CP-odd scalar with vector-like fermions at the LHC

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Abstract

Many theories beyond the standard model (BSM) contain CP-odd scalars (A) and new vector-like fermions (ψVL). The couplings of the A to two standard model gauge bosons (i.e. AVV couplings) cannot occur from renormalizable operators in a CP-conserving sector, but can be induced at the quantum loop level. We compute the AVV effective couplings at the 1-loop level induced by the SM fermions and vector-like fermions, present analytical expressions for them, and plot them numerically. Using the 8 TeV Large Hadron Collider (LHC) γγ, τ+τ− and t ¯t channel data, we derive constraints on the effective couplings of the A to standard model (SM) gauge bosons and fermions, present the gluon-fusion channel cross-sections at the 8 and 14 TeV LHC, and present the branching-ratio of the A into SM fermion and gauge-boson pairs. We present our results first model-independently, and then also for some simple models containing A and ψVL in the singlet and doublet representations of SU(2). In the doublet case, we focus on the two-Higgs-doublet (2HDM) Type-II and Type-X models.

1 Introduction

A long series of experiments culminating in the Large Hadron Collider (LHC) discovery of the Higgs boson at a mass of about 125 GeV have firmly established the standard model (SM) as the correct description of Nature up to an energy scale of a few hundred GeV. With this discovery, the theoretical puzzle as to why the Higgs boson remains this light when quantum effects should correct it to the highest scales present in the theory (such as the Planck scale) comes to the fore. This problem of the stability of the electroweak (EW) scale is the well known hierarchy problem of the SM. This could be a clue that some new physics beyond the standard model (BSM) is present near the EW scale which renders it stable against quantum corrections, making it natural. Many theoretical proposals have been made for this new physics (for reviews see Ref. [1]), and they usually contain new particles at the TeV energy scale. We are poised at a very interesting time when the LHC is probing this energy scale and can tell us if one of these proposals is realized in Nature.

Among the possibilities of BSM physics that makes the EW scale natural are models in which the Higgs-doublet of the SM is a pseudo-Nambu-Goldstone boson (PNGB). Concrete realizations of this idea, for example, are in models of little-Higgs, composite-Higgs and warped extra dimensions (for reviews see Refs. [2, 3, 4] respectively). In such models, in addition to the CP-even Higgs boson, there could be CP-odd scalars (denoted as A) that are also PNGBs, and therefore as “light”. Also, new heavy vector-like fermions (VL, denoted as ψVL) are usually required, that along with the SM fermions, complete some representation of a bigger group containing SU(2)⊗U(1). By vector-like fermions we mean that fermions in a representation of the SM gauge-group and in its conjugate representation both appear in the theory (for more details see for example Ref. [5]). Some supersymmetric models also include vector-like matter, and thus have A and ψVL both present, along with many superpartners. In this work, our goal is to capture the phenomenology of CP-odd scalars in such BSM theories, allowing A to couple to vector-like fermions and to SM fermions.

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The phenomenology of a CP-odd scalar at the LHC can be quite distinct as compared to a CP-even scalar (such as the SM Higgs boson). \(AW^+W^-\), \(AZZ\), \(A\gamma\gamma\) and \(AZ\gamma\) couplings (collectively called \(AVV\) couplings) cannot arise from renormalizable operators since we require that CP-invariance be not spontaneously broken by an A vacuum expectation value (VEV), i.e. we require \(\langle A\rangle = 0\). Also, the last two do not arize from renormalizable operators because of unbroken electromagnetic (EM) gauge invariance, the same reason why \(h\gamma\gamma\) and \(h\gamma Z\) are zero at the renormalizable level. These can then only result from higher-dimensional operators generated at loop-level. In contrast, for the CP-even SM Higgs boson (denoted as \(h\)), the \(hW^+W^-\) and \(hZZ\) couplings are generated at tree-level from dimension-four operators after electroweak symmetry breaking (EWSB), i.e. with \(\langle h\rangle = v/\sqrt{2}\). Therefore, generically speaking, the \(AW^+W^-\) and \(AZZ\) effective couplings, generated at loop-level, are much smaller in magnitude compared to the tree-level \(hW^+W^-\) and \(hZZ\) couplings; the \(A\gamma\gamma\) and \(h\gamma\gamma\) effective couplings are both loop suppressed and small, and similarly the \(A\gamma Z\) and \(h\gamma Z\) are also both loop suppressed. Thus, similar to the \(h\), the \(gg \rightarrow A\) "gluon-fusion" channel is important at the LHC, while compared to the \(h\), the vector-boson fusion channel of \(A\) is much suppressed. Turning next to \(A\) couplings to fermions, we include \(A\) couplings to new vector-like fermions at the tree-level. Furthermore, if \(A\) is part of a doublet, it couples also to SM fermions at the tree-level (similar to \(h\)). We consider the situation when \(A\) couples significantly only to third generation SM fermions, a situation common in many BSM extensions. Thus, the relevant couplings to SM fermions are \(Abb\), \(A\tau^+\tau^-\) and \(Att\). If the \(Abb\) coupling is sizable, \(bb \rightarrow A\), \(bg \rightarrow hA\) and \(gg \rightarrow bbA\) can be important production channels of the \(A\). However, we do not include these production channels in this work, but restrict ourselves only to the gluon-fusion channel.

We restrict ourselves to the situation when \(m_A < 2MVL\) so that \(A\) cannot decay to a pair of VLFs. If the \(\psi_{VL}\) is light enough they can also be studied directly at the LHC as discussed for instance in Ref. [5] and references therein. However, if they are too heavy to be directly produced at the LHC, but the \(A\) (or \(h\) as studied in Ref. [6]) can be directly produced and its coupling measured, the VLF contributions to the \(A\) effective couplings we derive here can be useful in probing the \(\psi_{VL}\) indirectly.

In order to capture many different BSM models, we perform a model-independent effective theory analysis of the \(A\) coupled to SM fields. As a function of the \(A\) effective-couplings (denoted \(\kappa\)), we present the constraints from the recent 8 TeV LHC run using the \(\gamma\gamma, \tau^+\tau^-\) and \(tt\) channels, and present the signal cross section (c.s., \(\sigma\)) at the LHC and branching ratio (\(BR\)) into the \(\gamma\gamma, \gamma Z, ZZ, W^+W^-\) and fermion final states. We do not focus much on the \(ZZ\) and \(W^+W^-\) decay channels as the branching-ratio into these modes are much smaller than the other modes due to phase-space suppression. These are some of the main results of this work. We also present many simple models containing \(A\) and \(\psi_{VL}\) in SU(2) singlet and doublet representations. For \(A\) in a doublet, we restrict ourselves to the two-Higgs-doublet model (2HDM) Type-II and Type-X. Another important result of this work is 1-loop analytical expressions for the \(AVV\) effective couplings \(\kappa_{AVV}\) induced by SM fermions and VLFs in each of these models; as a function of the model parameters, we plot numerically these effective couplings and the \(BR\) into SM gauge bosons and fermions listed above.

In previous studies, one of us has considered the implications of models with vector-like quarks (VLQ) and vector-like leptons (VLL) coupled to the CP-even Higgs boson \(h\) in Ref. [6], and the direct LHC signatures of VLQs in Refs. [5]; this work complements them by considering aspects of CP-odd scalars \(A\). From the vast literature, we give a sampling below of studies that deal with a CP-odd scalar, namely \(A\) and \(H\), are studied in Ref. [22], where the LHC 8 TeV exclusion and 14 TeV reach from the processes \(gg \rightarrow H \rightarrow AZ\) and \(gg \rightarrow A \rightarrow HZ\) are presented. Ref. [23]
calculates the loop factors for the AVV couplings in the MSSM and the 2HDM with a heavy chiral fourth generation. Ref. [24] studies \( A \rightarrow WW, ZZ \) decays and compares this with the corresponding CP-even scalar decays in 2HDM-II, and also with a chiral fourth generation or additional heavy vector-like quarks (VLQ) added. In addition to these, here we also include the effects of VLFs on \( A \rightarrow \gamma\gamma, Z\gamma \) decays. These studies are done with specific models in mind while we present the LHC limits and signal c.s. in a model-independent manner, and using these, derive results for the models we introduce and also for some of the models above.

The paper is organized as follows: In Sec. 2 we present a model-independent analysis of the CP-odd scalar \( A \), present constraints on its effective couplings from the 8 TeV LHC run, the LHC gluon-fusion c.s., and \( BR \) into SM fermion and gauge boson decay modes. In Sec. 3 we present many simple models containing \( A \) and \( \psi_{VL} \) as \( SU(2) \) singlets or doublets. For each of these models, we work out the 1-loop effective couplings of the \( A \). One can read-out the current constraints and gluon-fusion c.s of the \( A \) at the LHC for each of these models in conjunction with the results in Sec. 2. The models considered include \( A \) as an \( SU(2) \) singlet, or contained in the 2HDM, with correspondingly the \( \psi_{VL} \) also in singlet or doublet representations. We offer our conclusions in Sec. 4. We compile the expressions for the 1-loop effective couplings and mixing angles in App. A, and for \( \psi_{VL} \) mixing with the SM top in App. B.

## 2 Model-independent Analysis

In this section, we define an effective Lagrangian with couplings of the CP-odd scalar \( A \) to SM gauge bosons and fermions. For the \( A \), we show the constraints from the 8 TeV LHC, signal c.s. \( \sigma \) at the 8 and 14 TeV LHC, and \( BR \) into various SM final states, as a function of the effective couplings and \( m_A \). For any given new physics model, one can obtain this effective Lagrangian by integrating out heavier fields, following which the results of this section can then be used to obtain the LHC limits and gluon-fusion cross-section in that model.

CP invariance requires the CP-odd scalar \( A \) coupling to SM gauge bosons to be only via higher dimensional operators. Showing only the new physics terms, the effective Lagrangian is

\[
\mathcal{L}_{eff} = \frac{1}{2} \partial_{\mu} A \partial^\mu A - \frac{1}{2} m_A^2 A^2 - y_{Af} A f_i \gamma_5 f_i - \frac{1}{16\pi^2 M} \kappa_{A\gamma\gamma} A F_{\mu\nu} \tilde{F}^{\mu\nu} - \frac{1}{16\pi^2 M} \kappa_{A\gamma Z} A F_{\mu\nu} \tilde{Z}^{\mu\nu} - \frac{1}{16\pi^2 M} \kappa_{A\gamma W} A W_{\mu\nu} \tilde{W}^{\mu\nu},
\]

where \( \kappa_{Aij} \) s contain fermion loop contributions, and \( \tilde{F}_{\mu\nu} = \epsilon_{\mu\nu\sigma\tau} F^{\sigma\tau} \). We have defined the dimensionless effective couplings \( \kappa \) by pulling out a new-physics mass-scale \( M \) in the effective AVV terms. For the numerical results we show, we set \( M = 1 \) TeV from now on and show only \( \kappa \), but for other values of \( M \), the \( \kappa \) can easily be rescaled. Although we have defined the effective couplings \( \kappa \) by extracting a heavy new-physics mass scale \( M \), SM fermion contributions are to be included when present. Eq. (1) is an effective Lagrangian at a scale just above \( m_A \). Heavy BSM fermions and the SM fermion contributions are to be included in \( \kappa \) before comparing with the plots we show in this section. For various simple SM extensions detailed in Sec. 3 we compute the \( \kappa \)’s and present them in App. A. If SM fermions contribute and can go onshell, the \( \kappa \) are complex. In this case, the \( \kappa_{AVV} \) that appear in our plots in this section should be read as \( |\kappa_{AVV}| \). We assume \( y_{Af} \) to be real in this work.

The CP-odd scalar \( A \) can decay to SM gauge bosons or fermions. In terms of the \( \kappa \) and \( y \) defined above, the decay rates to different final states are

\[
\Gamma(A \rightarrow Z\gamma) = \frac{1}{32\pi} \left( \frac{\kappa_{A\gamma\gamma}}{16\pi^2 M} \right)^2 m_A^3 (1 - \lambda_Z)^3, \quad \Gamma(A \rightarrow gg) = \frac{1}{8\pi} \left( \frac{\kappa_{Aga}}{16\pi^2 M} \right)^2 m_A^3, \quad \Gamma(A \rightarrow \gamma\gamma) = \frac{1}{16\pi} \left( \frac{\kappa_{A\gamma\gamma}}{16\pi^2 M} \right)^2 m_A^3.
\]

Using these expressions, one can work out the \( BR \) of the \( A \) into these final states in any new physics model.

We turn next to discussing limits from the 8 TeV LHC and the gluon-fusion cross-section at 14 TeV. To obtain the limits on the effective couplings \( \kappa \) and \( y \), we use upper-limits (UL) from recent LHC analyses on \( \sigma(pp \rightarrow \phi) \times BR(\phi \rightarrow XX) \), where \( \phi \) is either \( h \) or \( A \), and the currently relevant constraints are \( XX = \{\gamma\gamma, \tau^+\tau^-, \ell\ell\} \). We take the limits on the \( \gamma\gamma \) channel from the CMS analysis Ref. [25] which
has an upper limit up to $M_\phi$ of 850 GeV, on the $\tau^+\tau^-$ channel from the ATLAS analysis Ref. [26] up to $M_\phi$ of 1000 GeV, and from the ATLAS analysis Ref. [27] for the $tt'$ channel. Using these we constrain the effective couplings of Eq. (1).

At the LHC, the $A$ can be produced by $gg \to A$ (called gluon-fusion channel), which starts at the 1-loop level when $A$ couples to colored fermions. In addition to the above production channel, if $A$ couples to $b$-quarks, there are additional production channels, namely, $bb \to A$ (called $bb$-fusion), $bg \to bA$ and $gg \to b\bar{b}A$ (called $b$-quark associated production) channels; how these compare with the gluon-fusion channel depends on how large the $b\bar{b}A$ coupling is in a given model. For instance, for $y_{bA} = 0.5$, we find that the production rate via $bb$-fusion and $b$-quark associated production channels becomes comparable to the gluon-fusion channel with $\kappa_{Agg} \approx 20$. We include only the gluon-fusion channel in this study, but in models with a large $b\bar{b}A$ coupling, the $bb$ fusion and $b$-quark associated production channels may have to be included, which we do not do here. For a study involving the $b$-quark associated production channels of the $h$ including $gg \to b\bar{b}h$, see Ref. [32]. One can separately study the $b$-quark associated production channels by tagging on the final state $b$-jet as discussed in Ref. [26]. Although there are some LHC limits using $b$-tagged events to which the $bb$ decay mode and the $b$-quark associated production channels contribute, we do not include them in our analysis here. So far these results have been presented for $m_A < 350$ GeV (see Refs. [28, 29, 30]).

Rather than compute the $A$ production rate at the LHC ourselves, we relate it to the $h$ production rate at the same mass, and make use of the vast literature on $h$ production rate. Denoting $\phi = \{A, h\}$, and since $\sigma(gg \to \phi) \propto \Gamma(\phi \to gg)$, we can write the $\sigma \times BR$ for $A$ production followed by decay into the final-state $XX$ as

$$\sigma(gg \to A) = \frac{\Gamma(A \to gg)}{\Gamma(h \to gg)} \times \sigma(gg \to h).$$  \hspace{1cm} (3)

We compute $\Gamma(A \to gg)$ and $BR(A \to XX)$ as a function of the effective couplings and apply the upper-limit $UL$ from the 8 TeV LHC quoted above using Eq. (3). For our numerical work, we take $\sigma(gg \to h)$ from Ref. [31]. We assume here that the dependence on the PDF, and the acceptance at the LHC for $A$ and $h$ are not very different, which should be reasonable assumptions. For the decay $A \to XX$, the final-states $XX$ we consider are $\gamma\gamma$, $\tau^+\tau^-$ and $tt'$ as these are currently the significant ones. We compute the $BR(A \to XX)$ using Eq. (2). If more than one state is fairly close in mass to the $A$, i.e. closer than the experimental resolution to separate them, we should include all of them into the $\sigma \times BR$ above.

In Fig. 1 we show $\sigma(gg \to A)$ at the 8 TeV LHC (left plot) and 14 TeV LHC (right plot) as a function of $\kappa_{Agg}$. $\sigma(gg \to A)$ is obtained using Eq. (3) and the $\sigma(gg \to h)$ from Ref. [31] as mentioned earlier. In a given new physics model, one can compute $\kappa_{Agg}$ and then use these plots to obtain the $\sigma(gg \to A)$. Using the $\sigma(gg \to A)$, we obtain constraints from the 8 TeV LHC data as a function of the BR into a particular mode. We show this in Fig. 2 obtained from the $\gamma\gamma$, $\tau^+\tau^-$ and $tt'$ channels. The regions to the top and right of the curves are excluded at the 95% CL level. In the $\gamma\gamma$ channel, the bound is strongest for $m_A = 200$ GeV since the experimental exclusion is tightest at that mass. We see that there is no constraint from this channel for $BR(A \to \gamma\gamma) \lesssim 10^{-4}$ for the range of $\kappa_{Agg}$ shown. From the
Figure 2: 8 TeV LHC constraints from the $\gamma \gamma$ channel (left), $\tau^+\tau^-$ channel (middle) and $t\bar{t}$ channel (right), for $m_A = 200$ GeV (red), 500 GeV (green), 800 GeV (blue) and 1000 GeV (yellow). The regions to the top and right of the curves are excluded at the 95% CL level.

$\tau^+\tau^-$ channel, we find the strongest limit for $m_A$ of about 500 GeV since the experimental exclusion is tightest at that mass. We show in Fig. 3 the total $\sigma(gg \rightarrow A) \times BR(A \rightarrow XX)$ contours (in pb) for $XX = \{\gamma \gamma, \tau^+\tau^-, t\bar{t}\}$ at the 14 TeV LHC, making use of the fact that the total $\sigma(gg \rightarrow A \rightarrow XX)$ $\propto \kappa_{Agg} \times BR(A \rightarrow XX)$, omitting kinematical factors independent of couplings. Thus, each mode $XX$ can be considered and presented independently of the others as we do here. The 95% CL LHC exclusion discussed above is also shown labeled as ‘8 TeV’.

As already mentioned, the model-independent results presented in this section can be used to obtain the LHC constraints and gluon-fusion c.s. in any particular model by computing first the effective couplings in that model. We compute the effective couplings in many simple models next.

3 Models

In this section we consider some specific models for the CP-odd scalar $A$ and study its LHC production and decays. The goal is to capture in simple models many of the features present in realistic BSM models as far as the LHC phenomenology of $A$ is concerned. We mostly focus on the situation when $m_A < 2M_{VL}$ and do not focus on the phenomenology due to the $A$ decaying to a pair of on-shell VLF. We first consider the models where $A$ is an $SU(2)$ singlet and couples to $SU(2)$ singlet or doublet VLF (minimal vector-like up-type-singlet or MVU model) and $SU(2)$ doublet VLF (minimal vector-like quark doublet or MVQ model). Following this, we consider some models with and without $SU(2)$ singlet and doublet VLFs present, and $A$ as part of an $SU(2)$ doublet in the 2HDM framework for Type-II ($MVQD,MVQU$ models) and Type-X ($MVQDX_{11}$ model).

3.1 $SU(2)$ singlet $A$

We consider some models with an $SU(2)$ singlet $A$ coupled to $SU(2)$ singlet or doublet VLFs.

3.1.1 MVU model (singlet $A$ with a singlet VLF)

We study a model with an $SU(2)$ singlet $A$, and a vector-like fermion ($\psi$) that is an $SU(2)$ singlet, $SU(3)$ triplet (i.e. a vector-like quark, VLQ) and has a hypercharge $Y_\psi$.\footnote{A model with only a vector-like lepton singlet is uninteresting for $A$ phenomenology since no LHC production channels are significant (note that the $Ab\bar{b}$ coupling is also not possible for this model).} Clearly, the electromagnetic charge $Q = Y_\psi$. To the SM Lagrangian we add

$$L \supset \frac{1}{2} \partial_{\mu} A \partial^{\mu} A - \frac{1}{2} m_A^2 A^2 + \bar{\psi} i\gamma_5 \psi + e QA_{\mu} \bar{\psi} \gamma^{\mu} \psi - g Q \frac{3}{c_W} Z_{\mu} \bar{\psi} \gamma^{\mu} \psi$$

$$+ \bar{\psi} i\gamma_5 \psi - i y_A A \bar{\psi} \gamma_5 \psi - m_\psi \bar{\psi} \psi - \frac{\lambda A}{6} A^2 H^\dagger H.$$
Figure 3: Contours of the 14 TeV LHC $\sigma \times BR$ (in pb) in the $\gamma\gamma$ channel (upper-row) and the $\tau^+\tau^+$ channel (middle-row) for $m_A = 200$ GeV (left), 500 GeV (middle), 800 GeV (right), and in the $t\bar{t}$ channel (bottom-row) for $m_A = 500$ GeV (left), 800 GeV (middle), 1000 GeV (right). The region to the right of the contour labeled ‘8 TeV’ is excluded at the 95% CL level from 8 TeV LHC result.
Figure 4: BR \((A \rightarrow \gamma\gamma)\) (black), BR \((A \rightarrow \gamma Z)\) (blue), BR\((A \rightarrow ZZ)\) (red), BR\((A \rightarrow WW)\) (cyan) as a function of \(m_A\) with \(y_A = 0.1\) and \(m_\psi = 1000\) GeV for MVU (left) and MVQ (right) models.

Figure 5: \(\kappa_{Agg}/y_A^2\) as a function of \(m_A\) for \(m_\psi = 800\) GeV (red) and 1200 GeV (blue) for MVU model.

Here we have not considered possible terms coupling the \(A\) to a SM fermion and a VLF for \(Y_\psi = 2/3, -1/3\) such as \(\bar{\psi}_L A u_R, \bar{\psi}_L A d_R, q H \bar{\psi}_R\), which may lead to nontrivial constraints from FCNCs if mixing is allowed with the first two generations. We will comment on the possibility of mixing with the third generation SM fermions in Sec. 3.5.

We restrict ourselves to \(m_A < 2M_{VL}\), so that \(A\) cannot decay to a VLF pair. The possible decay modes of \(A\) are to \(gg, \gamma\gamma, Z\gamma\) and \(ZZ\) through a VLF loop, but no decay to \(W^+W^-\). \(A\) cannot decay to a pair of SM fermions since such couplings are forbidden by gauge invariance. The effective \(AV^\mu V^\nu\) couplings induced by VLFs are given in App. A.1. From these we compute the partial widths and the BR into the above modes. In Fig. 4 we plot BR\((A \rightarrow \gamma\gamma)\), BR\((A \rightarrow Z\gamma)\) and BR\((A \rightarrow ZZ)\) where we chose \(Y_\psi = 2/3\) as an example. BR\((A \rightarrow gg)\) is almost constant at around 0.999.

In Fig. 5 we plot \(\kappa_{Agg}/y_A^2\) as a function of \(m_A\). From this, one can read-off the \(\sigma(gg \rightarrow A)\) at the 8 and 14 TeV LHC from Fig. 1 in Sec. 2. The peaks in Fig. 5 are due to the VLFs going onshell, although as mentioned earlier, we do not explore its consequences in this work. In this model, the gluon-fusion c.s. of \(A\) is induced only through loops of the heavy VLFs due to which the 8 TeV LHC exclusion limits on \(\sigma \times \text{BR}\) into the \(ZZ\) channel (see Ref. [33]) or the \(\gamma\gamma\) channel (see Ref. [25]) are rather weak, unless \(y_A\) becomes so large that perturbativity is lost.

If \(m_A < m_h/2\) (where \(h\) is the 125 GeV Higgs), then \(h \rightarrow AA\) becomes kinematically allowed and becomes a means of producing \(A\) in addition to the gluon-fusion channel discussed above. In Fig. 6 we plot BR\((h \rightarrow AA)\) for \(\lambda_A = 0.1, 0.05\) and 0.001. When this decay is allowed, it will contribute to the Higgs total width thereby modifying the BRs into the other channels. In particular, it will modify the signal strength \(\mu_{\gamma\gamma} = \Gamma(h \rightarrow \gamma\gamma)/\Gamma_{SM}(h \rightarrow \gamma\gamma)\), which is measured to about 10% precision (see for example Ref. [34]). We plot \(\mu_{\gamma\gamma}\) in Fig. 6. We thus see that the constraint on \(\lambda_A\) from the 8 TeV LHC is of the order of 0.01 if \(m_A < m_h/2\).
We consider a BSM extension with an $SU(2)$ singlet $A$, and one $SU(2)$ doublet vector-like fermion

$\psi = \psi_{L,R} = (\psi_{1L,R} \psi_{2L,R})/Y_\psi$. To the SM Lagrangian we add

$$L \supset \frac{1}{2} \partial^\mu A \partial_\mu A - \frac{1}{2} m_A^2 A^2 + \bar{\psi} i D \psi - i y_A A \bar{\psi} \gamma_5 \psi - m_\psi \bar{\psi} \psi - \frac{\lambda_A}{4!} A^4 - \frac{\lambda_A}{6} A^2 H^\dagger H,$$

where the gauge interactions of the $\psi$ are understood and are not explicitly shown. For $Y_\psi = 1/6$ one can add the terms $y'_A \bar{\psi}_L H^\dagger u_R + y'_d \bar{\psi}_L H d_R + iy_A A_q l_R^\dagger \psi_R + h.c$, which we will not consider here as they can induce nontrivial FCNC constraints if mixing is allowed to the first two generations. We comment on the mixing to third generation SM fermions in Sec. 3.5. As in MVU model, there are no decays to a pair of SM fermions, but unlike there, in this model $A \rightarrow W^+ W^-$ decay is also possible through the VLF loop, in addition to $gg, \gamma \gamma, Z \gamma$ and $ZZ$ modes. The expressions for $\kappa_{AVV}$ are given in App. A.2. We take $Y_\psi = 1/6$ as an example.

In Fig. 4 we plot the $BR$ of $A$ into $\gamma \gamma, Z \gamma, ZZ$ and $W^+ W^-$ modes. As in MVU model, the $BR$ into $gg$ remains almost constant at around 0.99 for $m_A \gtrsim 300$ GeV. As the $\psi_1 \psi_2 W$ coupling ($g$) is greater than the $\psi_i \psi_i Z$ couplings ($g/c_W(T_3 - Q s_W^2)$), the $BR$ into $WW$ is larger than into $ZZ$. Again, for the same reasons explained in MVU model, the exclusion limits from the 8 TeV LHC in the $\gamma \gamma, ZZ, WW$ channels are rather weak in this model also.

$\sigma(gg \rightarrow A)$ in this model is twice of what was obtained in MVU model because there are two degenerate VLFS in the loop. The VLFS are degenerate since no Yukawa terms involving the SM Higgs can be written down that can split the masses after EWSB. Since no couplings to a pair of SM fermions exist, there are no $b$-quark initiated production processes possible.

### 3.2 CP-odd scalar $A$ in 2HDM-II

Here, we consider a CP-odd scalar ($A$) as a part of a scalar doublet and find the allowed regions of parameter space from the exclusion-limit on $\sigma(A) \times BR(A \rightarrow \tau^+ \tau^-)$ presented by ATLAS [26, 35]. We focus on the $\tau^+ \tau^-$ channel as currently this is the most constraining. We do this first in the 2HDM Type-II ($2HDM-II$), then add VLFS to the 2HDM-II and study the $A$ production and $BR$ in the presence of the VLFS.

In the 2HDM, as we will discuss more in detail below, there are two CP-even scalars ($h$ and $H$) in addition to the CP-odd scalar ($A$). In some regions of parameter-space, $m_A \approx m_H$, i.e. their masses are within the experimental resolution to distinguish them. If so, we must add the contributions from both $A$ and $H$ to any given channel; their sum is incoherent due to the different CP quantum-numbers. For instance, the experimental invariant-mass resolution in the $\tau^+ \tau^-$ channel is about 30% (see for instance Ref. [35]). Therefore, we consider two cases, one when $m_A$ and $m_H$ are within 30% and add the contributions from the “degenerate” $A$ and $H$, and another when they are split by more than 30% and treat them separately. When they are degenerate, for the $\tau^+ \tau^-$ channel for instance, we have $BR(A \rightarrow \tau^+ \tau^-) \approx BR(H \rightarrow \tau^+ \tau^-)$ in the so-called alignment limit (as will be defined precisely later),
and we can use the constraints obtained in Sec. 2 if we interpret $\kappa_{Agg}$ shown there as $\sqrt{\kappa_{Agg}^2 + \kappa_{Hgg}^2}$ and $BR(A \to \tau^+\tau^-)$ as $BR(A \to \tau^+\tau^-) + BR(H \to \tau^+\tau^-)$. For the non-degenerate case, again one can make use of our results in Sec. 2 to obtain constraints either for the $H$ or $A$. For the $H$ alone, one should just read the $\kappa_{Agg}$ and $BR(A \to XX)$ as $\kappa_{Hgg}$ and $BR(H \to XX)$ respectively in Sec. 2 and obtain the c.s. and constraints. In this case, we assume that $m_H > m_A$ so that the $A \to HZ$ decay is not kinematically allowed.

### 3.2.1 2HDM-II:

In the 2HDM-II we have two scalar doublets, $\Phi_1$ with hypercharge $+1/2$ and $\Phi_2$ with hypercharge $-1/2$. The SM Yukawa couplings and the Higgs potential are replaced by

$$L = - \left( y_d q_d^T \Phi_1 d_R + y_u q_u^T \Phi_2 u_R + h.c \right) + (D_\mu \Phi_1)^2 + (D_\mu \Phi_2)^2 - V(\Phi),$$

where

$$V(\Phi_1, \Phi_2) = m_{12}^2 \Phi_1^\dagger \Phi_1 + m_{22}^2 \Phi_2^\dagger \Phi_2 - m_{12}^2 (\Phi_1^\dagger \Phi_2 + h.c) + \lambda_1 (\Phi_1^\dagger \Phi_1)^2 + \lambda_2 (\Phi_2^\dagger \Phi_2)^2 + \lambda_3 (\Phi_1^\dagger \Phi_1) (\Phi_2^\dagger \Phi_2) + \lambda_4 (\Phi_1^\dagger \Phi_2) (\Phi_2^\dagger \Phi_1) + \lambda_5 ((\Phi_1^\dagger \Phi_1)^2 + h.c).$$

In the limit when $m_{12} = 0$, the Lagrangian has a discrete $Z_2$ symmetry under which $\Phi_1 \to -\Phi_1$, $d_R \to -d_R$ (with all other fields unchanged), if the down-type right-handed fermions couple only to the $\Phi_1$ and the up-type right-handed fermions only couple to the $\Phi_2$ so that there are no tree-level FCNCs (see for example Ref. [36]). Nonzero $m_{12}$ softly breaks this $Z_2$ symmetry. We will not consider the hard $Z_2$ breaking terms $(\Phi_1^\dagger \Phi_1 \Phi_2^\dagger \Phi_2 + \Phi_2^\dagger \Phi_2 \Phi_1^\dagger \Phi_1 + h.c)^2$. There are eight free parameters in $V$. After we fix the minimum of the potential at $\langle \Phi_1 \rangle = v_1$ and $\langle \Phi_2 \rangle = v_2$, with the constraint $v_1^2 + v_2^2 = \bar{v}^2 = (246 \text{ GeV})^2$, the number of free parameters reduces to seven which we take to be $m_{\alpha, \beta}, m_b, m_H, m_{H^\pm}, \tan \beta, \alpha$ and $m_{12}^2$ in notation that is common in the literature (see Ref. [37]). We parametrize the scalar doublets as

$$\Phi_1 = \begin{pmatrix} \phi_1^+ \\ \frac{1}{\sqrt{2}} (v_1 + \rho_1 + i \eta_1) \end{pmatrix}, \quad \Phi_2 = \begin{pmatrix} \frac{1}{\sqrt{2}} (v_2 + \rho_2 + i \eta_2) \\ \phi_2^- \end{pmatrix}$$

with $v_1 = v \sin \beta$, $v_2 = v \cos \beta$. The physical mass eigenstates are: a neutral goldstone boson $G^0 = \eta_1 \cos \beta + \eta_2 \sin \beta$, a heavy scalar $H = \rho_1 \cos \alpha + \rho_2 \sin \alpha$, a light scalar $h = -\rho_1 \sin \alpha + \rho_2 \cos \alpha$, a CP-odd scalar $A = -\eta_1 \sin \beta + \eta_2 \cos \beta$, and charged scalars $H^\pm = -\phi_1^\pm \sin \beta + \phi_2^\pm \cos \beta$. The expressions of $\alpha, \beta$ in terms of the model parameters can be found, for example, in Ref. [13, 37]. It is this CP-odd scalar $A$ that we are studying in this work.

All the effective couplings, relevant BRs and the cross sections in the 2HDM can be found in Refs. [37, 38]. We plot the BR of $A$ into $VV$, $\tau^+\tau^-$ and $bb$ in Fig. 8. Our results for $\Gamma(A \to \gamma\gamma)$ and $\Gamma(A \to Z\gamma)$ match with that of the Ref. [38]. We see that the BRs into $\gamma\gamma$ and $Z\gamma$ are smaller compared to that of the corresponding loop induced SM Higgs branching ratios even for tan $\beta = 1$ when the couplings of $A$ to the SM fermions are equal to the Higgs Yukawa couplings. This is because the partial width $\Gamma(h \to \gamma\gamma, Z\gamma)$, being dominated by the $W$ loop, is larger than the partial width $\Gamma(A \to \gamma\gamma, Z\gamma)$ in which only the fermions contribute (see for example Fig. 2.10 of Ref. [38]). For larger tan $\beta$ the branching ratios are even smaller because of the increased $\Gamma(A \to bb)$ and $\Gamma(A \to \tau^+\tau^-)$ (recall that the $Abb$ and $Ar\tau^-\tau^-$ couplings are proportional to tan $\beta$). The discontinuity at $m_A = 2m_t$ in the BRs in Fig. 8 for tan $\beta = 1$ is because of the onset of $A \to t\bar{t}$ on-shell decay. For larger tan $\beta$, the discontinuity is smaller since the $Att$ coupling becomes smaller. The $h \to AA$ decay, possible for $m_A < m_h/2$, is studied in Ref. [12] and we will not discuss it here.

We are interested in the case where the lighter CP-even scalar ($h$) is the observed 125 GeV Higgs boson. For this, the $\cos(\alpha - \beta) \approx 0$ is the most favored region (see Fig. 18 of Ref. [12]). Only a small range of other values of $(\alpha - \beta)$ are allowed where the sign of the down-type coupling of the Higgs is reversed. For the 2HDM with exact $Z_2$ symmetry (i.e. $m_{12} = 0$), tan $\beta$ has an upper limit of 7 from perturbativity constraint (see Ref. [16]). We will work with a nonzero $m_{12}$ which allows for larger values of tan $\beta$ (see Ref. [14]). We also assume that the “alignment limit” $(\alpha - \beta = \pi/2)$ holds sufficiently

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2This is a natural choice since if these terms are zero to start with they will not be induced at the loop level even if the soft breaking terms are present.
Figure 7: Regions of the $m_A$-$\tan \beta$ parameter space (blue region) which is excluded at 95% confidence level from $\phi \to \tau^+\tau^-$ decay when only $A$ is present (left) and when $m_A$ and $m_H$ are degenerate (right) for 2HDM-II.

accurately so that the $h$ couplings are SM like to match with the properties of the observed 125 GeV state at the LHC as discussed in Ref. [11]. In this limit, the $H \to WW$ and $H \to ZZ$ decays do not give any significant constraints on the parameter space (see for example Ref. [33]).

Using the $\tau^+\tau^-$ channel constraints shown in Sec. 2 (Fig. 2) we obtain constraints on this model. In Fig. 7 we plot the 95% confidence level constraints on the $m_A$-$\tan \beta$ plane, when only $A$ is present (left), and for $m_A = m_H$ when both contribute (right). Ref. [35] has presented similar constraints in the $m_A$-$\tan \beta$ plane, but for the MSSM.

In the following subsections we add various combinations of $SU(2)$ singlet and doublet VLFs to the 2HDM-II. Our goal is to study how VLFs affect the $A$ LHC production rate and decay BRs. As in MVU and MVQ models we do not consider here possible mixing term of the VLFs with the SM fermions, so a singlet VLF cannot couple to the Higgs doublets without the doublet VLF. We will discuss this possibility in Sec. 3.5. There are eight different ways in which the $\Phi_1$ and the $\Phi_2$ can couple to the VLFs consistent with the symmetries of the 2HDM-II namely $\Phi_1 \to -\Phi_1$ and $d_R \to -d_R$ (with all other fields unchanged). Among these eight models we will discuss only three representative ones that also capture the effects in the others.

### 3.2.2 MVQD model

We introduce one doublet VLQ, $\psi = (\psi_1, \psi_2)$ with hypercharge $Y_\psi$ and one singlet VLQ ($\chi$) with hypercharge $(Y_\psi - 1/2)$ so that VLF couplings with $\Phi_1$ are allowed. The additional Lagrangian terms to the 2HDM-II are

$$L \supset -y_1 \sqrt{2} A \sin \beta (i y_1 \bar{\psi}_2 L \chi_R + i \tilde{y}_1 \bar{\psi}_2 R \Phi_1 \chi_L + h.c) - M_{\psi} \bar{\psi}_2 - M_\chi \bar{\chi}_R .$$

We can also write the terms $\bar{\psi}_L \Phi_1^C \chi_R$ and $\bar{\psi}_R \Phi_2^C \chi_L$, which we do not add here but will consider them subsequently as another model. These terms are forbidden if $\chi \to -\chi$ under the $Z_2$ symmetry of 2HDM-II. The terms involving $h$, $A$ and VLFs after EWSB are

$$L \supset \frac{1}{\sqrt{2}} A \sin \beta (i y_1 \bar{\psi}_2 L \chi_R + i \tilde{y}_1 \bar{\psi}_2 R \chi_L + h.c) - \frac{v}{\sqrt{2}} \cos \beta (y_1 \bar{\psi}_2 L \chi_R + \tilde{y}_1 \bar{\psi}_2 R \chi_L + h.c)$$

$$- \frac{y_1}{\sqrt{2}} h \sin \alpha (y_1 \bar{\psi}_2 L \chi_R + \tilde{y}_1 \bar{\psi}_2 R \chi_L + h.c) - M_{\psi} \bar{\psi}_2 - M_\chi \bar{\chi}_R .$$

Gauge interactions of the VLFs are present and not shown explicitly. $\psi_2$ and $\chi$ mix after EWSB, while $\psi_1$ is itself a mass eigenstate. We define the mass eigenstates $\zeta_1$ and $\zeta_2$ as

$$\psi_{2L,R} = \zeta_{1L,R} \cos \theta_{L,R} - \zeta_{2L,R} \sin \theta_{L,R}$$

$$\chi_{L,R} = \zeta_{1L,R} \sin \theta_{L,R} + \zeta_{2L,R} \cos \theta_{L,R} .$$

10
where the mixing angles $\theta_L$ and $\theta_R$ are defined in App. A.3. In terms of these mass eigenstates, the Lagrangian in Eq. (10) can be written as

$$\mathcal{L} \supset -y_{ij} (iA_{iL} \bar{\zeta}_j + h.c) - M_{\psi} \bar{\psi}_i \psi_1 + \kappa_{L,R} Z_{\mu} \bar{\zeta}_i \gamma_\mu \zeta_j + k_i \bar{\zeta}_i \gamma_\mu \zeta_i \mu \zeta_i$$

$$-y_{ij} (h \bar{\zeta}_i \gamma_\mu \zeta_j + h.c) , \quad (13)$$

where $i,j = 1,2$. We take the $y_1$ and $\tilde{y}_1$ to be real, enforcing CP invariance in the BSM sector. The relative sign between $y_1$ and $\tilde{y}_1$ in Eq. (9) is physical for the following reason. If we want to get rid of this relative sign we need to make the transformations $\chi_L \rightarrow -\chi_L$ and $\chi_R \rightarrow \chi_R$, or $\chi_L \rightarrow \chi_L$ and $\chi_R \rightarrow -\chi_R$. In either case, the $M_{\chi}$ changes its sign and is therefore a physical effect. For chiral fermions, the sign of the mass term is not physical since one can rotate it away by the above transformations.

Instead of the $\chi$ (with hypercharge $(Y_\psi - 1/2)$), if we consider a VLF (say $\xi$) of hypercharge $(Y_\psi + 1/2)$, we get a different model where the $\xi$ couples to the $\Phi^C$ instead of the $\Phi_1$. This model will have similar phenomenology as $MVQD_{11}$ model, which we discuss later.

The effective couplings for this model are given in App. A. When $y_1 = \tilde{y}_1$, in addition to CP invariance, the Lagrangian in Eq. (10) is also invariant under $P$ and $C$ individually, with $A$ transforming as $A^P \rightarrow A$, $A^C \rightarrow -A$. This implies that the VLF contribution to $\kappa_{AVV}$ is zero since $AV_{\mu\nu}V^{\mu\nu}$ is not $P$ invariant (although it is $CP$ invariant). Also, the VLF contributions are maximum for $M_\psi = M_\chi$ when the mixing between the VLFs ($\psi_2$ and $\chi$) is maximum. We will take $M_\psi$ and $M_\chi$ to be equal from now on.

In Fig. 8, we plot $\text{BR}(A \rightarrow \gamma\gamma)$, $\text{BR}(A \rightarrow B\gamma)$, $\text{BR}(A \rightarrow b\bar{b})$, and $\text{BR}(A \rightarrow \tau^+\tau^-, t\bar{t})$ for $Y_\psi = 1/6$ as an example, which is the SM quark-doublet hypercharge assignment. We see that for small values of $\tan \beta$ the VLF contribution to $\text{BR}(A \rightarrow \gamma\gamma)$ is small compared to the 2HDM-II. This is because $y_{ij}$s are proportional to $\sin \beta$. For large $\tan \beta$ and for large $m_A$, the VLF contributions to the $\text{BR}(A \rightarrow \gamma\gamma)$ become significant.
In Fig. 9, we plot contours of $\kappa_{Agg}$ for $M_\psi = M_\chi = 800$ GeV, 1700 GeV. For comparison we have also plotted the corresponding contours in 2HDM-II. Using this, one can read-off the $\sigma(gg \to A)$ at the 8 and 14 TeV LHC from Fig. 1 in Sec. 2. For comparison, we also show the corresponding contours in the 2HDM-II (without the VLFs). In Fig. 9 we also plot $y_{11}$ and $y'_{11}$ (defined in Eq. (13)) in the alignment limit ($\alpha - \beta = \pi/2$), which shows that the $h$ couplings to the VLFs become very small as $\tan \beta$ increases. Thus, the VLFs can modify $\sigma(gg \to A)$ and $\Gamma(A \to VV)$ significantly, while the $h$ remains SM-like as required by the LHC measurements of the 125 GeV state. We find that the VLF contributions partially cancel the SM fermion contributions for a range of low $\tan \beta$ values and for some ranges of the $m_A$, while for larger $\tan \beta$ the effective couplings always increase compared to the 2HDM-II. To illustrate this point more explicitly, we plot $\kappa_{Agg}$ as a function of $\tan \beta$ in Fig. 10 for $m_A = 300$ GeV and 600 GeV. The constraint on the 2HDM was nontrivial only for large $\tan \beta$ (see Fig. 7). Therefore, for large $\tan \beta$, since the $\kappa_{Agg}$ is bigger for this model compared to 2HDM (see Fig. 10), and the tree-level $\tau^+\tau^-$ BR from which the tightest constraint appears is almost unchanged, the constraint on this model will be tighter.

Figure 9: Contours of $\kappa_{Agg}$ for $M_\psi = M_\chi = 800$ GeV (left), 1700 GeV (middle), $y_1 = 0.5$, $\tilde{y}_1 = 1$ for MVQD$_{11}$ model, and corresponding contours for 2HDM-II (right). The bottom plot shows $y_{11}$ (red), $y'_{11}$ (blue) as a function of $\tan \beta$.

Figure 10: $\kappa_{Agg}$ as a function of $\tan \beta$ for $m_A = 300$ GeV (left) and 600 GeV (right), with $y_1 = 0.5$, $\tilde{y}_1 = 1$ and $M_\psi = 800$ GeV (green), 1000 GeV (blue) for MVQD$_{11}$ model and 2HDM-II (dashed-black).
Figure 11: Contours of $\kappa_{Hgg}$ for $y_1 = 0.5$, $\tilde{y}_1 = 1$, for $M_\psi = M_\chi = 800$ GeV (left), 1700 GeV (right) for $MVQD_{11}$ model and corresponding contours in 2HDM-II.

In Fig. 11, we plot contours of $\kappa_{Hgg}$ for $m_A = m_H$, in the alignment limit. Corresponding contours in 2HDM-II are also plotted for comparison. From this, one can also obtain $\sigma(gg \rightarrow H)$ from Fig. 1 by reading $\kappa_{Agg}$ there as $\kappa_{Hgg}$ as mentioned earlier.

### 3.2.3 MVQU$_{22}$ model

We introduce one doublet VLQ ($\psi$) with hypercharge $Y_\psi$ and one singlet VLQ ($\xi$) with hypercharge $Y_\xi + 1/2$, which couples only to $\Phi_2$. We add the following terms to the 2HDM-II Lagrangian

$$
\mathcal{L} \supset \bar{\psi}_L \Phi_2^C \xi_R + \bar{\psi}_R \Phi_2^C \xi_L + h.c. - M_\psi \bar{\psi}_L \psi - M_\xi \bar{\xi}_L \xi.
$$

Here we do not include the terms $\bar{\psi}_L \Phi_1^C \xi_R$ and $\bar{\psi}_R \Phi_1^C \xi_L$ as their effects have been considered in $MVQD_{11}$ model. As the BR($A \rightarrow VV$)s do not change much compared to the 2HDM-II case, we do not show them here. Instead of the $\xi$ (with hypercharge $Y_\xi + 1/2$) if we consider a VLF (say $\chi$) of hypercharge ($Y_\chi - 1/2$) we get a different model where the $\chi$ couples to the $\Phi_2^C$ instead of the $\Phi_2$. This will give similar effects to what we consider here.

The effective couplings for this model are given in App. A. As in $MVQD_{11}$ model, the $\kappa_{AVV}$ becomes zero when $y_2 = \tilde{y}_2$. In Fig. 12 we plot contours of $\kappa_{Agg}$ in $m_A$-$\tan \beta$ plane. For large $\tan \beta$ VLFs have very small contribution compared to the 2HDM-II which can be seen from Fig. 13. This is because $y_{1,2}$s are proportional to $\cos \beta$ which become small as $\tan \beta$ increases. Similar conclusions hold for $\kappa_{Hgg}$. In Fig. 14 we plot $\kappa_{Hgg}$ using which one can read-off the $\sigma(gg \rightarrow H)$ from Fig. 1 by reading $\kappa_{Agg}$ there as $\kappa_{Hgg}$ as mentioned earlier. Since $\kappa_{Agg}$ and $\kappa_{Hgg}$ do not change much compared to the 2HDM-II, constraints on the $m_A$-$\tan \beta$ plane will almost remain same as in the 2HDM-II case. Thus, VLFs if realized as in $MVQU_{22}$ model have little impact on the observables we consider here.

### 3.2.4 MVQU$_{12}$ model

We introduce one doublet VLQ ($\psi$) with hypercharge $Y_\psi$ and one singlet VLQ ($\xi$) with hypercharge ($Y_\xi + 1/2$). We consider the case where $\xi_R$ couples only to $\Phi_1$ and $\xi_L$ couples only to $\Phi_2$. To the 2HDM-II Lagrangian, we add

$$
\mathcal{L} \supset \bar{\psi}_L \Phi_1^C \xi - (y_1 \bar{\psi}_L \Phi_1^C \xi_R + \tilde{y}_1 \bar{\psi}_R \Phi_2 \xi_L + h.c.) - M_\psi \bar{\psi}_L \psi - M_\xi \bar{\xi}_L \xi.
$$

We get different models if instead of the couplings above, the $\psi$ and $\xi$ couples to $\Phi_2^C$ and $\psi$ couples to $\Phi_2$, or, if instead of $\xi$ we introduce a VLF singlet (say $\chi$) with hypercharge ($Y_\chi - 1/2$) with couplings to $\Phi_1$ and $\Phi_2$. All these models have similar phenomenology as $MVQU_{12}$ model.

The effective couplings for this model are given in App. A. In this model, the effective couplings do not reduce to zero for $y_1 = \tilde{y}_1$, unlike in $MVQD_{11}$ and $MVQU_{22}$ models, as there are no additional $P$ and $C$ symmetries in the VLF sector. In Fig. 15, we plot the BR($A \rightarrow VV$), BR($A \rightarrow b\bar{b}, \tau^+\tau^-, tt$) for an example choice of $Y_\psi = 1/6$. In Fig. 16 we plot contours of $\kappa_{Agg}$ for $y_1 = \tilde{y}_1 = 1$ and $M_\psi = M_\xi = 800$ GeV
Figure 12: Contours of $\kappa_{\text{Agg}}$ for $y_2 = 0.5$, $\tilde{y}_2 = 1$, for $M_\psi = M_\chi = 800$ GeV (left), 1700 GeV (right) for MVQU$_{22}$ model.

Figure 13: $\kappa_{\text{Agg}}$ with $\tan \beta$ for $m_A = 300$ GeV (left) and 600 GeV (right) with $y_2 = 0.5$, $\tilde{y}_2 = 1$ and $M_\psi = 800$ GeV (green), 1000 GeV (blue) for MVQU$_{22}$ model and 2HDM-II (dashed-black).

and 1700 GeV. From this, one can obtain $\sigma(gg \to A)$ at the 8 and 14 TeV LHC from Fig. 1 in Sec. 2. For low values of $\tan \beta$ the effective coupling increases compared to the 2HDM-II case, while for larger values of $\tan \beta$ the effective coupling decreases compared to the 2HDM-II. To show this more explicitly, we plot $\kappa_{\text{Agg}}$ with $\tan \beta$ in Fig. 17. The decreased coupling is due to a destructive interference between the contributions from SM fermions and the VLFs. If we reverse the sign of $y_1$ or $\tilde{y}_1$, we get the opposite effect; for low values of $\tan \beta$ the effective coupling decreases compared to the 2HDM-II while for larger values of $\tan \beta$ the effective coupling increases compared to the 2HDM-II. In Fig. 18 we plot contours of $\kappa_{Hgg}$ in the alignment limit. From this, one can also obtain $\sigma(gg \to H)$ from Fig. 1 by reading $\kappa_{\text{Agg}}$ there as $\kappa_{Hgg}$ as mentioned earlier.

In Fig. 19 we plot the regions of the $m_A$-$\tan \beta$ parameter-space which is excluded at 95% confidence level for two cases, when only $A$ is present, and when $A$ and $H$ are degenerate and both present. For comparison, we have also plotted the corresponding limit for the 2HDM-II case. We see that the constraints are loosened compared to the 2HDM-II due to the presence of VLFs. This happens because of the reduction of $\kappa_{\text{Agg}}$ ($\kappa_{Hgg}$) compared to the 2HDM-II.
3.3 CP-odd scalar $A$ in 2HDM Type-X:

In the 2HDM Type-X Model (2HDM-X) (see Ref. [20, 37] for a description of this model) all the SM quarks couple to $\Phi_2$ and all the leptons couple to $\Phi_1$. As a result, $A$ coupling to the quarks and leptons are proportional to $\cot \beta$ and $\tan \beta$ respectively. In the Type-X model, since all SM quarks couple very weakly to $A$ for large $\tan \beta$, $\sigma(gg \rightarrow A)$ becomes very small for large $\tan \beta$. As a consequence there are no constraints from $\sigma(pp \rightarrow A) \times BR(A \rightarrow \tau^+\tau^-)$. The Lagrangian for the model 2HDM-X is given by

$$L \supset - (y_d \bar{q}_L \Phi_2^C \bar{d}_R + y_u \bar{u}_L \Phi_2 \bar{u}_R + y_e \bar{e}_L \Phi_1 \bar{e}_R + h.c) + (D_{\mu} \Phi_1)^2 + (D_{\mu} \Phi_2)^2 - V(\Phi).$$

(16)

3.3.1 MVQDX$_{11}$ model

To the 2HDM Type-X model in Eq. (16), we introduce VLFs in a similar fashion as in MVQD$_{11}$ model as a representative case, and call it MVQDX$_{11}$ model. The other ways of coupling VLFs similar to MVQU$_{22}$ or MVQU$_{12}$ model will be qualitatively similar to our results here. We introduce a doublet VLQ $\psi = (\psi_1, \psi_2)$ with hypercharge $Y_{\psi}$, and a singlet VLQ $\chi$ with hypercharge $(Y_{\psi} - 1/2)$ which couples only to $\Phi_1$. To the 2HDM-X Lagrangian we add

$$L \supset \bar{\psi}_i D\psi_i + \chi_i D\chi_i - (y_1 \bar{\psi}_L \Phi_1 \chi_R + \bar{\psi}_R \bar{\Phi}_1 \chi_L + h.c) - M_{\psi} \bar{\psi}_i \psi_i - M_{\chi} \bar{\chi}_i \chi_i.$$  

(17)

The effective couplings of $A$ with VLFs are same as in MVQD$_{11}$ model and can be read off from App. A.1. The SM quark contribution to $\kappa_{AVV}$ for 2HDM-X can obtained from that of 2HDM-II (see Ref. [38]) by replacing $\tan \beta$ with $\cot \beta$ in the $A_{bb}$ coupling. In Fig. 20 we plot $BR(A \rightarrow \tau^+\tau^-)$ and $BR(A \rightarrow b\bar{b}, t\bar{t})$. We see that for large $\tan \beta$, $BR(A \rightarrow \gamma\gamma)$ and $BR(A \rightarrow Z\gamma)$ are similar to MVQDX$_{11}$ model, as expected. For $\tan \beta = 30$, $BR(A \rightarrow VV)$ is increased compared to 2HDM-II. This is because for large $\tan \beta$, $\Gamma(A \rightarrow b\bar{b})$ becomes much smaller in 2HDM-X.

In Fig. 21 we plot contours of $\kappa_{A_{gg}}$. For comparison, we also plot the corresponding contours in 2HDM-X (without VLFs). Using these plots, one can read off $\sigma(gg \rightarrow A)$ for 8 TeV and 14 TeV LHC from Fig. 1 in Sec. 2. As expected, for large $\tan \beta$, $\kappa_{A_{gg}}$ is significantly larger in this model compared to 2HDM-X since the VLFs contribute substantially while the SM quark contributions alone are very small.

In order to show explicitly how large the change is, we plot $\kappa_{A_{gg}}$ as a function of $\tan \beta$ for $m_A = 300$ GeV and 600 GeV in Fig. 22. The results for $\kappa_{A_{gg}}$, $\kappa_{H_{gg}}$ in 2HDM-X are also applicable for 2HDM-I as the SM quarks couple to $H,A$ in an identical fashion as in 2HDM-X. In Fig. 23 we plot contours of $\kappa_{H_{gg}}$ in $m_A\tan \beta$ plane in the alignment limit. From this, one can also obtain $\sigma(gg \rightarrow H)$ from Fig. 1 by reading $\kappa_{A_{gg}}$ there as $\kappa_{H_{gg}}$ as mentioned earlier.
Figure 15: $\text{BR}(A \rightarrow \gamma \gamma)$ (top-left and top-middle) and $\text{BR}(A \rightarrow Z\gamma)$ (top-right) with $M_\psi = M_\chi = 1000$ GeV (solid-black) for $\tan \beta = 1$ and 30 for $MVQU_{12}$ model, and the corresponding variation in the 2HDM-II (dashed-black); $\text{BR}(A \rightarrow \tau^+\tau^-, b\bar{b})$ for $\tan \beta = 1, 5, 10, 15, 30$ (middle-middle and middle-right) and $\text{BR}(A \rightarrow t\bar{t})$ for $\tan \beta = 1, 5, 10, 15$ (bottom).

Figure 16: Contours of $\kappa_{Agg}$ for $y_1 = 1, \tilde{y}_1 = 1$, for $M_\psi = M_\chi = 800$ GeV (left) and 1700 GeV (right) for $MVQU_{12}$ model.
Figure 17: $\kappa_{Agg}$ with $\tan\beta$ for $m_A = 300$ GeV (left) and 600 GeV (right) with $y_1 = 1$, $\tilde{y}_1 = 1$ and $M_\psi = 800$ GeV (green), 1000 GeV (blue) for $MVQU_{12}$ model and 2HDM-II (dashed-black).

Figure 18: Contours of $\kappa_{Hgg}$ for $y_1 = 1$, $\tilde{y}_1 = 1$, for $M_\psi = M_\chi = 800$ GeV (left), 1700 GeV (right) for $MVQU_{12}$ model.

Figure 19: For $MVQU_{12}$ model, regions of the $m_A$-$\tan\beta$ parameter-space excluded at the 95% CL from $\phi \to \tau^+\tau^-$ decay when only $A$ is present (left), and when $A$ and $H$ are degenerate and both present (right), with $y_1 = \tilde{y}_1 = 1$, $M_\psi = M_\chi = 800$ GeV (dark-blue region), 1000 GeV (light-blue and dark-blue regions). All shaded regions are excluded in the 2HDM-II.
Figure 20: \( \text{BR}(A \to \gamma \gamma) \) (top-left and top-middle) and \( \text{BR}(A \to Z \gamma) \) (top-right and top-middle) with \( M_\psi = M_\chi = 1000 \text{ GeV} \) (solid-black) for \( \tan \beta = 1 \) and 30 for \( MVQDX_{11} \) model, and the corresponding variation in the 2HDM-X (dashed-black); \( \text{BR}(A \to \tau^+ \tau^-) \) for \( \tan \beta = 1, 5, 10, 15, 30 \), \( \text{BR}(A \to b \bar{b}) \) for \( \tan \beta = 1, 5 \) (middle-middle and middle-right) and \( \text{BR}(A \to t \bar{t}) \) for \( \tan \beta = 1, 5, 10, 15, 30 \) (bottom).

Figure 21: Contours of \( \kappa_{A_{gg}} \) for \( y_1 = 0.5, \tilde{y}_1 = 1 \), \( M_\psi = 800 \text{ GeV} \) (left), 1700 GeV (middle) for \( MVQDX_{11} \) model and corresponding contours in 2HDM-X (right).
\( \kappa_{A\gamma\gamma} \) with \( \tan \beta \) for \( m_A = 300 \) GeV (left) and 600 GeV (right) with \( y_1 = 0.5, \tilde{y}_1 = 1 \) and \( M_\psi = 800 \) GeV (green), 1000 GeV (blue) for \( \text{MVQDX}_{11} \) model and 2HDM-X (dashed-black).

Figure 22: \( \kappa_{A\gamma\gamma} \) with \( \tan \beta \) for \( m_A = 300 \) GeV (left) and 600 GeV (right) with \( y_1 = 0.5, \tilde{y}_1 = 1 \) and \( M_\psi = 800 \) GeV (green), 1000 GeV (blue) for \( \text{MVQDX}_{11} \) model and 2HDM-X (dashed-black).

\( \kappa_{H\gamma\gamma} \) for \( y_1 = 1, \tilde{y}_1 = 1 \), for \( M_\psi = M_\chi = 800 \) GeV (left), 1700 GeV (middle) for \( \text{MVQDX}_{11} \) model and the corresponding contours in 2HDM-X (right).

Figure 23: Contours of \( \kappa_{H\gamma\gamma} \) for \( y_1 = 1, \tilde{y}_1 = 1 \), for \( M_\psi = M_\chi = 800 \) GeV (left), 1700 GeV (middle) for \( \text{MVQDX}_{11} \) model and the corresponding contours in 2HDM-X (right).

### 3.4 \( \text{MVLE}_{11} \) (Model with VL-lepton)

VLLs do not contribute in \( gg \to A \), but can contribute in \( A \to VV \). We show the effect of VLLs in a simple model similar to \( \text{MVQD}_{11} \) model, but with VLLs instead of VLQs. We introduce one doublet VLL (\( \psi \)) with hypercharge, \( Y_\psi \), and one singlet VLL (\( \chi \)) with hypercharge, \( Y_\chi = -1/2 \). The Lagrangian we consider is exactly the same as in Eq. (9), except here the VLLs \( \psi \) and \( \chi \) do not couple to gluons. The effective couplings are the same as for \( \text{MVQD}_{11} \) model except for color factors. As an example, we choose \( Y_\psi = -1/2 \) and plot BR(A \( \to \) \( \gamma \gamma \)) as a function of \( m_A \) in Fig. 24, with \( M_\psi = M_\chi = 500 \) GeV, for \( \tan \beta = 1 \) and 30. We see that the effect of VLLs is qualitatively similar to vector-like quarks; for low \( \tan \beta \) the effect of VLLs is negligible while for large \( \tan \beta \) and large \( m_A \) VLL contributions are significant. Near \( m_A = 1000 \) GeV, the VLL contribution is quite large due to them going onshell for our choice of VLL mass of 500 GeV. BR(A \( \to \) Z\( \gamma \)) will show the same behavior.

### 3.5 VLF mixing with SM-fermions

In this section we will briefly discuss the effect of mixing between the SM quarks and the VLFs. Constraints on the mixing from EWPT and a \( t' \) (vector-like top) decaying to \( Wb, Zt, Ht \) are studied in Ref. [5, 6, 39, 40]. Constraints from flavor observables are studied in Ref. [39]. All of the studies mentioned above requires the mixing to be small. Even if we add a small mixing between \( t' \) and \( t \), \( \kappa_{AVV} \)'s obtained in Sec. 3 will receive only tiny corrections. As an example, we consider a simple model in App. B, incorporating the mixing between \( t' \) and \( t \) and show that this is indeed the case.
4 Conclusions

Many theories beyond the standard model (BSM) contain CP-odd scalars (A) and new vector-like fermions (ψVL). Our goal was to study the LHC phenomenology of A when it is lighter than ψVL and coupled to it.

In Sec. 2 we have written an effective Lagrangian with A coupled to standard model (SM) gauge-bosons and fermions. We consider the situation when A couples significantly only to third generation SM fermions, namely t, b, τ, a situation common in many BSM extensions. The couplings of the A to standard model gauge bosons (i.e. AVV couplings) cannot occur from renormalizable operators in a CP-conserving sector, but can be induced as loop-generated non-renormalizable operators. These operators are induced by SM fermions and also the heavy ψVL. In Sec. 2 we present model-independent results that are useful whatever be the origin of these effective couplings. In Fig. 1 we present the 8 TeV and 14 TeV LHC gg → A (gluon-fusion channel) cross-sections as a function of the effective couplings. We also obtain limits on the effective couplings from the 8 TeV LHC data on the γγ, τ⁺τ⁻ and tt modes. We do not include the b¯b decay mode and the b-quark associated production channels in this work.

We define some simple models in Sec. 3 that are representative of BSM constructions as far as the phenomenology of A is concerned. These models include A and ψVL in the singlet and doublet representations of SU(2). In the doublet case, we focus on the two-Higgs-doublet (2HDM) Type-II and Type-X models. We compute the Agg and AVV effective couplings induced by the SM fermions and vector-like fermions at the 1-loop level and present analytical expressions for them in App. A. For the various models we define, we present the effective couplings κAgg, κAVV for V = {γ, Z}, BR(A → VV) and BR(A → f ¯f) for f = {τ, b, t} as a function of the model parameters. From the κAgg, one can see if a point in parameter-space in a given model is allowed by the 8 TeV data from our plots in Sec. 2. One can also read-off the gluon-fusion cross-section at the 8 TeV and 14 TeV LHC from Fig 1. Interestingly, for some of the 2HDM cases we studied, we find that the addition of vector-like fermions loosens the constraint compared to the 2HDM alone, and allows more of the parameter-space. This can be seen for instance in Fig. 19. The 14 TeV LHC gluon-fusion c.s. and the BRs in the different models we present should be useful in identifying promising discovery channels for the A after including a study of the relevant backgrounds.

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A Effective couplings and mixing angles

Here, we give expressions for $\kappa_{AVV}$’s for the various models we have considered. We also provide the explicit expressions for $\kappa_{ij}$’s and $y_{ij}$’s defined in Eq. (13). App. A.1, A.2 contain expressions for $\kappa_{AVV}$’s for MVU, MVQ models respectively. Sections A.3, A.4 and A.5 contain explicit expressions for $y_{ij}$’s and $\kappa_{ij}$’s for $\text{MVQD}_{11}$, $\text{MVQU}_{22}$ and $\text{MVQU}_{12}$ models respectively. App. A.6 contains general expressions for $\kappa_{AVV}$ for $\text{MVQD}_{11}$, $\text{MVQU}_{22}$, $\text{MVQU}_{12}$ models.

A.1 $\kappa_{AVV}$’s in MVU model

The effective couplings $\kappa_{AVV}$ (defined in Eq. 1) for MVU model are given by
\[
\kappa_{A\gamma\gamma} = 4N_c g_A e_Q (cQ)^2 \left( \int_0^1 dy \int_0^{y-1} dx \frac{\lambda_{\psi}}{m_{\psi}} \right),
\]
\[(A.18)\]
\[
\kappa_{A\gamma Z} = 4N_c g_A (cQ)^2 \left( \frac{\lambda_{\psi}}{m_{\psi}} \right) \left( \int_0^1 dy \int_0^{y-1} dx \frac{\lambda_{\psi}}{m_{\psi}} \right),
\]
\[(A.19)\]
\[
\kappa_{Agg} = 4g_A g_s^2 \left( \int_0^1 dy \int_0^{y-1} dx \frac{\lambda_{\psi}}{m_{\psi}} \right),
\]
\[(A.20)\]
\[
\kappa_{AZZ} = 4N_c g_A \left( \frac{\lambda_{\psi}}{m_{\psi}} \right) \left( \int_0^1 dy \int_0^{y-1} dx \frac{\lambda_{\psi}}{m_{\psi}} \right),
\]
\[(A.21)\]

where $\lambda_{\psi} = \frac{m_c^2}{m_A^2}$, $\lambda_{Z} = \frac{m_c^2}{m_A^2}$.

A.2 $\kappa_{AVV}$’s in MVQ model

The effecting couplings $\kappa_{AVV}$ for MVQ model are given by
\[
\kappa_{A\gamma\gamma} = 4N_c g_A e_Q (cQ) \left( (Q_1)^2 + (Q_2)^2 \right) \left( \int_0^1 dy \int_0^{y-1} dx \frac{\lambda_{\psi}}{m_{\psi}} \right),
\]
\[(A.22)\]
\[
\kappa_{A\gamma Z} = \kappa_{1\gamma Z} + \kappa_{2\gamma Z}, \text{ with}
\]
\[(A.23)\]
\[
\kappa_{ijZ} = 4N_c g_A e_Q (cQ) \left( \frac{\lambda_{\psi}}{m_{\psi}} \right),
\]
\[(A.24)\]
\[
\kappa_{ijZ} = 4N_c g_A e_Q \left( \frac{\lambda_{\psi}}{m_{\psi}} \right),
\]
\[(A.25)\]
\[
\kappa_{ijZ} = \kappa_{1\gamma Z} + \kappa_{2\gamma Z}, \text{ with}
\]
\[(A.26)\]

where $\lambda_{W} = \frac{m_w^2}{m_A^2}$ and $Q_{1,2}$, $T_{3}^{1,2}$ refers to $Q$ and $T_{3}$ of $\psi_1$ and $\psi_2$ respectively.
The couplings $\kappa_{ij}$ defined in Eq. 13 for $MVQD_{11}$ model and also for $MVQDX_{11}$ model are given by

$$\kappa_{11} = \frac{g}{c_W} \left( \frac{s_W}{3}  - \frac{1}{2} \cos^2 \theta_L \right) + \frac{g}{c_W} \left( \frac{s_W}{3}  - \frac{1}{2} \cos^2 \theta_R \right),$$

(A.27)

$$\kappa_{11} = \frac{g}{c_W} \left( \frac{s_W}{3}  - \frac{1}{2} \sin^2 \theta_L \right) + \frac{g}{c_W} \left( \frac{s_W}{3}  - \frac{1}{2} \sin^2 \theta_R \right),$$

(A.28)

$$\kappa_{L,R12} = \frac{g}{2c_W} \sin \theta_{L,R} \cos \theta_{L,R},$$

(A.29)

The couplings $y_{ij}$ (in Eq. 13) are given by

$$y_{11} = \frac{1}{\sqrt{2}} \sin \beta(y_1 \cos \theta_L \sin \theta_R - \tilde{y}_1 \sin \theta_L \cos \theta_R),$$

(A.30)

$$y_{22} = -\frac{1}{\sqrt{2}} \sin \beta(y_1 \sin \theta_L \cos \theta_R - \tilde{y}_1 \cos \theta_L \sin \theta_R),$$

(A.31)

$$y_{12} = \frac{1}{\sqrt{2}} \sin \beta(y_1 \cos \theta_L \sin \theta_R + \tilde{y}_1 \sin \theta_L \cos \theta_R),$$

(A.32)

$$y_{21} = -\frac{1}{\sqrt{2}} \sin \beta(y_1 \sin \theta_L \sin \theta_R + \tilde{y}_1 \cos \theta_L \cos \theta_R).$$

(A.33)

The mass eigenvalues $M_{1,2}$ (in Eq. 13) are given by

$$M_{1,2} = \frac{1}{2} \sqrt{(M_\phi + M_\chi)^2 + \frac{1}{2} \cos \beta \tilde{v}^2 (y_1 - \tilde{y}_1)^2 \pm \sqrt{(M_\phi - M_\chi)^2 + \frac{1}{2} \cos \beta \tilde{v}^2 (y_1 + \tilde{y}_1)^2}}$$

(A.34)

and the mixing angles $\theta_{L,R}$ are given by

$$\tan 2 \theta_L = \frac{2 \sqrt{2} \tilde{v} \cos \beta(y_1 M_\chi + \tilde{y}_1 M_\phi)}{2(M_\phi^2 - M_\chi^2) - \tilde{v}^2 \sin \beta \tilde{v}^2 (y_1^2 - \tilde{y}_1^2)},$$

(A.35)

$$\tan 2 \theta_R = \frac{2 \sqrt{2} \tilde{v} \cos \beta(y_1 M_\chi + \tilde{y}_1 M_\phi)}{2(M_\phi^2 - M_\chi^2) + \tilde{v}^2 \sin \beta \tilde{v}^2 (y_1^2 - \tilde{y}_1^2)},$$

(A.36)

A.4 Couplings and mass eigenvalues for $MVQU_{22}$ model

The couplings $\kappa_{ij}$ for $MVQU_{22}$ model are given by

$$\kappa_{11} = \frac{g}{c_W} \left( \frac{s_W}{3}  - \frac{1}{2} \cos^2 \theta_L \right) + \frac{g}{c_W} \left( \frac{s_W}{3}  - \frac{1}{2} \cos^2 \theta_R \right),$$

(A.37)

$$\kappa_{11} = \frac{g}{c_W} \left( \frac{s_W}{3}  - \frac{1}{2} \sin^2 \theta_L \right) + \frac{g}{c_W} \left( \frac{s_W}{3}  - \frac{1}{2} \sin^2 \theta_R \right),$$

(A.38)

$$\kappa_{L,R12} = \frac{g}{2c_W} \sin \theta_{L,R} \cos \theta_{L,R},$$

(A.39)

The couplings $y_{ij}$ are given by

$$y_{11} = \frac{1}{\sqrt{2}} \cos \beta(y_1 \cos \theta_L \sin \theta_R - \tilde{y}_1 \sin \theta_L \cos \theta_R),$$

(A.40)

$$y_{22} = -\frac{1}{\sqrt{2}} \cos \beta(y_1 \sin \theta_L \cos \theta_R - \tilde{y}_1 \cos \theta_L \sin \theta_R),$$

(A.41)

$$y_{12} = \frac{1}{\sqrt{2}} \cos \beta(y_1 \cos \theta_L \sin \theta_R + \tilde{y}_1 \sin \theta_L \cos \theta_R),$$

(A.42)

$$y_{21} = -\frac{1}{\sqrt{2}} \cos \beta(y_1 \sin \theta_L \sin \theta_R + \tilde{y}_1 \cos \theta_L \cos \theta_R).$$

(A.43)
The mass eigenvalues are given by

\[ M_{1,2} = \frac{1}{2} \sqrt{(M_\psi + M_\chi)^2 + \frac{1}{2} \sin \beta^2 v^2 (y_1 - \tilde{y}_1)^2 \pm \sqrt{(M_\psi - M_\chi)^2 + \frac{1}{2} v^2 \sin \beta^2 (y_1 + \tilde{y}_1)^2}. \] (A.44)

and the mixing angles \( \theta_{L,R} \) are given by

\[
\begin{align*}
\tan 2\theta_L &= \frac{2\sqrt{2} v \sin \beta (y_1 M_\chi + \tilde{y}_1 M_\psi)}{2(M_\psi^2 - M_\chi^2) - v^2 \cos \beta^2 (\tilde{y}_1^2 - y_1^2)}, \\
\tan 2\theta_R &= \frac{2\sqrt{2} v \sin \beta (y_1 M_\chi + \tilde{y}_1 M_\psi)}{2(M_\psi^2 - M_\chi^2) + v^2 \cos \beta^2 (\tilde{y}_1^2 - y_1^2)}. 
\end{align*}
\] (A.45, A.46)

### A.5 Couplings and mass eigenvalues for MVQU\(_{12}\) model

The couplings \( \kappa_{ij} \) for MVQU\(_{12}\) model are given by

\[
\kappa_{11} = \frac{g}{c_W} \left( \frac{1}{3} s^2_W - \frac{1}{2} \cos^2 \theta_L \right) + \frac{g}{c_W} \left( \frac{1}{3} s^2_W - \frac{1}{2} \cos^2 \theta_L \right), \\
\kappa_{12} = \frac{g}{c_W} \left( \frac{1}{3} s^2_W - \frac{1}{2} \cos^2 \theta_L \right) + \frac{g}{c_W} \left( \frac{1}{3} s^2_W - \frac{1}{2} \cos^2 \theta_L \right), \\
\kappa_{L,R12} = \frac{g}{2c_W} \sin \theta_{L,R} \cos \theta_{L,R}. 
\] (A.47, A.48, A.49)

The couplings \( y_{ij} \) are given by

\[
\begin{align*}
y_{11} &= \frac{1}{\sqrt{2}} (y_1 \sin \beta \cos \theta_L \sin \theta_R - \tilde{y}_1 \cos \beta \sin \theta_L \cos \theta_R), \\
y_{22} &= -\frac{1}{\sqrt{2}} (y_1 \sin \beta \cos \theta_L \sin \theta_R - \tilde{y}_1 \cos \beta \sin \theta_L \cos \theta_R), \\
y_{12} &= \frac{1}{\sqrt{2}} (y_1 \sin \beta \cos \theta_L \cos \theta_R + \tilde{y}_1 \cos \beta \sin \theta_L \sin \theta_R), \\
y_{21} &= -\frac{1}{\sqrt{2}} (y_1 \sin \beta \sin \theta_L \sin \theta_R + \tilde{y}_1 \cos \beta \cos \theta_L \cos \theta_R). 
\end{align*}
\] (A.50, A.51, A.52, A.53)

The mass eigenvalues are given by

\[ M_{1,2} = \frac{1}{2} \sqrt{(M_\psi + M_\chi)^2 + \frac{1}{2} v^2 (y_1 \cos \beta - \tilde{y}_1 \sin \beta)^2 \pm \sqrt{(M_\psi - M_\chi)^2 + \frac{1}{2} v^2 (y_1 \cos \beta + \tilde{y}_1 \sin \beta)^2}. \] (A.54)

and the mixing angles \( \theta_{L,R} \) are given by

\[
\begin{align*}
\tan 2\theta_L &= \frac{2\sqrt{2} v (y_1 \cos \beta M_\chi + \tilde{y}_1 \sin \beta M_\psi)}{2(M_\psi^2 - M_\chi^2) - v^2 (\tilde{y}_1^2 \sin \beta^2 - y_1^2 \cos \beta^2)}, \\
\tan 2\theta_R &= \frac{2\sqrt{2} v (y_1 \cos \beta M_\chi + \tilde{y}_1 \sin \beta M_\psi)}{2(M_\psi^2 - M_\chi^2) + v^2 (\tilde{y}_1^2 \sin \beta^2 - y_1^2 \cos \beta^2)}. 
\end{align*}
\] (A.55, A.56)

### A.6 \( \kappa_{AVV} \)'s for MVQD\(_{11}\), MVQU\(_{22}\), MVQU\(_{12}\) and MVQDX\(_{11}\) models

This section contains general expressions for \( \kappa_{AVV} \) (defined in Eq. 1) applicable to all of the four models, MVQD\(_{11}\), MVQU\(_{22}\), MVQU\(_{12}\) and MVQDX\(_{11}\).

\[
\begin{align*}
\kappa_{AZ\gamma} &= 4N_c (\kappa_{1Z\gamma} + \kappa_{2Z\gamma} + \kappa_{AZ\gamma} + \kappa_{AZ\gamma}), \\
\kappa_{AY\gamma} &= 4N_c (\kappa_{1Y\gamma} + \kappa_{2Y\gamma}), \\
\kappa_{A\gamma g} &= 4(\kappa_{1\gamma g} + \kappa_{2\gamma g}).
\end{align*}
\] (A.57, A.58, A.59)
and another heavy quark
After EWSB the top quark mixes with the VLF giving rise to two mass eigen states, the physical top

generation quarks as, MVQU
an example we choose
cases are given in App. A.3, A.4, A.5.

In this section we describe a simple model where the SM top mixes with a top-like VLQ. We only wish
The SM top yukawa coupling is given by,

\[
\lambda = \kappa_1 \sin \theta_L + \kappa_2 \sin \theta_R.
\]

with, \( \lambda_Z = \frac{m_Z^2}{m_t^2}, \lambda_1 = \frac{m_1^2}{m_t^2}, k_i = Q_i \) (electric charge of \( \zeta_i \)) and the couplings \( \kappa_{ij}, y_{ij} \) for each of the three
cases are given in App. A.3, A.4, A.5.

B A model with \( t' \leftrightarrow t \) mixing

In this section we describe a simple model where the SM top mixes with a top-like VLQ. We only wish
to estimate the possible corrections to \( \kappa_{AVV} \) and \( \kappa_{HVV} \) obtained in Sec. 3, because of the mixing. As
an example we choose MVQU22 model with \( Y_{\psi} = 1/6 \) and allow a mixing of the VLFs and the third
generation quarks as,

\[
\mathcal{L} \supset - y_t \bar{Q}_L \Phi_2 \xi_R + h.c.
\]

After EWSB the top quark mixes with the VLF giving rise to two mass eigen states, the physical top \( t \)
and another heavy quark \( t' \). The mass eigenvalues are given by

\[
m_t = \frac{1}{\sqrt{2}} y_t v \sin \beta \cos \theta_L^t \cos \theta_R^t + \frac{1}{\sqrt{2}} y_t' v_2 \sin \theta_L^t \sin \theta_R^t + M_\xi \sin \theta_L^t \sin \theta_R^t
\]

\[
m_{t'} = \frac{1}{\sqrt{2}} y_t v \sin \beta \sin \theta_L^t \sin \theta_R^t - \frac{1}{\sqrt{2}} y_t' v_2 \sin \theta_L^t \cos \theta_R^t + M_\xi \cos \theta_L^t \cos \theta_R^t,
\]

where \( \theta_L \) and \( \theta_R \) are the mixing angels which are given by,

\[
\tan(\theta_L^t + \theta_R^t) = \frac{y_t'/\sqrt{2}}{y_t v/\sqrt{2} \sin \beta - M_\xi}
\]

\[
\tan(\theta_L^t - \theta_R^t) = -\frac{y_t'/\sqrt{2}}{y_t v/\sqrt{2} \sin \beta + M_\xi}
\]

The SM top yukawa coupling is given by,

\[
y_t|_{\text{physical}} = \frac{1}{v}(m_t - M_\xi \sin \theta_L^t \sin \theta_R^t).
\]
We must fix the physical top mass at 173 GeV. Also $y_t|_{\text{physical}}$ has to be reasonably close to the SM value to make the theory consistent with the observed $\sigma(h) \times BR(h \rightarrow \gamma\gamma)$, which again reinforces the smallness of the mixing angles.

Now, $At^L$ coupling in 2HDM-II is modified as,

$$y_{tA} = \frac{m_t}{v} \cot \beta - \left( \frac{M_2}{v} \cot \beta + y_t' \cos \beta \right) \sin \theta_L^t \sin \theta_R^t. \tag{B.73}$$

The modification is very small when the mixing is small. The contribution to $At't', Ht't'$ coupling due to the mixing is,

$$y_{vA}, y_{vH} = \frac{1}{\sqrt{2}} y_t' \sin \theta_L^t \cos \theta_R^t, \tag{B.74}$$

which are also small for small mixing. As a result the corrections to $\kappa_{AVV}, \kappa_{HVV}$, due to the mixing will be very small.

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