Color-octet scalar decays to a gluon and an electroweak gauge boson in the Manohar-Wise model

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\textbf{Abstract:} We evaluate the one loop amplitudes giving rise to couplings between a scalar color octet, a gluon, and an electroweak gauge boson. These one loop amplitudes could give rise to new physics signals in $\gamma$ jet, $Z$ jet and $W$ jet production at the LHC. Branching ratios for color octet scalar decay into these modes can reach the 10\% ($\gamma$ jet), and a few percent ($Z$ jet) level for masses below $2m_t$. In a narrow kinematic window, the charged scalar can decay to $W$ jet with a substantial branching fraction.

\textbf{Keywords:} Vector boson plus jet, new physics, color octet scalars

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1 Introduction

The search for new physics at the Large Hadron Collider (LHC) continues to constrain the parameters of many possible new particles that appear in a variety of extensions of the standard model (SM). A channel that has not received much attention is the production of a single jet in conjunction with one electroweak gauge boson ($\gamma, Z, W^\pm$), which occurs in extensions of the SM with additional colored particles. The modes $g\gamma$ and $gZ$ have been recently considered in the phenomenological study of a pseudoscalar color octet $\pi_8$ [1], where it is argued that $g\gamma$, in particular, is a clean channel due to the presence of an energetic photon. An effective vertex in a model independent Lagrangian containing these modes appeared in Ref. [2]. Early studies of an apparent di-jet anomaly reported by CDF [3] considered decays into $\gamma, Z, W^\pm$ with multiple jets originating from transitions between two scalars with slightly different masses [4, 5]. In this case the lighter scalar decays into two jets, a different mechanism from the one considered in this paper.

In this paper we present the necessary one-loop results to study the constraints that these channels can place on a scalar color octet. In section 2 we review the relevant features of the Manohar-Wise (MW) model with an electroweak doublet, color octet scalar [6]. In section 3 we present explicit one-loop results for the $SVg$ vertices. Finally in section 4 we study the scalar widths into the $Vg$ channels in the regions of parameter space and kinematic variables where they can be largest.

2 The Model

The MW model contains fourteen new parameters in the scalar potential and an additional four in the Yukawa sector giving rise to a rich phenomenology. In particular three of these can modify substantially the loop level Higgs production and decay [6] as studied in the context of effective Higgs couplings [7–12]. The same three parameters are also constrained by precision electroweak measurements [6, 13, 14]. Flavor physics [15–17] constrains the new parameters in the Yukawa sector, and theoretical considerations such as unitarity and vacuum stability have been applied to all the new parameters [18–21]. Taking these constraints into account there have been several LHC studies of the model as well [22–25, 34].

In the MW model, the new scalars form a color octet and an electroweak doublet, with the respective gauge interactions being responsible for its pair production at the LHC. The possible Yukawa couplings reduce to two complex numbers once minimal flavor violation is imposed [6],

$$\mathcal{L}_Y = -\eta_{U} e^{i\alpha_{U}} g^{U}_{ij} \bar{u}_{Ri} T^{A} Q_{j} S^{A} - \eta_{D} e^{i\alpha_{D}} g^{D}_{ij} \bar{d}_{Ri} T^{A} Q_{j} S^{A} + h.c.$$  \hspace{1cm} (1)

Here $Q_i$ are the usual left-handed quark doublets, $S^a$ the new scalars and the normalization of the generators is $\text{Tr}(T^a T^b) = \delta^{ab}/2$. The matrices $g^{U,D}_{ij}$ are the same as the Higgs couplings to quarks, and the overall strength of the interactions is given by $\eta_{U,D}$ along with their phases $\alpha_{U,D}$. The latter introduce Charge-Parity (CP) violation beyond the SM and contribute for example to the Electric Dipole Moment (EDM) and Chromo-Electric Dipole Moment (CEDM) of quarks [6, 14, 17, 24].

The most general renormalizable scalar potential is given in Ref. [6], but our study will only
depend on the following terms,
\[ V = \lambda \left( H^{ij} H_i - \frac{v^2}{2} \right)^2 + 2m_s^2 \text{Tr} S^{ij} S_i + \lambda_1 H^{ij} H_i \text{Tr} S^{ij} S_j + \lambda_2 H^{ij} H_j \text{Tr} S^{ij} S_i \]
\[ + \left( \frac{\lambda_3}{3} H^{ij} H_j \text{Tr} S^{ij} S_j + \lambda_4 e^{i\alpha_4} H^{ij} \text{Tr} S^{ij} S_j S_i + \lambda_5 e^{i\phi_5} H^{ij} \text{Tr} S^{ij} S_j S_i + \text{h.c.} \right) \] (2)

where \( v \approx 246 \) GeV. The number of parameters in Eq.(2) can be further reduced by theoretical considerations: first \( \lambda_3 \) can be chosen to be real by a suitable definition of \( S \); custodial \( SU(2) \) symmetry implies the relations \( 2\lambda_3 = \lambda_2 \) (and hence \( M_{S^+} = M_{S^0} \)) \[6\] and \( \lambda_4 = \lambda_5^* \[14\]; and CP conservation removes all the phases, \( \alpha_U, \alpha_D, \) and \( \phi_4 \).

After symmetry breaking, the Higgs vev in Eq.(2) splits the octet scalar masses as,
\[ m_{S^\pm}^2 = m_3^2 + \frac{\lambda_1 v^2}{4}, \quad m_{S_{R,I}}^2 = m_3^2 + (\lambda_1 + \lambda_2 \pm 2\lambda_3) \frac{v^2}{4}, \] (3)

The following can be used as a convenient set of independent input parameters:
\[ M_{S^+, S^0}, \lambda_2, \lambda_4, \eta_U, \eta_D \] (4)

\( \lambda_2 \) then controls the mass splitting between the two neutral resonances \( S_{I,R} \) while \( \lambda_4 \) determines the magnitude of scalar loop contributions to single neutral scalar production via gluon fusion.

Finally, \( \eta_{U,D} \) control respectively the strength of the \( Stt \) and \( Sbb \) interactions.

The effective one-loop coupling \( SVg \) can be written in terms of two factors \( F(\bar{F}) V_g \) as,
\[ L_{Sg}^\alpha = \frac{\alpha_s}{8\pi v} \left[ \left( F_{2g}^R G_{\mu\nu} G_{\mu\nu} + \bar{F}_{2g}^R \tilde{G}_\mu G_{\mu\nu} \right) S_{R}^C + \left( F_{4g}^R G_{\mu\nu} G_{\mu\nu} + \bar{F}_{4g}^R \tilde{G}_\mu G_{\mu\nu} \right) S_{I}^C \right] \delta^{ABC}, \]
\[ L_{Sg}^\gamma = \frac{\sqrt{\alpha_s}}{3\pi v} \left[ \left( F_{2g}^\gamma G^{A\mu\nu} + \bar{F}_{2g}^\gamma \tilde{G}^{A\mu\nu} \right) S_{R}^B + \left( F_{4g}^\gamma G^{A\mu\nu} + \bar{F}_{4g}^\gamma \tilde{G}^{A\mu\nu} \right) S_{I}^B \right] \delta^{AB}, \]
\[ L_{Sg}^{Zg} = \frac{\sqrt{\alpha_s}}{12\pi v} \left[ \left( F_{2g}^{Zg} G^{A\mu\nu} Z_{\mu\nu} + \bar{F}_{2g}^{Zg} \tilde{G}^{A\mu\nu} Z_{\mu\nu} \right) S_{R}^B + \left( F_{4g}^{Zg} G^{A\mu\nu} Z_{\mu\nu} + \bar{F}_{4g}^{Zg} \tilde{G}^{A\mu\nu} Z_{\mu\nu} \right) S_{I}^B \right] \delta^{AB}, \]
\[ L_{SWg} = \frac{\sqrt{\alpha_s}}{12\pi v} \left[ \left( F_{Wg}^{\pm} G^{A\mu\nu} W^{\pm\mu} + \bar{F}_{Wg}^{\pm} \tilde{G}^{A\mu\nu} W^{\pm\mu} \right) S_{R}^\pm \delta^{AB} \right] \] (5)

where \( G_{\mu\nu}^A, A_{\mu\nu}, Z_{\mu\nu} \) and \( W_{\mu\nu} \) are the gluon, photon, \( Z \) and \( W \) field strength tensors respectively and \( \tilde{G}^{A\mu\nu} = (1/2)\epsilon^{\mu\nu\alpha\beta}G_{AB}^{A\alpha\beta} \).

Explicit one-loop results for these factors in the MW model are presented in the next section. Of these, only \( L_{Sg} \) exists in the literature and we find a sign difference described below. These effective vertices receive their main contributions from top-quark and color-octet scalar loops. The bottom-quark loop is important only for regions of parameter space where \( |\eta_D| >> |\eta_U| \).

3 Explicit one-loop results in the MW model

We perform the calculation with the aid of a number of software packages. We first implement the model in FeynRules \[26, 27\] to generate FeynArts \[28\] output where LoopTools \[29\] is used to
compute the one loop diagrams and simplification is assisted with FeynCalc [30, 31], FeynHelpers [32] and Package-X [33]. For completeness we present the result without assuming custodial or CP symmetries. For the numerical results, however, we will take the custodial and CP symmetry limits of these expressions.

3.1 \( S_{R,I} \rightarrow g \, g \)

![Diagram](image)

Figure 1: One-loop diagrams contributing to the factors appearing in Eq. (5) for the \( S_{R,1g} \) coupling.

The diagrams are shown in Figure 1 and the resulting form factors are given by

\[
F^g_{gR} = \left\{ \eta_{UC} I_q^G \left( \frac{m_t^2}{m_R^2} \right) + \eta_{DC} I_q^G \left( \frac{m_t^2}{m_I^2} \right) \right. \\
+ \frac{9}{4} v_2 \left( \lambda_4 c_4 + \lambda_5 c_5 \right) \frac{1}{2} \left[ I_s^G (1) + \frac{1}{3} I_s^G \left( \frac{m_t^2}{m_R^2} \right) + \frac{2}{3} I_s^G \left( \frac{m_{S \pm}^2}{m_R^2} \right) \right] \left\}
\]

\[
F^g_{gI} = \left\{ -\eta_{US} I_q^G \left( \frac{m_t^2}{m_I^2} \right) + \eta_{DS} I_q^G \left( \frac{m_t^2}{m_I^2} \right) \right. \\
- \frac{9}{4} v_2 \left( \lambda_4 s_4 + \lambda_5 s_5 \right) \frac{1}{2} \left[ I_s^G (1) + \frac{1}{3} I_s^G \left( \frac{m_t^2}{m_I^2} \right) + \frac{2}{3} I_s^G \left( \frac{m_{S \pm}^2}{m_I^2} \right) \right] \left\}
\]

\[
\tilde{F}^g_{gR} = \left[ -\eta_{US} \frac{m_t^2}{m_R^2} f \left( \frac{m_t^2}{m_R^2} \right) - \eta_{DS} \frac{m_t^2}{m_I^2} f \left( \frac{m_t^2}{m_I^2} \right) \right]
\]

\[
\tilde{F}^g_{gI} = \left[ -\eta_{UC} \frac{m_t^2}{m_I^2} f \left( \frac{m_t^2}{m_I^2} \right) + \eta_{DC} \frac{m_t^2}{m_I^2} f \left( \frac{m_t^2}{m_I^2} \right) \right]
\]

Imposing custodial symmetry, the scalar loops only contribute to \( F^g_{gR} \). Imposing CP symmetry \( S_R \) (\( S_I \)) are pure scalar (pseudo-scalar) and therefore only the factors \( F^g_{gR} \) and \( \tilde{F}^g_{gI} \) are not-zero. The bottom-quark loops are much suppressed with respect to the top-quark loops unless \( \eta_D >> \eta_U \).

In the limit of CP conservation and \( m_b = 0 \), these results agree with [13] except for the sign in front of the factor \( \frac{9}{4} \). This sign, however, is of no consequence for phenomenology as \( \lambda_{4,5} \) can have either sign. ¹

¹When CP violation is included we find the following errors in [24]: the factors \( \tilde{F}^g_{R,I} \) are a factor of two too large in [24]; the function \( I_s(z) \) in Eq. 2.6 of [24] contains an incorrect overall factor of \( z \) which is inconsequential in the limit of degenerate scalars. There is also a typo in Eq. 6 of [34], where there should be a minus sign in the term with \( \eta_U \) in \( F_I \), corresponding to \( \tilde{F}^g_I \) here.
3.2 $S_{R,I} \rightarrow \gamma g$

![Diagram](image)

Figure 2: One-loop diagrams contributing to the factors appearing in Eq. (5) for the $S_{R,I} \gamma g$ coupling.

The diagrams for this process are shown in Figure 2 and the resulting form factors are given by

$$
F_R^{\gamma g} = \left[ \eta_{UC} I_q \left( \frac{m_t^2}{m_R^2} \right) - \eta_{DC} \frac{1}{2} I_q \left( \frac{m_t^2}{m_R^2} \right) + \frac{9}{4} \frac{v^2}{m_R^2} \left( \lambda_4 c_4 - \lambda_5 c_5 \right) \frac{1}{2} I_s \left( \frac{m_{S\pm}^2}{m_R^2} \right) \right] \quad (10)
$$

$$
F_I^{\gamma g} = \left[ -\eta_{US} I_q \left( \frac{m_t^2}{m_I^2} \right) - \eta_{SD} \frac{1}{2} I_q \left( \frac{m_t^2}{m_I^2} \right) + \frac{9}{4} \frac{v^2}{m_I^2} \left( \lambda_4 s_4 - \lambda_5 s_5 \right) \frac{1}{2} I_s \left( \frac{m_{S\pm}^2}{m_I^2} \right) \right] \quad (11)
$$

$$
\tilde{F}_R^{\gamma g} = \left[ -\eta_{US} \frac{m_t^2}{m_R^2} f \left( \frac{m_t^2}{m_R^2} \right) + \eta_{SD} \frac{1}{2} \frac{m_t^2}{m_R^2} f \left( \frac{m_t^2}{m_R^2} \right) \right] \quad (12)
$$

$$
\tilde{F}_I^{\gamma g} = \left[ -\eta_{CU} \frac{m_t^2}{m_I^2} f \left( \frac{m_t^2}{m_I^2} \right) - \eta_{DC} \frac{1}{2} \frac{m_t^2}{m_I^2} f \left( \frac{m_t^2}{m_I^2} \right) \right] \quad (13)
$$

Unlike the $S_R \rightarrow gg$ case, the scalar loop contribution to $S_R \rightarrow \gamma g$ vanishes in the custodial symmetry limit. In the custodial and CP symmetry limits, the scalar loops do not contribute to these processes at all.

3.3 $S_{R,I} \rightarrow Z g$

![Diagram](image)

Figure 3: One-loop diagrams contributing to the factors appearing in Eq. (5) for the $S_{R,I} Zg$ coupling.

The diagrams for this process are shown in Figure 3 and the resulting form factors are given by...
by

\[
F_{R}^{Zg} = \frac{1}{(m_{R}^2 - m_{Z}^2)^2} \left\{ -\eta_{U} c_{U} \left( \frac{3c_{W}^2 - 5s_{W}^2}{s_{2W}} \right) m_{t}^2 \left[ (m_{R}^2 - m_{Z}^2 - 4m_{t}^2) \left[ f \left( \frac{m_{b}^2}{m_{R}^2} \right) - f \left( \frac{m_{b}^2}{m_{Z}^2} \right) \right] - 2m_{Z}^2 \left[ I_{q}^Z \left( \frac{m_{t}^2}{m_{R}^2} \right) - I_{q}^Z \left( \frac{m_{t}^2}{m_{Z}^2} \right) \right] - 2(m_{R}^2 - m_{Z}^2) \right] + \eta_{DCD} \left( \frac{3c_{W}^2 - s_{W}^2}{s_{2W}} \right) m_{b}^2 \left[ (m_{R}^2 - m_{Z}^2 - 4m_{b}^2) \left[ f \left( \frac{m_{s}^2}{m_{R}^2} \right) - f \left( \frac{m_{s}^2}{m_{Z}^2} \right) \right] - 2m_{Z}^2 \left[ I_{q}^Z \left( \frac{m_{b}^2}{m_{R}^2} \right) - I_{q}^Z \left( \frac{m_{b}^2}{m_{Z}^2} \right) \right] - 2(m_{R}^2 - m_{Z}^2) \right] \right\} \tag{14}
\]

\[
F_{I}^{Zg} = \frac{1}{(m_{I}^2 - m_{Z}^2)^2} \left\{ \eta_{USU} \left( \frac{3c_{W}^2 - 5s_{W}^2}{s_{2W}} \right) m_{t}^2 \left[ (m_{I}^2 - m_{Z}^2 - 4m_{t}^2) \left[ f \left( \frac{m_{t}^2}{m_{I}^2} \right) - f \left( \frac{m_{t}^2}{m_{Z}^2} \right) \right] - 2m_{Z}^2 \left[ I_{q}^Z \left( \frac{m_{t}^2}{m_{I}^2} \right) - I_{q}^Z \left( \frac{m_{t}^2}{m_{Z}^2} \right) \right] - 2(m_{I}^2 - m_{Z}^2) \right] + \eta_{DS} \left( \frac{3c_{W}^2 - s_{W}^2}{s_{2W}} \right) m_{b}^2 \left[ (m_{I}^2 - m_{Z}^2 - 4m_{b}^2) \left[ f \left( \frac{m_{b}^2}{m_{I}^2} \right) - f \left( \frac{m_{b}^2}{m_{Z}^2} \right) \right] - 2m_{Z}^2 \left[ I_{q}^Z \left( \frac{m_{b}^2}{m_{I}^2} \right) - I_{q}^Z \left( \frac{m_{b}^2}{m_{Z}^2} \right) \right] - 2(m_{I}^2 - m_{Z}^2) \right] \right\} \tag{15}
\]
\[ F_R^{Zg} = \frac{1}{(m_R^2 - m_Z^2)} \left\{ -\eta_{USU} \frac{(3c_W^2 - 5s_W^2)}{s_{2W}} m_t^2 \left[ f \left( \frac{m_t^2}{m_R^2} \right) - f \left( \frac{m_t^2}{m_Z^2} \right) \right] 
+ \eta_{DSD} \frac{(3c_W^2 - s_W^2)}{s_{2W}} m_b^2 \left[ f \left( \frac{m_b^2}{m_I^2} \right) - f \left( \frac{m_b^2}{m_Z^2} \right) \right] \right\} \] (16)

\[ F_I^{Zg} = \frac{1}{(m_I^2 - m_Z^2)} \left\{ -\eta_{UCU} \frac{(3c_W^2 - 5s_W^2)}{s_{2W}} m_t^2 \left[ f \left( \frac{m_t^2}{m_I^2} \right) - f \left( \frac{m_t^2}{m_Z^2} \right) \right] 
+ \eta_{DCD} \frac{(3c_W^2 - s_W^2)}{s_{2W}} m_b^2 \left[ f \left( \frac{m_b^2}{m_I^2} \right) - f \left( \frac{m_b^2}{m_Z^2} \right) \right] \right\} \] (17)

As in the $S \rightarrow \gamma g$ case, all scalar loop contributions to $S \rightarrow Zg$ vanish in the custodial and CP symmetry limit.

### 3.4 $S^+ \rightarrow W^+ g$

![One-loop diagrams contributing to the factors appearing in Eq. (5) for the $S^+ W^+ g$ coupling.](image)

Finally, the diagrams for $S^+ \rightarrow W^+ g$ are shown in Figure 4 and the resulting form factors are
given by

\[ P_{Wg} = \frac{1}{(m_f^2 - m_W^2)^2 s_W} \left( \frac{3}{2} V_{tb}^2 \left( \frac{m_t ((m_b - m_t)(m_b + m_t)^2 - m_t^2)}{m_W^2} + 2m_b m_W^2 \eta_U e^{i\alpha_U} \right) \right. \]

\[ + \frac{m_b ((m_b - m_t)(m_b + m_t)^2 - m_t^2)}{m_W^2} \left( \frac{m_b^2 - m_W^2}{2} \frac{I_D^W}{(m_b^2 + m_t^2 - m_W^2)^2} \right) \]

\[ + \left. \left( m_t^2 (m_W^2 (m_b - m_t)(m_b + m_t)^2 - 2m_t^2 m_b) + m_t^2 (m_b - m_t)(m_t^2 - 2(m_b + m_t)^2) \right) \eta_U e^{i\alpha_U} \right) \]

\[ + \frac{m_b (m_W^2 ((m_b - m_t)(m_b + m_t)^2 + 2m_t^2 m_t) + m_t^2 (m_b - m_t)(m_t^2 - 2(m_b + m_t)^2))}{m_W^2} \eta_D e^{-i\alpha_D} \right) \]

\[ \times \left( \frac{m_b^2 + m_t^2 - m_W^2}{2} \frac{I_D^W}{(m_b^2 + m_t^2 - m_W^2)^2} \right) \]

\[ - m_b m_t (m_W^2 - m_b^2) \left( 2m_b(m_b + m_t) + m_W^2 - m_t^2 \right) \eta_U e^{i\alpha_U} C_0(0, m_t^2, m_W^2, m_b^2, m_t^2, m_b^2) \]

\[ + m_b m_t (m_t^2 - m_W^2)(2m_t(m_b + m_t) + m_W^2 - m_t^2) \eta_D e^{-i\alpha_D} C_0(0, m_t^2, m_W^2, m_b^2, m_t^2, m_b^2) \]

\[ + \frac{m_t (m_b + m_t)(m_t^2 - m_W^2)}{2 m_t^2} \left( \frac{1}{(m_t^2 - m_W^2)(m_t^2 - m_b^2) + 2m_t^2 m_W^2 m_b^2} \eta_U e^{i\alpha_U} \log \left( \frac{m_b^2}{m_t^2} \right) \right) \]

\[ + \frac{m_b (m_b + m_t)(m_t^2 - m_W^2)}{2 m_B^2} \left( \frac{1}{(m_t^2 - m_W^2)(m_t^2 - m_b^2) + 2m_t^2 m_W^2 m_b^2} \eta_D e^{-i\alpha_D} \log \left( \frac{m_b^2}{m_t^2} \right) \right) \]

\[ + \frac{9}{4} v^2 \left[ \lambda_4 s_4 - \lambda_5 s_5 \left( \frac{m_W^2 I_{s_1}^W}{m_t^4} + 2m_t^2 f \left( \frac{m_t^2}{m_W^2} \right) + m_t^2 I_{s_2}^W \right) \right] \]

\[ + \left( \frac{(m_4^2 + m_5^2 - m_W^2)}{2} I_s^W \left( \frac{m_W^2 m_R^2}{m_t^2 + m_R^2 - m_W^2} \right)^2 \right) \]

\[ - \left( \lambda_4 c_4 - \lambda_5 c_5 \right) \frac{m_W^2 m_R^2}{m_t^2} \frac{m_R^2}{m_t^2} \log \left( \frac{m_R^2}{2m_t^2} \left( g \left( \frac{m_t^2}{m_R^2} + 1 \right) \right) \right) \]

\[ - \left( \lambda_4 c_4 - \lambda_5 c_5 \right) \frac{m_W^2 m_R^2}{m_t^2} \frac{m_R^2}{m_t^2} \log \left( \frac{m_W^2}{m_R^2} \right) \]

\[ - \left( \lambda_4 c_4 - \lambda_5 c_5 \right) \frac{m_W^2}{m_t^2} \lambda_4 c_4 - \lambda_5 c_5 \left( m_R^2 C_0(0, m_t^2, m_W^2, m_R^2, m_t^2, m_R^2) + m_R^2 C_0(0, m_t^2, m_W^2, m_t^2, m_R^2, m_R^2) \right) \]

\[ - \left( \lambda_4 c_4 - \lambda_5 c_5 \right) \frac{m_W^2}{m_t^2} \]
\[ \tilde{F}^{Wg} = \frac{1}{sW} \left\{ -\frac{3i}{2} |V_{tb}|^2 \left[ m_h^2 \eta U e^{i\alpha U} C_0(0, m_t^2, m_W^2, m_t^2, m_h^2) + m_b^2 \eta D e^{i\alpha D} C_0(0, m_t^2, m_W^2, m_b^2, m_t^2) \right] \right\} \]  

(19)

This rather cumbersome expression of \( F^{Wg} \) simplifies considerably in the custodial and CP symmetric limits. In that case, and working to leading order in \( m_b \), one finds

\[
F^{Wg} = 3\eta U e^{i\alpha U} |V_{tb}|^2 m_1^2 (m_1^2 - m_t^2) \left[ \frac{m_1^4 (m_1^2 - m_t^2)}{m_t^4 (m_1^2 - m_t^2)^2} \log \left( \frac{m_b m_t}{m_t^2 - m_1^2} \right) + \frac{i}{2} \right]
\]

(20)

The functions appearing in the above results are given by

\[
I_q^G(x) = I_\gamma(x) = 2x - x g(x)^2 f(x), \quad I_s^G(x) = I_s^V(x) = I_{s_1}^W(x) = -(1 + 2xf(x)),
\]

\[
I_q^Z(x) = I_s^Z(x) = I_q^W(x) = I_s^W(x) = \pm g(x) \sqrt{2f(x)}, (+ \text{ for } x > 1/4, - \text{ for } x < 1/4)
\]

\[
I_{s_2}^W(x) = 1 - g(x) \sqrt{2f(x)}
\]

(21)

where

\[
f(x) = \begin{cases} 
\frac{1}{2} \left( \ln \left( \frac{1 + g(x)}{1 - g(x)} \right) - i\pi \right)^2 & \text{for } x < \frac{1}{4} \\
-2 \left( \arcsin \left( \frac{1}{2\sqrt{x}} \right) \right)^2 & \text{for } x > \frac{1}{4}
\end{cases}
\]

(22)

and some specific values that are useful are

\[
I_s^G(1) = \frac{\pi^2}{9} - 1, \quad I_s^Z(1) = -\frac{\pi}{\sqrt{3}}
\]

(23)

In addition, \( C_0 \) is the well-known Passarino-Veltman function.

4 Decay widths and discussion

The decay widths for the one-loop induced modes are given by

\[
\Gamma(S_{R,I} \to gg) = \frac{5}{12\pi \alpha_s}{\left( \frac{\alpha_s}{8\pi v} \right)}^2 m_{R,I}^3 \left( |F_{gg}^{R,I}|^2 + |\tilde{F}_{gg}^{R,I}|^2 \right)
\]

\[
\Gamma(S_{R,I} \to \gamma g) = \frac{1}{8\pi}{\left( \frac{\alpha_s}{3\pi v} \right)}^2 m_{R,I}^3 \left( |F_{\gamma g}^{R,I}|^2 + |\tilde{F}_{\gamma g}^{R,I}|^2 \right)
\]

\[
\Gamma(S_{R,I} \to Z g) = \frac{1}{8\pi}{\left( \frac{\alpha_s}{12\pi v} \right)}^2 \frac{(m_{R,I}^2 - m_Z^2)^3}{m_{R,I}^3} \left( |F_{Zg}^{R,I}|^2 + |\tilde{F}_{Zg}^{R,I}|^2 \right)
\]

\[
\Gamma(S^{\pm} \to W^{\pm} g) = \frac{1}{8\pi}{\left( \frac{\alpha_s}{12\pi v} \right)}^2 \frac{(m_{S^{\pm}}^2 - m_{W}^2)^3}{m_{S^{\pm}}^3} \left( |F_{Wg}^{R,I}|^2 + |\tilde{F}_{Wg}^{R,I}|^2 \right)
\]

(24)
Numerically, these result in small branching ratios that are negligible for phenomenology except in special cases. The region of parameter space where these modes can be most relevant is understood as follows

- The $S_{R,I}$ neutral resonances will decay predominantly into top pairs and $S^\pm$ will decay predominantly into top-bottom pairs if those channels are kinematically available. The loop induced modes become important for the mass ranges

$$100 \text{ GeV} \lesssim m_{S_{R,S_I}} \lesssim 350 \text{ GeV}, \quad 100 \text{ GeV} \lesssim m_{S^\pm} \lesssim 175 \text{ GeV}$$

where the lower limit is approximately the LEP exclusion for scalar pair production [14]. The decays into two jets through the lighter quarks are not suppressed in these ranges and in some instances dominate as shown below.

- The decays to $t\bar{t}$ or $t\bar{b}$ can also be suppressed with very small values of $\eta_U$. This would also suppress the production mechanism for a single scalar, but would not affect the cross-section for pair-production which depends only on the QCD coupling constant and at 13 TeV is approximately 0.2 pb. [13, 34]. However, reducing $\eta_U$ also reduces the top-quark contribution to the loop decays, which is dominant in most cases.

- The $S_{R,I}$ will also decay predominantly into bottom pairs unless $\eta_D$ is small.

- A large value of $\lambda_{4,5}$ will affect the $gg$ channel but not the other ones. If the sign is such that the $gg$ decay is suppressed, this may result in slightly larger branching ratios for the $\gamma g$ or $Zg$ modes.

- For our numerics below we choose parameter values $\lambda_2 = 0$ to have degeneracy between $S_I$ and $S_R$ thus preventing decays between two of the scalars. We will illustrate results with $\eta_U = 5$, below its unitarity constraint [19], and with $\eta_D = 1$ as we do not want to enhance the $b\bar{b}$ decay modes. Finally, we illustrate the effect of reducing the $gg$ mode with $\lambda_4 = \lambda_5 = -10$, which is below the tree level unitarity constraint [19] but in the range of next to leading order constraints [35].

We use the above arguments to show the relevant branching ratios for $S_R$ decay in Figure 5. The figure indicates that the $gg$, $qq$ (light-quarks) and $b\bar{b}$ modes are still dominant in this case, but these dijet final states would be very hard to see over QCD backgrounds. The right hand panel zooms in on the $\gamma g$ and $Zg$ modes showing up to 7 percent branching rations for $\gamma$-jet states are possible. In Figure 6 we illustrate the same values of $\eta_U,\eta_D$ but with $\lambda_{4,5} = -10$ which suppresses the $gg$ mode at $m_R$ close to the $t\bar{t}$ threshold. In this case the $\gamma g$ mode can reach branching ratios near 12% and the $Zg$ mode near 1%. Similarly, we show in Figure 7 the analogous situation for $S_I$ decay. In this case the parameters $\lambda_{4,5}$ do not enter in the CP conserving and custodial symmetry obeying case and the resulting branching ratios behave very similarly to those for $S_R$ decay.

Finally, we show in Figure 8 the situation for $S^\pm$ decay. In this case the $Wg$ mode (or $W$ plus jet final state) can dominate the decay for values of $m_{S^\pm}$ below the top mass. The $qq$ light-quark mode below $t\bar{b}$ threshold is dominated by $c\bar{b}$ quarks.
Figure 5: Branching ratios for $S_R$ decay below the $t\bar{t}$ threshold for parameter values $\eta_U = 5$, $\eta_D = 1$ and $\lambda_{4,5} = 0$.

Figure 6: Branching ratios for $S_R$ decay below the $t\bar{t}$ threshold for parameter values $\eta_U = 5$, $\eta_D = 1$ and $\lambda_{4,5} = -10$.

A $H \rightarrow gg$

To validate our mechanical calculation of the one-loop amplitudes we also compute the well known $H \rightarrow gg$ [36, 37] in the notation

$$\mathcal{L}(hgg) = \frac{\alpha_s}{4\pi v} (F_R^a R_{\mu\nu} G^{B\mu\nu} + F_R^b \tilde{G}^{A}_{\mu\nu} G^{B\mu\nu}) H \delta^{AB}$$

We find, in agreement with the well known result, the values

$$F_R^a = I_q^G \left( \frac{m_t^2}{m_H^2} \right) + 3 \frac{v^2}{4 m_H^2} \left[ (\lambda_1 + \lambda_2 - 2\lambda_3) I_s^G \left( \frac{m_t^2}{m_H^2} \right) \right. \left. + (\lambda_1 + \lambda_2 + 2\lambda_3) I_s^G \left( \frac{m_t^2}{m_H^2} \right) + 2\lambda_1 I_s^G \left( \frac{m_{S^\pm}^2}{m_H^2} \right) \right]$$

and $F_R^b = 0$.

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Figure 7: Branching ratios for $S_I$ decay below the $t\bar{t}$ threshold for parameter values $\eta_U = 5$ and $\eta_D = 1$.

Figure 8: Branching ratios for $S^\pm$ decay for parameter values $\eta_U = 5$ and $\eta_D = 1$.

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