

Diffraction as a CP and Lineshape Analyzer for MSSM Higgs Bosons at the LHC

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ABSTRACT

We study the production and decay of a coupled system of mixed neutral MSSM Higgs bosons in exclusive double-diffractive processes at the LHC, including non-vanishing CP phases in the soft supersymmetry-breaking gaugino masses and third-generation trilinear squark couplings. The three neutral Higgs bosons are naturally nearly degenerate, for large values of $\tan\beta$, when the charged Higgs boson weighs around 150 GeV. Large mixing between all three neutral Higgs bosons is possible when CP is violated, a three-way mixing scenario which we also term tri-mixing. A resolution in the Higgs mass of ~ 1 GeV, which may be achievable using the missing-mass method, would allow one to distinguish nearly degenerate Higgs bosons by studying the production lineshape. Measurements of the polarizations of the tau leptons coming from the Higgs-boson decays could offer a direct and observable signal of CP violation in the Higgs sector.

1 Introduction

Direct search experiments at LEP established that, within the Standard Model, the single physical Higgs boson must weigh more than about 114 GeV [1], and the mass range favoured by indirect measurements is close to this lower limit [2]. Both the direct and indirect limits must be re-examined in non-minimal Higgs scenarios. One popular example is the minimal supersymmetric extension of the Standard Model (MSSM) [3], which predicts the existence of at least one light neutral Higgs boson weighing less than about 135 GeV. This may be joined by two other light neutral Higgs bosons if the charged Higgs bosons H^\pm are also light. If CP is conserved, only the two CP-even neutral Higgs bosons h, H may mix with each other dynamically, through off-diagonal absorptive self-energy effects. Instead, if CP is violated, all the three neutral Higgs bosons, including the CP-odd neutral Higgs boson A , mix through CP-violating quantum effects [4–8] to mass eigenstates, $H_{1,2,3}$, of indefinite CP. These CP-mixed Higgs states give rise to a coupled system whose resonant dynamics is properly described by a 3-by-3 propagator matrix [9, 10]. One very interesting phenomenological feature of the CP-violating MSSM is that the lightest neutral Higgs boson could be considerably lighter than 114 GeV [11, 12].

Inclusive experiments at the LHC will be able to discover the Standard Model Higgs boson, whatever its mass. They should also be able to discover the lightest neutral MSSM Higgs boson — at least if CP is conserved — and will be able to explore significant mass ranges for the heavier MSSM Higgs bosons. Supplementing inclusive searches, interest has recently been growing in the search for Higgs bosons in diffractive events at the LHC [13, 14]. These may offer novel prospects for measuring the properties of light neutral Higgs bosons [15], and disentangling their CP properties [16].

The MSSM offers additional sources of CP violation beyond the single Kobayashi–Maskawa phase in the Standard Model. If the soft supersymmetry-breaking parameters $m_0, m_{1/2}$ and A are universal, two new physical CP-odd phases remain. These may, without loss of generality, be parametrized as one phase in the trilinear couplings A and one in the gaugino masses $m_{1/2}$. In addition to signatures of CP violation in sparticle production and decay at high-energy colliders [17–19], these phases may have observable radiative effects on the Higgs sector [4], on electric dipole moments [20–22], and in B decays [23, 24]. One of the principal motivations for studying such models is the prospect of electroweak baryogenesis in the MSSM [25].

In this paper, we consider the prospects for studying light neutral MSSM Higgs bosons in diffractive events at the LHC, particularly in scenarios where CP is violated [4–8, 26–30], and the three neutral Higgs bosons mix strongly. This work continues previous studies of the

masses, couplings, production and decays of the mixed-CP Higgs bosons $H_{1,2,3}$, with a view to searches at LEP [11], the LHC [11,12,15,16,31–33], the ILC [34], a $\mu^+\mu^-$ collider [35] and a $\gamma\gamma$ collider [36–40]. As in our previous works [10,40], we include a complete treatment of loop-induced CP violation and three-way mixing, including off-diagonal absorptive effects in the resummed Higgs-boson propagator matrix [9].

Higgs-boson production in an exclusive diffractive collision $p + p \rightarrow p + H_i + p$, where the outgoing protons remain intact and scatter through small angles, offers a unique environment for investigating the MSSM Higgs sector, in particular when $\tan\beta$ is large and $M_{H^\pm}^{\text{pole}} \sim 150$ GeV. In such scenarios, all three neutral Higgs bosons have similar masses and there is strong three-way mixing. We call a scenario with these properties the *three-way mixing* or *tri-mixing* scenario of the CP-violating MSSM.

Moreover, the production cross section of the Higgs boson can be much enhanced in the MSSM for large values of $\tan\beta$, as compared with the Standard Model. Furthermore, good Higgs-mass resolution of the order of 1 GeV may be achievable [13] by precise measurements of the momenta of the outgoing protons in detectors a long way downstream of the interaction point. This enables one to disentangle nearly degenerate Higgs bosons by examining the production lineshape of the coupled system of neutral Higgs bosons.

The layout of this paper is as follows. Section 2 provides basic formulae for the luminosity of the exclusive double-diffractive process, based on [41,42]. Section 3 presents the formalism for the production, mixing and decay of a coupled system of CP-violating neutral MSSM Higgs bosons in diffraction, based on [10]. It also presents numerical example in a couple of CP-violating scenarios, considering two final states: $\bar{b}b$ and $\tau^+\tau^-$. Our conclusions are given in Section 4.

2 Luminosity for the Exclusive Double Diffractive Process

The effective luminosity for producing via double diffraction a system of invariant mass M and rapidity y can be written as [41]

$$M^2 \frac{\partial^2 \mathcal{L}}{\partial y \partial M^2} = \hat{S}^2 L, \quad (1)$$

where the ‘soft’ survival factor \hat{S}^2 is quite model- and process-dependent. Denoting the longitudinal momenta of two gluons that fuse into a system with invariant mass M by $x_1 p_1$

and $x_2 p_2$, we have

$$x_1 = \frac{M}{\sqrt{s}} e^y, \quad x_2 = \frac{M}{\sqrt{s}} e^{-y}, \quad (2)$$

where $s \equiv (p_1 + p_2)^2$ is the collider centre-of-mass energy squared, see Fig. 1.

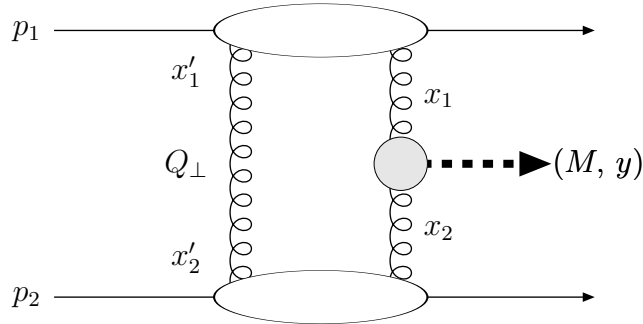


Figure 1: *The mechanism contributing to the exclusive double-diffractive process.*

Assuming $J_z = 0$, a parity-even central system and negligibly small perpendicular momenta of the outgoing protons: $p_\perp \ll 1$ GeV, the expression for L becomes [41, 42]

$$L = \left[\frac{\pi}{(N_C^2 - 1)b} \int_{Q_{\min}^2}^{\mu^2} \frac{dQ_T^2}{Q_T^4} f_g(x_1, x'_1, Q_T^2, \mu^2) f_g(x_2, x'_2, Q_T^2, \mu^2) \right]^2, \quad (3)$$

where $N_C = 3$ and Q_T^2 is the virtuality of the soft gluon needed for colour screening. The hard scale μ and the t -slope parameters are chosen as follows: $\mu = M/2$ and $b = 4$ GeV⁻². Formally, we also have introduced a non-vanishing cutoff $Q_{\min} \lesssim 1$ GeV to avoid encountering the Landau pole, but the sensitivity of L to Q_{\min} turns out to be not large. Furthermore, if $x' \ll x$ (which is actually the case for the exclusive diffractive process [41, 42]), the skewed or off-diagonal unintegrated gluon density $f_g(x, x', Q_T^2, \mu^2)$ may take on the factorizable form:

$$f_g(x, x', Q_T^2, \mu^2) \simeq R_g \tilde{f}_g(x, Q_T^2, \mu^2), \quad (4)$$

where R_g is a constant. The simplified form (4) is estimated to have an accuracy of 10 to 20%, where the function $\tilde{f}_g(x, Q_T^2, \mu^2)$ is defined by

$$\begin{aligned} \tilde{f}_g(x, Q_T^2, \mu^2) &\equiv \frac{\partial}{\partial \ln Q_T^2} \left[\sqrt{T(Q_T, \mu)} x g(x, Q_T^2) \right] \\ &= \frac{1}{2} \sqrt{T(Q_T, \mu)} x Q_T \left[\frac{dg(x, Q_T^2)}{dQ_T} - \frac{g(x, Q_T^2)}{2} \frac{dS(Q_T, \mu)}{dQ_T} \right]. \end{aligned} \quad (5)$$

We note that $\tilde{f}_g(x, Q_T^2, \mu^2)$ consists of derivatives of the gluon distribution function $g(x, Q_T^2)$ and the Sudakov factor $T(Q_T, \mu) = e^{-S(Q_T, \mu)}$. Collecting all the factors, the effective luminosity can be rewritten as

$$M^2 \frac{\partial^2 \mathcal{L}}{\partial y \partial M^2} = \hat{S}^2 \left[\frac{\pi R_g^2}{8b} \int_{\ln Q_{\min}}^{\ln \mu} F_g(x_1, x_2, Q_T, \mu) d \ln Q_T \right]^2, \quad (6)$$

where the integrand functions $F_g(x_1, x_2, Q_T, \mu)$ is defined by

$$F_g(x_1, x_2, Q_T, \mu) \equiv 2 \frac{\tilde{f}_g(x_1, Q_T^2, \mu^2) \tilde{f}_g(x_2, Q_T^2, \mu^2)}{Q_T^2}, \quad (7)$$

for the double-diffractive process.

For our numerical analysis, we also need to know the parton distribution functions (PDFs) for $g(x, Q_T^2)$ in $\tilde{f}_g(x, Q_T^2, \mu^2)$. Specifically, we take the PDFs given by CTEQ6M [44] and MRST2004NNLO [45]. The function $S(Q_T, \mu)$ that determines the Sudakov factor $T(Q_T, \mu)$ can be calculated as

$$S(Q_T, \mu) = \frac{1}{2\pi} \int_{Q_T^2}^{\mu^2} \frac{\alpha_s(k_t^2)}{k_t^2} \left[\mathcal{F}_g(k_t) + \sum_{q=u,d,s,c,\bar{u},\bar{d},\bar{s},\bar{c}} \mathcal{F}_q(k_t) \right] dk_t^2, \quad (8)$$

where

$$\begin{aligned} \mathcal{F}_g(k_t) &= \int_0^{1-\Delta(k_t)} z P_{gg}(z) dz = -11/2 - 6 \ln \Delta + 12\Delta - 9\Delta^2 + 4\Delta^3 - 3\Delta^4/2, \\ \mathcal{F}_q(k_t) &= \int_0^{1-\Delta(k_t)} P_{qg}(z) dz = 1/3 - \Delta/2 + \Delta^2/2 + \Delta^3/3, \end{aligned} \quad (9)$$

with $\Delta = \Delta(k_t) \equiv k_t/(\mu + k_t)$.

Equations (1)–(9) provide a complete basis for our numerical evaluation of the effective luminosity for exclusive double-diffractive Higgs production. In particular, we find that the effective luminosity may conveniently be computed as follows:

$$M^2 \frac{\partial^2 \mathcal{L}}{\partial y \partial M^2} = 4.0 \times 10^{-4} \left[\frac{\int_{\ln Q_{\min}}^{\ln \mu} F_g(x_1, x_2, Q_T, \mu) d \ln Q_T}{\text{GeV}^{-2}} \right]^2 \left(\frac{\hat{S}^2}{0.02} \right) \left(\frac{4}{b \text{ GeV}^2} \right)^2 \left(\frac{R_g}{1.2} \right)^4. \quad (10)$$

We note that the luminosity is very sensitive to the choice of R_g , in particular, but less to the hard scale μ (see also our discussion below).

In Fig. 2, we show $xg(x, Q_T^2)$ (upper left), the Sudakov factor $T(Q_T, \mu)$ (upper right), $\tilde{f}_g(x, Q_T^2, \mu^2)$ (lower left), and $F_g(x_1, x_2, Q_T)$ (lower right), see (7), as functions of Q_T or $\ln Q_T$ when the rapidity $y = 0$, $M = 120$ GeV, and $\sqrt{s} = 14$ TeV. The solid (dashed) lines

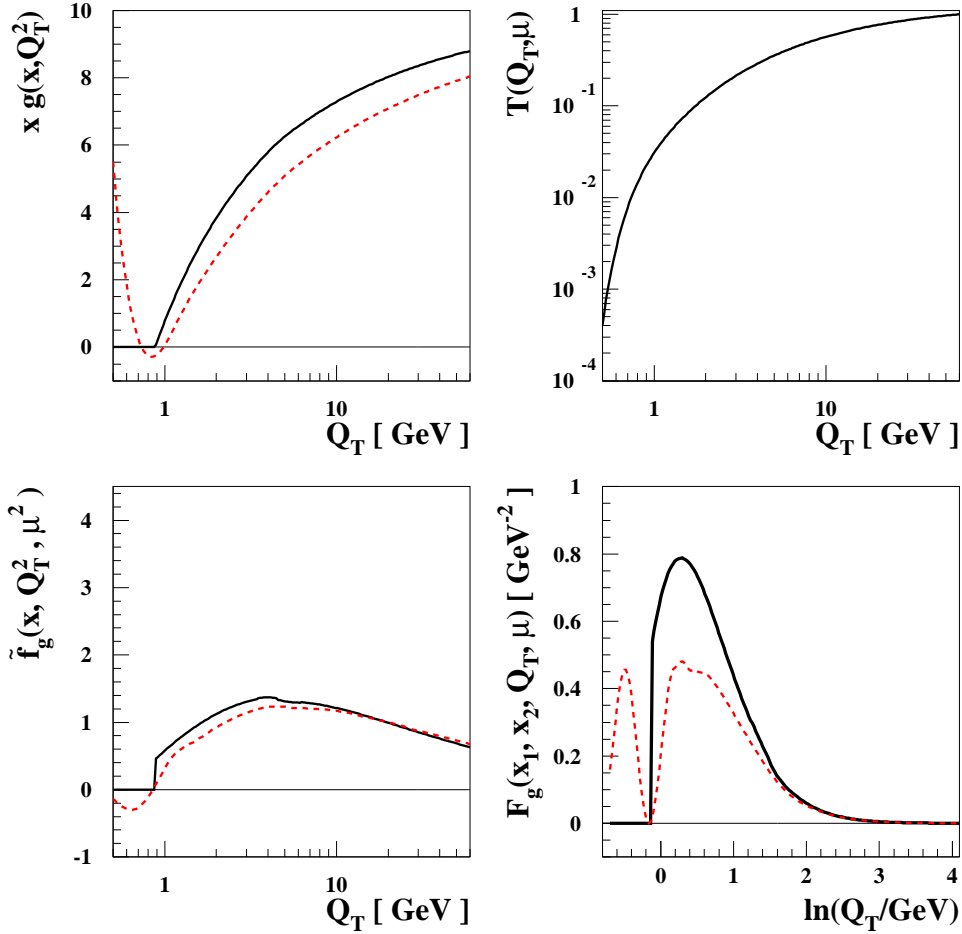


Figure 2: The quantities $xg(x, Q_T^2)$ (upper left), the Sudakov factor $T(Q_T, \mu)$ (upper right), $\tilde{f}_g(x, Q_T^2, \mu^2)$ (lower left), and $F_g(x_1, x_2, Q_T, \mu)$ (lower right), as functions of Q_T or $\ln Q_T$ when the rapidity $y = 0$, $M = 120$ GeV, and $\sqrt{s} = 14$ TeV. The solid (dashed) lines are the outputs of CTEQ6M (MRST2004NNLO).

are outputs of CTEQ6M (MRST2004NNLO). We observe significant differences between the functions $xg(x, Q_T^2)$, which results in a strong dependence on the parton distribution function used. Specifically, MRST2004 gives a non-vanishing and partly negative $xg(x, Q_T^2)$ when $Q_T \lesssim 0.9$ GeV, where $\tilde{f}_g(x_{1,2}, Q_T^2, \mu^2)$ becomes negative due to the rapidly decreasing $xg(x, Q_T^2)$. On the other hand, CTEQ6M returns $g(x, Q_T^2) = 0$ in this region. For $0.9 \text{ GeV} < Q_T < 10 \text{ GeV}$, MRST2004 returns smaller values for $F_g(x_1, x_2, Q_T, \mu)$ than CTEQ6M does. When $\ln(Q_T/\text{GeV}) > 3$ ($Q_T \gtrsim 20$ GeV), the contribution from this large- Q_T region to the luminosity is negligible, making the prediction insensitive to the choice of $\mu \gtrsim 20$ GeV. We note that the most significant contribution comes from the region around $Q_T = 1.3$ GeV, as seen in the lower-right frame for $F_g(x_1, x_2, Q_T, \mu)$.

In fact, the calculations of $g(x, Q_T^2)$ in both the CTEQ6 and MRST2004 parameterizations are not reliable when $Q_T \lesssim 1$ GeV. There are some phenomenological prescrip-

tions for making the low- Q_T^2 behaviour more sensible [43]. For Higgs bosons with masses $M \sim 120$ GeV, however, the luminosity does not depend strongly on the variation between the prescriptions, and we take $Q_{\min} = 1$ GeV to avoid unphysical effects associated with the inapplicability of the $g(x, Q_T^2)$ calculation for $Q_T \lesssim 1$ GeV.

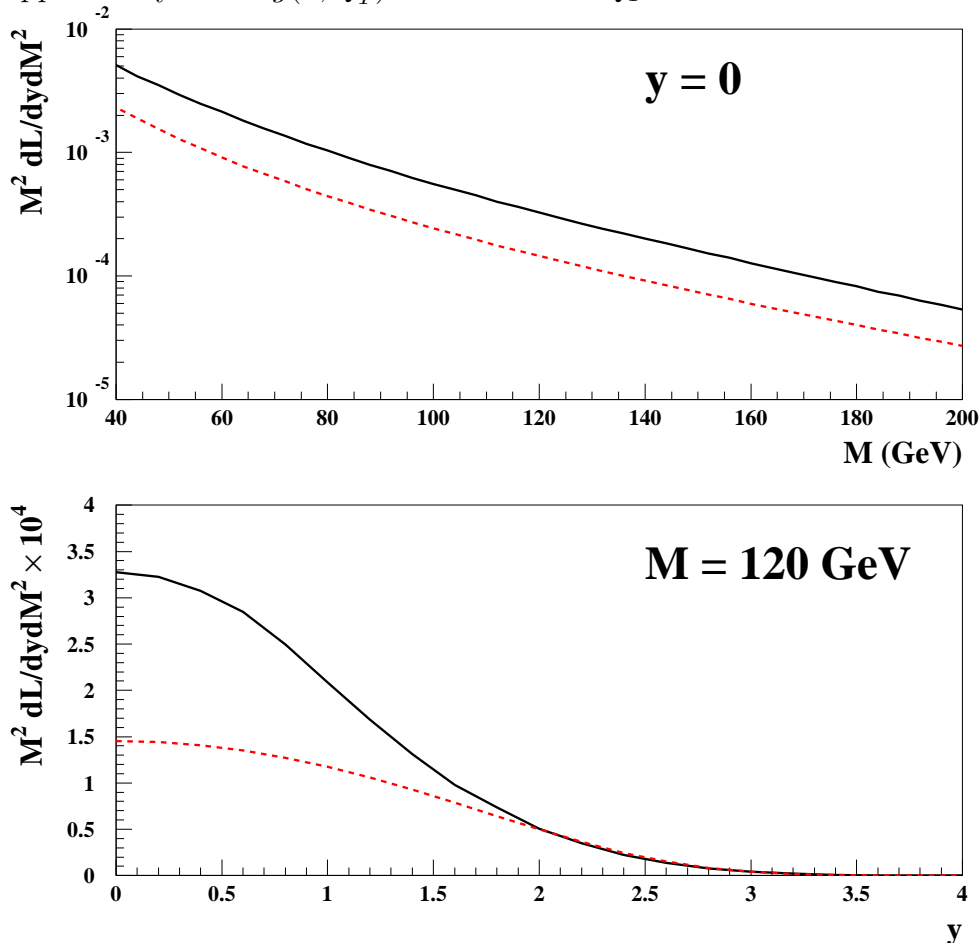


Figure 3: The effective luminosity for the exclusive double-diffractive process for $\sqrt{s} = 14$ TeV as a function of M ($y = 0$) and y ($M = 120$ GeV) in the upper (lower) frame, respectively. Solid (dashed) lines are obtained using CTEQ6M (MRST2004NNLO), taking $Q_{\min} = 1.0$ GeV in each case. Predictions using other PDFs lie between the above two extreme cases, according to the analysis in [42].

In Fig. 3, we show the effective luminosity for the exclusive double-diffractive process at the LHC as a function of M for $y = 0$ (upper frame) and as a function of y for $M = 120$ GeV (lower). The solid (black) lines are obtained using CTEQ6M and the dashed (red) lines by MRST2004NNLO. We take $Q_{\min} = 1.0$ GeV, $\hat{S}^2 = 0.02$, and $R_g = 1.2$ [41], see (10). We note large differences between the two predictions over large regions of M and y , which should be regarded as two extreme cases. For example, at $(y, M) = (0, 120$ GeV), the CTEQ6M prediction is about twice as large as that of MRST2004NNLO, and we

find that the prediction of CTEQ6M has a stronger dependence on y than does that of MRST2004NNLO. We remark that the predictions using other PDFs lie between the above two extreme cases, as discussed in [42].

3 The Process $pp \rightarrow p + H_i + p \rightarrow p + [f(\sigma)\bar{f}(\bar{\sigma})] + p$

The helicity amplitude for the process $g_1^a(\lambda_1)g_2^b(\lambda_2) \rightarrow H_i \rightarrow f(\sigma)\bar{f}(\bar{\sigma})$ in the double-diffractive production of Higgs bosons H_i is given by

$$\mathcal{M}(\sigma\bar{\sigma}; \lambda_1\lambda_2) = \frac{g\alpha_s m_f \sqrt{\hat{s}} \delta^{ab}}{8\pi v M_W} \langle \sigma \rangle_f \delta_{\sigma\bar{\sigma}} \delta_{\lambda_1\lambda_2}, \quad (11)$$

where the amplitude $\langle \sigma \rangle_f$ is defined as

$$\langle \sigma \rangle_f \equiv \sum_{i,j=1,2,3} S_i^g(\sqrt{\hat{s}}) D_{ij}(\hat{s}) (\sigma\beta_f g_{H_j\bar{f}f}^S - i g_{H_j\bar{f}f}^P). \quad (12)$$

For the $J_z^P = 0^+$ process with $p_\perp \approx 0$ that we consider, the pseudoscalar form factor of the g - g - H_i vertex does not contribute, making the helicity amplitude independent of the helicities of gluons. For the definitions of the couplings, the threshold corrections that are enhanced for large values of $\tan\beta$ for $f = b, \tau$, and the full 3×3 propagator matrix $D_{ij}(\hat{s})$, we refer to [10, 46].

Similarly as in [10], one can define the parton-level cross section as:

$$\hat{\sigma}_i^f \equiv 2(N_C^2 - 1) \frac{N_f \beta_f}{512\pi \hat{s}} \left(\frac{g\alpha_s m_f \sqrt{\hat{s}}}{8\pi v M_W} \right)^2 C_i^f \quad (13)$$

where the enhancement factor $2(N_C^2 - 1)$ for the exclusive process and color factors $N_f : N_l = 1$ and $N_q = 3$ have been included. The coefficients C_i^f are given in terms of the amplitudes $\langle \sigma \rangle_f$:

$$\begin{aligned} C_1^f &= \frac{1}{2}(|\langle + \rangle_f|^2 + |\langle - \rangle_f|^2), & C_2^f &= \frac{1}{2}(|\langle + \rangle_f|^2 - |\langle - \rangle_f|^2), \\ C_3^f &= -\Re(\langle + \rangle_f \langle - \rangle_f^*), & C_4^f &= \Im(\langle + \rangle_f \langle - \rangle_f^*). \end{aligned} \quad (14)$$

For our numerical results, we make the following choices of parameters:

$$\begin{aligned} \tan\beta &= 50, & M_{H^\pm}^{\text{pole}} &= 155 \text{ GeV}, \\ M_{\tilde{Q}_3} &= M_{\tilde{U}_3} = M_{\tilde{D}_3} = M_{\tilde{L}_3} = M_{\tilde{E}_3} = M_{\text{SUSY}} &= 0.5 \text{ TeV}, \\ |\mu| &= 0.5 \text{ TeV}, & |A_{t,b,\tau}| &= 1 \text{ TeV}, & |M_2| &= |M_1| = 0.3 \text{ TeV}, & |M_3| &= 1 \text{ TeV}, \\ \Phi_\mu &= 0^\circ, & \Phi_A &= \Phi_{A_t} = \Phi_{A_b} = \Phi_{A_\tau} = 90^\circ, & \Phi_1 &= \Phi_2 = 0^\circ, \end{aligned} \quad (15)$$

and we consider two values for the phase of the gluino mass parameter M_3 : $\Phi_3 = -10^\circ, -90^\circ$. For $\Phi_3 = -10^\circ$, CPsuperH [46] yields for the masses and widths of the neutral Higgs bosons:

$$\begin{aligned} M_{H_1} &= 120.2 \text{ GeV}, & M_{H_2} &= 121.4 \text{ GeV}, & M_{H_3} &= 124.5 \text{ GeV}, \\ \Gamma_{H_1} &= 1.19 \text{ GeV}, & \Gamma_{H_2} &= 3.42 \text{ GeV}, & \Gamma_{H_3} &= 3.20 \text{ GeV}, \end{aligned} \quad (16)$$

and for $\Phi_3 = -90^\circ$:

$$\begin{aligned} M_{H_1} &= 118.4 \text{ GeV}, & M_{H_2} &= 119.0 \text{ GeV}, & M_{H_3} &= 122.5 \text{ GeV}, \\ \Gamma_{H_1} &= 3.91 \text{ GeV}, & \Gamma_{H_2} &= 6.02 \text{ GeV}, & \Gamma_{H_3} &= 6.34 \text{ GeV}, \end{aligned} \quad (17)$$

respectively. In the above MSSM scenario, which has originally been introduced in [10, 40], all the three neutral CP-violating Higgs bosons are nearly degenerate, with masses of ~ 120 GeV, and mix strongly through off-diagonal absorptive self-energy effects. We call a scenario with these properties, which can only be realized in a CP-violating MSSM, the *three-way mixing* or *tri-mixing* scenario.

For $f = b$, the polarization of the b quark and the \bar{b} anti-quark cannot be measured, and the only observable is the total cross section, which is given by

$$M^2 \frac{\partial^2 \sigma_{\text{tot}}^b}{\partial y \partial M^2} = 4K \hat{\sigma}_1^b M^2 \frac{\partial^2 \mathcal{L}}{\partial y \partial M^2}, \quad (18)$$

where we include the perturbative QCD correction $K \equiv 1 + \frac{\alpha_s(\hat{s})}{\pi}(\pi^2 + 11/2)$ [47]. The factor 4 comes from the sum over the b - and \bar{b} -quark polarizations [15].

Figure 4 displays two numerical examples of three-way mixing scenarios with $\Phi_3 = -90^\circ$ (solid lines) and $\Phi_3 = -10^\circ$ (dashed lines). In this case, we have taken CTEQ6M and $Q_{\text{min}} = 1$ GeV. The differential cross section becomes as large as ~ 13 fb for $\Phi_3 = -90^\circ$, in which case it exhibits a single peak located between the pole masses, that are indicated by the solid vertical lines. This peak is broader than the expected missing-mass resolution $\delta M \sim 1$ GeV. On the other hand, twin peaks are discernible when $\Phi_3 = -10^\circ$, thanks to the expected good resolution in the Higgs mass: $\delta M \sim 1$ GeV. The twin peaks appear close to (but not at) the outer pair of pole masses, indicated by the vertical dashed lines, and the cross section actually exhibits a dip at the second pole mass. These examples demonstrate that the unique sensitivity to CP-conserving observables provided by the good mass resolution in double-diffraction events would in turn provide sensitivity to the CP-violating phase Φ_3 .

For $f = \tau$, we have four observables which can be constructed from $\hat{\sigma}_i^\tau$ with $i = 1 - 4$.

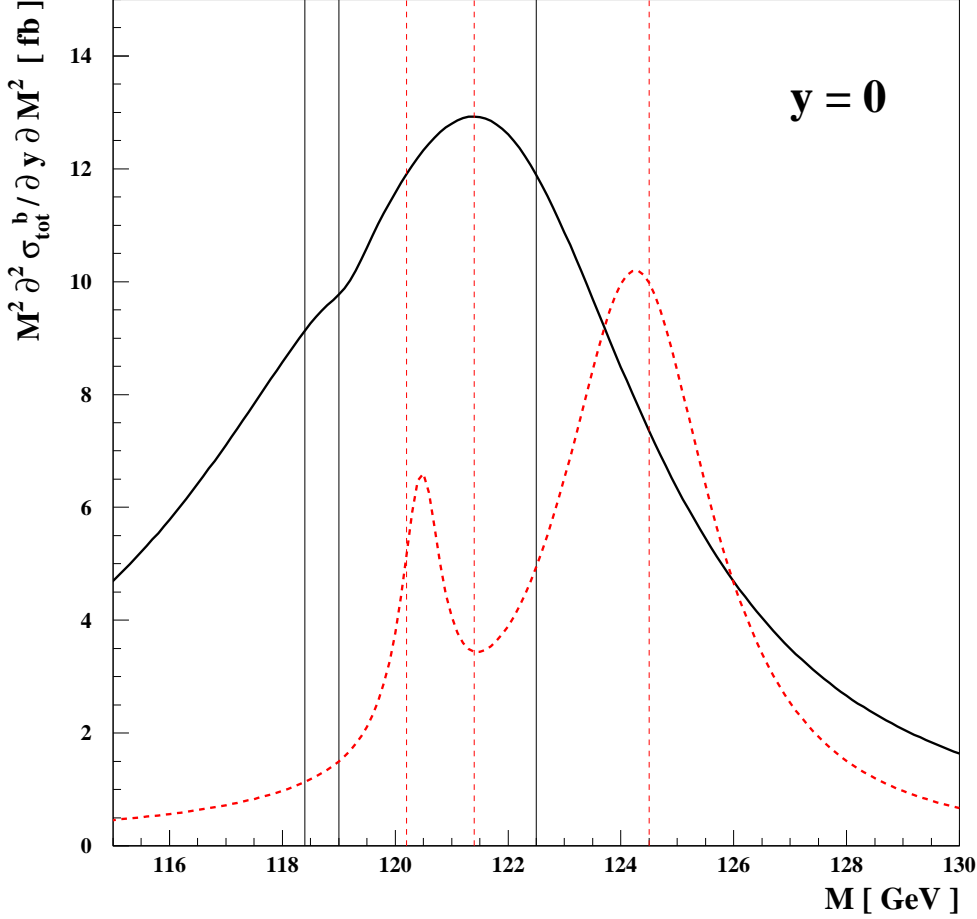


Figure 4: The hadronic level cross section $M^2 \frac{\partial^2 \sigma_{\text{tot}}^b}{\partial y \partial M^2}$ ($y = 0$) when the produced Higgs bosons decay into b quarks, calculated using CTEQ6M PDFs. Tri-mixing scenarios have been taken with $\Phi_3 = -90^\circ$ (solid lines) and $\Phi_3 = -10^\circ$ (dotted lines). The vertical lines indicate the three Higgs-boson pole-mass positions.

In particular, to analyze the signatures of CP violation in the production of longitudinally-polarized τ leptons, we define

$$\Delta\sigma_{\text{CP}}^\tau \equiv \sigma(pp \rightarrow p H_i p; H_i \rightarrow \tau_R^+ \tau_L^-) - \sigma(pp \rightarrow p H_i p; H_i \rightarrow \tau_L^+ \tau_R^-). \quad (19)$$

We then have the total and CP-violating cross sections given by

$$\begin{aligned} M^2 \frac{\partial^2 \sigma_{\text{tot}}^\tau}{\partial y \partial M^2} &= 4K \hat{\sigma}_1^\tau M^2 \frac{\partial^2 \mathcal{L}}{\partial y \partial M^2}, \\ M^2 \frac{\partial^2 \Delta\sigma_{\text{CP}}^\tau}{\partial y \partial M^2} &= 4K \hat{\sigma}_2^\tau M^2 \frac{\partial^2 \mathcal{L}}{\partial y \partial M^2}. \end{aligned} \quad (20)$$

Finally, we define the CP-violating asymmetry a_{CP}^τ as

$$a_{\text{CP}}^\tau \equiv \frac{M^2 \frac{\partial^2 \Delta\sigma_{\text{CP}}^\tau}{\partial y \partial M^2}}{M^2 \frac{\partial^2 \sigma_{\text{tot}}^\tau}{\partial y \partial M^2}} = \frac{\hat{\sigma}_2^\tau}{\hat{\sigma}_1^\tau}. \quad (21)$$

Figure 5 displays numerical examples in tri-mixing scenarios with $\Phi_3 = -90^\circ$ (solid lines) and $\Phi_3 = -10^\circ$ (dotted lines). We have again taken CTEQ6M and $Q_{\min} = 1$ GeV, and find cross sections as large as $\sim 1.5 - 3$ fb. The peaks and dips in the total cross section are located relative to the pole masses in the same way as for the $\bar{b}b$ final state, again offering sensitivity to the CP-violating phase. Moreover, comparing Figs. 4 and 5, we see that the relative sizes of the peaks are different for the two values of Φ_3 , thus providing more sensitivity to this CP-violating parameter.

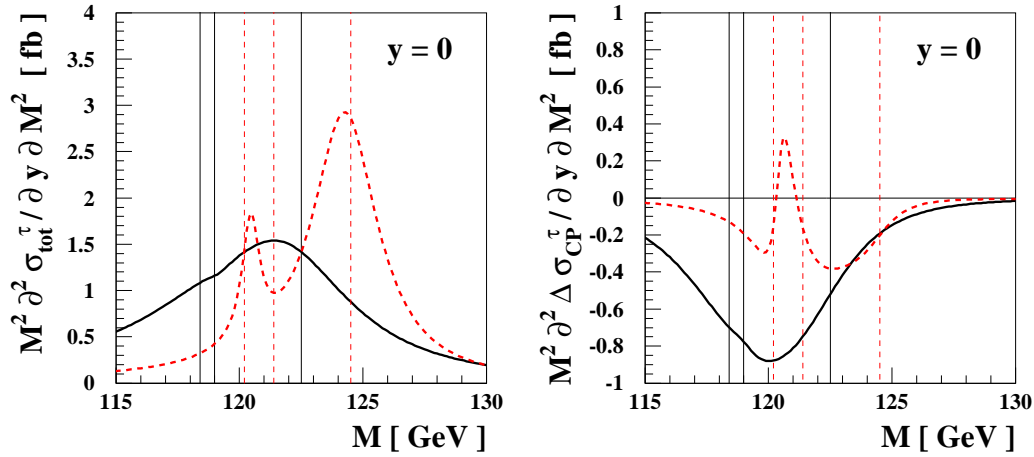


Figure 5: The hadron-level CP-conserving and CP-violating cross sections when Higgs bosons decay into τ leptons: $M^2 \frac{\partial^2 \sigma_{\text{tot}}^\tau}{\partial y \partial M^2}$ (left) and $M^2 \frac{\partial^2 \Delta \sigma_{\text{CP}}^\tau}{\partial y \partial M^2}$ (right), calculated using CTEQ6M PDFs. We have considered tri-mixing scenarios with $\Phi_3 = -90^\circ$ (solid lines) and $\Phi_3 = -10^\circ$ (dotted lines). The vertical lines indicate the three Higgs-boson pole-mass positions.

The CP-violating cross-section difference observable in the $\tau^+ \tau^-$ final state is shown in the right panel of Fig. 5. In the case $\Phi_3 = -90^\circ$, the cross-section difference always has the same sign and is maximized between the pole masses, whereas in the case $\Phi_3 = -10^\circ$ it is generally smaller and exhibits two sign changes. The expected missing-mass resolution $\delta M \sim 1$ GeV should be sufficient to resolve some of these structures. We expect that a CP asymmetry a_{CP}^τ larger than 10 % may be detected with an integrated luminosity $\gtrsim 100 \text{ fb}^{-1}$. As we see in Fig. 6, the CP-violating observable a_{CP}^τ attains values considerably larger than 10% in both the scenarios studied. Whereas the individual cross sections depend on the PDFs used, the CP asymmetry shown in Fig. 6 is insensitive to this choice.

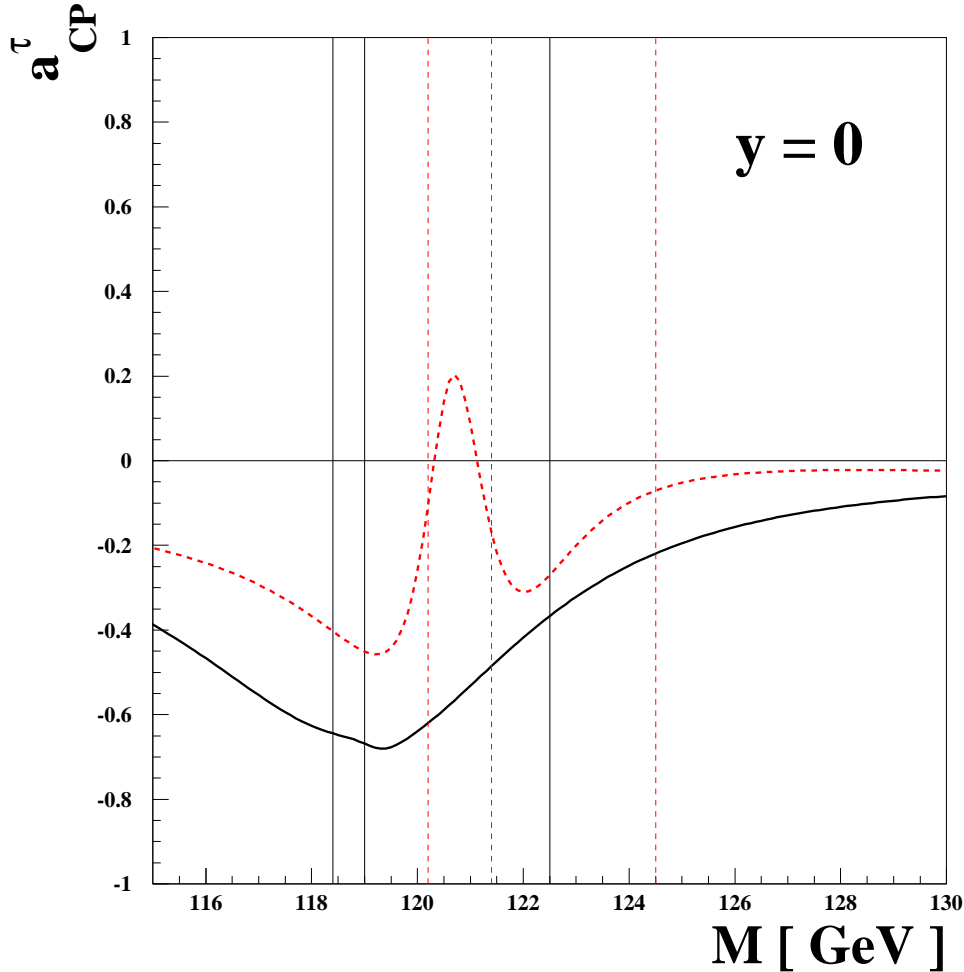


Figure 6: The CP-violating asymmetry a_{CP}^{τ} observable in tri-mixing scenarios when Higgs bosons decay into τ leptons. The line styles are the same as in Fig. 5.

4 Conclusions

We have extended our previous studies of CP-violating MSSM Higgs scenarios with large tri-mixing [10, 40] to diffractive production at the LHC. The production cross sections are large compared to those in the Standard Model, and the good Higgs-mass resolution obtainable via the missing-mass method should enable one to disentangle the different adjacent resonant peaks. Although it is difficult to construct CP-violating observables in $H_i \rightarrow \bar{b}b$ decays without tagging the final protons [16] or analyzing the b -quark decay products, observations of τ^{\pm} helicities would permit a CP asymmetry to be measured in $H_i \rightarrow \tau^+\tau^-$ decays. This information may also be used to further resolve the underlying resonant dynamics of a strongly mixed Higgs-boson system, thereby offering a sensitive window into CP violation and new physics due to an extended Higgs sector, such as the MSSM Higgs sector.

This example shows that exclusive double diffraction may offer unique possibilities for exploring Higgs physics in ways that would be difficult or even impossible in inclusive Higgs production. In particular, we have shown that exclusive double diffraction constitutes an efficient CP and lineshape analyzer of the resonant Higgs-boson dynamics in multi-Higgs models. In the specific case of CP-violating MSSM Higgs physics discussed here, which is potentially of great importance for electroweak baryogenesis, diffractive production may be the most promising probe at the LHC.

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