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Observability of the Lightest MSSM Higgs Boson with Explicit CP Violation via Gluon Fusion at the LHC

S.Y. Choi[1], Kaoru Hagiwara[2] and Jae Sik Lee[2]

[1] Department of Physics, Chonbuk National University, Chonju 561–756, Korea

[2] Theory Group, KEK, Tsukuba, Ibaraki 305–0801, Japan

Abstract

We investigate the observability of the lightest Higgs boson in the gluon–fusion channel at the CERN Large Hadron Collider (LHC) in the minimal supersymmetric Standard Model with explicit CP–violating mixing among three neutral Higgs bosons. The lightest Higgs boson with its mass less than 130 GeV can be detected at the LHC via its gluon–fusion production followed by the decay into two photons. The explicit CP violation can suppress both the production cross section and the two–photon decay branching fraction so significantly that the signal cross section may be more than ten times smaller than the SM signal. This reduction factor can be as small as 1/40 if the lightest Higgs boson mass is 115 GeV and its production cross section at LEP2 is more than 90 % that of the SM case.

The soft CP violating Yukawa interactions in the minimal supersymmetric standard model (MSSM) cause the CP–even and CP–odd neutral Higgs bosons to mix through loop corrections [1, 2]. The loop–induced CP violation in the MSSM Higgs sector can be large enough to affect Higgs phenomenology at present and future colliders significantly [1, 3, 4, 5, 6, 7]. In this letter, we study the effects of the CP–violating mixing on the production of the lightest neutral MSSM Higgs boson through gluon fusion and its decay into a photon pair, which is of crucial importance for detecting the lightest Higgs boson with its mass less than 130 GeV at the LHC [8]. After briefly reviewing the loop–induced CP–violating mixing [2] of three neutral Higgs bosons in the MSSM, we estimate the effects of the CP phases on the production of the lightest Higgs boson though gluon fusion and its branching fraction of the two–photon decay mode, respectively. Finally, combining both the production and decay of the lightest Higgs boson, we discuss the observability of the lightest MSSM Higgs boson at the LHC in the presence of explicit CP violation in the Higgs sector.

The loop–induced CP–violating neutral Higgs boson mixing is determined by the Higgs boson mass matrix obtained by taking all the second derivatives of the effective MSSM Higgs potential [9, 10]

$$
\begin{aligned}
V_{\rm Higgs} &= \tfrac{1}{2}m_1^2\,(\phi_1^2+a_1^2)+\tfrac{1}{2}m_2^2\,(\phi_2^2+a_2^2)-|m_{12}^2|\,(\phi_1\phi_2-a_1a_2)\cos(\xi+\theta_{12})\\
&\quad+|m_{12}^2|\,(\phi_1a_2+\phi_2a_1)\sin(\xi+\theta_{12})+\tfrac{\hat g^2}{8}\mathcal{D}^2+\tfrac{1}{64\,\pi^2}{\rm Str}\left[\mathcal{M}^4\left(\log\tfrac{\mathcal{M}^2}{Q^2}-\tfrac{3}{2}\right)\right]\,, \qquad (1)
\end{aligned}
$$

where $\mathcal{D}=\phi_2^2+a_2^2-\phi_1^2-a_1^2$, $\hat g^2=(g^2+g'^2)/4$ with the ${\rm SU(2)}_L$ and ${\rm U(1)}_Y$ gauge couplings g and g', and ϕ_i and a_i $(i=1,2)$ are the neutral components of the two Higgs doublet fields:

$$
H_1^0=\frac{1}{\sqrt{2}}\,(\phi_1+ia_1)\;,\qquad H_2^0=\frac{{\rm e}^{i\xi}}{\sqrt{2}}\,(\phi_2+ia_2)\;. \qquad (2)
$$

All the tree–level parameters of the Higgs potential (1) such as m_1^2, m_2^2 and $m_{12}^2=|m_{12}^2|\,{\rm e}^{i\theta_{12}}$ are the running parameters evaluated at the renormalization scale Q, rendering the Higgs potential (almost) independent of Q up to two–loop–order corrections. The super–trace is to be taken over all the bosons and fermions that couple to the Higgs fields.

The matrix $\mathcal{M}$ in Eq. (1) is the field–dependent mass matrix of all modes that couple to the Higgs bosons. The dominant contributions in the MSSM come from the third generation quarks and squarks because of their large Yukawa couplings. The field–dependent masses of the bottom and top quarks are given by $m_b^2=|h_b|^2|H_1^0|^2$ and $m_t^2=|h_t|^2|H_2^0|^2$ with the bottom and top Yukawa couplings h_b and h_t, and the bottom– and top–squark mass matrices read

$$
\begin{aligned}
\mathcal{M}_{\tilde t}^2 &= \left(\begin{array}{cc}
m_{\widetilde Q}^2+m_t^2-\frac{1}{8}\left(g^2-\frac{g'^2}{3}\right)\mathcal{D} & -h_t^*\left[A_t^*\,(H_2^0)^*+\mu H_1^0\right]\\
-h_t\left[A_tH_2^0+\mu^*\,(H_1^0)^*\right] & m_{\widetilde U}^2+m_t^2-\frac{g'^2}{6}\mathcal{D}
\end{array}\right)\,,\\
\mathcal{M}_{\tilde b}^2 &= \left(\begin{array}{cc}
m_{\widetilde Q}^2+m_b^2+\frac{1}{8}\left(g^2+\frac{g'^2}{3}\right)\mathcal{D} & -h_b^*\left[A_b^*\,(H_1^0)^*+\mu H_2^0\right]\\
-h_b\left[A_bH_1^0+\mu^*\,(H_2^0)^*\right] & m_{\widetilde D}^2+m_b^2+\frac{g'^2}{12}\mathcal{D}
\end{array}\right)\,. \qquad (3)
\end{aligned}
$$

where $m_{\widetilde Q}^2$, $m_{\widetilde U}^2$ and $m_{\widetilde D}^2$ are the real soft SUSY–breaking squark mass parameters, A_b and A_t are the complex soft SUSY–breaking trilinear parameters, and μ is the complex supersymmetric Higgsino mass parameter.

The second derivatives of the potential, giving the mass matrix of the Higgs bosons (at vanishing external momenta), are then evaluated at its minimum point $(\phi_1, \phi_2, a_1, a_2) = (v\cos\beta, v\sin\beta, 0, 0)$ with $v \simeq 246$ GeV and $\tan\beta = \langle\phi_2\rangle/\langle\phi_1\rangle$. After absorbing a Goldstone mode $G^0 = a_1\cos\beta - a_2\sin\beta$ into the Z boson, we are left with a real and symmetric 3×3 mass–squared matrix $\mathcal{M}_H^2$ of three physical states, $a\,(= a_1\sin\beta + a_2\cos\beta)$, ϕ_1 and ϕ_2. The two CP–violating entries of the symmetric matrix, which mix a with ϕ_1 and ϕ_2, are given by

$$\mathcal{M}_H^2\Big|_{a\phi_1} = \frac{3}{16\pi^2}\left\{\frac{m_t^2\Delta_{\tilde t}}{\sin\beta}\,F_t + \frac{m_b^2\Delta_{\tilde b}}{\cos\beta}\,G_b\right\}, \quad \mathcal{M}_H^2\Big|_{a\phi_2} = \frac{3}{16\pi^2}\left\{\frac{m_t^2\Delta_{\tilde t}}{\sin\beta}\,G_t + \frac{m_b^2\Delta_{\tilde b}}{\cos\beta}\,F_b\right\}. \quad (4)$$

The explicit forms of the dimensionless quantities $F_{t,b}$ and $G_{t,b}$ and all the CP–preserving entries of the mass–squared matrix $\mathcal{M}_H^2$ can be found in Ref. [2]. The rephasing–invariants

$$\Delta_{\tilde t} = \frac{\Im\mathrm{m}(A_t\mu \mathrm{e}^{i\xi})}{m_{\tilde t_2}^2 - m_{\tilde t_1}^2}\,, \qquad \Delta_{\tilde b} = \frac{\Im\mathrm{m}(A_b\mu \mathrm{e}^{i\xi})}{m_{\tilde b_2}^2 - m_{\tilde b_1}^2}, \quad (5)$$

measure the amount of CP violation in the top and bottom squark–mass matrices and vanish in the CP–invariant theories, leading to $|m_{12}^2|\,\sin(\xi+\theta_{12}) = 0$ in the potential (1). The matrix $\mathcal{M}_H^2$ can be diagonalized by an orthogonal matrix O;

$$O^T\mathcal{M}_H^2\,O = \mathsf{diag}\,(m_{H_1}^2, m_{H_2}^2, m_{H_3}^2), \quad (6)$$

where the three mass-eigenvalues are ordered as $m_{H_1} < m_{H_2} < m_{H_3}$.

The loop–corrected neutral–Higgs–boson sector depends on many parameters in the Higgs and squark sectors; a loop–corrected pseudoscalar mass m_A, $\tan\beta$, μ, A_t, A_b, the scale Q, and the soft–breaking masses, $m_{\tilde Q}$, $m_{\tilde U}$, and $m_{\tilde D}$, as well as on the complex gluino–mass parameter $M_{\widetilde g}$ through one–loop corrections to the top and bottom quark masses [11]. However, the CP violation in the Higgs sector is determined essentially by the rephasing invariant combinations $A_t\mu \mathrm{e}^{i\xi}$ and $A_b\mu \mathrm{e}^{i\xi}$, see Eq. (5), and is dominantly by the top–squark sector if $\tan\beta \leq 10$. Therefore, we take in our numerical analysis the following parameter set:

$$\begin{aligned}
&|A_t| = |A_b| = \kappa\, M_{\rm SUSY}\,, && |\mu| = 2\,|A_t|\,,\\
&m_{\widetilde Q,\widetilde U,\widetilde D} = |M_{\widetilde g}| = M_{\rm SUSY} = 0.5 \text{ TeV}\,, && \mathrm{Arg}(M_{\widetilde g}) = 0\,,\\
&\Phi \equiv \mathrm{Arg}(A_t\mu \mathrm{e}^{i\xi}) = \mathrm{Arg}(A_b\mu \mathrm{e}^{i\xi})\,. &&
\end{aligned} \quad (7)$$

Then, we vary the dimensionless parameter κ, the common phase Φ and $\tan\beta$ in the numerical analysis, for which the pseudoscalar mass parameter m_A is chosen to fix the the lightest Higgs boson mass m_{H_1}. Clearly, a large κ implying large values of $|A_{t,b}|$ leads to large CP–violating effects as clearly seen from Eq. (5). However, κ cannot be too large, because it generates an unacceptably large value of the electron and neutron electric dipole moments (EDM's)* at the two–loop level through the one–loop effective CP–odd couplings

*It is possible that the stringent two–loop EDM constraints may be satisfied by a cancellation among various contributions [13, 14, 5].

of the Higgs boson to the gauge bosons [12]. Moreover, in order to avoid a color and electric–charge breaking minimum deeper than the electroweak vacuum, κ cannot be significantly larger than the unity [15, 16].

In the presence of the CP–violating neutral Higgs–boson mixing, the amplitude for the resonance production $gg \to H_i$ ($i = 1, 2, 3$) can be written as

$$\mathcal{M}_{ggH_i} = \frac{m_{H_i}\alpha_s\delta_{ab}}{4\pi}\left\{S_i^g(m_{H_i})\left(\epsilon_1\cdot\epsilon_2 - \frac{2k_1\cdot\epsilon_2\, k_2\cdot\epsilon_1}{m_{H_i}^2}\right) - P_i^g(m_{H_i})\frac{2}{m_{H_i}^2}\epsilon_{\mu\nu\rho\sigma}\,\epsilon_1^\mu\epsilon_2^\nu k_1^\rho k_2^\sigma\right\}, \quad (8)$$

where a, b = (1 to 8) are the color indices for the eight gluon fields, and $k_{1,2}$ and $\epsilon_{1,2}$ are the momenta and polarization vectors of two colliding gluons, respectively. The scalar and pseudo–scalar form factors are then given by

$$\begin{aligned} S_i^g(m_{H_i}) &= \sum_f \left\{ g_{sf}^i \frac{m_{H_i}}{m_f} F_{sf}(\tau_{if}) + \frac{1}{4}\sum_{j=1,2} g^i_{\tilde{f}_j\tilde{f}_j}\frac{m_{H_i}}{m^2_{\tilde{f}_j}} F_0(\tau_{i\tilde{f}_j})\right\}, \\ P_i^g(m_{H_i}) &= \sum_f g^i_{pf}\frac{m_{H_i}}{m_f} F_{pf}(\tau_{if}), \end{aligned} \quad (9)$$

where $\tau_{ix} = m_{H_i}^2/4m_x^2$, g_{sf}^i and g_{pf}^i are the couplings of the Higgs boson H_i to the scalar and pseudo–scalar fermion bilinears $\bar{f}f$ and $i\,\bar{f}\gamma_5 f$, respectively. The CP–violating Higgs mixing leads to a simultaneous existence of these two couplings. On the other hand, $g^i_{\tilde{f}_j\tilde{f}_j}$ is the coupling of H_i to a diagonal sfermion pair. We refer to Ref. [6] for the explicit forms of the couplings as well as the form factors F_{sf}, F_{pf}, and F_0. Note that in the minimal SM only the scalar form factor due to the top–quark and bottom–quark contributions survives.

The production cross section of a neutral Higgs boson H_i in gg fusion is given by

$$\sigma(gg\to H_i) = \frac{\alpha_s^2}{256\pi m_{H_i}^2}\left[|S_i^g|^2 + |P_i^g|^2\right]\delta(1-\frac{m_{H_i}^2}{\hat{s}}) \equiv \hat{\sigma}_{\rm LO}(gg\to H_i)\,\delta(1-\frac{m_{H_i}^2}{\hat{s}}), \quad (10)$$

with $\sqrt{\hat{s}}$ the two–gluon c.m. energy. Figure 1 shows the leading–order (LO) parton–level cross section $\hat{\sigma}_{\rm LO}(gg \to H_1)$ as a function of the phase Φ for m_{H_1} = 80 GeV (solid line), 90 GeV (dashed line), 100 GeV (dotted line), 110 GeV (dash–dotted line), 115 GeV (thick dashed line), and 120 GeV (thick solid line) with the parameter set (7) for $\kappa = 1.6$ (upper) and 2.0 (lower) and for $\tan\beta = 4$ (left) and 10 (right), respectively. Note that the cases with $m_{H_1} \leq 100$ GeV are also considered because the lightest Higgs boson H_1 could be undetected at LEP2 with the ZZH_1 coupling suppressed for non–vanishing Φ.

The SM LO parton-level cross section at $m_{H_{\rm SM}} = m_{H_1}$ is 45.0 fb $\leq \hat{\sigma}^{\rm SM}_{\rm LO} \leq$ 46.6 fb for $m_{H_{\rm SM}}$ between 80 GeV and 120 GeV, implying that the SM cross section does not depend on the Higgs boson mass significantly. On the contrary, the MSSM cross section is very sensitive to Φ and m_{H_1} for both $\kappa = 1.6$ and 2.0 and for both $\tan\beta = 4$ and 10. In particular, the cross section is significantly smaller than the SM one for small Φ and $\tan\beta$. This is due to the suppression of the coupling of the lightest MSSM Higgs boson to top quarks and to the cancellations between the fermionic and bosonic contributions. The cancellation is more significant when the top–squark mass splitting is larger, for smaller Φ, larger κ and smaller

$\tan\beta$. For $\kappa = 1.6$, a significant cancellation between the fermionic and bosonic contributions occurs for all the Higgs mass cases at $\tan\beta = 4$ when $\Phi \lesssim 70^o$, and for $m_{H_1} = 115$ GeV (thick dashed line) at $\tan\beta = 10$ when $\Phi \lesssim 50^o$. For $\kappa = 2.0$ and $\tan\beta = 4$, a significant cancellation occurs for $\Phi \lesssim 90^o$, while for $\kappa = 2.0$ and $\tan\beta = 10$ it occurs when $m_{H_1} = 110$ GeV and 115 GeV for $\Phi \lesssim 70^o$. In all cases, the cancellation between fermionic and bosonic contributions is suppressed for $\Phi \gtrsim 120^o$, reflecting the suppressed sfermion mass splitting. We find that for $\tan\beta = 4$ and $m_{H_1} \leq 100$ GeV the lightest Higgs boson has a large CP–odd component when $70^o \lesssim \Phi \lesssim 110^o$ ($\kappa = 1.6$) and $80^o \lesssim \Phi \lesssim 130^o$ ($\kappa = 2.0$). Finally, since the scalar and pseudoscalar couplings of H_1 to the top quarks are $g^1_{pt} \sim O_{11}/\tan\beta$ and $g^1_{st} \sim O_{31}/\sin\beta$, respectively, the top–quark loop contribution for the lightest Higgs boson of a large CP–odd mixture is suppressed by $\cos\beta$ as compared to the case of a pure CP–even lightest Higgs boson at $\Phi = 0^o/180^o$.

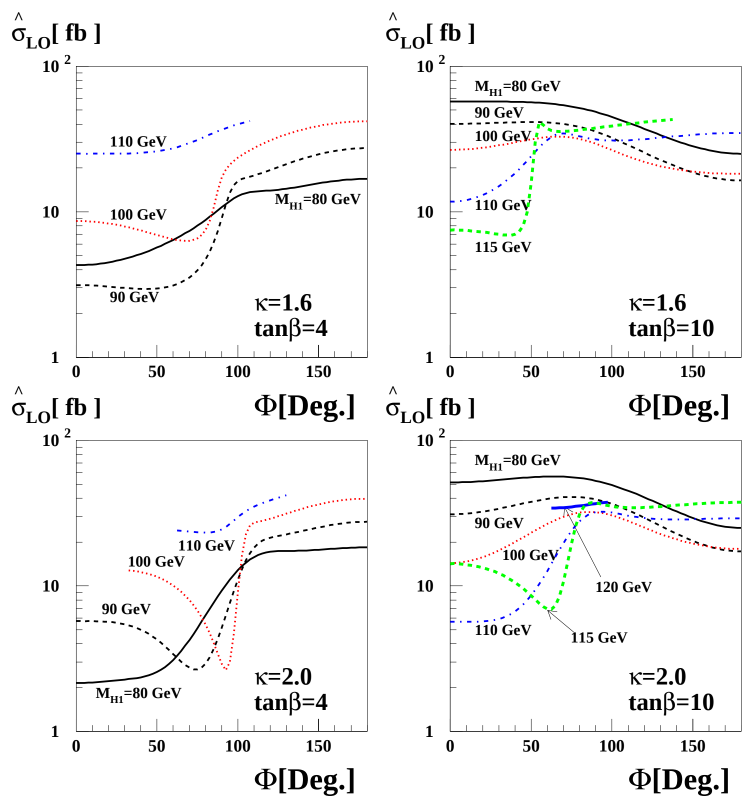

Figure 1: *The LO parton–level cross section of the lightest Higgs boson as a function of the phase Φ for $m_{H_1} = 80$ GeV (solid line), 90 GeV (dashed line), 100 GeV (dotted line), 110 GeV (dash–dotted line), 115 GeV (thick dashed line), and 120 GeV (thick solid line). We take the parameter set (7) with $\kappa = 1.6$ (upper) and $\kappa = 2.0$ (lower) and two values of* $\tan\beta = 4$ *(left) and 10 (right).*

For a realistic estimate of the production cross section it is necessary to include the next–to–leading–order (NLO) QCD loop correction, denoted by the $\tan\beta$–dependent K factor to a good approximation [17]; for small $\tan\beta$, it is 1.5–1.7 and for large $\tan\beta$ it is in general close to unity except when the lightest Higgs boson approaches the SM limit, for which $K \approx 1.5$. In addition to the QCD NLO correction, we need to fold the parton–level cross section with the gluon distribution function to obtain the hadronic level cross section as

$$\sigma(pp \to H_1) = K\,\hat{\sigma}_{LO}(gg \to H_1)\;\tau\frac{\mathrm{d}\mathcal{L}^{gg}_{LO}}{\mathrm{d}\tau}\,, \tag{11}$$

where $\tau = m^2_{H_i}/s$ with $\sqrt{s}$ the hadron collider c.m. energy. At the LHC with $\sqrt{s} = 14$ TeV, the size of the gluon fusion luminosity factor $(\tau\frac{\mathrm{d}\mathcal{L}^{gg}_{LO}}{\mathrm{d}\tau})$ is between 0.6×10^3 and 0.3×10^3 for $m_{H_1} = 80 - 130$ GeV [18].

In the presence of the radiatively induced CP–violating neutral Higgs boson mixing, the amplitude for the decay $H_i \to \gamma\gamma$ $(i = 1, 2, 3)$ is written as

$$\mathcal{M}_{\gamma\gamma H_i} = \frac{m_{H_i}\alpha}{4\pi}\left\{S^\gamma_i(m_{H_i})\left(\epsilon^*_1\cdot\epsilon^*_2 - \frac{2\,k_1\cdot\epsilon^*_2\,k_2\cdot\epsilon^*_1}{m^2_{H_i}}\right) - P^\gamma_i(m_{H_i})\frac{2}{m^2_{H_i}}\epsilon_{\mu\nu\rho\sigma}\,\epsilon^{*\mu}_1\epsilon^{*\nu}_2 k^\rho_1 k^\sigma_2\right\}, \tag{12}$$

where $k_{1,2}$ and $\epsilon_{1,2}$ are the momenta and polarization vectors of the two photons, respectively. The scalar and pseudoscalar form factors due to the (s)quark, $W^\pm$ and charged Higgs loops read

$$\begin{aligned}
S^\gamma_i(m_{H_i}) &= 2N_C\sum_f Q^2_f\left\{g^i_{sf}\frac{m_{H_i}}{m_f}F_{sf}(\tau_{if}) + \frac{1}{4}\sum_{j=1,2} g^i_{\tilde{f}_j\tilde{f}_j}\frac{m_{H_i}}{m^2_{\tilde{f}_j}}F_0(\tau_{i\tilde{f}_j})\right\}\\
&\quad + \frac{g m_{H_i}}{2m_W}(c_\beta O_{2\,i} + s_\beta O_{3\,i})F_1(\tau_{iW}) + \frac{v m_{H_i} C_i}{2M^2_{H^\pm}}F_0(\tau_{iH^\pm})\,,\\
P^\gamma_i(m_{H_i}) &= 2N_C\sum_{f=t,b} Q^2_f g^i_{pf}\frac{m_{H_i}}{m_f}F_{pf}(\tau_{if})\,,
\end{aligned} \tag{13}$$

where $N_C = 3$, Q_f is the electric charge of the (s)fermion $f(\tilde{f})$ in the unit of the positron charge, and C_i is the coupling of H_i to the charged Higgs boson pair: $\mathcal{L}_{H_iH^+H^-} = v\,C_i\,H_iH^+H^-$. We refer again to Ref. [6] for the explicit forms of the form factors and the couplings C_i's. The possible chargino contributions are neglected by assuming that the chargino states are very heavy. Note that the SM pseudoscalar form factor $P^\gamma_{\rm SM}$ vanishes and the SM scalar form factor $S^\gamma_{\rm SM}$ has only the top–quark and and $W^\pm$–boson loop contributions.

The decay width $\Gamma(H_i \to \gamma\gamma)$ is then given by

$$\Gamma(H_i \to \gamma\gamma) = \frac{m_{H_i}\alpha^2}{256\pi^3}\left[\,|S^\gamma_i(m_{H_i})|^2 + |P^\gamma_i(m_{H_i})|^2\right], \tag{14}$$

in terms of the scalar and pseudoscalar form factors in Eq. (13). The main contribution to the decay of the lightest MSSM Higgs boson into two photons is from the $W^\pm$–boson loop giving rise to the scalar form factor $S^\gamma_1(m_{H_1})$, which is determined by the coupling of H_1 to W^+W^-. We find that the $H_1W^+W^-$ coupling is very sensitive to the CP–violating neutral Higgs boson mixing. For example, if H_1 is a pure CP–odd state, the coupling vanishes. As a result, the partial decay width and its branching fraction can be significantly suppressed in the presence of the CP–violating phases. Figure 2 shows the the branching fraction $\mathcal{B}(H_1 \to \gamma\gamma)$ as a function of the phase Φ for $m_{H_1} = 80$ GeV (solid line), 90 GeV (dashed line), 100 GeV (dotted line), 110 GeV (dash–dotted line), 115 GeV (thick dashed line), and 120 GeV (thick solid line) for the parameter set (7) with $\kappa = 1.6$ (upper) and $\kappa = 2.0$ (lower) and with $\tan\beta = 4$ (left) and 10 (right). For reference, the SM branching fraction, which turns out to be between 1×10^{-3} and 3×10^{-3}, is also shown in each frame with the same line convention, fixing $m_{H_{\rm SM}} = m_{H_1}$. The MSSM branching fraction is so sensitive to the phase Φ that it is suppressed by a factor of 10^3 around $\Phi = 100^{\rm o}$ ($\tan\beta = 4$) and 10^4 around $\Phi = 50^{\rm o} - 75^{\rm o}$ ($\tan\beta = 10, \kappa = 1.6$) and around $\Phi = 70^{\rm o} - 95^{\rm o}$ ($\tan\beta = 10, \kappa = 2.0$) where H_1 is a dominantly CP–odd state.

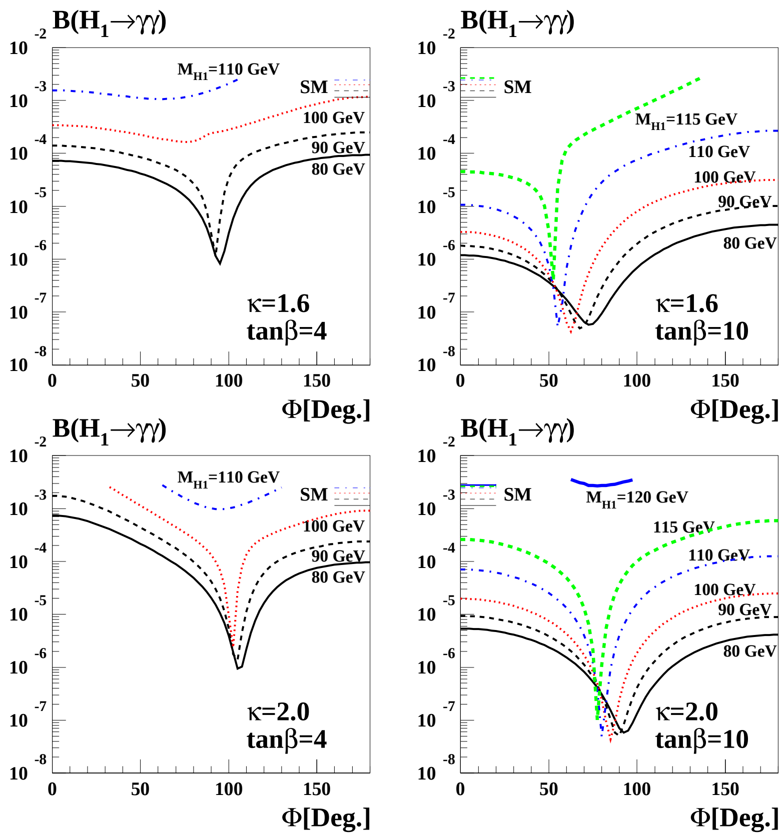

Figure 2: *The branching fraction $\mathcal{B}(H_1 \to \gamma\gamma)$ as a function of Φ for $m_{H_1} = 80$ GeV (solid line), 90 GeV (dashed line), 100 GeV (dotted line), 110 GeV (dash–dotted line), 115 GeV (thick dashed line), and 120 GeV (thick solid line). The parameter set (7) is with $\kappa = 1.6$ (upper) and $\kappa = 2.0$ (lower) and with $\tan\beta = 4$ (left) and 10 (right). The SM branching fraction is also shown with the same line convention at $m_{H_{\rm SM}} = m_{H_1}$.*

Since gluon fusion is the main production mode for H_1 and the decay $H_1 \to \gamma\gamma$ is the major signal mode for H_1 at the LHC for $m_{H_1} \leq 130$, it is crucial to investigate the observability of the lightest MSSM Higgs boson with explicit CP violation through this channel at the LHC. For this purpose, we consider the ratio of the signal cross sections:

$$R_{g\gamma}^{H_i} \equiv \frac{[\hat{\sigma}_{\rm LO}(gg \to H_i) \times \mathcal{B}(H_i \to \gamma\gamma)]_{\rm MSSM}}{[\hat{\sigma}_{\rm LO}(gg \to H_{\rm SM}) \times \mathcal{B}(H_{\rm SM} \to \gamma\gamma)]_{\rm SM}\Big|_{m_{H_{\rm SM}}=m_{H_i}}} . \tag{15}$$

It measures the H_i signal cross section of the MSSM as compared to the SM Higgs boson signal cross section at the same mass. Figure 3 shows the ratio as a function of the phase Φ for $m_{H_1} = 80 - 120$ GeV, $\kappa = 1.6$ and 2.0, and $\tan\beta = 4$ and 10, as in the previous figures.

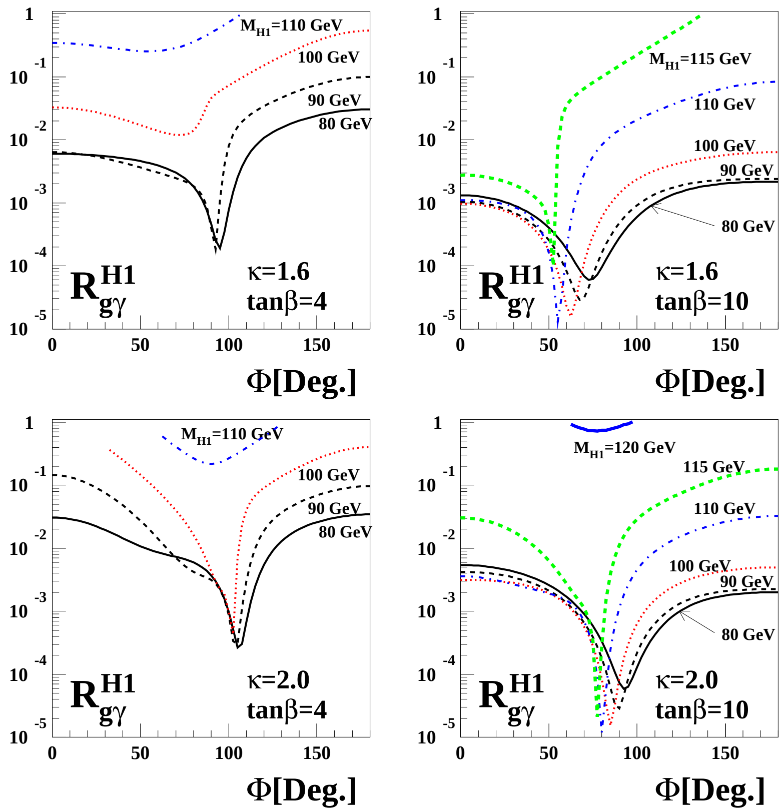

Figure 3: *The ratio $R_{g\gamma}^{H_1}$ as a function of the phase Φ for $m_{H_1} = 80$ GeV (solid line), 90 GeV (dashed line), 100 GeV (dotted line), 110 GeV (dash–dotted line), 115 GeV (thick dashed line), and 120 GeV (thick solid line). The parameter set (7) is taken with $\kappa = 1.6$ (upper) and 2.0 (lower) and with $\tan\beta = 4$ (left) and 10 (right).*

The ratio $R_{g\gamma}^{H_1}$ can be highly suppressed for a wide range of Φ, which is mainly due to the significant decrease of $\mathcal{B}(H_1 \to \gamma\gamma)$ caused by the suppressed coupling of H_1 to W^+W^- for non–vanishing Φ. The ratio is strongly suppressed for a wide range of Φ around $90^{\rm o}$, although otherwise it is not suppressed and can be as large as the unity for some cases. For the cases studied in this letter, we find that the ratio is almost always less than unity.

On the other hand, the case for the heavy neutral Higgs bosons is converse to that for the lightest Higgs boson; the sum rule for the couplings $g_{H_iVV} = c_\beta\, O_{2\,i} + s_\beta\, O_{3\,i}$ of the neutral Higgs bosons to a gauge boson pair $V\ (= W^\pm, Z)$,

$$\sum_{i=1}^{3} g^2_{H_iVV} = 1\,, \tag{16}$$

implies that if g_{H_1VV} is suppressed, either g_{H_2VV} or g_{H_3VV} for $H_{2,3}$ should be enhanced. Based on the sum rule (16), it has been argued [13] that the tantalizing hints for the Higgs boson(s) with its mass around 115 GeV at the LEP experiments [19] might be due to the intermediate or the heaviest Higgs boson instead of the lightest Higgs boson which can be CP–odd. In this light, it is worthwhile to simultaneously consider the ratios $R^{H_{2,3}}_{g\gamma}$ with $m_{H_{2,3}}$ around 115 GeV as well as $R^{H_1}_{g\gamma}$. Let us examine the dependence of $R^{H_i}_{g\gamma}$ on the parameter $\kappa = |A_{t,b}|/M_{\rm SUSY}$ in the parameter set (7) for $m_{H_i} = 115$ GeV and $\tan\beta = 10$. We find that it is impossible for H_3 to be the 115 GeV Higgs boson for the LEP2 excess events in our parameter set (7) with $\kappa \geq 1.2$. Although, the coupling g_{H_3VV} may dominate the other couplings for $\kappa < 1.2$, we do not consider this case in the present work because we do not find significant CP–violating mixing for $\kappa < 1.2$. We find that the coupling g_{H_2VV} is larger than the coupling g_{H_1VV} in the following parameter space:

$$\begin{aligned}
30^{\rm o} \leq \Phi \leq 60^{\rm o} \qquad & \text{for} \quad \kappa = 1.6\,,\\
40^{\rm o} \leq \Phi \leq 80^{\rm o} \qquad & \text{for} \quad \kappa = 1.8\,,\\
60^{\rm o} \leq \Phi \leq 100^{\rm o} \qquad & \text{for} \quad \kappa = 2.0\,.
\end{aligned} \tag{17}$$

In this regard, we present the ratio $R^{H_2}_{g\gamma}$ with $m_{H_2} = 115$ GeV instead of $R^{H_1}_{g\gamma}$ in the parameter space (17). Figure 4 shows the ratio $R^{H_1,H_2}_{g\gamma}$ with $m_{H_1,H_2} = 115$ GeV as a function of the phase Φ for $\kappa = 1.3$ (solid line), $\kappa = 1.4$ (dashed line), $\kappa = 1.6$ (dotted line), $\kappa = 1.8$ (dash–dotted line), and $\kappa = 2.0$ (thick solid line). Note that the ratio $R^{H_1}_{g\gamma}$ with $m_{H_1} = 115$ GeV depends strongly on the phase Φ independently of the value of κ. The ratio $R^{H_1,H_2}_{g\gamma}$ becomes smaller for larger κ. The ratio $R^{H_2}_{g\gamma}$ with $m_{H_2} = 115$ GeV is about 0.01 in the range (17).

It is known [8] that for the integrated luminosity of 100 fb^{-1} the signal significance for the discovery of the SM Higgs boson with $m_{H_{\rm SM}} \leq 130$ GeV through the process $gg \to H_{\rm SM} \to \gamma\gamma$ is less than 10 per experiment at the LHC. It means that the 5σ–level discovery of the lightest MSSM Higgs boson may not be possible at the LHC through this channel if the ratio $R^{H_1}_{g\gamma}$ is significantly less than a quarter. As shown in Figs. 3 and 4 the lightest Higgs boson of the MSSM can escape detection if $\kappa \geq 1.7$. On the other hand, it may be possible to discover the lightest Higgs boson at the LHC if κ is less than 1.7 and the phase Φ is sufficiently large.

Let us elaborate on the ratios $R^{H_2}_{g\gamma}$ a little more. For a fixed m_{H_1} there should be an anti–correlation between $R^{H_1}_{g\gamma}$ and $R^{H_2}_{g\gamma}$ due to the sum rule (16) for $g^2_{H_1VV} + g^2_{H_2VV} \approx 1$; if the coupling g_{H_1WW} is suppressed, the coupling g_{H_2WW} is enhanced and the mass difference $m_{H_2} - m_{H_1}$ is also reduced. Nevertheless, the ratio $R^{H_2}_{g\gamma}$ for $m_{H_2} \leq 150$ GeV is always less than 0.1 in spite of the anti–correlation for the parameter space under consideration. It is therefore possible that LHC discovers neither H_1 nor H_2 at the LHC for a wide range of Φ.

On the other hand, all the cases shown in Fig. 4 give either $g^2_{H_1ZZ}$ or $g^2_{H_2ZZ}$ greater than 0.5 so that one of their production cross sections at LEP2 is not suppressed significantly. If we require that $\max\{g^2_{H_iZZ}\} > 0.9$, then we find that the minimum of the ratio is around 1/40.

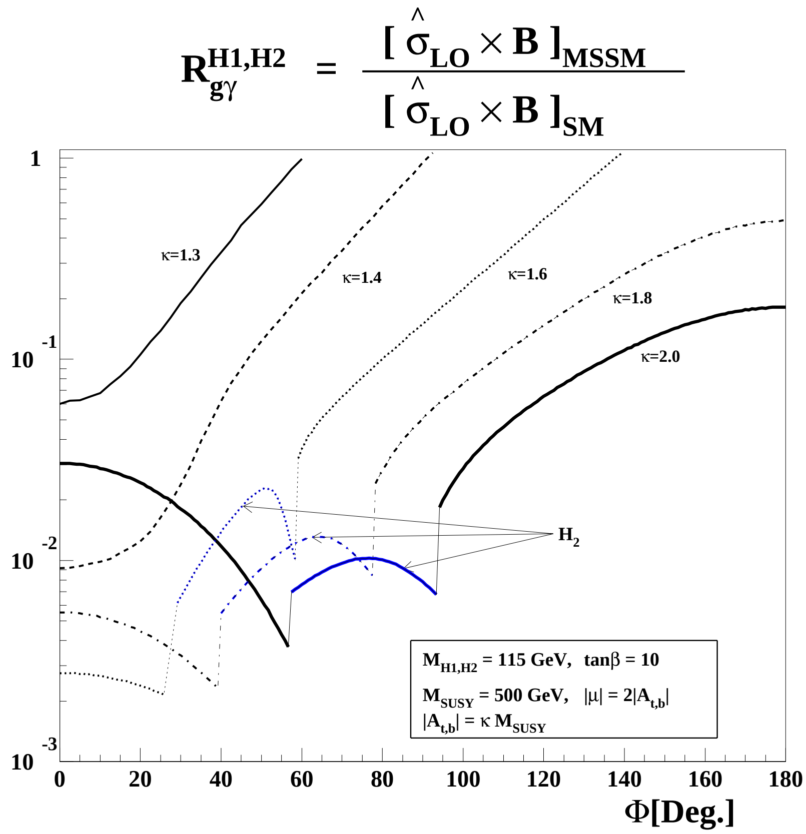

Figure 4: *The ratio* $R^{H_1,H_2}_{g\gamma}$ *as a function of the phase* Φ *for* $\kappa = 1.3$ *(solid line),* $\kappa = 1.4$ *(dashed line),* $\kappa = 1.6$ *(dotted line),* $\kappa = 1.8$ *(dash–dotted line), and* $\kappa = 2.0$ *(thick solid line). The mass of the lightest or the second lightest Higgs boson is fixed with* $m_{H_1,H_2} = 115$ *GeV. The parameter set (7) with* $\tan\beta = 10$ *is taken.*

To summarize, we have investigated the observability of the lightest Higgs boson at the LHC by studying its production through gluon fusion and its decay into two photons in the MSSM where the tree-level CP invariance in the Higgs sector is explicitly broken by the loop effects of third–generation squarks with CP–violating phases. We find that both the production cross section and the decay branching fraction can be strongly suppressed for non–trivial phase Φ and for large κ, while the maximal signal cross section is always for $\Phi = 180^o$. *Consequently, it is possible that the lightest MSSM Higgs boson escapes detection through the gluon fusion and its decay into two photons at the LHC if the CP–violating mixing is significant.* It is therefore important to study seriously the vector–boson fusion signal at the LHC [20].

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