Radiative decays of a singlet scalar boson through vector-like quarks

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Abstract: If the standard model Higgs boson were much heavier, it would appear as a broad resonance since its decay into a pair of longitudinally polarized gauge bosons is highly enhanced. We study whether the same enhancement happens at loop level in a simple extension of the standard model with a singlet scalar boson $S$. In order to focus on the loop effects, we assume that $S$ does not interact with the standard model particles at tree level. The singlet scalar $S$ is linked to the standard model world by vector-like quarks running in the loop. We introduce three vector-like quark multiplets, an $SU(2)_L$ doublet, an up-type singlet, and a down-type singlet. There are two kinds of loop effects in the $S$ phenomenology, the mixing with the Higgs boson and the radiative decays into $hh$, $WW$, $ZZ$, $gg$, and $\gamma\gamma$ through the triangle loops. We show that the crucial condition for enhancing loop effects including the longitudinal polarization enhancement is the large mass differences among vector-like quarks. The current LHC constraints on $S$ from the heavy scalar searches and the Higgs precision data are shown to be very significant: the mixing angle with the Higgs boson should be smaller than about 0.1 for $m_S = 750$ GeV.
1 Introduction

The Standard Model (SM) has been more solidified by the Large Hadron Collider (LHC) data at $\sqrt{s} = 13$ TeV. Supersymmetry models, composite Higgs models, and other SM extensions are all strongly constrained. Despite the absence of new signals, we believe that new physics beyond the SM must exist since the issue of naturalness and the existence of dark matter cannot be explained within the SM. There are two kinds of strategies to confront no signals, pushing new particles out of the LHC reach [1] or introducing a hidden sector [2, 3]. If either is the case, we turn to radiative corrections mediated by new particles or the linking to the hidden sector [4] in order to find a clue about new physics.

An interesting question is how significantly the radiative correction or the linking changes tree level results. They generally cause subleading corrections, but there exist extreme cases also. We may see entirely new signals which are absent at tree level, such as the flavor changing neutral current processes through loops and the invisible Higgs decay modes through the mixing with a hidden sector. In this work, we investigate a possibility that such new signals are so significant that they could be useful in the search for a new heavy scalar boson.

Among various tree level results, the decay of a heavy scalar boson shows an intriguing and unique feature, the longitudinal polarization enhancement in its decay into a massive gauge boson pair. It is well-known that if the SM Higgs boson $h$ were much heavier, it would have decayed dominantly into a heavy gauge boson pair, $VV$ ($V = W, Z$). The $t\bar{t}$ channel, the next dominant one, has the maximum branching ratio about 19%. The extraordinarily large $\Gamma(h \to VV)$ when $m_h \gg m_V$ is due to its decay into the longitudinal modes, $V_LV_L$ [5, 6]. The longitudinal polarization vector of $V$ is proportional to $p_V^L/m_V$ in the high energy limit, which leads to $\Gamma(h \to V_LV_L) \propto m_h^3$. The heavier the Higgs boson
is, the larger the decay rate into $V_LV_L$ becomes. For instance, $\Gamma(h \to V_LV_L)$ is about 99% of $\Gamma(h \to VV)$ when $m_h \simeq 440$ GeV. Accordingly its total decay rate is also enhanced so that a heavy Higgs boson becomes a broad resonance.

We wonder whether the same thing happens when a new heavy scalar boson decays only radiatively. To answer this question, we consider a simple extension of the SM where a singlet scalar $S$ [7] and vector-like quarks (VLQs) [8] are introduced. This can be considered as a simplified model. Theoretically, a singlet scalar has drawn a lot of interest in the context of Higgs portal models [2]. Its phenomenological signatures have been extensively studied [9–14]. Heavy VLQs are also interesting as they appear around the electroweak scale in many new physics models [15, 16]. The VLQs are compatible with the current experimental results while new heavy SM-like chiral quarks are excluded by the measurement of Higgs production rates [17, 18]. Moreover, the combination of a singlet scalar and vector-like quarks is attractive: it can shift the metastability of the electroweak vacuum in the SM [19–23]; it is crucial to construct a model where all of the gauge and Yukawa couplings remain asymptotically safe up to infinite energy [24, 25].

In order to focus on the role of radiative corrections exclusively, we consider a limiting scenario where $S$ does not couple to the SM particles at tree level. The VLQs play the role of messengers as connecting the SM particles and $S$ at loop level, as interacting with the singlet scalar $S$, the SM Higgs boson, and the SM gauge bosons. In order to allow the Yukawa interactions of VLQs with the Higgs boson, we introduce three VLQ multiplets, an $SU(2)_L$ doublet, and two (up-type and down-type) $SU(2)_L$ singlets. The presence of multiple VLQs shall be shown crucial in the $S$ phenomenology. The VLQ loops have two kinds of implications. First $S$ decays into $gg$, $\gamma\gamma$, $WW$, $ZZ$, and $hh$ through triangle VLQ loops. The singlet scalar $S$ can be produced and probed at high energy colliders. Secondly, $S$ is radiatively mixed with the Higgs boson. Naive expectation is that the heavier the VLQs are, the smaller the loop corrections become. We shall show that this is not true. Large mass differences in the VLQ mass spectrum induce the longitudinal polarization enhancement and increase the $S$-$h$ radiative mixing. The obtained condition for the enhancement at loop level shall help to study the physical properties of new particles running in the loop. These are our main results.

The paper is organized as follows. In Sec. 2, we provide the general helicity amplitude framework for the decay of a scalar boson into a massive gauge boson pair and into a Higgs boson pair. Section 3 summarizes our new physics model with a singlet scalar $S$ and VLQs. The gauge and Yukawa couplings of the VLQs in terms of the mass eigenstates are given. In Sec. 4, we present our main analytic results of loop calculations. The radiatively generated $S$-$h$ mixing and the loop induced decay rates of $S \to hh, VV$ are to be shown. In particular, the asymptotic behaviors of the loop functions are very useful to understand the enhancement of $\Gamma(S \to hh, VV)$ by large mass differences of the VLQs. Section 5 is devoted to our numerical results in a simple benchmark scenario. The general physical properties of $S$ such as its branching ratio and total decay rate are studied. We also calculate the exclusion limits from the current LHC data of the heavy scalar searches and the Higgs precision observation. The future prospect at the 13 TeV LHC is also discussed. Section 6 contains our conclusions.
2 Decays of a scalar boson into $VV$ and $hh$

We consider a $J^{PC} = 0^{++}$ scalar particle $S$ which has a mass $m_S$ and a momentum $p^\mu$. In the CP-conserving framework, the most general coupling of $S$ to a pair of gauge bosons and that to a pair of the Higgs bosons can be parameterized by

$$S(p)V_\mu(p_1)V'_\nu(p_2) : m_S \left[ A g_{\mu\nu} + B \frac{p_{2\mu}p_{1\nu}}{m_S^2} \right],$$

$$S hh : m_S C,$$

where $A$, $B$, and $C$ are dimensionless.

We write the helicity amplitudes for the decay $S \rightarrow VV'$ as

$$\langle V_\mu(p_1, \lambda_1)V'_\nu(p_2, \lambda_2) | S(P) \rangle \equiv m_S T_{\lambda_1, \lambda_2},$$

where $\lambda_1$ and $\lambda_2$ are the helicities of the outgoing gauge bosons. The dimensionless amplitudes $T_{\lambda_1, \lambda_2}$’s are then written in terms of $A$ and $B$ in Eq. (2.1) as [26]

$$T_{++} = T_{--} = -A,$$

$$T_{00} = \begin{cases} \frac{m_S^2}{4m_V^2}(2A + B) - (A + B), & \text{if } m_V \equiv m_{V_1} = m_{V_2} \neq 0; \\ 0, & \text{if } m_{V_1} = 0 \text{ or } m_{V_2} = 0, \end{cases}$$

and the other helicity amplitudes are zero. The partial decay rates are

$$\Gamma(S \rightarrow VV') = \frac{1}{S} \frac{\beta_{VV'}}{16\pi} m_S \sum_{\lambda_1, \lambda_2} |T_{\lambda_1, \lambda_2}|^2,$$

$$\Gamma(S \rightarrow hh) = \frac{\beta_{hh}}{32\pi} m_S |C|^2,$$

where the symmetric factor $S$ is $1/2$ for two identical outgoing particles, and $\beta_{ij} = \sqrt{1 - 2m_i^2 + m_j^2}/m_S^2 + (m_i^2 - m_j^2)/m_S^4$.

When a scalar particle is heavy enough, its decay into a massive gauge boson pair $VV$ ($V = W^\pm, Z$) has a special feature. The condition $m_S \gg m_V$ makes $T_{00}$ greatly enhanced if $(2A + B) \neq 0$. The SM Higgs boson, if its mass is greater than $2m_V$, has

$$A^{SM} = \frac{2m_V^2}{v m_h}, \quad B^{SM} = 0.$$

The partial decay rate of $h_{SM} \rightarrow V_L V_L$ is proportional to the cube of $m_h$ while that of $h_{SM} \rightarrow V_T V_T$ is inversely proportional to $m_h$. The heavier the Higgs boson is, the more dominant $h \rightarrow V_L V_L$ will become. Another significant decay rate $\Gamma(h \rightarrow t\bar{t})$ is linearly proportional to $m_h$. The Higgs boson decay into $V_L V_L$ is dominant in the total decay rate. This is called the longitudinal polarization enhancement.

The partial decay rate of $S$ into a pair of SM Higgs bosons is sizable if $C \sim \mathcal{O}(1)$. In the MSSM, an obvious scalar boson candidate which decays into $hh$ is the heavy CP-even
Higgs boson $H$. However, the decay into a pair of light Higgs bosons is suppressed in the alignment limit since $C$ is [27]

$$C_{\text{MSSM}}^{H} = -\frac{3g_2^2 \sin 4\beta}{8} \frac{v}{M_H} [1 + \mathcal{O}(\cos(\beta - \alpha))],$$

(2.6)

where $g_2 = g / \cos \theta_W$ and $\theta_W$ is the weak mixing angle. The partial decay rate $\Gamma(H \to hh)$ is inversely proportional to the heavy Higgs mass: there is no enhancement in the $hh$ decay channel.

### 3 Model with a singlet scalar and vector-like quarks

We consider a simple extension of the SM by introducing a CP-even singlet scalar boson $S_0$, a VLQ doublet $Q_{L/R}$, two VLQ singlets $U_{L/R}$ and $D_{L/R}$:

$$Q_{L/R} = \left( \begin{array}{c} U' \\ D' \end{array} \right)_{L/R}, \quad U_{L/R}, \quad D_{L,R}.$$  

(3.1)

The $SU(3)_c \times SU(2)_L \times U(1)_Y$ quantum numbers of $Q_{L/R}$, $U_{L/R}$, and $D_{L/R}$ are $(3, 2, 1/3)$, $(3, 1, 4/3)$, and $(3, 1, -2/3)$, respectively. The hypercharges of VLQs are set to be the same as the SM quarks. Different assignment shall affect the decays of $S$ into $ZZ$ and $\gamma\gamma$.

The most general scalar potential of the SM Higgs doublet $H$ and a real singlet scalar $S_0$ is [28]

$$V(H, S_0) = -\mu^2 H^\dagger H + \lambda (H^\dagger H)^2 + \frac{a_1}{2} S_0 H^\dagger H + \frac{a_2}{2} S_0^2 H^\dagger H + b_1 S_0 + \frac{b_2}{2} S_0^2 + \frac{b_3}{3} S_0^3 + \frac{b_4}{4} S_0^4.$$  

(3.2)

Note that we do not assume any discrete $Z_2$ symmetry for $S_0$. When defining the neutral component of $H$ as $\phi_0 = (v_0 + h_0)/\sqrt{2}$ and the VEV of the singlet field as $\langle S_0 \rangle = x$, the extrema of the potential satisfy

$$\frac{\partial V(v_0/\sqrt{2}, x)}{\partial v_0} = 0, \quad \frac{\partial V(v_0/\sqrt{2}, x)}{\partial x} = 0.$$  

(3.3)

Although there exist many possible extrema, the true minimum of $H$ should generate proper EWSB, i.e., $v_0 = v = 246$ GeV. On the other hand, the VEV of $S$ is free to choose since the shift of the singlet field, $S \to S + \Delta S$, corresponds to redefining the parameters of $a_{1,2}$ and $b_{1,\ldots,4}$. There is no change in physics. Without loss of generality we take $(v_0, x) = (v, 0)$. Note that the choice of vanishing VEV for $S_0$ eliminates the tadpole term of $S_0$. The minimization conditions in Eq. (3.3) become

$$\mu^2 = \lambda v^2, \quad b_1 = -\frac{v^2}{4} a_1.$$  

(3.4)

The Yukawa terms of VLQs with the singlet $S_0$ and the SM Higgs doublet $H$ as well as their mass terms are

$$\mathcal{L}_Y = S_0 \left[ y_Q \bar{Q} Q + y_R \bar{R} U + y_D \bar{D} D \right] + M_Q \bar{Q} Q + M_R \bar{R} U + M_D \bar{D} D$$

$$+ \left[ Y_Q \bar{Q} L H D_R + Y_D' \bar{Q} R H D_L + Y_U \bar{U} \tilde{H} U_R + Y_D' \bar{D} \tilde{H} D_L + H.c. \right],$$  

(3.5)
where $\tilde{H} = i \tau_2 H^*$. For simplicity, we assume $y_Q = y_U = y_D \equiv y_S$, $Y_U = Y_{U^c}$, and $Y_D = Y_{D^c}$ in what follows.

The VLQ mass matrix $M_F$ in the basis of $(F', F)$ where $F = U, D$ is

$$M_F = \left( \begin{array}{cc} M_Q & Y_F \sqrt{2} \\ Y_{F'} \sqrt{2} & M_F \end{array} \right),$$

which is diagonalized by the mixing matrix of

$$R_{\theta_F} = \left( \begin{array}{cc} c_{\theta_F} & -s_{\theta_F} \\ s_{\theta_F} & c_{\theta_F} \end{array} \right).$$

Here we adopt simplifying notations of $c_x = \cos x$ and $s_x = \sin x$. The $Y_U$ and $Y_D$ terms generate the mixings between VLQ doublet and VLQ singlets. If $Y_F \sim O(1)$, VLQ mixing angles are small since VLQs are expected to be heavy. The mass eigenvalues and the mixing angle are then

$$M_{F_1, F_2} = \frac{1}{2} \left[ M_Q + M_F \mp \sqrt{(M_F - M_Q)^2 + 2Y_F^2 v^2} \right],$$

$$s_{2\theta_F} = \frac{\sqrt{2}Y_F v}{M_{F_2} - M_{F_1}},$$

where $M_{F_1} < M_{F_2}$.

The Yukawa terms of the VLQ mass eigenstates become

$$-L_{\text{Yukawa}} = y_S S_0 \sum_i \left[ U_i U_i + D_i D_i \right] + h_0 \sum_i y_{h_{F_i F_j}} \hat{F}_i F_j,$$

where $F = U, D, i, j = 1, 2$, and $y_{h_{F_i F_j}}$ are

$$y_{h_{F_1 F_1}} = -y_{h_{F_2 F_2}} = -\frac{Y_F}{\sqrt{2}} s_{2\theta_F}, \quad y_{h_{F_1 F_2}} = y_{h_{F_2 F_1}} = -\frac{Y_F}{\sqrt{2}} c_{2\theta_F}.$$

The gauge interaction Lagrangian in terms of the VLQ mass eigenstates is

$$L_{\text{gauge}} = e A_\mu \sum_F \sum_i Q_F \hat{F}_i \gamma^\mu F_i + g Z Z_\mu \sum_F \sum_{i,j} \hat{g}_{Z_{F_i F_j}} \hat{F}_i \gamma^\mu F_j$$

$$+ \frac{g}{\sqrt{2}} \left[ W^{+\mu} \sum_{i,j} \hat{g}_{W_{U_i D_j}} U_i \gamma^\mu D_j + H.c. \right].$$

Here $Q_F$ is the electric charge of the fermion $F$ and the effective gauge couplings $\hat{g}_{V_{F'F}}$ are

$$\hat{g}_{Z_{F_1 F_1}} = g_{Q}^2 c_{\theta_F}^2 + g_{F}^2 s_{\theta_F}^2, \quad \hat{g}_{Z_{F_2 F_2}} = g_{Q}^2 s_{\theta_F}^2 + g_{F}^2 c_{\theta_F}^2,$$

$$\hat{g}_{Z_{F_1 F_2}} = \left( g_{Q}^2 - g_{F}^2 \right) s_{\theta_F} c_{\theta_F},$$

$$\hat{g}_{W_{U_1 D_1}} = c_{\theta_F} c_{\theta_D}, \quad \hat{g}_{W_{U_1 D_2}} = c_{\theta_F} s_{\theta_D},$$

$$\hat{g}_{W_{U_2 D_1}} = s_{\theta_F} c_{\theta_D}, \quad \hat{g}_{W_{U_2 D_2}} = s_{\theta_F} s_{\theta_D},$$

- 5 -
where \( \hat{g}_{VFF} = \hat{g}_{V'F} \) and \( \bar{g}_V^F = \frac{1}{2} T_3^F - s_W^2 Q_F \) for \( F = Q, U, D \). There is a big difference between \( h-F-F' \) couplings and \( V-F-F' \) couplings. In the limit of \( \theta_{U,D} \ll 1 \), the gauge couplings to different mass eigenstates of VLQs (e.g. \( \hat{g}_{VFI_1I_2} \)) are suppressed by \( s_{\theta_F} \). On the contrary, the VLQ couplings to the Higgs boson are suppressed for the same mass eigenstates.

Without the \( Z_2 \) symmetry, the \( S_0 \) field can couple to the SM particles at tree level. Since the singlet scalar \( S_0 \) is neutral under all quantum numbers of the SM gauge group, the only possible renormalizable couplings of \( S_0 \) to the SM particles at tree level are to the Higgs boson through \( a_1 \) and \( a_2 \) terms in Eq. (3.2). However, a nonvanishing \( a_1 \) term will generate the \( S-h \) mixing with the mixing angle \( \eta \), which shall change the Higgs coupling modifiers of \( \kappa_V \) and \( \kappa_f \) into \( c_\eta \). According to the global fit analysis of the LHC Higgs precision data [29–31], \( c_\eta \) is very close to 1. Nonzero \( a_1 \) builds up some tension with the Higgs boson constraints. Moreover, our main question is whether the unique characteristic of a heavy scalar boson such as the longitudinal polarization enhancement remains even at loop level. Therefore, we consider a limiting scenario in which the singlet scalar has no tree level couplings with the Higgs boson:

\[
a_1^{\text{tree}} = 0 = a_2^{\text{tree}}. \tag{3.13}
\]

### 4 The effects of the VLQ loops

In the previous section, we suggested a rather extreme scenario where \( S_0 \) does not interact with the SM Higgs boson at tree level. The singlet field \( S_0 \) could be considered as a field in a hidden sector. In the model, the visible sector and the hidden sector are connected via VLQ loops: the VLQs play the role of messengers. There are two phenomenological implications: (i) the singlet-Higgs mixing and (ii) the radiative decays of \( S \) into SM particles. We study the effects at one loop level.

#### 4.1 \( S-h \) mixing and Higgs Modifiers

![Figure 1](image_url)

**Figure 1**: Feynman diagrams for the loop induced \( S-h \) mixing.

First, the \( S-h \) mixing is radiatively generated through the VLQ loops as shown in Fig. 1. The scalar-mass-squared matrix in the basis of \( (h_0, S_0) \) becomes

\[
M_{hS}^2 \equiv \begin{pmatrix}
2\lambda v^2 & \delta M^2_{Sh} \\
\delta M^2_{Sh} & M^2_{SS}
\end{pmatrix}, \quad \tag{4.1}
\]
where $M_{SS}^2 = b_2$ since we have used the conditions in Eq. (3.4) for our choice of the vacuum $(v_0, x) = (v, 0)$. At one loop level, we have

$$
\delta M_{Sh}^2 = -\frac{y_S N_c}{4\pi^2} \sum_F \sum_i y_{h,F_i} M_F^2 \left[ 4(\tau_{F_i}^S - 1)g(\tau_{F_i}^S) - 4\tau_{F_i}^S + 5 \right],
$$

where $N_c = 3$ is the color factor of the VLQ, $F = U, D$, $i = 1, 2$, $\tau_j^F = m_j^2/(4m_j^2)$, and $y_{h,F_i}$’s are in Eq. (3.10). The loop function $g(\tau)$ is given by

$$
g(\tau) = \begin{cases} \sqrt{\tau^{-1} - 1} \arcsin \sqrt{\tau} & \text{if } \tau \leq 1; \\
\frac{\sqrt{1-\tau^{-1}}}{2} \left[ \log \frac{1+\sqrt{1-\tau^{-1}}}{1-\sqrt{1-\tau^{-1}}} - i\pi \right] & \text{if } \tau > 1.
\end{cases}
$$

Note that $\delta M_{Sh}^2$ vanishes if $M_{F_1} = M_{F_2}$ since $y_{h,F_1}F_1 = -y_{h,F_2}F_2$; see Eq. (3.10). Significant $S$-$h$ mixing requires sizable mass differences of $F_1$ and $F_2$.

The mass eigenvalues and the $S$-$h$ mixing angle $\eta$ are

$$
m_{h,S}^2 = \frac{1}{2} \left( M_{hh}^2 + M_{SS}^2 \mp \sqrt{(M_{SS}^2 - M_{hh}^2)^2 + 4(\delta M_{Sh}^2)^2} \right),
$$

$$s_{2\eta} = -\frac{2\delta M_{Sh}^2}{m_S^2 - m_h^2},
$$

where we use the $S$-$h$ mixing matrix $\mathbb{R}_{\eta}$ in Eq. (3.7). Since $\delta M_{Sh}^2$ is radiatively generated, we expect $s_\eta \ll 1$. We take the mass eigenstate $h = c_\eta h_0 - s_\eta S_0$ to be the observed Higgs boson with a mass of 125 GeV, and assume $S$ to be heavy such as $m_S \gtrsim 500$ GeV.

The nonzero $S$-$h$ mixing changes the Higgs coupling modifiers of $\kappa_Z$, $\kappa_W$, $\kappa_t$, $\kappa_\tau$, and $\kappa_b$ to be $c_\eta$. The loop induced decays of the Higgs boson into $gg$ and $\gamma\gamma$ have additional loop contributions from VLQs. We define $\kappa_g$ and $\kappa_\gamma$ as

$$
\mathcal{L}_{\text{Higgs}} = \kappa_g c_g^{SM} \frac{h}{v} G^{\mu\nu} G_{\mu\nu}^{SM} + \kappa_\gamma c_\gamma^{SM} \frac{h}{v} F^{\mu\nu} F_{\mu\nu}.
$$

The SM values $c_g^{SM}$ and $c_\gamma^{SM}$ are

$$
c_g^{SM} = \frac{\alpha_s}{16\pi} A_{hgg}^{SM}, \quad c_\gamma^{SM} = \frac{\alpha_e}{8\pi} A_{h\gamma\gamma}^{SM},
$$

where $A_{hgg}^{SM} = \sum_{f=t,b} A_{1/2}^{SM}(\tau_f^h)$, $A_{h\gamma\gamma}^{SM} = A_1(\tau_W^h) + \sum_{f=t,b,\tau} N_{Q_f}^2 Q_f^2 A_{1/2}^{SM}(\tau_f^h)$, and $A_{1/2}(\tau)$ and $A_1(\tau)$ are referred to Ref. [32]. The modifiers $\kappa_g$ and $\kappa_\gamma$ receive two kinds of new contributions. One is from the modified couplings of $h$ to the SM particles through the $S$-$h$ mixing. The other is from the triangle VLQ loops, parameterized by

$$
A_{hgg}^{VLQ} = \sum_F \sum_i y_{h,F_i} \frac{v}{M_F} A_{1/2}(\tau_F^h),
$$

$$
A_{h\gamma\gamma}^{VLQ} = \sum_F \sum_i N_{Q_f} Q_f^2 y_{h,F_i} \frac{v}{M_F} A_{1/2}(\tau_F^h),
$$

1 There is UV divergence in the one loop calculation of $\delta M_{Sh}^2$ which must be properly renormalized. Detailed description on the renormalization of the whole model is in preparation [33].
\begin{align*}
\kappa_{g,\gamma} &= c_\eta A_{hgg,h\gamma\gamma}^{SM} + A_{hgg,h\gamma\gamma}^{VLQ}.
\end{align*}

Figure 2: Feynman diagrams of $S \rightarrow hh$ and $S \rightarrow WW$ from the VLQ loops.

where $F = U, D$, $i = 1, 2$, $\tau_j^i = m_j^2/(4m_i^2)$. Then $\kappa_g$ and $\kappa_\gamma$ are

\begin{align*}
\kappa_{g,\gamma} &= c_\eta A_{hgg,h\gamma\gamma}^{SM} + A_{hgg,h\gamma\gamma}^{VLQ}.
\end{align*}

Since $y_{hF_1F_i}$ and $y_{hF_2F_j}$ in Eq. (3.10) have opposite signs, both $\delta M_S^2$ as well as $A_{hgg,h\gamma\gamma}^{VLQ}$ are suppressed when $M_{F_1} \simeq M_{F_2}$.

4.2 Radiative Decays of S

Another important effect of the VLQ loops is the radiative decay of $S$ into the SM particles, which occurs through the $S$-$h$ mixing as in Fig. 1 and/or through the triangle VLQ loops into a gauge boson pair or a Higgs boson pair as in Fig. 2. Since we consider the case of $m_S \gtrsim 500$ GeV, the main decay modes are into $t\bar{t}$, $gg$, $\gamma\gamma$, $WW$, $ZZ$, and $hh$.

The decay of $S$ into a top quark pair is only through the $S$-$h$ mixing. The partial decay rate is

\begin{align*}
\Gamma(S \rightarrow t\bar{t}) &= s_\eta^2 \Gamma(h_{SM} \rightarrow t\bar{t}) \bigg|_{m_{h_{SM}} = m_S}.
\end{align*}

Another important decay channel is $S \rightarrow hh$ shown in Fig. 2(a). The vertex $C$ in Eq. (2.1) at one loop level is

\begin{align*}
C &= \frac{ys_N}{4\pi^2} \sum_F \sum_{i,j} y_{hF_1F_j}^2 C_T(m_h, m_S, M_{F_i}, M_{F_j}) + \frac{3m_h^2}{v m_S} s_\eta,
\end{align*}

where $y_{hF_iF_j}$ are given in Eq. (3.10). The first term is due to the triangle diagrams while the second one is from the $S$-$h$ mixing. The asymptotic expression\footnote{Full expressions of form factors such as $A_T$, $B_T$, and $C_T$ are to be presented in Ref. [33].} of $C_T$ when $m_h \ll m_S$ and $\Delta F \ll M_F$, where $\Delta F = M_{F_j} - M_{F_i}$ and $M_F = (M_{F_i} + M_{F_j})/2$, is very useful to understand the enhancement of $\Gamma(S \rightarrow hh)$ in some parameter space:

\begin{align*}
\sqrt{\tau} C_T &= 2 + (1 - 2\tau^{-1})f(\tau) - 2g(\tau) \\
&\quad + \left(\frac{\Delta F^2}{M_F^2}\right) \left[8\tau^2 + 49\tau - 48 \right] + \left(\frac{(\tau^3 + 12\tau^2 - 26\tau + 16)g(\tau)}{4\tau(1-\tau)^2} \right] \\
&\quad + \left(\frac{m_h^2}{m_S^2}\right) \left[\frac{2(6 - \tau)}{3} + \frac{2(2\tau - 2)f(\tau)}{\tau} \right] + O\left(\frac{\Delta F^4}{M_F^4}\right) + O\left(\frac{m_h^4}{m_S^4}\right).
\end{align*}
where \( \tau = m_2^2/(4M_F^2) \) and \( f(\tau) \) is referred to Ref. [32]. Note that the odd power terms in \((\Delta F/M_F)\) are neglected since they cancel each other after the summation in Eq. (4.10). If \( y_S, Y_{U,D} \sim O(1) \), \( \mathcal{C} \) is not suppressed by large \( m_S \), contrary to the case of a heavy CP-even scalar \( H \) of the MSSM in Eq. (2.6). Another important result is that the partial decay rate \( \Gamma(S \to hh) \) increases with \( \Delta F \), the mass difference between \( M_{F_i} \) and \( M_{F_j} \). Since \( \Delta F \) is proportional to the Higgs VEV from the SM-like Yukawa couplings of the VLQs to the Higgs boson, the enhancement of \( S \to hh \) can be considered as non-decoupling effects.

The VLQ loops also allow the decay of \( S \) into a massive gauge boson pair \( VV \) \((V = W, Z)\) as shown in Fig. 2(b). The dimensionless parameters \( \mathcal{A} \) and \( \mathcal{B} \) in Eq. (2.1) are

\[
\mathcal{A}_{WW} = \frac{g^2 y_S N_c}{8\pi^2} \sum_{i,j} \left[ \hat{g}_{WU,V}^2 \mathcal{A}_T(m_W, m_S, M_{U_i}, M_{D_j}) + \{U \leftrightarrow D\} \right] + \frac{2m_W^2}{v m_S} s_\eta, \\
\mathcal{B}_{WW} = \frac{g^2 y_S N_c}{8\pi^2} \sum_{i,j} \left[ \hat{g}_{WU,V}^2 \mathcal{B}_T(m_W, m_S, M_{U_i}, M_{D_j}) + \{U \leftrightarrow D\} \right], \\
\mathcal{A}_{ZZ} = \frac{g^2 y_S N_c}{4\pi^2} \sum_{i,j} \left[ \hat{g}_{ZU,V}^2 \mathcal{A}_T(m_Z, m_S, M_{U_i}, M_{D_j}) + \{U \leftrightarrow D\} \right] + \frac{2m_Z^2}{v m_S} s_\eta, \\
\mathcal{B}_{ZZ} = \frac{g^2 y_S N_c}{4\pi^2} \sum_{i,j} \left[ \hat{g}_{ZU,V}^2 \mathcal{B}_T(m_Z, m_S, M_{U_i}, M_{D_j}) + \{U \leftrightarrow D\} \right],
\]

where \( i, j = 1, 2 \). \( \mathcal{A}_{VV} \) consists of two parts, one from the triangle VLQ loops and the other from the \( S-h \) mixing, while \( \mathcal{B}_{VV} \) is only from the triangle loops.

Our main question is whether the longitudinal polarization enhancement in \( S \to VV \) remains significant at loop level, which happens when \( 2\mathcal{A} + \mathcal{B} \neq 0 \) as shown in Eq. (2.3). The \( S-h \) mixing induced terms, proportional to \( s_\eta \) in Eq. (4.12), appear only in \( \mathcal{A} \) and thus generate the longitudinal polarization enhancement. The condition that the triangle VLQ loops induce the enhancement is easy to see through the asymptotic behaviors of \( \mathcal{A}_T \) and \( \mathcal{B}_T \) in the limit of \( \Delta_F \ll M_F \) and \( m_V \ll m_S \), given by

\[
\sqrt{\tau} \mathcal{A}_T = 1 + (1 - \tau^{-1}) f(\tau) + \left( \frac{\Delta_F^2}{M_F^2} \right) \left[ -1 + \frac{3\tau - 4}{4\tau^2} f(\tau) - \frac{(\tau^2 + 4\tau - 8)g(\tau)}{4\tau(\tau - 1)} \right] + 2 \left( \frac{m_V^2}{m_S^2} \right) \left[ 3 - \tau - \tau^{-1} f(\tau) - 2g(\tau) \right] + O\left( \frac{\Delta_F^4}{M_F^4} \right) + O\left( \frac{m_V^4}{m_S^4} \right),
\]

\[
\sqrt{\tau} \mathcal{B}_T = -2 - 2(1 - \tau^{-1}) f(\tau) + \left( \frac{\Delta_F^2}{M_F^2} \right) \left[ \frac{5}{2} + \frac{(8 - 5\tau)f(\tau)}{2\tau^2} - \frac{(\tau^2 + 12\tau - 16)g(\tau)}{2\tau(\tau - 1)} \right] + 4 \left( \frac{m_V^2}{m_S^2} \right) \left[ \tau - 4 + (2 - \tau)\tau^{-1} f(\tau) + 2g(\tau) \right] + O\left( \frac{\Delta_F^4}{M_F^4} \right) + O\left( \frac{m_V^4}{m_S^4} \right),
\]

where \( \tau = m_2^2/(4M_F^2) \). Equations (4.13) and (4.14) show that \( 2\mathcal{A} + \mathcal{B} \sim O\left( \frac{m_V^2}{m_S^2} \right) \) if \( \Delta_F = 0 \). Sizable mass differences of VLQs are crucial for the longitudinal polarization enhancement through the triangle VLQ loops.
The last category of the radiative decays of $S$ is into $gg$, $\gamma\gamma$, and $Z\gamma$. When at least one of the outgoing gauge bosons is massless, there is no longitudinal polarization mode as shown in Eq. (2.3). The $A$’s are

$$A_{\gamma\gamma} = \frac{e^2y_SN_c}{4\pi^2} \sum_F \sum_i Q_{F_i}^2 \frac{1}{\sqrt{\tau_{F_i}}} \left[ 1 + (1 - \tau_{F_i}^{-1})f(\tau_{F_i}) \right],$$

$$A_{gg} = \delta^{ab}g_s^2y_S \sum_F \sum_i \frac{1}{\sqrt{\tau_{F_i}}} \left[ 1 + (1 - \tau_{F_i}^{-1})f(\tau_{F_i}) \right],$$

$$A_{Z\gamma} = \frac{eg_Zy_SN_c}{2\pi^2} \sum_F \sum_i Q_{F_i} \hat{g}_{ZF_i} \frac{1}{\sqrt{\tau_{F_i}}} \left[ -1 - (1 - \tau_{F_i}^{-1})f(\tau_{F_i}) + O \left( \frac{m_Z^2}{m_S} \right) \right],$$

where $a, b$ are color indices of the outgoing gluons, $F = U, D$, $i = 1, 2$, and $\tau_{F_i} = m_S^2/(4M^2_{F_i})$. The $B$’s can be obtained by using Ward identity as follows

$$B_{\gamma\gamma} = -2A_{\gamma\gamma}, \quad B_{gg} = -2A_{gg}, \quad B_{Z\gamma} = -2 \left( 1 - \frac{m_Z^2}{m_S^2} \right)^{-1} A_{Z\gamma}.$$  

The final comment in this section is the importance of the VLQ Yukawa couplings with the Higgs boson in enhancing the radiative decay rates of $S$. If we do not allow the $Y_U$ and $Y_D$ terms, which happens for example when we introduce only one VLQ multiplet, the $S-h$ mixing and the $S \rightarrow hh$ decay will be absent. In addition, the VLQs running in the triangle VLQ loops for the decay of $S \rightarrow WW, ZZ$ have the same masses because of no VLQ mixing. There is no longitudinal polarization enhancement and thus the signal rates of the radiative decays have typical loop suppression [34]. In summary, the presence of the VLQ doublet and the VLQ singlets are crucial for the enhanced radiative decays of $S$.

5 Numerical Results

The phenomenological characteristics of the singlet scalar $S$ depend on the model parameters of $y_S$, $m_S$, $Y_{U,D}$, $M_Q$, $M_U$ and $M_D$. The $y_S$ contributes equally to all of the partial decay rates of $S$ by the common factor of $y_S^2$ since $S$ decays only radiatively through VLQ loops in our model. The branching ratios of $S$ are independent of $y_S$. The $m_S$ dependence on the branching ratios is also weak for the heavy $S$. The $Y_{U,D}$, $M_Q$, and $M_{U,D}$ specify the VLQ mass matrices and thus the mass difference $\Delta_F$. Since $Y_U$ and $Y_D$ also quantify the VLQ couplings with the Higgs boson, they are the most crucial parameters.

Therefore, we consider a simple benchmark parameter line, given by

$$M_Q = M_U = M_D, \quad Y_U = 0, \quad Y_D \text{ varies.} \quad (5.1)$$

We found that the results in this simple case display the main characteristic features of the radiative decays of $S$. The VLQ mass spectra become

$$M_{U_1} = M_{U_2} = M_Q, \quad M_{D_{1,2}} = M_Q \mp \frac{1}{\sqrt{2}} |Y_D| v. \quad (5.2)$$
Figure 3: Branching ratios of the radiative decays of the singlet scalar $S$ with mass $m_S = 500, 750$ GeV as functions of $\Delta M_{U_1,D_1} (\equiv M_{U_1} - M_{D_1})$. For the VLQ masses we set the lightest VLQ mass as $M_{D_1} = 0.6 m_S$ and assume $M_Q = M_{U_1} = M_D$ and $Y_{U_1} = 0$ with varying $Y_D$.

Note that $D_1$ becomes the lightest VLQ and $\Delta M_{U_1,D_1} = \Delta M_{D_2,U_1} = (1/2) \Delta M_{D_2,D_1}$, where $\Delta M_{ij} \equiv M_i - M_j$. Our setting of $Y_{U_1} \neq Y_D$ generates a sizable mass difference $\Delta M_{U_1,D_1}$ which is essential for the longitudinal polarization enhancement of $S \rightarrow WW$.

Brief comments on the VLQ masses are in order here. The mass bounds on the VLQs from the direct searches at the Tevatron and the LHC depend sensitively on the decay channels of the VLQs. If the main decay mode includes the third generation quarks, the bounds are strong: $M_{VLQ} \lesssim 400–600$ GeV [35]. If VLQs mix only with lighter generations, the mass bounds become much less than 400 GeV [35], which is adopted here.

In Fig. 3, we present the branching ratios of the singlet scalar $S$ as functions of $\Delta M_{U_1,D_1}$, or equivalently of $Y_D$, along the benchmark parameter line. We consider two cases, $m_S = 500$ GeV and $m_S = 750$ GeV with $M_{D_1} = 0.6 m_S$. When $Y_D = 0$ ($Y_{U_1} = 0$ by setting), the dominant decay mode is into $gg$ with almost 100% branching ratio. The radiative decay of $S$ into $hh$ is certainly prohibited. In addition there is no radiatively generated $S$-$h$ mixing, i.e., $s_{\eta} = 0$, which forbids the decay of $S \rightarrow t \bar{t}$. The mixing induced decays in $S \rightarrow WW, ZZ$ are closed and only the triangle VLQ loop contributions become relevant. The next dominant decay mode is into $WW$ with very small branching ratio of the order of $10^{-3}$. This is because of the suppression of the longitudinal polarization enhancement since the $Y_D = 0$ condition makes all of the VLQ masses degenerate and thus $2\mathcal{A} + \mathcal{B} \sim \mathcal{O}(m_{Y_1}^2/m_S^2)$: see Eqs. (4.11), (4.13), and (4.14). The reason why $B(S \rightarrow WW)$ is much larger than $B(S \rightarrow ZZ)$ when $Y_{U_1,D} = 0$ is that the gauge couplings of VLQs to the $Z$ boson are smaller than those to the $W$ boson with our choice of the electric charges of VLQs. Note that $\Gamma(S \rightarrow WW) \gg \Gamma(S \rightarrow \gamma\gamma, ZZ)$ is generic in the view of high dimensional operators in the effective field theory [36].

As $Y_D$ increases, the decay modes into $hh$, $WW$, $ZZ$ and $t\bar{t}$ all become significant. For both $m_S = 500$ GeV and $m_S = 750$ GeV cases, the $hh$ mode is as important as the
**Figure 4:** Total decay rate of the singlet scalar $S$ as a function of $\Delta M_{U_1D_1}$ or $Y_D$ for $m_S = 500$ GeV and $m_S = 750$ GeV. We take the benchmark parameter line in Eq. (5.1).

For the $gg$ mode when $Y_D \simeq 0.8$, and dominant when $Y_D \gtrsim 0.9$, followed by the $WW$, $ZZ$, and $t\bar{t}$ modes. We found that the little hierarchy among $hh$, $WW$, $ZZ$ and $t\bar{t}$ modes is quite generic with more general parameter setup other than our benchmark scenario. In some extreme corners of the parameter space such as small $Y_{U,D}$ but large $\Delta_F$, the $WW$ decay mode is dominant.

In Fig. 4, we show the total decay rate of $S$ as a function of $\Delta M_{U_1D_1}$ for $m_S = 500$ GeV and $m_S = 750$ GeV. When $\Delta M_{U_1D_1} = 0$, $\Gamma_S^{\text{tot}} \sim 0.1$ GeV for both mass cases. The singlet scalar is a very narrow resonance. With increasing $\Delta M_{U_1D_1}$, $\Gamma_S^{\text{tot}}$ starts decreasing, which is expected since $U_{1,2}$ and $D_2$ become heavier with the fixed $M_{D_1}$ and thus make smaller loop corrections. When $\Delta M_{U_1D_1}$ is large enough, however, $\Gamma_S^{\text{tot}}$ turns to increase, reaching about $10$ GeV when $\Delta M_{U_1D_1} = 300$ GeV. The enhancement compared to the $\Delta M_{U_1D_1} = 0$ case is almost by two orders of magnitude. This is unexpected since the VLQ masses for $\Delta M_{U_1D_1} = 300$ GeV are much heavier than those for $\Delta M_{U_1D_1} = 0$. This shows how dramatical the enhancement of the radiative decays of $S$ can be when there exist sizable mass differences of the VLQs.

Figure 5 presents the 95% C.L. exclusion region in the $(\Delta M_{U_1D_1}, y_S)$ parameter plane by the LHC Higgs precision data as well as the heavy Higgs search results in the $WW$, $ZZ$, and $hh$ channels. We also show the contours for $s_\eta$ by dashed (orange) lines. For the Higgs precision data, we adopt the global fit results from the ATLAS/CMS combined analysis for $\kappa_V \leq 1$ [31]: $\kappa_V = 0.97 \pm 0.060$, $\kappa_g = 0.81^{+0.13}_{-0.10}$, and $\kappa_\gamma = 0.90^{+0.10}_{-0.09}$. Note that $\kappa_f = 0.87^{+0.12}_{-0.11}$ and $\kappa_b = 0.57^{+0.16}_{-0.16}$ are consistent within $2\sigma$ but $\kappa_t = 1.42^{+0.23}_{-0.22}$ shows some deviation. For heavy scalar boson searches with mass $m_S = 500$ (750) GeV, the observed 95% C.L. upper bounds on $\sigma \cdot B$ at $\sqrt{s} = 8$ TeV are $200 \text{ fb}$ ($40 \text{ fb}$) for $WW$ [37, 38], $43 \text{ fb}$ ($12 \text{ fb}$) for $ZZ$ [39], and $107.6 \text{ fb}$ ($34 \text{ fb}$) for $hh$ [40–42]. We found that the heavy scalar search channels of dijet [43, 44] and $W\gamma/Z\gamma$ [45] provide weaker constraints. We do not
consider the $t\bar{t}$ channel [46, 47] because the current bound ignores the interference with the continuum background, which can be very significant [48–50].

The Higgs precision data exclude large $\Delta M_{U_1,D_1}$, almost independently of $y_S$: $\Delta M_{U_1,D_1} \lesssim 200\,(300)$ GeV for $m_S = 500\,(750)$ GeV is allowed. This exclusion mainly comes from the constraint on $\kappa_g$ of which the deviation from the SM value is generated from the $S$-$h$ mixing or the triangle VLQ loops. When $\Delta M_{U_1,D_1}$ is small, or equivalently when all of the VLQ masses are almost degenerate, the opposite signs between $y_{hF_1}$ and $y_{hF_2}$ cause significant cancellation of the $F_1$ and $F_2$ contributions. Therefore, $\kappa_g$ is within the allowed value. As the VLQ mass difference increases, the VLQ loop corrections become more important. The Higgs precision data put an upper bound on $\Delta M_{U_1,D_1}$. The $\kappa_\gamma$ is less sensitive since the dominant contribution to $\kappa_\gamma$ comes from the $W$ loop. The $S$-$h$ mixing effect, mainly on $\kappa_V$, is minor since we adopt the Higgs precision data at $2\sigma$ level such that $s_\eta \lesssim 0.5$ [30].

Figure 5 shows that the $ZZ$ channel in the heavy scalar searches puts the strongest bound for both mass cases. This is attributed to compatible branching ratios of $WW$, $hh$, and $ZZ$ modes but much smaller LHC upper bounds on $\sigma \cdot B$ for the $ZZ$ mode because of its clean signal. The parameter space with large $y_S$ and large $Y_D$ is excluded. We also present the contours of $s_\eta$ by dashed (orange) lines. It is clear to see that the current heavy Higgs searches put stronger bounds on the $S$-$h$ mixing angle than the Higgs precision data. In most parameter space, $s_\eta$ should be less than about 0.01 (0.05) for $m_S = 500\,(750)$ GeV. The radiatively generated $S$-$h$ mixing is significantly constrained by the current LHC data.

Finally, we show in Fig. 6 the cross section times branching ratio $\sigma(pp \to S) \times B(S \to$
Figure 6: Cross sections of production and decay of $S$ for the main decay channels with $m_S = 500 \text{ GeV}$ and $\sqrt{s} = 13 \text{ TeV}$ at the LHC. The cross sections are normalized by $y_S^2$.

$XY)$ as a function of $\Delta M_{U_1 D_1}$ with $m_S = 500 \text{ GeV}$ at the 13 TeV LHC. The decay of $S$ into $gg$ is not considered because of the overwhelming QCD background. We normalize $\sigma \cdot B$ by $y_S^2$. Incorporating the current Higgs precision constraint on $\Delta M_{U_1 D_1}$, we present the results for $\Delta M_{U_1 D_1}$ up to 200 GeV. In the whole parameter space, the $WW$ mode is leading or next-to-leading, having $\sigma \cdot B \sim O(100\text{–}1000) \text{ fb}$. The cleanest search mode, the $ZZ$ one, also has sizable signal rate about $100 \text{ fb}$ if $\Delta M_{U_1 D_1} \gtrsim 50 \text{ GeV}$. The $hh$ channel is also promising with sufficient VLQ mass differences.

6 Conclusions

In a simple extension of the SM with an additional singlet scalar field $S$, we answer the question whether a unique feature of a heavy scalar boson, the longitudinal polarization enhancement in its decay into a massive gauge boson pair, remains at loop level. In order to focus on the loop induced effects, we consider a limiting scenario where $S$ does not interact with the SM Higgs boson at tree level. Since $S$ decouples from the SM world at tree level, we introduced vector-like quarks (VLQs) as messengers between $S$ and the SM particles. In order for the Higgs boson to interact also with the VLQs, one VLQ doublet and two VLQ singlets are suggested. There are two up-type VLQs and two down-type VLQs, $U_{1,2}$ and $D_{1,2}$. Through the Yukawa couplings of VLQs with $S$ and the Higgs boson, the VLQs generate radiatively the $S$-$h$ mixing as well as the decays of $S$ into $gg$, $WW$, $ZZ$, and $hh$.

We found that the most required condition for enhancing the radiative decay rates of $S$ into $WW$, $ZZ$ and $hh$ is the large mass differences of VLQs. This is contrary to the common expectation since large mass differences with the fixed lightest VLQ mass mean heavy VLQs and thus smaller loop corrections. First the radiatively generated $S$-$h$ mixing is proportional to the coupling of $h$-$F_i$-$F_1$ ($F = U, D$ and $i = 1, 2$). When $M_{F_1} = M_{F_2}$, the opposite signs between $h$-$F_1$-$F_1$ and $h$-$F_2$-$F_2$ couplings cancel the contributions of $F_1$.
and $F_2$. As $\Delta M_{F_2} (\equiv M_{F_2} - M_{F_1})$ increases, the $S$-$h$ mixing angle is enhanced. The mixing induced decays of $S$ into $WW$, $ZZ$, $hh$, and $t\bar{t}$ become significant. Another kind of the VLQ contribution to the radiative decay of $S$ is through the triangle VLQ loops. We showed that the longitudinal polarization enhancement in $S \to WW, ZZ$ through the triangular VLQ loops happens also when the mass differences of the VLQs become large.

In order to illustrate the phenomenological features, we considered a simple benchmark scenario where $Y_D$ controls the VLQ mass differences with the fixed lightest VLQ mass. Two cases of $m_S = 500$ GeV and $m_S = 750$ GeV are studied. When $\Delta M_{FF'} = 0$, both the $S$-$h$ mixing and the longitudinal polarization enhancement in $S \to VV$ vanish, which makes $S \to gg$ dominant. The total decay rate is of the order of 0.1 GeV for $m_S \sim 500$ GeV. If $\Delta M_{FF'}$ is sizable such as $Y_D \simeq 0.8$, the decay of $S$ into $hh$ becomes as important as that into $gg$. For $Y_D \gtrsim 0.8$, $B(S \to gg)$ drops rapidly, and the decays into $hh$, $WW$, $ZZ$, and $t\bar{t}$ become similarly dominant. The enhancement of the total decay rate of $S$ is huge, by one order of magnitude when $Y_D = 1$. This is contrary to the naive prediction that heavier VLQs running in the loop would cause smaller loop corrections.

We also presented the 95\% C.L. exclusion regions of ($Y_D, y_S$) from the current LHC bounds including the Higgs precision data and the heavy scalar searches in the channel of $WW$, $ZZ$, and $hh$. Among various Higgs precision data, $\kappa_g$ puts the strongest bound on $Y_D$: $Y_D \gtrsim 1.1$ for $m_S = 500$ GeV and $Y_D \gtrsim 1.7$ for $m_S = 750$ GeV are excluded. The heavy scalar searches also put additional constraints. In particular the $ZZ$ channel data severely limit the $S$-$h$ mixing angle $\eta$, more than the Higgs precision data: $s_\eta \lesssim 0.05$ for $m_S = 500$ GeV and $s_\eta \lesssim 0.1$ for $m_S = 750$ GeV are allowed. In conclusions, our loop calculation in a UV model with a singlet scalar and three VLQ multiplets shows that the radiative decays of $S$ can be very enhanced when the mass spectrum of the VLQs shows diversity. Note that the presence of multiple VLQs is crucial for the enhanced radiative decays of $S$ since sizable mass differences among VLQs are required. Therefore, the persistent searches for a heavy scalar boson at the future LHC are of great importance in constraining new particles that appear at loop level.

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