constants $\tilde{y}_f^{h,H,A}$ 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 match with the observed $126$ GeV resonance observed $h$ at LHC, then $h^0 = h$, and the $\tilde{y}_f^{h}$ measure the deviation at tree level in the coupling between the Higgs and the SM fermions. The other neutral scalars are heavy than the corresponding Higgs boson and the coupling constants $\tilde{y}_f^{H,A}$ are determined by the Type II-1 model \cite{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 flauvor-physics and the unconstrained \cite{106}. We summarise the values of the Yukawa couplings in both submodels in the Table 1.

Considering the matching between the coupling constants in Eqs. (13) and (22), i.e. $y_{f,p}^{h} = m_f \tilde{y}_f^{h}/v$, 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

Fig. 2. (a) Variation of the Higgs boson branching ratio $\delta B_{h^{4F}}$ due to the 4-fermion interaction as a function of the new physics scale $\Lambda$. The dashed line denotes $h \rightarrow \tau^+\tau^-$ decay and the solid one $h \rightarrow b\bar{b}$ decay channel. (b) Expected Significance level ($S_{ILC}$) at ILC.
Table 1. Yukawa couplings to up-type quark, to down-type quark, and to charged leptons for both submodels.

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

are shown in Fig. 3. The differences in the Higgs decay width variation between the SM and the 2HDM frameworks are small for all $\Lambda$ (less than 1%), because the deviations in the yukawa coupling between the two cases are negligible.

![Graph showing variation of Higgs decay width into $b\bar{b}$ and $\tau\bar{\tau}$ at one-loop due to the 4-fermion interaction](image)

Fig. 3. Variation of the Higgs decay width 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 $\Lambda$.

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

Next, we estimate the corrections to the decay width to the heavy neutral scalars $(H,A)$. We exemplify in the unconstrained Type II-1 model, whose Yukawa couplings are summarised in Table 2.

Table 2. 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 $\delta\Gamma^{4F}$ between the scalar and pseudo-scalar case at any value of $\Lambda$, except near the threshold. This is because at lower scalar masses there is a bigger suppresion in the pseudoscalar variation (see (18) and (19)), therefore making the scalar variation visible at lower scalar masses. Note that always the corrections are valid below $\Lambda$, which is the effective cut-off theory, and at $M_\phi = \Lambda$ the curves in the plot fall steeply due to the logarithm behaviour in the correction. Complementary, Fig. 4b 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 than before. As it is expected, the $\Lambda$ quadratic 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 for masses near the threshold, for any $\Lambda$, the branching fraction corrections in the scalar case can become as big as one order of magnitude than the pseudoscalar case one.

![Graph](image)

(a)

![Graph](image)

(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 to point out that although our results for SM Higgs boson are less sensitive than the ones presented in Refs. [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 to SM Higgs decays, we have focus on $h \rightarrow b\bar{b}$ and $h \rightarrow \tau^+\tau^-$ 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, $\delta Br^{4F}(h \rightarrow b\bar{b})$ is roughly speaking an order of magnitude bigger than $\delta Br^{4F}(h \rightarrow \tau^+\tau^-)$ for any scale energy $\Lambda$, 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 $h \rightarrow b\bar{b}$ is $\sim 1\%$, meanwhile $h \rightarrow \tau^+\tau^-$ 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 $h \rightarrow b\bar{b}$ and $2.2 \times 10^{-2}\%$ for $h \rightarrow \tau^+\tau^-$. 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, and therefore having more precision in some observables. Additionally, it is expected higher precision measurements in the Higgs sector at both ILC and CLIC than LHC, allowing to explore more deeply into the quantitative information of the couplings and Higgs decay widths, and therefore the being 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 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 if 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 $\delta \Gamma^P_{4FI}$ and $\delta \Gamma^S_{4FI}$ can be distinguished in the lower mass threshold, when we have provided $y_s = 2.69$ and $y_p = 2.77$. However, it is important to note that in general (arbitrary values of $y_s$ and $y_p$) the condition of distinguishability is

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

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