Physics Letters B 708 (2012) 133–137

Contents lists available at SciVerse ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

Higgs search and flavor-safe fermion mass generation

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ARTICLE INFO

Article history:
Received 14 December 2011
Received in revised form 10 January 2012
Accepted 13 January 2012
Available online 20 January 2012
Editor: M. Cvetič

ABSTRACT

The combined analysis of ATLAS and CMS Collaborations has shrunk the possible mass range of the standard model Higgs boson to a 25 GeV window above 114 GeV, if it is not particularly heavy. Very recently, both collaborations have reported 2–3 σ excess of events in di-photon and four-lepton channels with invariants mass at 125 GeV. We study a scenario of electroweak symmetry breaking of allowing a fermiophobic sector to contribute to the weak gauge boson mass generation. Yukawa couplings are enhanced and couplings to W/Z bosons are suppressed. Consequently, the W-loop and top quark loop in $h \to \gamma\gamma$ significantly cancel each other. We study the phenomenological implications to the Higgs search in such scenario. Also discussed are implications of enhanced Yukawa couplings in flavor physics, in particular, the charged scalar contribution to the $B_d^0 - \bar{B}_d^0$ mixing.

1. Introduction

The origin of fundamental particle masses is closely related to the physics mechanism of spontaneous electroweak symmetry breaking (EWSB). It is one of the open questions in the standard model (SM). There are good reasons to expect this to be answered by current collider experiments. The simplest and the most economic mechanism of EWSB is via the minimal Higgs boson model. Recently, a combined analysis of ATLAS and CMS Collaborations provided a 95% C.L. exclusion of the SM Higgs boson in the mass range 141–476 GeV. Leading discovery channels for a light SM Higgs in the remaining window of about 25 GeV are $gg \to h \to \gamma\gamma$, associated productions Wh or Zh and weak boson fusions (WBF) with $h \to \tau^+\tau^-$. Very recently, both collaborations have reported 2–3 σ excess of events in di-photon and four-lepton channels [1]. The loop-induced $h \to \gamma\gamma$ decay is interesting since the SM background is relatively clean and understandable [2]. Similar to flavor constraints in low energy, the loop-induced $h \to \gamma\gamma$ decay and gluon fusion production $gg \to h$ probe virtual particles in loops. They are sensitive to many new physics. For models that were proposed to cancel the quadratic divergence, $gg \to h$ production is always reduced [3]. On the other hand, there exist two categories of models that enhance the gluon fusion $gg \to h$ production. In the standard model with fourth family (SM4), the additional fourth generation heavy quarks also contribute and enhance the $gg \to h$ [4]. In this Letter, we focus on the second category of models that can enhance top Yukawa coupling. The $h \to \gamma\gamma$ has both W-loop and top quark loop contributions but the two contributions are of opposite signs. In SM, the $h \to \gamma\gamma$ is completely dominated by the W-loop while in models with $gg \to h$ enhancement, the cancellation may become significant.

At hadron colliders, the search of SM Higgs heavily relies on the fact that the Higgs boson couples to weak gauge bosons and fermions simultaneously. Associated productions $W^{\pm}h$ or Zh with $h \to b\bar{b}$ result from both types of couplings. In the gluon fusion production $gg \to h$ with $h \to \gamma\gamma$, the production is due to Yukawa couplings of heavy quarks while the decay is dominated by the W-loop. WBF productions with $h \to \tau^+\tau^-$ depend also crucially on Yukawa couplings. For heavy SM Higgs with $m_h > 140$ GeV, $gg \to h$ with $h \to W^+W^-$ (WW^*) or ZZ become leading discovery channels. These channels play important role not only in discovering the Higgs boson but also in confirming the role of Higgs in EWSB.

Both masses of electroweak gauge bosons and SM fermions break electroweak symmetry. In addition, fermion masses break the chiral symmetry by flipping left-handed fields into right-handed ones. The Lagrangian for massless SM fermions is invariant under the chiral symmetry $U(3)_Q \otimes U(3)_u \otimes U(3)_d \otimes U(3)_\ell \otimes U(3)_e$ even after EWSB takes place. Thus, even though the sector that give rise to SM fermion masses must break EWSB, it is not necessarily the dominant source for generating weak gauge boson masses. In this Letter, we investigate implications of such a scenario in low energy measurements and collider physics.

EWSB models are constrained by various precision measurements. The S and T parameters are directly related to properties of weak gauge bosons [6]. Flavor changing neutral current (FCNC) effects constrain the mechanism of the SM fermion mass generation.

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0370-2693/$ – see front matter
doi:10.1016/j.physletb.2012.01.039

For instance, in extended TechniColor (ETC) models [7], in order to generate top quark mass, the ETC scale must be low. However, low ETC scale usually results in large FCNC effects since the ETC sector also couples to other fermions. Even for scenarios with SM fermion mass generation through scalars, there can be unacceptable FCNC at the tree level [5]. One has only two categories that satisfy the criterion for minimal flavor violation (MFV). In one case, one scalar doublet couples to u-type quarks, and another scalar doublet couples to d-type quarks. In the other case, the full $[U(3)]^5$ flavor symmetry is broken by Yukawa couplings to a single scalar doublet. In both cases, neutral scalars automatically have diagonal Yukawa couplings in the basis in which the quark mass matrix is diagonal.

In this Letter, we will use one doublet scalar to generate SM fermion masses via Yukawa couplings. It serves as the simplest scenario without large flavor violation at tree level. But the scalar is not solely responsible for weak gauge boson masses. There are additional sources for EWSB which contribute significantly to gauge boson masses. So Yukawa couplings are enhanced. To be concrete, we use the bosonic Technicolor (BTC) model [8–17] as the extra EWSB source. In this context, we study phenomenological implications of scenarios where EWSB and SM fermion mass generation come from two sources.

In Section 2, we briefly discuss spectra of BTC models. There exist now additional charged scalar states which couple to SM fermions with enhanced couplings. These enhanced Yukawa couplings may result in large FCNC mediated by charged scalars at loop levels. In Section 3, we use the FCNC, especially neutral meson mixings, to constrain Yukawa couplings. We then discuss possible implications in Higgs searches at colliders and conclude in the last section.

2. Theory realization

Since SM fermion masses break the electroweak symmetry, the fermion mass generation sector inevitably contributes to weak gauge boson masses. We assume that there is another fermiophobic sector that contributes to weak gauge boson masses [18]. Fermiophobic can be achieved through either SM gauge symmetries or additional symmetries.

Within the SM gauge symmetries, if the $SU(2)_L \times U(1)_Y$ symmetry is broken by the *vev* of a Higgs other than a doublet $(2,1)_1$ [19], the SM gauge symmetries automatically forbid the coupling between such Higgs and SM fermions thus fermiophobic. On the other hand, non-doublet Higgs in general breaks the $SU(2)_{L+R}$ custodial symmetry and generates unacceptable contribution to ΔT. Their *vevs* are usually constrained to be smaller than $\mathcal{O}(1\text{ GeV})$. The Georgi–Machacek model is an exception [20], where two $SU(2)$ triplet Higgs are introduced and the ΔT is under control even with larger *vev*. Such fermiophobic scalars ϕ have very unique phenomenology since they can only be produced via $W\phi$, $Z\phi$ associate production and WBF. For light scalars $m_\phi < 130$ GeV, with the absence of $\phi b\bar{b}$ coupling, $\phi \to \gamma\gamma$ becomes the dominate decay channel [18,21]. The heavy scalar decay would behave similarly to the SM-like Higgs except that there is no $t\bar{t}$ mode. In addition, the gg fusion production is absent and the search of such heavy fermiophobic Higgs is through WWW channel (as in [22]).

In a BTC theory, the electroweak symmetry is broken by the condensation of techni-quarks and the WW scattering amplitude is unitarized by both techni-pion π_{TC} and techni-rho ρ_{TC}. By assignment, the strongly interacting sector does not couple to SM fermions. There exists a scalar doublet in the theory which does not develop *vev* but the scalar couples to both SM fermions and the techni-fermions. The SM fermion masses arise via techni-fermions confinement and the strength of Yukawa couplings determines the mass of such fermion.

In both fermiophobic Higgs models and the BTC model, the Yukawa couplings between the Higgs-like scalar and the SM fermions are in general enhanced since the EWSB is dominated by a fermiophobic source. In this Letter, we use the BTC model to illustrate the general feature and phenomenological implications of enhanced Yukawa couplings.

Using the non-linear representation, one can define Σ [17]

$$\Sigma = \exp(2i\Pi/f), \quad \Pi = \begin{pmatrix} \pi^0/2 & \pi^+/\sqrt{2} \\ \pi^-/\sqrt{2} & -\pi^0/2 \end{pmatrix}, \tag{1}$$

where Π represents an isotriplet of techni-pions and f is their decay constant. The theory also contains a scalar doublet Φ that couples to SM fermions.

$$\Phi = \begin{pmatrix} \overline{\phi^0} & \phi^+ \\ -\phi^- & \phi^0 \end{pmatrix}. \tag{2}$$

Φ mixes with the techni-pion state through Yukawa couplings to techni-fermions. For convenience, one can rewrite the Φ field in the non-linear form. By expanding around the true vacuum, one obtains

$$\Phi = \frac{\sigma + f'}{\sqrt{2}}\Sigma', \quad \Sigma' = \exp\big(2i\Pi'/f'\big), \tag{3}$$

where f' is the *vev* of ϕ and Π' represents its isotriplet components.

Leading terms in the Lagrangian for the Φ and Σ fields are

$$\begin{aligned}\mathcal{L}_{KE} &= \frac{1}{2}\partial_\mu\sigma\partial^\mu\sigma + \frac{f^2}{4}\operatorname{Tr}\big(D_\mu\Sigma^\dagger D^\mu\Sigma\big) \\ &\quad + \frac{(\sigma+f')^2}{4}\operatorname{Tr}\big(D_\mu\Sigma'^\dagger D^\mu\Sigma'\big),\end{aligned} \tag{4}$$

where the covariant derivative is given by

$$D^\mu\Sigma = \partial^\mu\Sigma - igW_a^\mu\frac{\tau^a}{2}\Sigma + ig'B^\mu\Sigma\frac{\tau^3}{2}. \tag{5}$$

There exist specific linear combinations of pion fields:

$$\pi_A = \frac{f\Pi + f'\Pi'}{\sqrt{f^2+f'^2}}, \tag{6}$$

which mix with gauge fields in quadratic terms with derivatives. Such states are unphysical and can be gauged away. Physical states π_p arise from orthogonal linear combinations,

$$\pi_p = \frac{-f'\Pi + f\Pi'}{\sqrt{f^2+f'^2}}. \tag{7}$$

The physical pion mass can be obtained by potential terms involving couplings between Φ and Σ, which originate from techni-fermion Yukawa couplings and techni-fermion condensation. The mass of physical pions are model dependent. In the unitary gauge, weak gauge boson masses arise from the remaining quadratic terms

$$m_W^2 = \frac{1}{4}g^2v^2, \qquad m_Z^2 = \frac{1}{4}\big(g^2+g'^2\big)v^2, \tag{8}$$

where v is the electroweak scale $v \equiv \sqrt{f^2+f'^2} = 246$ GeV.

For simplicity, we define a mixing angle θ

$$\sin\theta = \frac{f'}{v}, \qquad \cos\theta = \frac{f}{v}. \tag{9}$$

As long as SM fermions get masses, f' is non-zero and $\cos\theta$ never reaches 1. In the limit of $\sin\theta = 1$, σ corresponds to the SM Higgs. For other θ, σ behaves like the SM Higgs but of different couplings. Expanding the third term in Eq. (4), we find couplings between σ and gauge bosons

$$\mathcal{L}_{\sigma WZ} = 2\sin\theta \frac{m_W^2}{v}\sigma W^{+\mu} W_\mu^- + \sin\theta \frac{m_Z^2}{v}\sigma Z^\mu Z_\mu, \tag{10}$$

which are reduced by a factor of $\sin\theta$, compared with those in the SM.

Couplings of Φ to quarks are given by regular Yukawa couplings. Couplings of σ to fermions are

$$\mathcal{L}_{\sigma\bar{f}f} = -\sum_{\text{fermions}} \frac{1}{\sin\theta}\frac{\sqrt{2}m_f}{v}\sigma\bar{f}f. \tag{11}$$

σ is mostly Φ-like. The σ Yukawa couplings are larger than the SM Yukawa couplings by a factor of $\csc\theta = v/f'$. Similarly, the charged physical pion which is a mixing state of $\phi^\pm$ and $\pi^\pm$ also couples to SM fermions. The couplings are slightly different from those of σ by a factor of $\cos\theta = f/v$

$$\mathcal{L}_{\pi_p^\pm} = -i\cot\theta\frac{m_{u_i}}{v}\pi_p^-\bar{d}_{iL}u_{iR} - i\cot\theta\frac{m_{d_j}}{v}\pi_p^+\bar{u}_L^i V_{CKM}^{ij} d_R^j + \text{h.c.} \tag{12}$$

Here, indices i, j stand for flavor. In the limit of $\sin\theta = 1$, these couplings vanish and the charged degree of freedom basically disappears.

3. Flavor constraints from enhanced Yukawa coupling

Tree level FCNC mediated by the neutral scalar can be avoided if quark mass matrices and Yukawa couplings are diagonalized simultaneously [5]. Couplings between the charged scalar π_p^+ and SM fermions are flavor violating. The $\csc\theta$ enhancement of couplings renders constraints more severe. We now examine bounds on Yukawa couplings due to FCNC processes involving the charged scalar.[1]

At tree level, the charged scalar contributes to flavor violating rare decays like $B_u \to \tau\nu_\tau$, $B \to D\tau\nu_\tau$. The SM prediction is $\text{BR}(B_u \to \tau\nu_\tau)_{\text{SM}} = (0.95 \pm 0.27) \times 10^{-4}$ while the current experimental value is $\text{BR}(B_u \to \tau\nu_\tau)_{\text{exp}} = (1.65 \pm 0.34) \times 10^{-4}$. The $W^\pm$-mediated SM contribution is suppressed due to helicity suppression. Even though the charged scalar may be much heavier, the no-helicity-suppression makes its contribution comparable to the SM contribution. The decay amplitude is proportional to

$$V_{ub}\frac{m_b}{f'}\frac{m_\tau}{f'}\frac{f^2}{v^2}. \tag{13}$$

In the type-II 2HDM, a similar contribution mediated by the charged Higgs only becomes relevant when the $\tan\beta \gtrsim 10$ for very light charged Higgs of $M_{H^+} = 100$ GeV but the bound for charged Higgs of 700 GeV is $\tan\beta < 60$ [23]. Translating such bound into the current model, the top quark Yukawa coupling is greater than 4π and become strongly coupled. Therefore, in the range that we are interested in, the bound from $B_u \to \tau\nu_\tau$ is irrelevant.

Recent search of $\mu \to e\gamma$ at the MEG experiment will soon reach $\text{BR}(\mu \to e\gamma) \simeq 1 \times 10^{-13}$. The one-loop contribution from charged state to $\mu \to e\gamma$ is suppressed by small lepton masses

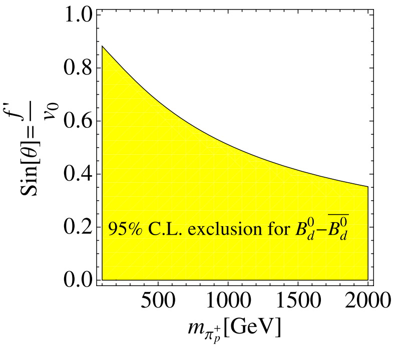

Fig. 1. Constraints on the f'/v and charged scalar mass $m_{\pi_p^+}$ from ΔM_{B_d}. The yellow region is excluded at 95% C.L. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this Letter.)

and additional helicity-flip. The largest contribution in Higgs mediated $\mu \to e\gamma$ is usually the Barr–Zee two-loop effects involving the charged scalar coupling to a top-bottom loop. The two-loop involving Z and neutral scalar does not exist since there is no tree level flavor violating vertices in the neutral scalar coupling to SM fermions. Again, similar study in type-II 2HDM has shown the contribution only reach the sensitivity for $\tan\beta$ of 60 [24]. The current model is not expected to receive constraints from lepton flavor violating experiments like $\mu \to e\gamma$.

Given the large top quark mass, the most severe constraints come from neutral B_d-mixing or B_s-mixing since these processes do not have suppression from m_b. The charged Higgs mediated B_d-mixing had been calculated for 2HDM [25]

$$\Delta M_{B_d} = \frac{G_F^2 m_t^2 f_{B_d}^2 \hat{B}_d M_B |V_{td}^* V_{tb}|^2 \eta_b}{24\pi^2} \times \left[I_{WW}\left(y^W\right) + I_{W\Pi}\left(y^W, y^\Pi, x\right) + I_{\Pi\Pi}\left(y^\Pi\right)\right], \tag{14}$$

where

$$y_W = \frac{m_t^2}{m_W^2}, \qquad y_\Pi = \frac{m_t^2}{m_{\pi_p^+}^2}, \qquad x = \frac{m_{\pi_p^+}^2}{m_W^2}, \tag{15}$$

and

$$\begin{aligned} I_{WW} &= 1 + \frac{9}{1-y^W} - \frac{6}{(1-y^W)^2} - \frac{6}{y^W}\left(\frac{y^W}{1-y^W}\right)^3 \ln y^W, \\ I_{W\Pi} &= \lambda_{tt}^2 y^\Pi \left[\frac{(2x-8)\ln y^\Pi}{(1-x)(1-y^\Pi)^2} + \frac{6x\ln y^W}{(1-x)(1-y^W)^2} - \frac{8-2y^W}{(1-y^W)(1-y^\Pi)}\right], \\ I_{\Pi\Pi} &= \lambda_{tt}^4 y^\Pi \left[\frac{1+y^\Pi}{(1-y^\Pi)^2} + \frac{2y^\Pi \ln y^\Pi}{(1-y^\Pi)^3}\right]. \end{aligned} \tag{16}$$

The QCD running of the Wilson coefficient is given by $\eta_b = 0.552$ [26]. The non-perturbative decay constant f_{B_d} and the bag parameter $\hat{B}_d$ from lattice QCD calculation is $f_{B_d}\hat{B}_d^{1/2} = 216 \pm 15$ MeV [27]. We take the conservative value $f_{B_d}\hat{B}_d^{1/2} = 201$ MeV. Since the theoretical uncertainty in $f_B\hat{B}^{1/2}$ dominates experimental uncertainties and they are correlated in B_s and B_d decays, we use only ΔM_{B_d}. It has the smallest total uncertainty and the current experiment measurement for B_d mixing is

$$\Delta M_{B_d} = (507 \pm 4) \times 10^9 \hbar\ \text{s}^{-1}. \tag{17}$$

Shown in Fig. 1 are correlated constraints on the mixing angle $\sin\theta$ and charged scalar mass $m_{\pi_p^+}$ from ΔM_{B_d}. $\sin\theta$ represents

[1] The mixing is also constrained by precision electroweak measurement like S-parameter. However, that depends on particular model of strong dynamics which is not focus of this Letter. Therefore, we don't discuss its implications.

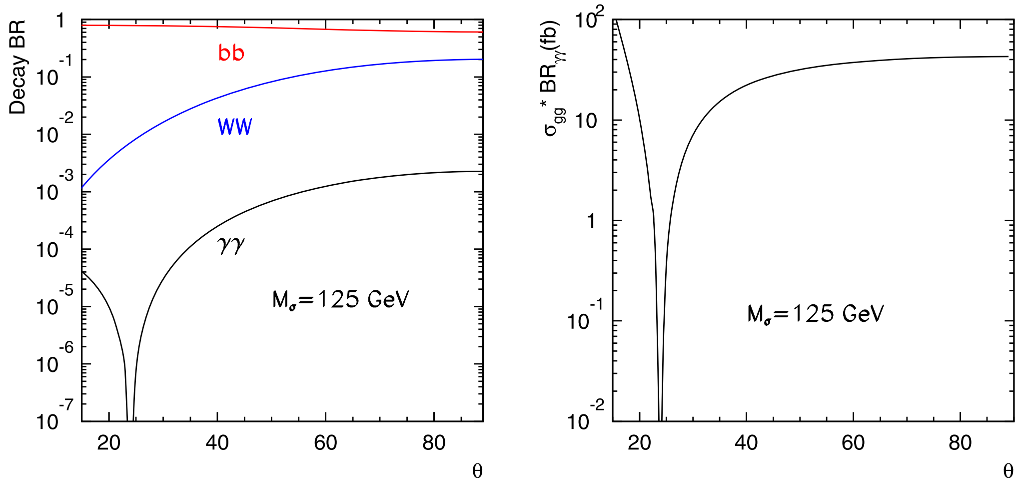

Fig. 2. (a) Decay BR of σ; (b) $\sigma(gg \to h) \cdot \mathrm{BR}(h \to \gamma\gamma)$ production rate.

the ratio of top quark mass effects in the total EWSB. In the limit of $\sin\theta = 1$ or $f' = v$, the charged scalar state disappears and σ corresponds to the SM Higgs boson.

The correlation in Fig. 1 constrains charged scalar masses and the charged scalar coupling to SM fermions.

4. Phenomenological implications

In BTC models, couplings of the lightest neutral scalar σ to weak gauge bosons are suppressed by $\sin\theta$ while those to SM fermions enhanced by $1/\sin\theta$, compared with the SM Higgs boson. This feature significantly changes Higgs phenomenologies. An immediate consequence is that $W\sigma$ associated production rate as well as the weak boson fusion production rate are both reduced. We use the recently reported excess region $m_h \simeq 125$ GeV to illustrate the phenomenology of σ particle.

The most important channel for light SM Higgs discovery is the di-photon search. Similar to the SM Higgs, the $\sigma \to \gamma\gamma$ has both WW and fermion loop contribution. Now the latter is enhanced and the former is reduced. The $\sigma \to \gamma\gamma$ partial width is

$$\Gamma(\sigma \to \gamma\gamma) = \frac{G_\mu \alpha^2 M_H^3}{128\sqrt{2}\pi^3} \left| \sum_f N_c Q_f^2 A_{1/2}^H(\tau_f)/\sin\theta + A_1^H(\tau_W)\sin\theta \right|^2 \tag{18}$$

where

$$A_{1/2}^H(\tau) = 2\big[\tau + (\tau - 1) f(\tau)\big]\tau^{-2},$$
$$A_1^H(\tau) = -\big[2\tau^2 + 3\tau + 3(2\tau - 1) f(\tau)\big]\tau^{-2},$$
$$f(\tau) = \begin{cases} \arcsin^2\sqrt{\tau}; & \tau \leqslant 1, \\ -\frac{1}{4}\big[\log\frac{1+\sqrt{1-\tau^{-1}}}{1-\sqrt{1-\tau^{-1}}}\big]^2; & \tau > 1. \end{cases} \tag{19}$$

Notice that the W-loop and heavy quark loop interfere destructively. In the case of SM Higgs boson, the W-loop dominates. With the reduced coupling in σWW and the enhanced coupling in $\sigma f\bar{f}$, the destructive interference effect become significant in the di-photon channel. $\Gamma(\sigma \to \gamma\gamma)$ can be much smaller than that of the SM Higgs with the same mass. Shown in Fig. 2 are (a) the BR of σ to various important searching modes and (b) the total width of σ, for $\theta = \pi/4$. The cancellation between the W-loop and top quark loop contribution maximizes at $\theta = 25$ degree. Given the enhancement in both top quark Yukawa and total width of σ, the production rate of $\sigma(gg \to h) \cdot \mathrm{BR}(h \to \gamma\gamma)$ is always smaller than the case of SM in $\theta \ni [25, 90]$.

The σ production via gluon fusion with $\sigma \to b\bar{b}$ decay would encounter tremendously irreducible QCD $b\bar{b}$ background. Other standard searches through associate productions $W\sigma$, $Z\sigma$ with $\sigma \to b\bar{b}$ are suppressed by a factor of $\sin^2\theta$. The light σ search is then much more challenging than the SM Higgs boson. The only all-enhanced channel is through $t\bar{t}\sigma$ with $\sigma \to \tau^+\tau^-$ or $h \to b\bar{b}$ and more realistic studies are on going [28].

5. Conclusions

In this Letter, we have used the minimal BTC model to illustrate some general features of scenarios, where weak gauge boson mass generation is dominated by a fermiophobic sector and the SM fermion mass generation due to a light scalar doublet. Yukawa couplings of the light Higgs to SM fermions are enhanced while couplings to weak gauge bosons reduced. The extra charged scalar state inevitably mediate flavor changing neutral current processes with enhanced Yukawa couplings. For them, ΔM_{B_d} provides the most severe bound. We use the bench-mark mass $m_\sigma = 125$ GeV to illustrate the phenomenology of σ. In the di-photon channel, the destructive interference between the top quark loop and the W-boson loop becomes significant, which may render the partial width negligible in some regions of model parameter space.

Acknowledgements

We thank Guohuai Zhu for useful discussions. The work is supported in part, by the Fundamental Research Funds for the Central Universities 2011QNA3017, the National Science Foundation of China (10875103, 11135006), National Basic Research Program of China (2010CB833000), and Zhejiang University Group Funding (2009QNA3015).

References

[1] ATLAS Collaboration, ATLAS-CONF-2011-163;
CMS Collaboration, CMS-PAS-HIG-11-011.

[2] C. Balazs, E.L. Berger, P.M. Nadolsky, C.-P. Yuan, Phys. Lett. B 637 (2006) 235, hep-ph/0603037;
T. Aaltonen, et al., CDF Collaboration, Phys. Rev. Lett. 107 (2011) 102003, arXiv:1106.5123 [hep-ex].

[3] M.S. Carena, S. Heinemeyer, C.E.M. Wagner, G. Weiglein, Eur. Phys. J. C 26 (2003) 601, hep-ph/0202167.

[4] Q. Li, M. Spira, J. Gao, C.S. Li, Phys. Rev. D 83 (2011) 094018, arXiv:1011.4484 [hep-ph].
[5] S.L. Glashow, S. Weinberg, Phys. Rev. D 15 (1977) 1958.
[6] M.E. Peskin, T. Takeuchi, Phys. Rev. Lett. 65 (1990) 964;
M.E. Peskin, T. Takeuchi, Phys. Rev. D 46 (1992) 381.
[7] S. Dimopoulos, L. Susskind, Nucl. Phys. B 155 (1979) 237.
[8] A. Kagan, S. Samuel, Phys. Lett. B 252 (1990) 605;
A. Kagan, S. Samuel, Phys. Lett. B 270 (1991) 37;
A. Kagan, S. Samuel, Phys. Lett. B 284 (1992) 289.
[9] E.H. Simmons, Nucl. Phys. B 312 (1989) 253.
[10] C.D. Carone, H. Georgi, Phys. Rev. D 49 (1994) 1427, hep-ph/9308205.
[11] C.D. Carone, E.H. Simmons, Nucl. Phys. B 397 (1993) 591, hep-ph/9207273.
[12] C.D. Carone, M. Golden, Phys. Rev. D 49 (1994) 6211, hep-ph/9312303.
[13] C.D. Carone, E.H. Simmons, Y. Su, Phys. Lett. B 344 (1995) 287, hep-ph/9410242.
[14] V. Hemmige, E.H. Simmons, Phys. Lett. B 518 (2001) 72, hep-ph/0107117.
[15] C.X. Yue, Y.P. Kuang, G.R. Lu, L.D. Wan, Mod. Phys. Lett. A 11 (1996) 289;
X.L. Wang, B. Huang, G.R. Lu, Y.D. Yang, H.B. Li, Commun. Theor. Phys. 27 (1997) 325;
Y.G. Cao, Z.K. Jiao, Commun. Theor. Phys. 38 (2002) 47.
[16] M. Antola, M. Heikinheimo, F. Sannino, K. Tuominen, arXiv:0910.3681 [hep-ph].
[17] C.D. Carone, J. Erlich, J.A. Tan, Phys. Rev. D 75 (2007) 075005, hep-ph/0612242.
[18] A.G. Akeroyd, Phys. Lett. B 368 (1996) 89, hep-ph/9511347.
[19] T. Han, B. Mukhopadhyaya, Z. Si, K. Wang, Phys. Rev. D 76 (2007) 075013, arXiv:0706.0441 [hep-ph];
P. Fileviez Perez, T. Han, G.-y. Huang, T. Li, K. Wang, Phys. Rev. D 78 (2008) 015018, arXiv:0805.3536 [hep-ph];
P. Fileviez Perez, H.H. Patel, M.J. Ramsey-Musolf, K. Wang, Phys. Rev. D 79 (2009) 055024, arXiv:0811.3957 [hep-ph].
[20] H. Georgi, M. Machacek, Nucl. Phys. B 262 (1985) 463.
[21] T. Aaltonen, et al., CDF Collaboration, Phys. Rev. Lett. 103 (2009) 061803, arXiv:0905.0413 [hep-ex];
V.M. Abazov, et al., D0 Collaboration, Phys. Rev. Lett. 101 (2008) 051801, arXiv:0803.1514 [hep-ex].
[22] T. Han, H.S. Liu, M.x. Luo, K. Wang, W. Wu, Phys. Rev. D 80 (2009) 095010, arXiv:0908.2186 [hep-ph].
[23] F. Mahmoudi, O. Stal, Phys. Rev. D 81 (2010) 035016, arXiv:0907.1791 [hep-ph].
[24] J. Hisano, S. Sugiyama, M. Yamanaka, M.J.S. Yang, Phys. Lett. B 694 (2011) 380, arXiv:1005.3648 [hep-ph].
[25] C.Q. Geng, J.N. Ng, Phys. Rev. D 38 (1988) 2857;
P. Ball, R. Fleischer, Eur. Phys. J. C 48 (2006) 413, hep-ph/0604249.
[26] G. Buchalla, A.J. Buras, M.E. Lautenbacher, Rev. Mod. Phys. 68 (1996) 1125, hep-ph/9512380.
[27] E. Gamiz, et al., HPQCD Collaboration, Phys. Rev. D 80 (2009) 014503, arXiv:0902.1815 [hep-lat].
[28] K.X. Fu, M.X. Luo, L.Y. Shan, K. Wang, in preparation.