Search for a light fermiophobic Higgs boson produced via gluon fusion at Hadron Colliders

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In this study, we propose new Higgs production mechanisms with multi-photon final states in the fermiophobic limit of the Two Higgs Doublet Model. The processes are: $gg \rightarrow hh$, $gg \rightarrow Hh$ followed by $H \rightarrow hh$ and $gg \rightarrow Ah$ followed by $A \rightarrow hZ$. In the fermiophobic limit, $gg \rightarrow hh$ and $gg \rightarrow Ah \rightarrow hhZ$ would give rise to $4\gamma$ signature while $gg \rightarrow Hh \rightarrow hh$ can give a $6\gamma$ final state. We show that both the Fermilab Tevatron and CERN’s Large Hadron Collider can probe a substantial slice of the parameter space in this fermiophobic scenario of the Two Higgs Doublet Model. If observed the above processes can give some information on the triple Higgs couplings involved.

I. INTRODUCTION

There are several extensions of the Standard Model (SM) with an enlarged scalar sector. Some of these extensions allow for Higgs with reduced or even no couplings to the fermions. They are referred to as fermiophobic Higgs scenarios in the literature [1]. The D0 collaboration has recently presented new results on fermiophobic Higgs searches [2, 3]. In [2] they have searched for a fermiophobic Higgs produced in association with a charged Higgs. The full process is $p\bar{p} \rightarrow hH^\pm \rightarrow hhW^{\pm*} \rightarrow 4\gamma + X$ and was proposed in [4, 5, 6]. D0 required at least three photons in the final state for maximizing signal efficiency. For each pair of values of ($\tan \beta$, $m_{H^\pm}$) a bound on the fermiophobic Higgs mass was set. In [3], D0 has performed a search for the inclusive production of di-photon final states via the Higgsstrahlung and vector boson fusion processes: $p\bar{p} \rightarrow hV \rightarrow \gamma\gamma + X$ and $p\bar{p} \rightarrow VV \rightarrow h + \gamma\gamma + X$, respectively, with a total integrated luminosity of 1.10 ± 0.07 fb$^{-1}$. A lower $m_h$ bound of 100 GeV was obtained in a benchmark scenario that assumes $hVV$ ($V=W, Z$) couplings to be exactly the same as in the SM and all fermion branching ratios to be exactly zero.

All LEP collaborations have searched for a fermiophobic Higgs in the channel $e^+e^- \rightarrow h(\rightarrow \gamma\gamma)Z$ [7, 8, 9, 10]. The combination of all results [11] yielded the lower bound for the fermiophobic Higgs mass of 109.7 GeV, at 95% confidence level, which again is valid only in the above benchmark scenario. Searches in the $e^+e^- \rightarrow hA$ channel were also performed at LEP with lower bounds derived for $m_h + m_A$ (see references [7, 8] for details). The new D0 bound [3] on $m_h$ is weaker but a larger region of the model’s parameter space is covered. The channel $qq' \rightarrow V^* \rightarrow h(\rightarrow \gamma\gamma)V$ had already been used at the Tevatron by the CDF [12] and by the D0 [13] collaborations to set limits of 78.5 GeV and 82 GeV at 95% confidence level, respectively, on the fermiophobic Higgs mass. It is expected [14, 15, 16] that all Tevatron bounds will improve once the data collected at 2 fb$^{-1}$ luminosity is analyzed.

There are however other ways of producing a four photon final state in a fermiophobic scenario. In this letter we will consider all fermiophobic Higgs production processes with at least three photons in the final state produced via gluon fusion. As shown in [2], a signal with at least three photons is very easy to extract at the Tevatron. The most relevant process for the analysis is $gg \rightarrow hh \rightarrow 4\gamma$ but we can also have $gg \rightarrow hH \rightarrow hhh \rightarrow 6\gamma$ and $gg \rightarrow hA \rightarrow hhZ \rightarrow 4\gamma + X$. We will show that both the Tevatron and the Large Hadron Collider (LHC) can probe a substantial region of the parameter space. We will discuss...
as well the complementarity between the different production modes. This paper is structured as follows: we will review the fermiophobic model in section II and then proceed to look at the all available theoretical and experimental bounds in section III. In section IV, we will discuss in detail the production process and in section V the signal. Analysis of the results and conclusions will be presented in section VI.

II. THE FERMIOPHOBIC THDM

To define our notation we start with a brief review of the two-Higgs doublet potential used here. The potential chosen is the most general, renormalizable, CP-conserving potential, invariant under \( SU(2) \otimes U(1) \) that one can build with two complex Higgs doublets. It can be written as

\[
V(\Phi_1, \Phi_2) = \lambda_1(|\Phi_1|^2 - v_1^2)^2 + \lambda_2(|\Phi_2|^2 - v_2^2)^2 + \lambda_3((|\Phi_1|^2 - v_1^2) + (|\Phi_2|^2 - v_2^2))^2 + \\
\lambda_4(|\Phi_1|^2|\Phi_2|^2 - |\Phi_1^*\Phi_2|^2) + \lambda_5|\text{Re}(\Phi_1^*\Phi_2) - v_1v_2|^2 + \lambda_6|3m(\Phi_1^*\Phi_2)|^2
\]  

(1)

where \( \Phi_1 \) and \( \Phi_2 \) have weak hypercharge \( Y = 1 \), \( v_1 \) and \( v_2 \) are respectively the vacuum expectation values of \( \Phi_1 \) and \( \Phi_2 \) and the \( \lambda_i \) are real-valued parameters. Note that this potential violates the discrete symmetry \( \Phi_i \rightarrow -\Phi_i \) only softly by the dimension two term \( \lambda_5 \text{Re}(\Phi_1^*\Phi_2) \). The hard breaking terms (dimension four) of the discrete symmetry have been set to zero. As in all other THDM, we end up with two CP-even Higgs states usually denoted by \( h \) and \( H \), one CP-odd state, \( A \) and two charged Higgs bosons, \( H^\pm \). The potential in eq. (1) has 8 parameters (including \( v_1 \) and \( v_2 \)). The combination \( v^2 = v_1^2 + v_2^2 \) is fixed as usual by the electroweak breaking scale through \( v^2 = (2\sqrt{2}G_F)^{-1} \). We are thus left with 7 independent parameters; namely \( (\lambda_i)_{i=1,...,6} \), and \( \tan\beta = v_2/v_1 \). Equivalently, we can take instead

\[
m_h , \ m_H , \ m_A , \ m_{H^\pm} , \ \tan\beta , \ \alpha \text{ and } \lambda_5.
\]  

(2)

as the 7 independent parameters. The angle \( \beta \) is the rotation angle from the group eigenstates to mass eigenstates in the CP-odd and charged sector. The angle \( \alpha \) is the corresponding rotation angle for the CP-even sector. In a general THDM it is possible to couple just one doublet to all fermions by choosing an appropriate symmetry for both the fermions and the scalars. This model is known as THDM type I in the literature. Like in the SM, where just one doublet couples to all fermions, each scalar couples to the different fermions with the same coupling constant. However, unlike the SM, the couplings are now proportional to the rotation angles \( \alpha \) and \( \beta \). For instance, the lightest CP-even Higgs couples to the fermions as \( \cos\alpha/\sin\beta \, g^\text{SM}_{hff} \). By choosing \( \cos\alpha = 0 \), the lightest CP-even Higgs decouples from all fermions. It is usually referred to as a fermiophobic Higgs scalar [1]. This way the heavy CP-even scalar will acquire larger couplings to the fermions than the corresponding SM couplings. The remaining scalars are not affected by this choice as they do not couple proportionally to \( \alpha \). However, \( h \) can still decay to two fermion pairs via \( h \rightarrow W^+W(Z^*Z) \rightarrow 2f\bar{f} \) or \( h \rightarrow W^*W^*(Z^*Z^*) \rightarrow 2f\bar{f} \). We will include these decays in our analysis. It is worth pointing out that these processes occur near the \( W(Z) \) threshold. Decays of \( h \) to two fermions can also be induced by scalar and gauge boson loops (see e.g. fig. [1]). In the THDM, the angle \( \alpha \) has to be renormalized to render \( h \rightarrow f\bar{f} \) finite. However, at \( \alpha = \pi/2 \), all one-loop decays \( h \rightarrow f\bar{f} \) are finite. Thus we can impose the following condition for the renormalization constant \( \delta\alpha \): the renormalized one-loop decay width for \( h \rightarrow f\bar{f} \) is equal to the finite unrenormalized decay width. This condition is equivalent to setting \( [\delta\alpha]_{\alpha=\pi/2} = 0 \). In [17] we have checked that this condition holds for all fermions. The only relevant one-loop decay is \( h \rightarrow b\bar{b} \) due to a large contribution of the Feynman diagram.

![FIG. 1: Feynman diagram of the largest contribution to $h^0 \rightarrow b\bar{b}$](image-url)
shown in fig. 10 to the total decay width \[^{28}\]. Thus, on one hand, \(h\) is not completely fermiophobic at \(\alpha = \pi/2\), and on the other hand, all decays \(h \to f \bar{f}\) but \(h \to b \bar{b}\) are almost zero even at one-loop level. Nevertheless it is possible, although hard, to find regions of the parameter space where \(h \to b \bar{b}\) has a sizeable effect in the fermiophobic Higgs signature. With this in mind it is obvious that we are working with a version of the model with one less parameter than the general CP-conserving THDM.

### III. THEORETICAL AND EXPERIMENTAL BOUNDS

In our analysis we take into account the following experimental and theoretical constraints. Note however that we will use the theoretical constraints as guides and explain in each case how they would affect the results given in the plots. We believe that if Nature did not provide a fermiophobic Higgs it should be disallowed ultimately by experiment.

- Experimental bounds from LEP; the LEP collaborations have set bounds on the mass of a fermiophobic Higgs as a function of \(\tan \beta\). The most stringent bound comes from the combination of all LEP experiments given in \[^{11}\]. The 95% CL limit on \(\tan \beta\) is about 18 for \(m_h = 20\) GeV. From \(m_h = 20\) GeV until \(m_h = 70\) GeV, the bound oscillates about \(\tan \beta = 10\). For \(m_h > 70\) GeV, \(\tan \beta > 10\) is already a conservative bound. In the plots \(\tan \beta = 10\) is used because much higher values of \(\tan \beta\) would violate perturbativity constraints. Note however that as we will show, \(gg \to hh\) is independent of \(\tan \beta\) and for \(\tan \beta > 10\) the dependence of all other relevant processes is negligible.

There is another bound on the fermiophobic Higgs coming from \(hA\) production at LEP which constrains severely the value of \(\tan \beta\) especially for small \(h\) masses. If however we take the \(A\) mass to be above 150 GeV, the \(hA\) production is no longer a constraint for all values of \(\tan \beta\).

Finally, if \(\tan \beta\) is large, the \(ZZh\) coupling is suppressed while the non-fermiophobic CP-even Higgs \(H\) will couple to the \(Z\) bosons with almost the SM strength. Therefore, in the plots shown, the minimum value of the \(H\) mass is 100 GeV and most cases presented are for an \(H\) mass above 120 GeV.

- As already stated in the introduction, the D0 and CDF experiments recently reported searches for a fermiophobic Higgs in two different channels. D0 searched \[^{3}\] for the inclusive production of di-photon final states via the Higgsstrahlung and vector boson fusion. The bound with the model benchmark described in the introduction is weaker than the LEP bound but it spans a larger region of the parameter space. The search \[^{2}\] in the \(p \bar{p} \to hH^{\pm} \to hhW^{\pm} \to 4\gamma + X\) channel sets a bound on the fermiophobic Higgs mass for each pair of values of \((\tan \beta, m_{H^{\pm}})\). We will take all these bounds into consideration in our analysis.

- The extra contributions to the \(\delta \rho\) parameter from the Higgs scalars \[^{18}\] should not exceed the current limits from precision measurements \[^{19}\]: \(|\delta \rho| \lesssim 10^{-3}\). Such an extra contribution to \(\delta \rho\) vanishes in the limit \(m_{H^{\pm}} = m_A\). To ensure that \(\delta \rho\) will be within the allowed range whenever possible we allow only a small splitting between \(m_{H^{\pm}}\) and \(m_A\).

- Recently, it has been shown in Ref. \[^{20}\] that for THDM models of the type II, data on \(B \to Xs\gamma\) imposes a lower limit of \(m_{H^{\pm}} \gtrsim 290\) GeV. In THDM type I, there is no such constraint on the charged Higgs mass. Therefore, in our numerical analysis which is valid for THDM type I we will ignore the limit on the charged Higgs.

- The scalar sector can also be constrained using perturbativity constraints on \(\lambda_i\) \[^{21, 22}\]. In the present study we will not impose those constraints in order to quantify the optimal cross sections and scan over all parameters space. In fact, in the fermiophobic limit, the process \(gg \to hh\) for example depends only on \(m_b, m_H\) and \(\lambda_5\). As we will explain later, the \(\tan \beta\) dependence in \(gg \to hh\) drops out. As a result, for a given \(m_b, m_H\) and \(\lambda_5\), one can tune \(\tan \beta, m_A\) and \(m_{H^{\pm}}\) in order to satisfy perturbativity constraints.

- From the requirement of perturbativity for the top and bottom Yukawa couplings \[^{23}\], \(\tan \beta\) is constrained to lie in the range \(0.3 \leq \tan \beta \leq 100\). But it turns out that from perturbativity argument on \(\lambda_i\), moderate values of \(\tan \beta\) less than about 10 are preferred.
In order to respect perturbativity constraints we will use moderate values for $M_h$, $M_H$, $M_A$ and $M_{H\pm}$.

IV. PRODUCTION CROSS SECTIONS

The process $pp(\bar{p}) \rightarrow hh$ has both tree level contributions mediated by Higgs exchange from $q\bar{q} \rightarrow H^*, h^* \rightarrow hh$ and one loop contributions from gluon fusion $gg \rightarrow hh$. The tree level contribution is proportional to the quark masses which will be neglected for the Tevatron energies. We have also checked that even in the large tan $\beta$ limit the production process $q\bar{q} \rightarrow hh$ is also negligible for the LHC when compared to $gg \rightarrow hh$. The process $gg \rightarrow hh$ occurs only at the one-loop level. As we will see, even if loop suppressed, this process can still be enhanced by the strong QCD coupling as well as by the heavy Higgs $H$ resonant effect when it can decay to two light CP even $h$ scalars. There are two types of diagrams that participate in the process $gg \rightarrow hh$. The box diagram, as shown in fig. 2 represents just a generic contribution to the process. We have also included all other quarks and because both the initial and the final state have identical particles, there is a total of six diagrams for each flavor. The second type are the vertex diagrams which again are six for each flavor. In fig. 3 we show just the representative diagrams with a generic fermion in the loop. In the fermiophobic limit, the top loop is always the dominant contribution. The other two processes, $pp(\bar{p}) \rightarrow Hh$ and $pp(\bar{p}) \rightarrow Ah$ have a similar structure (vertex and box) to $pp \rightarrow hh$, except for one additional contribution to the vertex diagrams in $pp(\bar{p}) \rightarrow Ah$ which is the $Z$ boson and the Goldstone boson s-channel exchange. However, in the fermiophobic limit, all box contributions vanish and the same is true for all vertex contribution with s-channel Higgs ($h$) exchange due to the fact that the fermiophobic Higgs coupling to fermions is zero. Since the box contributions drop out in the fermiophobic limit, the two processes $pp(\bar{p}) \rightarrow hh, Hh$ are directly proportional to the pure scalar couplings $Hhh$ and $HHz$. The third process $pp(\bar{p}) \rightarrow Ah$ is sensitive both to pure scalar couplings $AAh$ and $AGh$ as well as to the gauge coupling
ZAh. Hereafter we list the pure scalar coupling in the fermiophobic limit (FL) needed for our study:

\[ \lambda_{Hhh}^{FL} = \frac{e\lambda_5 v^2 \sin \beta}{4m_W s_W} \propto \lambda_5 \frac{\tan \beta}{\sqrt{1 + \tan^2 \beta}} \]

\[ \lambda_{Hh}^{FL} = \frac{e\lambda_5 v^2 \cos \beta}{4m_W s_W} \propto \lambda_5 \frac{1}{\sqrt{1 + \tan^2 \beta}} \]

\[ \lambda_{Ahh}^{FL} = \frac{-e}{4m_W s_W \cos \beta} \left[ -2 \sin^2 \beta m_h^2 + \lambda_5 v^2 - 4 \cos^2 \beta m_A^2 \right] \]

\[ \lambda_{Agh}^{FL} = \frac{e \sin \beta}{2m_W s_W} \left( m_A^2 - m_h^2 \right) \]

\[ \lambda_{H'h'h'}^{FL} = \frac{e}{2m_W s_W \sin 2\beta} \left[ 2 \sin^2 \beta m_h^2 - \lambda_5 v^2 \sin \beta - 2 \sin 2\beta \cos \beta m_{H^+}^2 \right] \]

Note that the \( Hhh \) and \( H\bar{H}h \) couplings are directly proportional to \( \lambda_5 \) while the \( AAh \) coupling depend both on \( \lambda_5, m_A \) as well as on \( m_h \). It is clear that in the case of exact discrete \( Z_2 \) symmetry (\( \lambda_5 = 0 \)), both \( Hhh \) and \( H\bar{H}h \) would vanish. In the fermiophobic limit, all fermionic couple to each scalar with the same strength. The \( f f H \) coupling is proportional to \( 1/\sin \beta \). Therefore from eq. 3 we conclude that the \( \beta \) angle dependence cancels out in the cross section, that is, \( \sigma_{gg\rightarrow hh} \) depends only on \( m_h, m_H \) and \( \lambda_5 \). The \( f f A \) coupling is proportional to \( \cos \beta/\sin \beta \) which means that in the large \( \tan \beta \) limit the \( \beta \) angle dependence is also very mild for very large \( h \) and/or \( A \) masses. Those cases will not be included in our study. Finally, note that not only \( \sigma_{gg\rightarrow Ah} \) vanishes in the limit \( \lambda_5 = 0 \) but also \( \sigma_{gg\rightarrow Ah} \), in the high \( \tan \beta \) limit becomes negligible when \( \lambda_5 = 0 \) due to to the smallness of the values of the masses involved.

The one-loop amplitudes were generated and calculated with the packages FeynArts \cite{24} and FormCalc \cite{25}. The scalar integrals were evaluated with LoopTools \cite{26}.

1. **Numerical results**

In this section, we present our numerical results. As stated earlier, to avoid the LEP bounds on the fermiophobic Higgs we will fix \( \tan \beta \) to be of the order 10 for the processes \( gg \rightarrow Ah \) and \( gg \rightarrow Hh \). This way we suppress the \( ZZh \) coupling while keeping perturbativity bounds on \( \lambda_i \) within the allowed range. The first consequence of this choice of \( \tan \beta \approx 10 \) is that the coupling \( Hhh \) is enhanced (\( \sin \beta \approx 1 \)) while \( H\bar{H}h \) is suppressed (see eqs. 3, 4). As we have discussed, in the fermiophobic limit, the process \( pp \rightarrow hh \) has only vertex contribution through s-channel heavy Higgs (\( H \)) exchange. Therefore, the cross section for \( pp(\bar{p}) \rightarrow hh \) depend both on \( m_H \) as well as on \( Hhh \) coupling which is proportional to \( \lambda_5 \). There are two sources of enhancements for \( pp \rightarrow hh \). The first one is to take \( Hhh \) (or equivalently \( \lambda_5 \)) as large as possible. The second one is when \( hh \) production is resonant, that is, \( m_H \approx 2m_h \).

We first discuss \( pp(\bar{p}) \rightarrow hh \) production. In Fig. 4(left) we illustrate the cross section of \( pp(\bar{p}) \rightarrow hh \) as a function of the fermiophobic Higgs mass \( m_h \) for two representative values of \( \lambda_5 = 4\pi \) and \( 8\pi \). The other heavy CP even mass is taken to be \( m_H = 2m_h \) such that the resonant channel \( H \rightarrow hh \) is open. The other parameters are: \( m_A = m_{H_+} = 300 \) GeV. On the \( H \) resonance, the cross section is enhanced and can reach a few hundreds of picobarn for a very light fermiophobic Higgs \( m_h = 50 \) GeV. Even for \( m_h \) of the order of 100 GeV, the cross section is still larger than \( 0.1 \) pb which would give thousands of produced events for the planned Tevatron luminosity of \( 10 fb^{-1} \). As can be seen from the analytic expression of the coupling \( Hhh \), the large \( \lambda_5 \) is the large is the coupling \( Hhh \) and so is the \( pp \rightarrow hh \) cross section. The observed kinks at \( m_H = 160 \) and 182 GeV in the left plot are due to the opening of the decay channels \( H \rightarrow WW \) and \( H \rightarrow ZZ \). On the right-hand side of Fig. 4 we show the same cross section, the only difference being the \( H \) mass which is now fixed at 120 GeV. It is clear that even away from the \( H \) resonance a significant set of fermiophobic Higgs masses and \( \lambda_5 \) values can still be probed at the Tevatron. The behavior with \( \lambda_5 \) changes at threshold due to the \( H \) width effect on the cross section. Above threshold the \( H \rightarrow hh \) channel is closed, the \( H \) width is then very small. Therefore in that region the cross section is just proportional to \( \lambda_5^2 \). Below threshold, the \( H \rightarrow hh \) channel is open and the width in the \( H \) propagator starts to play a role which makes the dependence with \( \lambda_5 \) no longer trivial.

In Fig. 3 we show the same plots as in Fig. 4 but for the LHC. As expected the plots are re-scaled by more than one order of magnitude. For the values shown, all masses between 10 and 100 GeV can be probed for most values of \( \lambda_5 \). Even for \( \lambda_5 = \pi/2 \) the smallest cross section value is of the order of 10 fb.
FIG. 4: $\sigma(gg \to hh)$ for the Tevatron in units of pb as a function of the fermiophobic Higgs mass with $\lambda_5 = 4\pi$ and $8\pi$. On the left the mass of the heavier CP-even Higgs boson is twice the fermiophobic Higgs mass, i.e., on threshold, and on the right it is fixed (120 GeV).

FIG. 5: $\sigma(gg \to hh)$ for the LHC in units of pb as a function of the fermiophobic Higgs mass with $\lambda_5 = \pi/2$, $4\pi$ and $8\pi$. On the left the mass of the heavier CP-even Higgs boson is twice the fermiophobic Higgs mass, i.e., on the threshold and on the right it is fixed (120 GeV).

In Fig. 6 we show the cross section of $pp \to hh$ as a function of the resonant Higgs mass $m_H$ for $\lambda_5 = 8\pi$ and for a fermiophobic Higgs $m_h = 60$ GeV, for the Tevatron (left) and for the LHC (right). As one can see, away from the resonance $m_H \approx 2m_h$, the cross section is of the order a few fb for the Tevatron and a few pb for the LHC. Once we cross the resonance, one can see a spectacular pick for $m_H = 120$ GeV which is due to the opening of the decay channel $H \to hh$. The peak is very sharp because the total width of $H$ is very narrow for $m_H \lesssim 2m_h$, less than $10^{-2}$ GeV. Once the decay channel $H \to hh$ is open for $m_H \gtrsim 2m_h$ the decay width of $H$ increase suddenly to more than 100 GeV. This is manifestly seen in the plot by a dramatic decrease of the cross section from few hundred pb to $10^{-2}$ pb.

Finally we present the $pp(\bar{p}) \to Hh$ and $pp(\bar{p}) \to Ah$ reactions. It is clear from both plots presented in Fig. 7 that the cross section for $Hh$ production is negligible for most of the parameter space. On the contrary, the cross section for $hA$ production can be very large and still within the Tevatron reach. The cross section $gg \to Ah$ can be several orders of magnitude larger than the corresponding $gg \to Z^* \to Ah$ (see ref. [5]). We will show in the next section that taking into account the pseudo-scalar branching ratios, the decay $A \to Zh$ can be the dominant one, and so the channel $gg \to Ah \to Zhh \to Z\gamma\gamma$ is still
FIG. 6: $\sigma(gg \to hh)$ for the Tevatron (left) and for the LHC (right) in units of $pb$ as a function of the heavier CP-even Higgs mass with $\lambda_5 = 8\pi$ and a fermiophobic Higgs mass of 60 GeV.

FIG. 7: $\sigma(gg \to hA)$ and $\sigma(gg \to hH)$ for the Tevatron (left) and for the LHC (right) in units of $pb$ as a function of the fermiophobic Higgs mass with $\tan\beta = 10$, $\lambda_5 = 4\pi, 8\pi$, $M_H = 150$ GeV and $M_A = 150$ GeV.

V. HIGGS SIGNATURE

Having established that $gg \to hh$ and $gg \to Ah$ are worth studying both at the Tevatron and at the LHC we now turn to the experimental signatures for the fermiophobic Higgs and for the pseudo-scalar in the parameter space under study. In the fermiophobic limit there is a dramatic change in the fermiophobic Higgs signatures. For smaller fermiophobic Higgs masses, the main decay is to two photons (through W and charged Higgs loops) until the WW channel starts to dominate. The crossing point of the branching ratios depends on the parameters of the scalar potential which enters the game through the charged Higgs contribution to $h \to \gamma\gamma$. Those parameters are mainly the charged Higgs mass, $\lambda_5$, $\tan\beta$ and the fermiophobic Higgs mass (see eq. (7)). Recently a detailed study of $h \to \gamma\gamma$ appeared in [27]. Due to all experimental and theoretical constraints, for a fermiophobic Higgs with mass between 10 and 100 GeV, $Br(h \to \gamma\gamma) \approx 100\%$ except in a tiny neighborhood of $\lambda_5 = 0$. In this neighborhood, all cross...
sections are extremely small. Therefore there will always be a tiny region around \( \lambda_5 = 0 \) that will not be probed with the processes proposed here. For illustration, we show in fig. 8 (left) the fermiophobic Higgs branching ratio as a function of \( \lambda_5 \). We have checked that the larger \( \tan \beta \) is the higher the value of the \( \gamma \gamma \) branching ratio. \( \text{Br}(h \rightarrow \gamma \gamma) \) decreases with the fermiophobic Higgs mass, but for \( m_h = 100 \text{ GeV} \) the plot looks almost the same. Regarding the pseudo-scalar decays, as one can see in fig. 8 (right), one can have \( \text{Br}(A \rightarrow hZ) \approx 100\% \) if the decay \( A \rightarrow H^\pm W^\mp \) is kinematically forbidden. We have checked that changes in \( \tan \beta, \lambda_5 \) and the \( A \) mass produce negligible changes in the branching ratio provided the charged Higgs channel \( A \rightarrow H^\pm W^\mp \) is kept closed. Therefore, the 4 \( \gamma \) final state is by far the dominant one.

### VI. DISCUSSION AND CONCLUSIONS

In this work we have shown that there are alternative channels to search for fermiophobic Higgs with a multi-photon signature. A vast region of the parameter space of the fermiophobic THDM can be probed at the Tevatron and the analysis can easily be extended for the LHC. In the fermiophobic limit, the angle \( \beta \) is already very constrained. LEP [11] has set a limit of \( \tan \beta > 10 \) for almost all values of the fermiophobic Higgs masses up to 100 GeV. For some masses the bound is even stronger. On the other hand theoretical constraints tell us that these values cannot be too high. This implies that \( gg \rightarrow hh \) and \( gg \rightarrow hA \) will be large while \( gg \rightarrow hH \) will be negligible. Note however that the \( gg \rightarrow hh \) cross section does not depend on \( \tan \beta \) and that the \( gg \rightarrow hA \) dependence on \( \tan \beta \) for values above 10 is negligible. Regarding the \( \lambda_5 \) dependence, we have shown that there is a "tiny to small" region around \( \lambda_5 = 0 \) that can not be probed. This is especially true for \( hh \) productions but \( gg \rightarrow Ah \) decreases with \( \lambda_5 \) as well.

The process \( p\bar{p} \rightarrow hH^\pm \rightarrow hhW^{\pm*} \rightarrow 4\gamma + X \) proposed in [4,5] and studied in [2] is complementary to the processes we propose in this work. We probe the region of small \( H \) and/or \( A \) masses while [2] probes the small charged Higgs mass region. The advantage in our case is that our study is independent of \( \tan \beta \) while their process does not depend on \( \lambda_5 \). Finally if all Higgs scalars besides the fermiophobic Higgs are very heavy, only the two photon search can exclude a fermiophobic Higgs.

As expected and as it was shown in [2] the background for a 3 or 4 photon final state is easy to control. In the case of the 3 photon final state the main background contribution comes from the direct tri-photon production (see [2] for details).

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The coupling $[H^+ b]_i$ is proportional to the $t$-quark mass.