Associated production of $Z$ and neutral Higgs bosons at the CERN Large Hadron Collider

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Abstract

We study the hadroproduction of a $CP$-even or $CP$-odd neutral Higgs boson in association with a $Z$ boson in the minimal supersymmetric extension of the standard model (MSSM) We include the contributions from quark-antiquark annihilation at the tree level and those from gluon-gluon fusion, which proceeds via quark and squark loops, and list compact analytic results. We quantitatively analyze the hadronic cross sections at the CERN Large Hadron Collider assuming a favorable supergravity-inspired MSSM scenario.

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1 Introduction

The search for Higgs bosons is among the prime tasks of the CERN Large Hadron Collider (LHC) [1]. While the standard model (SM) of elementary-particle physics contains one complex Higgs doublet, from which one neutral $CP$-even Higgs boson $H$ emerges in the physical particle spectrum after the spontaneous breakdown of the electroweak symmetry, the Higgs sector of the minimal supersymmetric extension of the SM (MSSM) consists of a two-Higgs-doublet model (2HDM) and accommodates a quintet of physical Higgs bosons: the neutral $CP$-even $h^0$ and $H^0$ bosons, the neutral $CP$-odd $A^0$ boson, and the charged $H^\pm$-boson pair. At the tree level, the MSSM Higgs sector has two free parameters, which are usually taken to be the mass $m_{A^0}$ of the $A^0$ boson and the ratio $\tan \beta = v_2/v_1$ of the vacuum expectation values of the two Higgs doublets.

In the following, we focus our attention on the $h^0$, $H^0$, and $A^0$ bosons, which we collectively denote by $\phi$. A recent discussion of $H^\pm$-boson production at the LHC may be found in Refs. [2–4] and the references cited therein. The dominant source of $\phi$ bosons is their single production by $gg$ fusion, $gg \to \phi$, which is mediated by heavy-quark [5] and squark [6] loops. Another, less important mechanism of single $\phi$-boson production is $b\bar{b} \to \phi$ [7]. The $\phi$ bosons thus produced have essentially zero transverse momentum ($p_T$). In order for the $\phi$ bosons to obtain finite $p_T$, they need to be produced in association with one or more other particles or jets ($j$). In leading order (LO), $j\phi$ associated production proceeds through the partonic subprocesses $gg \to g\phi$, $gq \to q\phi$, and $q\bar{q} \to g\phi$, which again involve heavy-quark [8,9] and squark [9,10] loops. Alternatively, the $\phi$ bosons can be produced, with interesting rates, in association with (i) a dijet via intermediate-boson fusion, $qq' \to qq'V^*V^* \to qq'\phi$, where $q$ and $q'$ stand for any light flavor of quark or antiquark, $V=W^\pm, Z$, and virtual particles are marked by an asterisk [11]; (ii) a quark-antiquark pair of heavy flavor $Q=t, b$ via $gg, qq \to QQ\phi$ [12]; (iii) an intermediate boson via $qq' \to W^\pm\phi$ [13–16], $q\bar{q} \to Z\phi$ [13,14,16–21], and $gg \to Z\phi$ [17–19,22–25]; or (iv) another, possibly different boson via $q\bar{q} \to \phi1\phi2$ [16,26,27] and $gg \to \phi1\phi2$ [27–29]. Note that, due to the absence of $A^0VV$ couplings at the tree level, $qq' \to qq'V^*V^* \to qq'\phi$ and the Drell-Yan processes $qq' \to W^{\pm*} \to W^{\pm}\phi$ and $qq' \to Z^* \to Z\phi$ are not possible for $\phi = A^0$. The partonic subprocesses $gg \to Z\phi$ and $gg \to \phi1\phi2$ are mediated by heavy-quark [17–19,22–25,27–29] and squark loops [18,27,29]. Comprehensive reviews of quantum corrections to Higgs-boson production within the SM and MSSM may be found in Refs. [30,31], respectively.

In this paper, we revisit $Z\phi$ associated hadroproduction via gluon fusion in the MSSM. In the SM case, mutual agreement between three independent calculations [17,22,25] has been established, and compact formulae for the partonic cross section are available [17]. In the MSSM, the status is much less advanced and, perhaps, somewhat unsatisfactory. As for $gg \to Z\phi$ with $\phi = h^0, H^0$, there exists only one analysis so far [19], which has not yet been verified by other authors. The analytic expressions presented in Ref. [19] are rather complicated; they involve 14 form factors. In order to translate the SM results [17,22,25] to the case of $gg \to Z\phi$ with $\phi = h^0, H^0$ in the MSSM, it is not sufficient to adjust the $HZZ$ and $Hqq$ couplings; it is also necessary to include certain quark triangle
diagrams with an $A^0$ boson in the $s$ channel [see Fig. 1(a)]. On the other hand, the squark loop contributions to these two MSSM processes vanish [19] for reasons explained below. As for $gg \to ZA^0$, the quark loop contributions were first studied on the basis of a numerical evaluation [23], which was recently employed for phenomenological signal-versus-background analyses taking into account the subsequent $Z \to l^+l^-$ and $A^0 \to b\bar{b}$ decays [20,24]. An independent analysis, including also the squark loops, was reported in Ref. [18], which does not contain an analytic expression for the partonic cross section either. Unfortunately, comparisons between Refs. [18,20,23,24] are not discussed in these papers. Recently, the helicity amplitudes of $gg \to Z\phi$ with $\phi = h^0, H^0, A^0$ were analyzed in Ref. [32] with regard to the asymptotic helicity conservation property of supersymmetry, and their real and imaginary parts were graphically presented at a specific scattering angle as functions of the center-of-mass (c.m.) energy.

With the Higgs hunt at the LHC being in full swing, it is an urgent matter to consolidate our knowledge of $Z\phi$ associated hadroproduction via gluon fusion in the MSSM, which is the motivation of this paper. Specifically, we present compact analytic expressions, also for the partonic cross sections of $q\bar{q}, b\bar{b} \to Z\phi$ [18–21], and perform a detailed numerical analysis using up-to-date input.

The importance of $b\bar{b}$-initiated subprocesses for Higgs-boson production has been variously emphasized in the literature, in particular, in connection with the final states $\phi$ [7], $\phi_1\phi_2$ [27], $H^+H^-$ [33], and $W^+H^\mp$ [2,34]. These subprocesses receive contributions from Feynman diagrams involving $b$-quark Yukawa couplings, which are generally strong for large values of $\tan\beta$. (The $btH^-$ and $tbH^+$ couplings are also strong for small values of $\tan\beta$.) If the two final-state particles couple to a $Z$ boson (or photon), as is the case for the final states $h^0A^0$, $H^0A^0$, $Zh^0$, $ZH^0$, and $H^+H^-$, then there are additional contributions from Drell-Yan-type diagrams, which are already present for the light flavors $q = u, d, s, c$. However, diagrams of the latter type are absent for the final states $h^0h^0$, $h^0H^0$, $H^0H^0$, $A^0A^0$, $ZA^0$, and $W^\pm H^\mp$, which can still be produced through $b\bar{b}$ annihilation.

As for $b\bar{b}$ annihilation, it should be noted that the treatment of bottom as an active flavor inside the colliding hadrons leads to an effective description, which comprises contributions from the higher-order subprocesses $gb \to Z\phi b$, $g\bar{b} \to Z\phi\bar{b}$, and $gg \to Z\phi b\bar{b}$. If all these subprocesses are to be explicitly included along with $b\bar{b} \to Z\phi$, then it is necessary to employ a judiciously subtracted $b$-quark PDF in order to avoid double counting [7,35]. The evaluation of $b\bar{b} \to Z\phi$ with an unsubtracted $b$-quark PDF is expected to slightly overestimate the true cross section [7,35]. For simplicity, we shall nevertheless adopt this effective approach in our analysis, keeping in mind that a QCD-correction factor below unity is to be applied. In fact, such a behavior has recently been observed for $b\bar{b} \to ZA^0$ [21].

In order to reduce the number of unknown supersymmetric input parameters, we adopt a scenario where the MSSM is embedded in a grand unified theory (GUT) involving supergravity (SUGRA) [36]. The MSSM thus constrained is characterized by the following parameters at the GUT scale, which come in addition to $\tan\beta$ and $m_{A^0}$: the universal scalar mass $m_0$, the universal gaugino mass $m_{1/2}$, the trilinear Higgs-sfermion coupling $A$, the bilinear Higgs coupling $B$, and the Higgs-higgsino mass parameter $\mu$. Notice that $m_{A^0}$
is then not an independent parameter anymore, but it is fixed through the renormalization group equation. The number of parameters can be further reduced by making additional assumptions. Unification of the $\tau$-lepton and $b$-quark Yukawa couplings at the GUT scale leads to a correlation between $m_\tau$ and $\tan \beta$. Furthermore, if the electroweak symmetry is broken radiatively, then $B$ and $\mu$ are determined up to the sign of $\mu$. Finally, it turns out that the MSSM parameters are nearly independent of the value of $A$, as long as $|A| \lesssim 500$ GeV at the GUT scale.

This paper is organized as follows. In Sec. 2, we list the LO cross sections of $q\bar{q} \to Z\phi$, where $\phi = h^0, H^0, A^0$, in the MSSM. We work in the parton model of QCD with $n_f = 5$ active quark flavors $q = u, d, s, c, b$, which we take to be massless. However, we retain the $b$-quark Yukawa couplings at their finite values, in order not to suppress possibly sizable contributions. We adopt the MSSM Feynman rules from Ref. [37]. The couplings of the $Z$ and $\phi$ bosons to quarks, $v_{Zqq}$, $a_{Zqq}$, and $g_{\phi qq}$, are given in Eq. (5) of Ref. [33] and Eq. (A3) of Ref. [27], respectively. As for the $\phi ZZ$ couplings, $g_{\phi^0 ZZ}$ and $g_{H^0 ZZ}$ are given by Eq. (A3) in Appendix A, while the $A^0 ZZ$ coupling vanishes at tree level. The $h^0 A^0 Z$ and $H^0 A^0 Z$ couplings, $g_{h^0 A^0 Z}$ and $g_{H^0 A^0 Z}$, may be found in Eq. (A2) of Ref. [27]. For each quark flavor $q$ there is a corresponding squark flavor $\tilde{q}_i$, which comes in two mass eigenstates $i = 1, 2$. The masses $m_{\tilde{q}_i}$ of the squarks and their trilinear couplings to the $\phi$ bosons, $g_{\tilde{q}_i \tilde{q}_j \tilde{q}_j}$, are listed in Eqs. (A5), (A7), and (A8) and in Table 1 of Ref. [38]1, Eq. (A.2) of Ref. [33], and Eq. (A4) of Ref. [27], respectively.

Considering the generic partonic subprocess $ab \to Z\phi$, we denote the four-momenta of the incoming partons, $a$ and $b$, and the outgoing $Z$ and $\phi$ bosons by $p_a, p_b, p_Z$, and $p_\phi$, respectively, and define the partonic Mandelstam variables as $s = (p_a + p_b)^2$, $t = (p_a - p_Z)^2$, and $u = (p_b - p_Z)^2$. The on-shell conditions read $p_a^2 = p_b^2 = 0$, $p_Z^2 = m_Z^2 = z$, and $p_\phi^2 = m_\phi^2 = h$. Four-momentum conservation implies that $s + t + u = z + h$. Furthermore, we have $sp_T^2 = tu - zh = N$, where $p_T$ is the absolute value of transverse momentum common to the $Z$ and $\phi$ bosons in the c.m. frame.

The tree-level diagrams for $b\bar{b} \to Z\phi$ with $\phi = h^0, H^0$ and $\phi = A^0$ are depicted in Fig. 1(a) and (b), respectively. As already mentioned above, there is no Drell-Yan diagram in Fig. 1(b) because of the absence of a $A^0 ZZ$ coupling at the tree level. The differential cross sections for the first class of partonic subprocesses may be generically

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1In Ref. [38], $m_{\tilde{q}_i}$ and $g_{\tilde{q}_i \tilde{q}_j \tilde{q}_j}$ are called $M_{Qa}$ and $\tilde{V}_{Qab}/g$, respectively.
written as

\[
\frac{d\sigma}{dt} (b\bar{b} \rightarrow Z\phi) = \frac{G_F^2 c_w^4 z}{3\pi s} \left[ (2z + p_T^2) g_{\phi ZZ}^2 (v_{Zbb}^2 + a_{Zbb}^2) |P_z(s)|^2 + \lambda |P|^2 \right. \\
- 4 s p_T^2 \left( \frac{1}{t} + \frac{1}{u} \right) g_{\phi bb} a_{Zbb} \text{Re} P + \left. g_{\phi bb}^2 \left( v_{Zbb}^2 T_+ + a_{Zbb}^2 T_- \right) \right],
\]

where \( G_F \) is Fermi’s constant, \( c_w = m_W/m_Z \) is the cosine of the weak mixing angle, \( \lambda = s^2 + z^2 + h^2 - 2(sz + zh + hs) \), and

\[
P = g_{\phi A^0 Z} g_{A^0 bb} P_{A^0}(s),
\]

\[
T_\pm = 2 \pm 2 + 2 p_T^2 \left[ z \left( \frac{1}{t} \pm \frac{1}{u} \right) \mp \frac{2s}{tu} \right].
\]

Here,

\[
P_X(s) = \frac{1}{s - m_X^2 + im_X \Gamma_X}
\]
is the propagator function of particle \( X \), with mass \( m_X \) and total decay width \( \Gamma_X \). For the second class of partonic subprocesses, we have

\[
\frac{d\sigma}{dt} (b\bar{b} \rightarrow Z A^0) = \frac{G_F^2 c_w^4 z}{3\pi s} \left[ |S|^2 - 4 s p_T^2 \left( \frac{1}{t} + \frac{1}{u} \right) g_{A^0 bb} a_{Zbb} \text{Re} S \right. \\
+ \left. g_{A^0 bb}^2 \left( v_{Zbb}^2 T_+ + a_{Zbb}^2 T_- \right) \right],
\]

where

\[
S = g_{h^0 A^0 Z} g_{h^0 bb} P_{h^0}(s) + g_{H^0 A^0 Z} g_{H^0 bb} P_{H^0}(s).
\]

As for \( Zh^0 \) and \( ZH^0 \) production, there are also sizable contributions from \( q\bar{q} \) annihilation via a virtual \( Z \) boson for the quarks of the first and second generations, \( q = u, d, s, c \), whose Yukawa couplings are negligibly small. The corresponding Drell-Yan cross sections are obtained from Eq. (1) by putting \( P = T_\pm = 0 \) and substituting \( b \rightarrow q \). The resulting expression agrees with Eq. (2.8) of Ref. [17], appropriate for \( q\bar{q} \rightarrow ZH \) in the SM, after adjusting the \( HZZ \) coupling. The full tree-level cross sections are then obtained by complementing the \( b\bar{b} \)-initiated cross sections of Eq. (1) with the Drell-Yan cross sections for \( q = u, d, s, c \).

The non-vanishing one-loop diagrams pertinent to \( gg \rightarrow Z\phi \), with \( \phi = h^0, H^0 \) and \( \phi = A^0 \) are depicted in Figs. 2(a) and (b), respectively. As already mentioned in the Introduction, the presence of the quark triangle diagrams involving an \( s \)-channel \( A^0 \)-boson exchange in Fig. 2(a) represents a qualitatively new feature of the MSSM as compared to the SM. Furthermore, similarly to Fig. 1(b), quark triangle diagrams with an \( s \)-channel \( Z \)-boson exchange do not appear in Fig. 2(b). In the following, we refer to a squark loop diagram involving an \( s \)-channel propagator as a triangle diagram. The residual squark loop diagrams are regarded to be of box type. The squark triangle and box diagrams for \( gg \rightarrow Z\phi \) with \( \phi = h^0, H^0 \) vanish, and so do the squark box diagrams for \( gg \rightarrow ZA^0 \). This may be understood as follows. (i) The \( g_{q\bar{q} \tilde{q} \bar{q}} \), \( g_{ggq\bar{q}} \), \( g_{gZq\bar{q}} \), and \( g_{Zq\bar{q}} \) couplings
are symmetric in \( i \) and \( j \), while the \( g_{A^0q_i\bar{q}_j} \) coupling is antisymmetric \[39\]. Thus, squark loops connecting gluons and \( Z \) bosons with an odd number of \( A^0 \) bosons vanish upon summation over \( i \) and \( j \). (ii) The \( g_{gq_i\bar{q}_j} \) and \( g_{Zq_i\bar{q}_j} \) couplings are linear in the squark four-momenta, while the \( g_{gg_i\bar{q}_j} \), \( g_{gZq_i\bar{q}_j} \), and \( g_{g\phi_i\bar{q}_j} \) couplings are momentum independent \[39\]. Thus, a squark loop connecting gluons, \( Z \) bosons, and \( \phi \) bosons vanishes upon adding its counterpart with the loop-momentum flows reversed if the total number of gluons and \( Z \) bosons is odd.

As in Refs. \[3,4\], we express the quark and squark loop contributions in terms of helicity amplitudes. We label the helicity states of the two gluons and the \( Z \) boson in the partonic c.m. frame by \( \lambda_a = \pm 1/2 \), \( \lambda_b = \pm 1/2 \), and \( \lambda_Z = 0, \pm 1 \). We first consider \( gg \rightarrow Z\phi \) with \( \phi = h^0, H^0 \). The helicity amplitudes of the quark triangle contribution read

\[
\mathcal{M}_{\lambda_a,\lambda_b,0}^\Delta = -2\sqrt{z} (\lambda_a + \lambda_b) \sum_q \left[ \frac{z-s}{z} a_{Zqq} g_{\phi ZZ} \mathcal{P}_Z(s) \left( F_\Delta \left( s, m_q^2 \right) + 2 \right) - \frac{s}{m_q} g_{A^0qq} g_{A^0\phi Z} \mathcal{P}_{A^0}(s) F_\Delta \left( s, m_q^2 \right) \right].
\]

The quark triangle form factor, \( F_\Delta \), is given in Eq. \(8\). As for the quark box contribution, all twelve helicity combinations contribute. Due to Bose symmetry, they are related by

\[
\mathcal{M}_{\lambda_a,\lambda_b,\lambda_Z}^\boxtimes (t, u) = (-1)^{\lambda_Z} \mathcal{M}_{\lambda_a,\lambda_b,\lambda_Z}^\boxtimes (u, t),
\]

\[
\mathcal{M}_{\lambda_a,\lambda_b,\lambda_Z}^\boxtimes (t, u) = \mathcal{M}_{\lambda_a-\lambda_b-\lambda_Z}^\boxtimes (t, u).
\]

Keeping \( \lambda_Z = \pm 1 \) generic, we thus only need to specify four expressions. These read

\[
\mathcal{M}_{\lambda_a,\lambda_b,\lambda_Z}^\boxtimes (t, u) = \mathcal{M}_{\lambda_a,\lambda_b,\lambda_Z}^\boxtimes (u, t),
\]

\[
\mathcal{M}_{\lambda_a,\lambda_b,\lambda_Z}^\boxtimes (t, u) = \mathcal{M}_{\lambda_a-\lambda_b-\lambda_Z}^\boxtimes (t, u).
\]

The quark box form factors, \( F_{\lambda_Z}^{\lambda_Z} \), are listed in Eq. \(8\). For the reasons explained above, we have \( \mathcal{M}_{\lambda_a,\lambda_b,\lambda_Z}^\Delta = \mathcal{M}_{\lambda_a,\lambda_b,\lambda_Z}^\boxtimes = 0 \) for the squark-induced helicity amplitudes.

We now turn to \( gg \rightarrow ZA^0 \). The helicity amplitudes of the quark and squark triangle contributions read

\[
\mathcal{M}_{\lambda_a,\lambda_b,0}^\Delta = -8\sqrt{z} (1 + \lambda_a \lambda_b) \sum_q m_q (g_{h^0A^0Zg_{h^0qq} P_{h^0}(s) + g_{H^0A^0Zg_{H^0qq} P_{H^0}(s) F_\Delta \left( s, m_q^2 \right)}}.
\]
\[
\tilde{M}_{\lambda_a \lambda_b \lambda_c}^\Delta = 2 \sqrt{\frac{\lambda}{z}} (1 + \lambda_a \lambda_b) \sum_{\tilde{q}_i} (g_{h^0 A^0 Z} g_{h^0 \tilde{q}_i \tilde{q}_i} P_{h^0} (s) + g_{H^0 A^0 Z} g_{H^0 \tilde{q}_i \tilde{q}_i} P_{H^0} (s)) \tilde{F}_\Delta \left( s, m_{\tilde{q}_i}^2 \right),
\]
(9)

respectively. The quark and squark triangle form factors, \( F_\Delta \) and \( \tilde{F}_\Delta \), may be found in Eq. (B.9). Again, the helicity amplitudes of the quark box contribution satisfy the Bose symmetry relations of Eq. (7). We find

\[
\begin{align*}
M_{\lambda_a \lambda_b \lambda_c}^{\Box +0} &= -\frac{8}{\sqrt{2 \lambda}} \sum_q g_{A^0 q q} a_{Z q q} m_q \left[ F_{++}^0 + (t \leftrightarrow u) \right], \\
M_{\lambda_a \lambda_b \lambda_c}^{\Box +0} &= -\frac{8}{\sqrt{2 \lambda}} \sum_q g_{A^0 q q} a_{Z q q} m_q \left[ F_{+-}^0 + (t \leftrightarrow u) \right], \\
M_{\lambda_a \lambda_b \lambda_c}^{\Box + + \lambda_Z} &= -4 \sqrt{2 N_s} \sum_q g_{A^0 q q} a_{Z q q} m_q \left[ F_{++}^1 - (t \leftrightarrow u) \right], \\
M_{\lambda_a \lambda_b \lambda_c}^{\Box + - \lambda_Z} &= -4 \sqrt{2 N_s} \sum_q g_{A^0 q q} a_{Z q q} m_q \left[ F_{+-}^1 - (t \leftrightarrow u, \lambda_Z \rightarrow -\lambda_Z) \right].
\end{align*}
\]
(11)

The quark box form factors, \( F_{\lambda_a \lambda_b \lambda_c}^{\Box |\lambda_Z|} \), are presented in Eq. (B.10). We recall that \( \tilde{M}_{\lambda_a \lambda_b \lambda_c}^\Delta = 0 \).

The differential cross section of \( gg \rightarrow Z \phi \) is then given by

\[
\frac{d\sigma}{dt}(gg \rightarrow Z \phi) = \alpha_s^2(\mu_r) G_F^2 m_W^4 \sum_{\lambda_a \lambda_b \lambda_c \lambda_d} \left| M_{\lambda_a \lambda_b \lambda_c}^\Delta + M_{\lambda_a \lambda_b \lambda_c}^{\Box \lambda_d} + \tilde{M}_{\lambda_a \lambda_b \lambda_c}^{\Delta \lambda_d} \right|^2,
\]
(12)

where \( \alpha_s(\mu_r) \) is the strong-coupling constant at renormalization scale \( \mu_r \). Due to Bose symmetry, the right-hand side of Eq. (12) is symmetric in \( t \) and \( u \).

The differential cross section of \( gg \rightarrow Z H \) in the SM is obtained from Eqs. (6)–(8) and (12), with \( \phi = h^0 \), by replacing \( h^0 \rightarrow H \), adjusting the \( h^0 ZZ \) and \( h^0 qq \) couplings, and discarding the contribution due to \( A^0 \)-boson exchange. In this way, we recover the result of Ref. [17], which is expressed in terms of Lorentz-invariant form factors rather than helicity amplitudes.

The kinematics of the inclusive reaction \( AB \rightarrow Z \phi + X \), where \( A \) and \( B \) are colliding hadrons, is described in Sec. II of Ref. [2]. Its double-differential cross section \( d^2 \sigma / dy dp_T \), where \( y \) and \( p_T \) are the rapidity and transverse momentum of the Z boson in the c.m. system of the hadronic collision, may be evaluated from Eq. (2.1) of Ref. [2].

3 Phenomenological Implications

We are now in a position to explore the phenomenological implications of our results. The SM input parameters for our numerical analysis are taken to be \( G_F = 1.16637 \times \)
$10^{-5}$ GeV$^{-2}$, $m_W = 80.399$ GeV, $m_Z = 91.1876$ GeV, $m_t = 172.0$ GeV, and $\overline{m}_b(\overline{m}_b) = 4.19$ GeV [40]. We adopt the LO proton PDF set CTEQ6L1 [41]. We evaluate $\alpha_s(\mu_t)$ and $m_0(\mu_t)$ from the LO formulas, which may be found, e.g., in Eqs. (23) and (24) of Ref. [42], respectively, with $n_f = 5$ quark flavors and asymptotic scale parameter $\Lambda_{QCD}^{(5)} = 165$ MeV [41]. We identify the renormalization and factorization scales with the $Z\phi$ invariant mass $\sqrt{s}$. We vary $\tan\beta$ and $m_{A^0}$ in the ranges $3 < \tan\beta < 32 \approx m_t/m_b$ and $180$ GeV $< m_{A^0} < 1$ TeV, respectively. As for the GUT parameters, we choose $m_{1/2} = 150$ GeV, $A = 0$, and $\mu < 0$, and tune $m_0$ so as to be consistent with the desired value of $m_{A^0}$. All other MSSM parameters are then determined according to the SUGRA-inspired scenario as implemented in the program package SUSPECT [43]. We do not impose the unification of the $\tau$-lepton and $b$-quark Yukawa couplings at the GUT scale, which would just constrain the allowed $\tan\beta$ range without any visible effect on the results for these values of $\tan\beta$. We exclude solutions which do not comply with the present experimental lower mass bounds of the sfermions, charginos, neutralinos, and Higgs bosons [40].

We now study the fully integrated cross sections of $pp \to Z\phi + X$ at the LHC, with c.m. energy $\sqrt{S} = 14$ TeV. Figures 3–5 refer to the cases $\phi = h^0, H^0, A^0$, respectively. In part (a) of each figure, the $m_\phi$ dependence is studied for $\tan\beta = 3$ and 30 while, in part (b), the $\tan\beta$ dependence is studied for $m_{A^0} = 300$ and 600 GeV. We note that the SUGRA-inspired MSSM with our choice of input parameters does not permit $\tan\beta$ and $m_{A^0}$ to be simultaneously small, due to the experimental lower bound on $m_{h^0}$ [40]. This explains why the curves for $\tan\beta = 3$ in Figs. 3–5(a) only start at $m_{A^0} \approx 280$ GeV, while those for $\tan\beta = 30$ already start at $m_{A^0} \approx 180$ GeV.

In Figs. 3 and 4, which refer to $\phi = h^0, H^0$, respectively, the total $q\bar{q}$-annihilation contributions (dashed lines), corresponding to the coherent superposition of Drell-Yan and Yukawa-enhanced amplitudes, and the $gg$-fusion contributions (solid lines), which arise only from quark loops, are presented separately. For a comparison with future experimental data, they should be added. For comparison, also the pure Drell-Yan contributions (dotted lines) are shown. As for $\phi = h^0$, we observe from Fig. 3 that the contribution due to $q\bar{q}$ annihilation is almost exhausted by the Drell-Yan process and greatly exceeds the one due to $gg$ fusion, by a factor of 3–5. The $q\bar{q}$-annihilation contribution falls off by a factor of two as $m_{h^0}$ runs from 82 GeV to 115 GeV and feebly depends on $\tan\beta$, except for the appreciable rise towards the lower edge of the considered $\tan\beta$ range. The $gg$-fusion contribution feebly depends on $m_{h^0}$, $m_{A^0}$, and $\tan\beta$. The situation is very different for $\phi = H^0$, as is obvious from Fig. 4. Here, $b\bar{b}$ annihilation is generally far more important than the Drell-Yan process, except for $m_{A^0} = 300$ GeV and $\tan\beta = 3$, where the latter gets close. The contribution due to $b\bar{b}$ annihilation monotonically increases with $\tan\beta$, while the one due to the Drell-Yan process decreases. Furthermore, $gg$ fusion competes with $q\bar{q}$ annihilation and even dominates for $\tan\beta \lesssim 7$.

As for $\phi = A^0$, the $b\bar{b}$-annihilation contribution (dashed lines) and the total $gg$-fusion contribution (solid lines), corresponding to the coherent superposition of quark and squark loop amplitudes, are presented separately in Fig. 5. For comparison, also the $gg$-fusion contribution due to quark loops only (dotted lines) is shown. As in the case of $\phi = H^0$,
$gg$ fusion competes with $b\bar{b}$ annihilation and even dominates for $\tan \beta \lesssim 7$. Again, the $b\bar{b}$-annihilation contribution monotonically increases with $\tan \beta$. The bulk of the $gg$-fusion contribution is due to the quark loops, especially at low values of $m_{A^0}$.

Finally, we compare our results with the literature. As already mentioned in Sec. 2, we recover the well-known SM result [17], for $\phi = H$, by taking the SM limit of our results for $\phi = h^0$ in Eqs. (6) and (8). The contribution due to $A^0$-boson exchange in Eq. (6), which is not probed in the SM limit, agrees with the analogous contribution to $gg \to W^- H^+$ given in Eq. (1) of Ref. [3] after appropriately adjusting the masses and couplings. On the other hand, the residual terms in the latter equation, which arise from the exchanges of $h^0$ and $H^0$ bosons, coincide with Eq. (9) after substituting the appropriate masses and couplings. Similarly, by adjusting masses and couplings in Eq. (10), we reproduce Eq. (2.3) in Ref. [4], which gives the squark triangle contribution to $gg \to W^- H^+$. In Ref. [23], numerical results for the cross section of $pp \to Z A^0$ via quark-loop-mediated $gg$ fusion were presented. Adopting the input parameters and proton PDF set specified in that reference, we nicely reproduce the separate contributions due triangle and box diagrams shown in Fig. 4 therein, while we fail to agree with their superposition. Furthermore, we find reasonable agreement with the cross section of $pp \to Z A^0 + X$ via $gg$ fusion represented graphically for different scenarios in Figs. 6 and 7 of Ref. [21] adopting the respective inputs from there.

4 Conclusions

We analytically calculated the cross sections of the partonic subprocesses $q\bar{q} \to Z\phi$ and $gg \to Z\phi$, where $\phi = h^0, H^0, A^0$, to LO in the MSSM. We included the Drell-Yan and Yukawa-enhanced contributions to $q\bar{q}$ annihilation (see Fig. 1) and the quark and squark loop contributions to $gg$ fusion (see Fig. 2). We presented these results as helicity amplitudes expressed in terms of standard scalar one-loop integrals.

We then quantitatively investigated the inclusive cross sections of $pp \to Z\phi + X$ at the LHC with $\sqrt{s} = 14$ GeV adopting a favorable SUGRA-inspired MSSM scenario, varying the input parameters $m_{A^0}$ and $\tan \beta$. Our results are presented in Figs. 3–5. The total cross section for $\phi = h^0$ is typically of order 1 pb, while those for $\phi = H^0, A^0$ are of order 100 fb (10 fb) for $m_{A^0} = 300$ GeV (600 GeV). Assuming design luminosity, $L = 10^{34}$ cm$^{-2}$s$^{-1}$, a cross section of 1 pb corresponds to $10^5$ events per year and experiment at the LHC (see Table I of Ref. [44]).

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A Feynman rules

In this appendix, we collect the Feynman rules used in this paper. The Feynman rules for the $Zq\bar{q}$ vertices are $ig\gamma^\mu(v_{Zq} - a_{Zq}g_5)$, with $g = e/s_w$, $e$ being the proton charge, $s_w^2 = 1 - c_w^2$, and

$$v_{Zq} = -\frac{I_q - 2s_w^2Q_q}{2c_w}, \quad a_{Zq} = -\frac{I_q}{2c_w}, \quad (A.1)$$

where $I_q = \pm 1/2$ and $Q_q = 2/3, -1/3$ are the weak hypercharge and electric charge of quark $q$, respectively. The Feynman rules for the $q\bar{q}g$ ($\phi = h^0, H^0$) and $A^0q\bar{q}$ vertices are $ig\phi_{qq}$ and $gg_{A^0qq}g_5$, respectively, with

$$g_{h^{0}tt} = \frac{m_t \cos \alpha}{2m_W \sin \beta}, \quad g_{h^{0}tt} = -\frac{m_t \sin \alpha}{2m_W \sin \beta}, \quad g_{A^{0}tt} = -\frac{m_t \cot \beta}{2m_W},$$

$$g_{h^{0}bb} = \frac{m_b \sin \alpha}{2m_W \cos \beta}, \quad g_{h^{0}bb} = -\frac{m_b \cos \alpha}{2m_W \cos \beta}, \quad g_{A^{0}bb} = -\frac{m_b \tan \beta}{2m_W}, \quad (A.2)$$

where $\alpha$ is the mixing angle that rotates the weak $CP$-even Higgs eigenstates into the mass eigenstates $h^0$ and $H^0$. The Feynman rules for the $\phi ZZ$ vertices are $ig\phi_{ZZ}g^{\mu\nu}$, with

$$g_{h^{0}ZZ} = -\frac{m_Z}{c_w} \sin (\alpha - \beta), \quad g_{H^{0}ZZ} = \frac{m_Z}{c_w} \cos (\alpha - \beta). \quad (A.3)$$

The Feynman rules for the $\phi A^{0}Z$ vertices are $gg_{\phi A^0Z}(p + p')^\mu$, where $p$ is the incoming four-momentum of the $\phi$ boson, $p'$ is the outgoing four-momentum of the $A^0$ boson, and

$$g_{h^{0}A^{0}Z} = \frac{\cos (\alpha - \beta)}{2c_w}, \quad g_{H^{0}A^{0}Z} = \frac{\sin (\alpha - \beta)}{2c_w}. \quad (A.4)$$

The Feynman rules for the $\phi q_i\bar{q}_j$ vertices are $ig\phi_{q_i\bar{q}_j}$, with

$$\begin{pmatrix}
g_{h^{0}i_1\bar{i}_1} & g_{h^{0}i_1\bar{i}_2} 
g_{h^{0}i_2\bar{i}_1} & g_{h^{0}i_2\bar{i}_2}
\end{pmatrix}
= \mathcal{M}^i \begin{pmatrix}
m_Z \sin (\alpha + \beta) \left( l_i^2 - s_w^2Q_i \right) - m_Z^2 \cos \alpha \frac{1}{2m_W \sin \beta} - m_t (\mu \sin \alpha + A, \cos \alpha) \frac{1}{2m_W \sin \beta} - m_b (\mu \cos \alpha + A, \sin \alpha) \frac{1}{2m_W \sin \beta} 
m_Z \sin (\alpha + \beta) \frac{1}{c_w} + m_Z^2 \sin \alpha \frac{1}{2m_W \cos \beta} + m_b (\mu \cos \alpha + A, \sin \alpha) \frac{1}{2m_W \cos \beta} + m_b (\mu \cos \alpha + A, \sin \alpha) \frac{1}{2m_W \cos \beta}
\end{pmatrix} (\mathcal{M}^i)^T,$$

$$\begin{pmatrix}
g_{h^{0}b_1\bar{b}_1} & g_{h^{0}b_1\bar{b}_2} 
g_{h^{0}b_2\bar{b}_1} & g_{h^{0}b_2\bar{b}_2}
\end{pmatrix}
= \mathcal{M}^b \begin{pmatrix}
m_Z \sin (\alpha + \beta) \left( l_i^2 - s_w^2Q_i \right) - m_Z^2 \cos \alpha \frac{1}{2m_W \sin \beta} - m_t (\mu \sin \alpha + A, \cos \alpha) \frac{1}{2m_W \sin \beta} - m_b (\mu \cos \alpha + A, \sin \alpha) \frac{1}{2m_W \sin \beta} 
m_Z \sin (\alpha + \beta) \frac{1}{c_w} + m_Z^2 \sin \alpha \frac{1}{2m_W \cos \beta} + m_b (\mu \cos \alpha + A, \sin \alpha) \frac{1}{2m_W \cos \beta} + m_b (\mu \cos \alpha + A, \sin \alpha) \frac{1}{2m_W \cos \beta}
\end{pmatrix} (\mathcal{M}^b)^T,$$
\begin{equation}
\begin{pmatrix}
g_{H^0i_1i_1} & g_{H^0i_1i_2} \\
g_{H^0i_2i_1} & g_{H^0i_2i_2}
\end{pmatrix}
- \mathcal{M}^\dagger \left( - \frac{m_Z \cos(\alpha + \beta)(t^2_s Q_t)}{c_w} - \frac{m_Z \sin(\alpha + A_t \sin(\alpha))}{2 m_W \sin(\beta)} \right. \\
\left. \frac{m_Z \sin(\alpha + A_t \sin(\alpha))}{c_w} - \frac{m_Z \cos(\alpha + \beta)(t^2_s Q_t)}{2 m_W \sin(\beta)} \right) (\mathcal{M}^\dagger)^T,
\end{equation}

\begin{equation}
\begin{pmatrix}
g_{H^0b_1b_1} & g_{H^0b_1b_2} \\
g_{H^0b_2b_1} & g_{H^0b_2b_2}
\end{pmatrix}
- \mathcal{M}^\dagger \left( - \frac{m_Z \cos(\alpha + \beta)(t^2_z Q_b)}{c_w} - \frac{m_Z \sin(\alpha + A_b \sin(\alpha))}{2 m_W \sin(\beta)} \right. \\
\left. \frac{m_Z \sin(\alpha + A_b \sin(\alpha))}{c_w} - \frac{m_Z \cos(\alpha + \beta)(t^2_z Q_b)}{2 m_W \sin(\beta)} \right) (\mathcal{M}^\dagger)^T, \quad (A.5)
\end{equation}

where

\begin{equation}
\mathcal{M}^\dagger = \begin{pmatrix}
\cos \theta_\tilde{q} & \sin \theta_\tilde{q} \\
-\sin \theta_\tilde{q} & \cos \theta_\tilde{q}
\end{pmatrix} \quad (A.6)
\end{equation}

are the squark mixing matrices, with \( \theta_\tilde{q} \) being the squark mixing angles.

### B Quark and squark loop form factors

In this appendix, we express the quark and squark triangle and box form factors, \( F_\Delta, \bar{F}_\Delta, \) and \( F_{\lambda\lambda|z|}, \) for \( \phi = h^0, H^0 \) and \( \phi = A^0, \) in terms of the standard scalar three- and four-point functions, which we abbreviate as \( C^{ab}_{ijk}(c) = C_0(a, b, c, m^2_A, m^2_B, m^2_C) \) and \( D^{abcd}_{ijkl}(e, f) = D_0(a, b, c, d, e, f, m^2_A, m^2_B, m^2_C, m^2_D), \) respectively. The definitions of the latter may be found in Eq. (5) of Ref. [3].

The quark triangle form factor for \( \phi = h^0, H^0 \) reads

\begin{equation}
F_\Delta(s, m^2_q) = 4 m^2_q C^{60}_{qqq}(s). \quad (B.7)
\end{equation}

The quark box form factors for \( \phi = h^0, H^0 \) read

\begin{equation}
F^{0+}_+ = 2s(t + u)C^{00}_{qqq}(s) + 2 \left( t + u + \frac{\lambda}{s} \right) \left[ (t - z)C^{20}_{qqq}(t) + (t - h)C^{00}_{qqq}(t) \right] \\
- \left[ N \left( t + u + \frac{\lambda}{s} \right) + 2m^2_q \lambda \right] D^{h00}_{qqqq}(t, u) - 4 \left( szh + m^2_q \lambda \right) D^{h00}_{qqqq}(s, t),
\end{equation}

\begin{equation}
F^{0-}_+ = \frac{(h - z - s)}{N} t(u - t) \left[ s(t + u)C^{00}_{qqq}(s) - \lambda C^{h0}_{qqq}(s) - 2m^2_q N D^{h00}_{qqqq}(t, u) \right] \\
+ 2(t + u)(t - z) \left[ 1 + \frac{t(t - u)(h - z - s)}{N(t + u)} \right] C^{20}_{qqq}(t) \\
+ \frac{2(t - h)}{N} \left[ z(u^2 - t^2 - \lambda) + (t + u)(t^2 - zh) \right] C^{h0}_{qqq}(t) \\
- (h - z - s) \left[ 2st - zh \right] + 4m^2_q(t - u) D^{h00}_{qqqq}(s, t),
\end{equation}
\[ F_{++}^1 = (t - u) \left[ \frac{z - h - s}{\sqrt{\lambda}} - \lambda_Z \right] \left[ \frac{s}{N} C_{qqq}^{00}(s) - \frac{1}{2} D_{qqq}^{h_{0z0}}(t, u) - \frac{s}{N} \left( t - \frac{2N}{t - u} \right) D_{qqq}^{h_{0z0}}(s, t) \right] \\
+ \frac{2(h - u)}{\sqrt{\lambda N}} \left( \lambda_Z \sqrt{\lambda} + t - u + \frac{2N}{h - u} \right) \left[ (h - t) C_{qqq}^{h0}(t) + (z - u) C_{qqq}^{z0}(u) \right], \]

\[ F_{+-}^1 = \frac{s}{N} \left( \frac{4s(t + u)}{\sqrt{\lambda}} + \sqrt{\lambda} - \lambda_Z(t - u) \right) C_{qqq}^{00}(s) - \frac{2s}{N} \left( \sqrt{\lambda} + \lambda_Z(t - u) \right) C_{qqq}^{hz}(s) \\
- \frac{2(t - h)}{\sqrt{\lambda N}} \left( -s(u + 3t) - 2N + (u - t)(t - z) + \lambda_Z(t - s - z) \sqrt{\lambda} \right) C_{qqq}^{h0}(t) \\
+ \frac{2(u - z)}{\sqrt{\lambda N}} \left( 3u(s - z) + th - 2z(h - 2u) - \lambda_Z(h - u) \sqrt{\lambda} \right) C_{qqq}^{z0}(u) \\
+ \frac{s}{\sqrt{\lambda N}} \left[ t(\lambda + 8zh - 4ts - 2(t + u)(z + h) \right. \\
+ \lambda_Z(-2h + 3t + u - 2z) \sqrt{\lambda} \left. \right] - 16m_q^2 N \right] D_{qqq}^{h_{0z0}}(s, t) \\
+ \frac{1}{2} \left( -\sqrt{\lambda} - \frac{16m_q^2 s}{\sqrt{\lambda}} + \lambda_Z(t - u) \right) D_{qqq}^{h_{0z0}}(t, u). \tag{B.8} \]

The quark and squark triangle form factors for \( \phi = A^0 \) read

\[ F_\Delta \left( s, m_q^2 \right) = 2 + \left( 4m_q^2 - s \right) C_{qqq}^{00}(s), \]

\[ \bar{F}_\Delta \left( s, m_q^2 \right) = 2 + 4m_q^2 C_{\bar{q}q\bar{q}}^{00}(s). \tag{B.9} \]

The quark box form factors for \( \phi = A^0 \) read

\[ F_{++}^0 = 2s(t + u)C_{qqq}^{00}(s) + 2 \left( t + u + \frac{\lambda}{s} \right) \left[ (t - z)C_{qqq}^{z0}(t) + (t - h)C_{qqq}^{h0}(t) \right] \\
- \left[ N \left( t + u + \frac{\lambda}{s} \right) + 2m_q^2 \lambda \right] D_{qqq}^{h_{0z0}}(t, u) - 4 \left( szh + m_q^2 \lambda \right) D_{qqq}^{h_{0z0}}(s, t), \]

\[ F_{+-}^0 = \frac{2 + (t - u)^2}{N} \left[ s(t + u)C_{qqq}^{00}(s) - \lambda C_{qqq}^{hz}(s) \right] \\
- 2 \left[ 3t - u + \frac{t}{N}(t - u)^2 \right] \left[ (t - z)C_{qqq}^{z0}(t) + (t - h)C_{qqq}^{h0}(t) \right] \\
- 2 \left[ zN - m_q^2 \lambda \right] D_{qqq}^{h_{0z0}}(t, u) \\
+ 2 \left[ st \left[ 3t - u + \frac{t}{N}(t - u)^2 \right] + 2m_q^2 \lambda \right] D_{qqq}^{h_{0z0}}(s, t), \]

\[ F_{++}^1 = \left( \frac{z - h - s}{\sqrt{\lambda}} - \lambda_Z \right) \left\{ (t - u) \left( \frac{s}{N} C_{qqq}^{00}(s) - \frac{1}{2} D_{qqq}^{h_{0z0}}(t, u) \right) \\
- s \left[ 2 + \frac{t}{N}(t - u) \right] D_{qqq}^{h_{0z0}}(s, t) \right\} \\
+ 2(t - z) \left\{ \lambda_Z \frac{h - t}{N} + \frac{1}{\sqrt{\lambda}} \left[ 2 + \frac{(t - u)(t - h)}{N} \right] \right\} C_{qqq}^{z0}(t) \]
\[-2(t-h)\left\{\lambda Z \frac{h-u}{N} + \frac{1}{\sqrt{\lambda}} \left[2 - \frac{(t-u)(u-h)}{N}\right]\right\} C_{qgq}^{h0}(t),\]

\[F_{++}^1 = \left(\lambda Z - \frac{t-u}{\sqrt{\lambda}}\right) \left[\frac{s}{N} (s-z+h) \left(C_{qqq}^{00}(s) - iD_{qqgq}^{h00}(s,t)\right) + \frac{s+z-h}{2} D_{qqgq}^{h00}(t,u)\right]\]

\[-2(t-z)\left\{\lambda Z \frac{t-h}{N} - \frac{1}{\sqrt{\lambda}} \left[2 + \frac{(t-u)(t-h)}{N}\right]\right\} C_{sgq}^{z0}(t)\]

\[-2(t-h)\left\{\lambda Z \frac{u-h}{N} + \frac{1}{\sqrt{\lambda}} \left[2 - \frac{(t-u)(u-h)}{N}\right]\right\} C_{qgq}^{h0}(t).\]  

(B.10)

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Figure 1: Tree-level Feynman diagrams for $q\bar{q} \rightarrow Z\phi$, with (a) $\phi = h^0, H^0$ and (b) $\phi = A^0$, in the MSSM.
Figure 2: One-loop Feynman diagrams for $gg \rightarrow Z\phi$, with (a) $\phi = h^0, H^0$ and (b) $\phi = A^0$, due to virtual quarks and squarks in the MSSM.
Figure 3: Total cross sections $\sigma$ (in pb) of $pp \rightarrow Zh^0 + X$ via $q\bar{q}$ annihilation (dashed lines) and $gg$ fusion (solid lines) at the LHC (a) as functions of $m_{h^0}$ for $\tan\beta = 3$ and 30; and (b) as functions of $\tan\beta$ for $m_{A^0} = 300$ GeV and 600 GeV. For comparison, also the Drell-Yan contribution to $q\bar{q}$ annihilation (dotted lines) is shown.
Figure 3 (Continued).

(b)
Figure 4: Total cross sections $\sigma$ (in pb) of $pp \rightarrow ZH^0 + X$ via $q\bar{q}$ annihilation (dashed lines) and $gg$ fusion (solid lines) at the LHC (a) as functions of $m_{H^0}$ for $\tan\beta = 3$ and $30$; and (b) as functions of $\tan\beta$ for $m_{A^0} = 300$ GeV and 600 GeV. For comparison, also the Drell-Yan contribution to $q\bar{q}$ annihilation (dotted lines) is shown.
Figure 4 (Continued).
Figure 5: Total cross sections $\sigma$ (in pb) of $pp \to ZA^0 + X$ via $b\bar{b}$ annihilation (dashed lines) and $gg$ fusion (solid lines) at the LHC (a) as functions of $m_{A^0}$ for $\tan\beta = 3$ and 30; and (b) as functions of $\tan\beta$ for $m_{A^0} = 300$ GeV and 600 GeV. For comparison, also the quark loop contribution to $gg$ fusion (dotted lines) is shown.
Figure 5 (Continued).