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# Bounds on Heavy Chiral Fermions <sup>1</sup>

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## Abstract

We derive the low-energy electroweak effective lagrangian for the case of additional heavy, unmixed, sequential fermions. Present data still allow for the presence of a new quark and/or lepton doublet with masses greater than  $M_Z/2$ , provided that these multiplets are sufficiently degenerate. Deviations of the effective lagrangian predictions from a full one-loop computation are sizeable only for fermion masses close to the threshold  $M_Z/2$ . Some of the constraints on new sequential fermions coming from accelerator results and cosmological considerations are presented. We point out that the new fermions can significantly affect the production and decay rate into  $\gamma\gamma$  of the intermediate Higgs at LHC.

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LEP precision data represent a step of paramount relevance in probing extensions of the Standard Model (SM). Through their virtual effects, the electroweak radiative corrections "feel" the presence of new particles running in the loops and the level of accuracy on the relevant observables is such that this set of tests is complementary to the traditional probes on virtual effects due to new physics (i.e. highly suppressed or forbidden flavour changing neutral current phenomena). In some cases, as that which we aim to discuss here, the electroweak precision tests represent the only indirect way to search for these new particles.

In this letter we will discuss electroweak radiative effects from extensions of the ordinary fermionic spectrum of the SM. The new fermions are supposed to possess the same colour and electroweak quantum numbers as the ordinary ones and to mix very thinly with the ordinary three generations.

The most straightforward realization of such a fermionic extension of the SM is the introduction of a fourth generation of fermions. This possibility has been almost entirely jeopardized by the LEP bound on the numbers of neutrinos species. Although there still exists the obvious way out of having new fermion generations with neutrinos of mass  $\geq M_Z/2$ , we think that these options are awkward enough not to deserve further studies. Rather, what we have in mind in tackling this problem are general frames discussing new physics beyond the SM which lead to new quarks and/or leptons classified in the usual chiral way with iso-doublets and iso-singlets for different chiralities.

Situations of this kind may be encountered in grand unified schemes where the ordinary fifteen Weyl spinors of each fermionic generation are only part of larger representations or where new fermions (possibly also mirror fermions) are requested by the group or manifold structure of the schemes. Chiral fermions with heavy static masses may also provide a first approximation of virtual effects in technicolor-like schemes when the dynamical behaviour of technifermion self-energies are neglected.

Although such effects have been extensively investigated in the literature [1], our presentation focuses mainly on two aspects, which have been only partially touched in the previous analyses: the use of effective lagrangians for a model-independent treatment of the problem and a discussion of the validity of this approach in comparison with the computation in the full-fledged theory.

While separate tests can be set up for each different extension of the SM, there may be some advantage in realizing this analysis in a model independent framework. The natural theoretical tool to this purpose is represented by an effective electroweak lagrangian where, giving up the renormalizability requirement, all  $SU(2)_L \otimes U(1)_Y$  invariant operators up to a given dimension are present with unknown coefficients, to be eventually determined from the experiments. Each different model fixes uniquely this set of coefficients and the effective lagrangian becomes in this way a common ground to discuss and compare several SM extensions. The introduction of the well known  $S$ ,  $T$  and  $U$  [2] or  $\epsilon$ 's [3] variables was much in the same spirit and the use of an effective lagrangian represents in a sense the natural extension of these approaches.

The use of an effective lagrangian for the electroweak physics has been originally advocated for the study of the large Higgs mass limit in the SM [4, 5, 6]. Subsequently, contributions from chiral  $SU(2)_L$  doublets have been considered in the degenerate case [7], for small splitting [8] and in the case of infinite splitting [9, 10]. In the present note we will deal with the general case of arbitrary splitting among the fermions in the doublet. Our results will be used to test the model with the latest available data.

The use of an effective lagrangian in precision tests has its own limitations, which are also discussed in the present note. In particular we are going to use an effective electroweak lagrangian organized in a derivative expansion which we truncate at the fourth order. When discussing two-point vector boson functions  $-i\Pi_{ij}^{\mu\nu}(q)$  ( $i, j = 0, 1, 2, 3$ ), this amounts to keep only the constant and linear terms in  $q^2$ :

$$\Pi_{ij}^{\mu\nu}(q) = \Pi_{ij}(q^2)g^{\mu\nu} + (q^\mu q^\nu \text{ terms}) \quad (1)$$

$$\Pi_{ij}(q^2) \equiv A_{ij} + q^2 F_{ij}(q^2) = A_{ij} + q^2 F_{ij}(0) + \dots \quad (2)$$

The next terms in the  $q^2$  expansion are suppressed by increasing powers of  $q^2/M^2$ ,  $M$  generically representing the mass of the particles running in the loop. One can ask how large has to be  $M$  to obtain a sensible approximation from the truncation of the full one-loop result. As we will see, for LEP I observables, the truncation is a very good approximation already for relatively light fermions, with masses around  $70 - 80 \text{ GeV}$ .

At the end of this note we will add some comments on the direct searches of new quarks and leptons and on the modifications induced by additional chiral fermions in the  $\gamma\gamma$  signature for an intermediate Higgs at LHC.

For new chiral fermions which do not mix with the ordinary ones, the virtual effects measurable at LEP 1 are all described by operators bilinear in the gauge vector bosons. Here, for completeness, we consider the standard list [5] of CP conserving operators containing up to four derivatives and built out of the gauge vector bosons  $W_\mu^i$  ( $i = 1, 2, 3$ ),  $B_\mu$  and the would be Goldstone bosons  $\xi^i$ :

$$\begin{aligned} \mathcal{L}_0 &= \frac{v^2}{4} [\text{tr}(TV_\mu)]^2 \\ \mathcal{L}_1 &= i\frac{gg'}{2} B_{\mu\nu} \text{tr}(T\hat{W}^{\mu\nu}) \\ \mathcal{L}_2 &= i\frac{g'}{2} B_{\mu\nu} \text{tr}(T[V^\mu, V^\nu]) \\ \mathcal{L}_3 &= g \text{tr}(\hat{W}_{\mu\nu}[V^\mu, V^\nu]) \\ \mathcal{L}_4 &= [\text{tr}(V_\mu V_\nu)]^2 \\ \mathcal{L}_5 &= [\text{tr}(V_\mu V^\mu)]^2 \\ \mathcal{L}_6 &= \text{tr}(V_\mu V_\nu) \text{tr}(TV^\mu) \text{tr}(TV^\nu) \\ \mathcal{L}_7 &= \text{tr}(V_\mu V^\mu) [\text{tr}(TV^\nu)]^2 \end{aligned}$$

$$\begin{aligned}
\mathcal{L}_8 &= \frac{g^2}{4} [\text{tr}(T\hat{W}_{\mu\nu})]^2 \\
\mathcal{L}_9 &= \frac{g}{2} \text{tr}(T\hat{W}_{\mu\nu}) \text{tr}(T[V^\mu, V^\nu]) \\
\mathcal{L}_{10} &= [\text{tr}(TV_\mu) \text{tr}(TV_\nu)]^2 \\
\mathcal{L}_{11} &= \text{tr}((\mathcal{D}_\mu V^\mu)^2) \\
\mathcal{L}_{12} &= \text{tr}(T\mathcal{D}_\mu \mathcal{D}_\nu V^\nu) \text{tr}(TV^\mu) \\
\mathcal{L}_{13} &= \frac{1}{2} [\text{tr}(T\mathcal{D}_\mu V_\nu)]^2 \\
\mathcal{L}_{14} &= ig\epsilon^{\mu\nu\rho\sigma} \text{tr}(\hat{W}_{\mu\nu} V_\rho) \text{tr}(TV_\sigma)
\end{aligned} \tag{3}$$

We recall the notation:

$$\begin{aligned}
T &= U\tau^3U^\dagger \quad , \\
V_\mu &= (D_\mu U)U^\dagger \quad ,
\end{aligned} \tag{4}$$

$$U = e^{i\frac{\vec{\xi} \cdot \vec{\tau}}{v}} \quad , \tag{5}$$

$$D_\mu U = \partial_\mu U - g\hat{W}_\mu U + g'U\hat{B}_\mu \quad , \tag{6}$$

$\hat{W}_\mu, \hat{B}_\mu$  are matrices collecting the gauge fields:

$$\begin{aligned}
\hat{W}_\mu &= \frac{1}{2i} \vec{W}_\mu \cdot \vec{\tau} \quad , \\
\hat{B}_\mu &= \frac{1}{2i} B_\mu \tau^3 \quad .
\end{aligned} \tag{7}$$

The corresponding field strengths are given by:

$$\begin{aligned}
\hat{W}_{\mu\nu} &= \partial_\mu \hat{W}_\nu - \partial_\nu \hat{W}_\mu - g[\hat{W}_\mu, \hat{W}_\nu] \quad , \\
\hat{B}_{\mu\nu} &= \partial_\mu \hat{B}_\nu - \partial_\nu \hat{B}_\mu \quad .
\end{aligned} \tag{8}$$

Finally the covariant derivative acting on  $V_\mu$  is given by:

$$\mathcal{D}_\mu V_\nu = \partial_\mu V_\nu - g[\hat{W}_\mu, V_\nu] \quad . \tag{9}$$

The effective electroweak lagrangian reads:

$$\mathcal{L}_{eff} = \mathcal{L}_{SM} + \sum_{i=0}^{14} a_i \mathcal{L}_i \quad , \tag{10}$$

where  $\mathcal{L}_{SM}$  is the SM lagrangian. Here we do not include the Wess-Zumino term [10]. For an extra doublet of fermions (quarks or leptons), we have determined the coefficients  $a_i$  ( $i = 0, \dots, 14$ ), by computing the corresponding one-loop contribution to a set of  $n$ -point gauge boson functions ( $n = 2, 3, 4$ ), in the limit of low external momenta and by matching

the predictions of the full and effective theories. By denoting with  $M$  and  $m$  the masses of the upper and lower weak isospin components, respectively, we obtain, in units of  $1/16\pi^2$ :

$$\begin{aligned}
a_0^q &= \frac{3M^2}{2v^2} \left( \frac{1 - r^2 + 2r \log r}{1 - r} \right) \\
a_1^q &= \frac{1}{12(-1 + r)^3} \left[ 3(1 - 15r + 15r^2 - r^3) + 2(1 - 12r - 6r^2 - r^3) \log r \right] \\
a_2^q &= \frac{1}{12(-1 + r)^3} \left[ 3(3 - 7r + 5r^2 - r^3) + 2(1 - r^3) \log r \right] \\
a_3^q &= \frac{1}{8(-1 + r)^3} \left[ 3(-1 + 7r - 7r^2 + r^3) + 6r(1 + r) \log r \right] \\
a_4^q &= \frac{1}{6(-1 + r)^3} \left[ 5 - 9r + 9r^2 - 5r^3 + 3(1 + r^3) \log r \right] \\
a_5^q &= \frac{1}{24(-1 + r)^3} \left[ -23 + 45r - 45r^2 + 23r^3 - 12(1 + r^3) \log r \right] \\
a_6^q &= \frac{1}{24(-1 + r)^3} \left[ -23 + 81r - 81r^2 + 23r^3 - 6(2 - 3r - 3r^2 + 2r^3) \log r \right] \\
a_7^q &= -a_6^q \\
a_8^q &= \frac{1}{12(-1 + r)^3} \left[ 7 - 81r + 81r^2 - 7r^3 + 6(1 - 6r - 6r^2 + r^3) \log r \right] \\
a_9^q &= -a_6^q \\
a_{10}^q &= 0 \\
a_{11}^q &= -\frac{1}{2} \\
a_{12}^q &= \frac{1}{8(-1 + r)^3} \left[ 1 + 9r - 9r^2 - r^3 + 6r(1 + r) \log r \right] \\
a_{13}^q &= 2a_{12}^q \\
a_{14}^q &= \frac{3}{8(-1 + r)^2} \left[ 1 - r^2 + 2r \log r \right] \tag{11}
\end{aligned}$$

for quarks and:

$$\begin{aligned}
a_i^l &= \frac{1}{3} a_i^q \quad (i = 0, \quad i = 3, \dots, 14) \\
a_1^l &= \frac{1}{12(-1 + r)^3} \left[ 1 - 15r + 15r^2 - r^3 - 2(1 + 6r^2 - r^3) \log r \right] \\
a_2^l &= \frac{1}{12(-1 + r)^3} \left[ -1 - 3r + 9r^2 - 5r^3 - 2(1 - r^3) \log r \right] \tag{12}
\end{aligned}$$

for leptons, where:

$$r = \frac{m^2}{M^2} \tag{13}$$

The coefficients  $a_i$  of the effective lagrangian  $\mathcal{L}_{eff}$  are related to measurable parameters. In particular, to make contact with the LEP data, we recall that, by neglecting higher

derivatives, the relation between the effective lagrangian  $\mathcal{L}_{eff}$  and the  $\epsilon$ 's parameters, is given by:

$$\begin{aligned}\delta\epsilon_1 &= 2a_0, \\ \delta\epsilon_2 &= -g^2(a_8 + a_{13}), \\ \delta\epsilon_3 &= -g^2(a_1 + a_{13}).\end{aligned}\tag{14}$$

The  $\epsilon$  parameters are obtained by adding to  $\delta\epsilon_i$  the SM contribution  $\epsilon_i^{SM}$ , which we regard as functions of the Higgs and top quark masses. From eqs. (11) and (12) one finds:

$$\delta\epsilon_1^q = 3\delta\epsilon_1^l = \frac{3M^2}{8\pi^2} \frac{G}{\sqrt{2}} \left[ \frac{1-r^2+2r\log r}{(1-r)} \right]\tag{15}$$

$$\delta\epsilon_2^l = 3\delta\epsilon_2^l = \frac{Gm_W^2}{12\pi^2\sqrt{2}} \left[ \frac{5-27r+27r^2-5r^3+(3-9r-9r^2+3r^3)\log r}{(1-r)^3} \right]\tag{16}$$

$$\delta\epsilon_3^q = \frac{Gm_W^2}{12\pi^2\sqrt{2}} [3 + \log r]\tag{17}$$

$$\delta\epsilon_3^l = \frac{Gm_W^2}{12\pi^2\sqrt{2}} [1 - \log r]\tag{18}$$

A recent analysis of the available precision data from LEP, SLD, low-energy neutrino scatterings and atomic parity violation experiments, leads to the following values for the  $\epsilon$  parameters [11]:

$$\begin{aligned}\epsilon_1 &= (3.6 \pm 1.5) \cdot 10^{-3} \\ \epsilon_2 &= (-5.8 \pm 4.3) \cdot 10^{-3} \\ \epsilon_3 &= (3.6 \pm 1.5) \cdot 10^{-3}\end{aligned}\tag{19}$$

Notice the relatively large error in the determination of  $\epsilon_2$ , mainly dominated by the uncertainty on the  $W$  mass. We illustrate our result in fig. 1, in the plane  $(\epsilon_1, \epsilon_3)$ , for the case of an extra quark doublet. The upper ellipsis represents the  $1\sigma$  experimentally allowed region, obtained by combining all LEP data.

If one also includes the SLD determination of the left-right asymmetry, then one gets the lower ellipsis. The predictions from an additional heavy quark doublet are given by the dashed line, obtained by fixing one of the masses to  $200\text{ GeV}$  and letting the other vary from  $200\text{ GeV}$  to  $300\text{ GeV}$ . One has in this way two branches, according to which mass,  $m$  or  $M$ , has been fixed. The top and Higgs masses has been fixed to  $175\text{ GeV}$  and  $100\text{ GeV}$ , respectively. As expected, it appears that only a small amount of splitting among the doublet components is allowed. For the chosen value of  $m_t$  and  $m_H$ , the SM prediction lies already outside the  $1\sigma$  allowed region and additional positive contributions to  $\epsilon_1$  tend to be disfavoured. On the contrary, the positive contribution to  $\epsilon_3$ , almost constant in the chosen range of masses, is still tolerated, and even preferred by the fit to the data which do not include the SLD result.

The result for a full extra generation of heavy quarks and leptons is shown in fig. 2, dashed line. We have assumed equal ratio  $r$  in the lepton and in the quark sector. As can be seen from eqs. (17-18), this makes the two branches of fig. 1 to degenerate in a unique line.

If new physics beyond the SM were modeled by additional heavy chiral fermions, of the kind we have considered, then, from the effective lagrangian of eqs. (10-13) we could draw informations on the future searches of anomalous trilinear couplings. Indeed, the anomalous magnetic and weak moments of the  $W$ ,  $\Delta k_\gamma$  and  $\Delta k_Z$ , can be expressed as combinations of the coefficients  $a_i$ . One finds [12, 8]<sup>1</sup>:

$$\begin{aligned}\Delta k_\gamma &= g^2(-a_1 + a_2 - a_3 + a_8 - a_9) \\ \Delta k_Z &= \frac{a_0}{(c^2 - s^2)} + \frac{g^2 s^2}{(c^2 - s^2)c^2}(a_1 + a_{13}) \\ &+ g^2 \left[ \frac{s^2}{c^2}(a_1 + a_{13} - a_2) - a_3 + a_8 - a_9 + a_{13} \right]\end{aligned}\quad (20)$$

$s$  and  $c$  denoting the sine and the cosine of the Weinberg angle. The contribution of a quark or lepton doublet to the anomalous moments can be readily evaluated by substituting in eq. (20) the explicit expressions of the coefficients  $a_i$  given in eqs. (10-13). For instance, for the anomalous magnetic moment  $\Delta k_\gamma$ , we obtain:

$$\Delta k_\gamma^q = \frac{Gm_W^2}{4\pi^2\sqrt{2}} \frac{1}{(1-r)^3} [-1 + 8r - 7r^2 + 2(r+2r)\log r] \quad , \quad (21)$$

$$\Delta k_\gamma^l = \frac{Gm_W^2}{12\pi^2\sqrt{2}} \frac{1}{(1-r)^3} [1 + 6r - 9r^2 + 2r^3 + 6r\log r] \quad (22)$$

for quarks and leptons, respectively. Considering in particular the case of degenerate doublets, as suggested by the small value of the  $\epsilon_1$  parameter, one finds:

$$\Delta k_\gamma^q = 3\Delta k_\gamma^l = -\frac{Gm_W^2}{4\pi^2\sqrt{2}} \sim -1.3 \cdot 10^{-3} \quad . \quad (23)$$

Similar contributions are also obtained for  $\Delta k_Z$  so that only very large multiplicities (many doublets) could push the predictions for the anomalous couplings to the level of observability available at LEP II. As a consequence, the contributions to  $\epsilon_3$  would be necessarily positive, large and hard to reconcile with the data.

We are thus lead to consider the possibility of relatively light ( $m, M \geq M_Z/2$ ) chiral fermions, both to check the agreement with the present data, and to test the reliability of our effective lagrangian approach. If the additional fermions are not sufficiently heavy, we do not expect that their one-loop effects are accurately reproduced by the coefficients  $a_i$  in eqs. (11-12). In this case we have to consider the full dependence on external momenta

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<sup>1</sup>The definitions of the anomalous couplings depend on the overall normalization of the trilinear  $WWN$  ( $N = \gamma, Z$ ) vertex, usually denoted by  $g_{WWN}$ . Here we are following the convention of ref. [8].

of the Green functions, not just the first two terms of the  $q^2$  expansion given in eq. (2). We recall that in this case the  $\epsilon$  parameters are given by [13]:

$$\begin{aligned}\delta\epsilon_1 &= e_1 - e_5 \\ \delta\epsilon_2 &= e_2 - s^2 e_4 - c^2 e_5 \\ \delta\epsilon_3 &= e_3 + c^2 e_4 - c^2 e_5\end{aligned}\tag{24}$$

where we have kept into account the fact that in our case there are no vertex or box corrections to four-fermion processes. In eq. (24)

$$\begin{aligned}e_1 &= \frac{A_{33} - A_{WW}}{M_W^2} \\ e_2 &= F_{WW}(M_W^2) - F_{33}(M_Z^2) \\ e_3 &= \frac{c}{s} F_{30}(M_Z^2) \\ e_4 &= F_{\gamma\gamma}(0) - F_{\gamma\gamma}(M_Z^2) \\ e_5 &= M_Z^2 F'_{ZZ}(M_Z^2)\end{aligned}\tag{25}$$

where the quantities  $A_{ij}$  and  $F_{ij}$  are defined in eqs. (1) and (2). The expressions for the quantities  $e_i$ , in the case of an ordinary quark or lepton doublet can be easily derived from the literature [14]. We plot the result for an extra quark doublet in fig. 1 (full line). We have taken  $m_t = 175 \text{ GeV}$  and  $m_H = 100 \text{ GeV}$ . One of the two masses is kept fixed at  $50 \text{ GeV}$ , and the other one runs from  $50 \text{ GeV}$  to  $170 \text{ GeV}$ . The small masses cause a substantial deviation from the asymptotic, effective lagrangian prediction. Nevertheless, as illustrated in fig. 1, such a large effect is still compatible with the present data. The case of a full new generation is shown in the solid line of fig. 2.

In particular, as it was observed in [13], a large negative contribution to both  $\epsilon_1$  and  $\epsilon_3$  is now possible, due to a formal divergence of  $F'_{ZZ}$  at the threshold which produces a large and positive  $e_5$ . Clearly, this behaviour cannot be reproduced by  $\mathcal{L}_{eff}$ , which, at the fourth order in derivatives, automatically sets  $F'_{ZZ} = 0$ . A relevant question is, then, when the asymptotical regime starts, i.e. how close to  $M_Z$  should be the masses of the new quarks or leptons for observing deviations due to the full expression of  $\Pi_{ij}(q^2)$  instead of the truncated expression given in eq. (2). A detailed analysis shows that already for masses of the new fermions above  $70 - 80 \text{ GeV}$  the difference between the values of the  $\epsilon_i$  obtained with the truncated and full expression of  $\Pi_{ij}(q^2)$  are as small as  $10^{-4}$ , i.e. below the present experimental level of accuracy. This is illustrated in fig. 3 where the asymptotical and full expression of  $\epsilon_3$  are compared as a function of  $r$ .

Beyond the indirect precision tests, the possibility of having new fermions carrying the usual  $SU(3)_C \times SU(2)_L \times U(1)_Y$  quantum numbers can clearly also be bounded by the direct searches.



Concerning the present searches, from LEP we have the lower bound of  $M_Z/2$  which applies independently from any assumption on the decay modes of the new fermions which couple to the  $Z$  boson. Much stronger limits on the new quarks masses can be inferred from the Tevatron results. However, as we know from the search for the top quark, these latter bounds rely on assumptions concerning the decay modes of the heavy quark. For instance, in the case of the top search it was stressed that if a new decay channel into the  $b$  quark and a charged Higgs were available to the top, then one could not use the CDF bounds on  $m_t$  [15] which came along these last years, before the final discovery of the top quark.

Now, it may be conceivable that the new physics related to the presence of extra-fermions can also affect their possible decay channels making the lightest of the new fermions unstable. Indeed, we stated in our assumption that the new fermions do not essentially mix with the ordinary ones, hence one has to invoke new physics if one wants to avoid the formation of stable heavy mesons made out of the lightest stable new fermion and of the ordinary fermions of the Standard Model. If the new fermions can decay within the detector, then the bounds on their masses, coming from Tevatron data, must be discussed in a model-dependent way and even the case of new quarks with masses close to  $M_Z/2$  is not fully ruled out.

If on the contrary the lightest new quark is stable, then searches for exotic heavy meson at CDF already ruled out the possibility of being near the threshold  $M_Z/2$ . The existence of coloured particle with charge  $\pm 1$  is strictly bounded over  $130 \text{ GeV}$  from CDF experiment [16]. Finally, note that for charged leptons the bound coming from CDF are much less stringent. A new stable charged lepton of mass of  $50 - 60 \text{ GeV}$  cannot be ruled out.

We are aware of the fact that apart from bounds coming from the direct production of the new fermions, there exist also cosmological limits which apply to the case of stable electrically charged and (or) coloured fermions. Cosmologically stable quarks of masses up to  $20 \text{ TeV}$  can be ruled out on the basis of the results concerning superheavy element searches [17]. As for stable charged hadronic superheavies in the  $20 - 10^5 \text{ TeV}$  range they seem to be in contrast with the bounds which are obtained requiring that heavy particles captured by neutrons star do not induce their collapse to a black hole. Clearly all these severe bounds apply only to the case of cosmologically stable new fermions, so that one can easily avoid them for fermions which have some (even very small) mixing with the ordinary fermions and (or) decay through particles which are related to the new physics beyond the SM.

Finally, concerning future searches, we comment on an effect due to the presence of new heavy quarks which may be potentially relevant for the LHC physics. We consider the production and decay into a photon pair of a Higgs of intermediate mass, i.e.  $100 \leq m_H \leq 150 \text{ GeV}$ . The presence of new quarks give rise to competing effects of opposite sign at the level of production and  $\gamma\gamma$  decay of this intermediate Higgs.

Indeed, as for production, the gluon-fusion amplitude increases due to the effect of the

new quarks which adds up to that of the top [18]. For a new doublet of quarks heavier than the Higgs, the SM gluon-fusion amplitude gets approximately multiplied by a factor 3, the total number of heavy quarks in the loop. This leads to an enhancement of a factor nine, roughly, for the cross section production.

However, in the decay into two photons the new quarks tend to decrease the rate. The dominant contribution to  $H \rightarrow \gamma\gamma$  comes from the loop where the  $W$  boson run, while the top exchange contribution yields an opposite sign. Again the new heavy fermions produce contribution which are analogous to those due to the top exchange and, hence, tend to reduce the decay rate. More precisely, when additional fermions are present, the partial width of the Higgs into a photon pair is given by [19]:

$$\Gamma(H \rightarrow \gamma\gamma) = \frac{\alpha^2}{128\pi^3\sqrt{2}} G_F m_H^3 \left| F_1(\tau) + \sum_i N_i Q_i^2 F_{1/2}(\tau_i) \right|^2 \quad (26)$$

where the sum extends over all fermions of masses  $M_i$ ,  $\tau = 4M_W^2/m_H^2$ ,  $\tau_i = 4M_i^2/m_H^2$ ,  $N_i = 3$  for quarks and  $N_i = 1$  for leptons. In the asymptotic regime, for  $2M_W > m_H$  and  $2M_i > m_H$ , one has:

$$F_1(\tau) = 7, \quad F_{1/2}(\tau_i) = -\frac{4}{3}, \quad (27)$$

When  $m_H \simeq 100 \text{ GeV}$ , from eqs. (26) and (27), one obtains, roughly:

$$\frac{[BR(H \rightarrow \gamma\gamma)]_{new}}{[BR(H \rightarrow \gamma\gamma)]_{SM}} \simeq \frac{1}{3} \quad (28)$$

when only a new doublet of heavy quarks is present, and

$$\frac{[BR(H \rightarrow \gamma\gamma)]_{new}}{[BR(H \rightarrow \gamma\gamma)]_{SM}} \simeq \frac{1}{10} \quad (29)$$

when a new complete generation of heavy quarks and leptons is considered. Notice that the above estimates for the  $BR$ 's are reliable as long as  $m_H \simeq 100 \text{ GeV}$ . For  $m_H$  approaching  $150 \text{ GeV}$  the ratio  $\tau$  gets close to one and, thus the asymptotic expression of  $F_1$  for the  $W$  exchange contribution given in eq. (27) can no longer be used. In this latter case, the  $W$  contribution can become sizeably larger than in the asymptotic situation, hence making the suppression of  $BR(H \rightarrow \gamma\gamma)$  due to the new fermions less severe. In the less favorable case of a full new heavy generation, the product of the production cross-section times the  $BR(H \rightarrow \gamma\gamma)$  reflects a conspicuous dependence on  $m_H$ . For  $m_H \simeq 100 \text{ GeV}$  we obtain:

$$\frac{[\sigma(gg \rightarrow H) \cdot BR(H \rightarrow \gamma\gamma)]_{new}}{[\sigma(gg \rightarrow H) \cdot BR(H \rightarrow \gamma\gamma)]_{SM}} \simeq 0.9 \quad (30)$$

while for  $m_H = 150 \text{ GeV}$  the above ratio becomes larger than one, reaching values close to 4.

In conclusion, we discussed the impact of the presence of new sequential fermions on the electroweak precision tests. We showed that the present data still allowed the presence

of a new quark and/or lepton doublet with masses greater than  $M_Z/2$ . Only for light new fermions which are close to the threshold  $M_Z/2$  one finds drastic departures of the effective lagrangian result from the full one-loop radiative corrections obtained in the SM. The presence of new fermions carrying usual  $SU(3)_C \times SU(2)_L \times U(1)_Y$  quantum numbers with mass as low as  $60 - 80 \text{ GeV}$  is severely limited both by accelerator results and cosmological constraints. Finally, the new fermions can significantly affect the production and decay rate into  $\gamma\gamma$  of the intermediate Higgs at LHC.

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## Figure Captions

Fig. 1 Predictions for  $\epsilon_1, \epsilon_3$  from an additional quark doublet. The lower (upper) dashed line represents the case  $m (M) = 200 \text{ GeV}$ ,  $M (m)$  varying between  $200 \text{ GeV}$  and  $300 \text{ GeV}$ , evaluated with  $\mathcal{L}_{eff}$ . The lower (upper) full line corresponds to  $m (M) = 50 \text{ GeV}$ ,  $M (m)$  varying between  $50 \text{ GeV}$  and  $170 \text{ GeV}$ , evaluated with a complete 1-loop computation. The SM point corresponds to  $m_t = 175 \text{ GeV}$  and  $m_H = 100 \text{ GeV}$ . The upper (lower) ellipses is the  $1 \sigma$  allowed region, obtained by a fit of the high energy data which excludes (includes) the SLD measurement.

Fig. 2 Same as for Fig. 1, in the case of a full extra generation of quarks and leptons.

Fig. 3 Comparison between the asymptotic (solid line) and full one-loop (dashed lines) computations of  $\epsilon_3$  versus  $r = m^2/M^2$ , for an additional quark doublet. For the SM contribution,  $m_t = 175 \text{ GeV}$  and  $m_H = 100 \text{ GeV}$  are assumed. The  $2 \sigma$  allowed region is also displayed.

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