The LHC di-photon Higgs signal predicted by little Higgs models (Postprint)

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Full Text

Preamble

The LHC di-photon Higgs signal predicted by little Higgs models

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Abstract

Little Higgs theory naturally predicts a light Higgs boson whose most important discovery channel at the LHC is the di-photon signal $pp \to \gamma\gamma$. In this work we perform a comparative study for this signal in some typical little Higgs models, namely the littlest Higgs model (LH), two littlest Higgs models with T-parity (named LHT-I and LHT-II) and the simplest little Higgs model (SLH). We find that compared with the Standard Model prediction, the di-photon signal rate is always suppressed and the suppression extent can be quite different for different models.

The suppression is mild (< 10%) in the LH model but can be quite severe (90%) in the other three models. This means that discovering the light Higgs boson predicted by little Higgs theory through the di-photon channel at the LHC will be more difficult than discovering the SM Higgs boson.

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INTRODUCTION

The little Higgs [1] is proposed as an elegant mechanism of electroweak symmetry breaking with a naturally light Higgs sector. So far various realizations of the little Higgs symmetry structure have been proposed [2–4], which can be categorized generally into two classes [5].

One class uses the product group, represented by the littlest Higgs model (LH) [3], in which the SM $SU(2)_L$ gauge group arises from the diagonal breaking of two (or more) gauge groups. The other class uses the simple group, represented by the simplest little Higgs model (SLH) [4], in which a single larger gauge group is broken down to the SM $SU(2)_L$. Further, to relax the constraints from electroweak precision tests [4, 6], a discrete symmetry called T-parity is proposed [7], which can also provide a candidate for cosmic dark matter.

For the LH there are two different implementations of T-parity in the fermion sector, called respectively LHT-I and LHT-II [8, 9]. A characteristic difference between LHT-I and LHT-II is that the top quark partner responsible for canceling the one-loop quadratic divergence of the Higgs mass contributed by the top quark is T-even for the former and T-odd for the latter. The implementation of T-parity in the SLH has also been attempted [10].

To test little Higgs theory at the LHC, Higgs phenomenology will play an important role [11]. At the LHC different search strategies will be applied for different mass ranges. For a light Higgs boson below about 140 GeV the di-photon signal $pp \to \gamma\gamma$ is the most important discovery channel because the narrow $\gamma\gamma$ peak can be reconstructed to distinguish the signal from backgrounds. In contrast, the dominant channel $pp \to b\bar{b}$ cannot be utilized for discovery because of overwhelming QCD backgrounds. Recently the ATLAS collaboration reported their di-photon search results with 209 pb ^{1} of data collected in early 2011 and excluded a signal rate of 4.2–15.8 times the SM prediction for 110 GeV $< m_h$ $<$ 140 GeV [12]. With a luminosity of 2 fb ^{1} the ongoing LHC will be able to use the di-photon signal to exclude a light SM Higgs boson. So the di-photon Higgs channel will be a sensitive probe for new physics models like little Higgs theory.

So far the di-photon signal has been studied in some new physics models [13–16]. Although some little Higgs models have also been discussed [14–16], these previous studies are performed separately in different frameworks. To show the difference of model predictions, it is necessary to perform a comparative study for various models. Further, the study for the SLH has not been reported in detail in the literature. In this work we consider all these models (LH, LHT-I, LHT-II and SLH) to perform a comparative study.

Our work is organized as follows. In Sec. II we recapitulate the models. In Sec. III we calculate the rate of $pp \to \gamma\gamma$ at the LHC in these models. Finally, we give our conclusion in Sec. IV.

II. LITTLE HIGGS MODELS

A. Littlest Higgs model (LH)

The LH model [3, 17] is based on a non-linear model in the coset space of $SU(5)/SO(5)$ with additional local gauge symmetry $[SU(2)\times U(1)]^2$. The vacuum expectation value (VEV) of an $SU(5)$ symmetric tensor field breaks $SU(5)$ to $SO(5)$ at the scale f. The top quark partner T-quark, heavy gauge bosons (W_H, Z_H, A_H) and triplet scalar (Φ^{++}, Φ^{+}, Φ^{0}, Φ^{P}) are respectively introduced to cancel the Higgs mass one-loop quadratic divergence contributed by the top quark, gauge bosons and Higgs boson in the SM.

The top quark and T-quark can give the dominant contributions to the effective coupling hgg. Their Higgs couplings are given by

$$L_t \simeq -\frac{\lambda_1}{\sqrt{2}} f \bar{u}_L u_R \cos\frac{\sqrt{2}(v+h)}{f} + \lambda_2 f \bar{U}_L U_R + \text{h.c.},$$

where $c_\Sigma \equiv \cos\frac{\sqrt{2}(v+h)}{f}$ and $s_\Sigma \equiv \sin\frac{\sqrt{2}(v+h)}{f}$, with h and v being the neutral Higgs boson field and its VEV, respectively. After diagonalization of the mass matrix in Eq. (1), we can obtain the mass eigenstates t and T as well as their couplings with the Higgs boson [14],

$$y_t \bar{t} t h + y_T \bar{T} T h,$$

where

$$y_t = \frac{1}{2}(c_t^2 - s_t^2) + \frac{c_t^2 s_t^2 v^2}{f^2}.$$

The parameter x is a free parameter of the Higgs sector proportional to the triplet VEV v' and defined as $x = 4fv'/v^2$. The c_t and s_t are the mixing parameters between t and T, with

$$c_t = \frac{r}{\sqrt{r^2+1}}, \quad s_t = \frac{1}{\sqrt{1+r^2}}.$$

In addition to the Higgs couplings with charged fermions, the Higgs couplings with charged bosons also contribute to the effective coupling $h\gamma\gamma$, which are given as

$$y_W W^+ W^- h + 2 y_{W_H} W_H^+ W_H^- h + y_{\Phi^+} \Phi^+ \Phi^- h + y_{\Phi^{++}} \Phi^{++} \Phi^{--} h,$$

where

$$m_{W_H} = \frac{gf}{2sc},$$

$$y_W = 1 + \frac{v^2}{3f^2}, \quad y_{W_H} = \frac{v^2}{f^2}, \quad y_{\Phi^+} = \frac{v^2}{f^2} + \mathcal{O}\left(\frac{v^4}{f^4}\right), \quad y_{\Phi^{++}} = \frac{v^2}{f^2} + \mathcal{O}\left(\frac{v^4}{f^4}\right).$$

The c and s are the mixing parameters in the gauge boson sector. Since the $h\Phi^{++}\Phi^{--}$ coupling is very small, the contributions of the doubly-charged scalar can be ignored. In the littlest Higgs model, the relation between G_F and v is modified from its SM form, which can induce [14]

$$v = v_{SM}\left[1 - \frac{v_{SM}^2}{f^2}(1-x^2)\right],$$

where $v_{SM} = 246$ GeV is the SM Higgs VEV.

B. Littlest Higgs models with T-parity (LHT)

The LHT-I and LHT-II have the same kinetic term of the Σ field where T-parity can be naturally implemented, requiring that the coupling constant of $SU(2)_1$ ($U(1)_1$) equals that of $SU(2)_2$ ($U(1)_2$). This makes the four mixing parameters in the gauge sector c, s, c', s' equal to $1/\sqrt{2}$, respectively. Under T-parity, the SM bosons are T-even and the new bosons are T-odd. Therefore, the coupling $H^\dagger\Phi H$ is forbidden, leading to the triplet VEV $v' = 0$ and $x = 0$. Since the correction of W_H to the relation between G_F and v is forbidden by T-parity, the Higgs VEV v is modified as [15, 16]

$$v = v_{SM}\left(1 + \frac{v_{SM}^2}{f^2}\right).$$

The Higgs couplings with charged bosons of LHT-I and LHT-II can be obtained from Eq.(5) and Eq.(6) by taking $c = s = 1/\sqrt{2}$ and $x = 0$.

For each SM quark (lepton), a heavy copy of mirror quark (lepton) with T-odd quantum number is added in order to preserve T-parity. In the LHT-I [8, 15, 18], T-parity is simply implemented by adding the T-parity images for the original top quark interaction to make the Lagrangian T-invariant, so that the top quark partner canceling the one-loop quadratic divergence of Higgs mass is still T-even. Inspired by the way that the quadratic divergence given by top quark is canceled in the SLH, ref. [9] takes an alternative implementation of T-parity in LHT-II, where all new particles including the heavy top partner responsible for canceling the SM one-loop quadratic divergence are odd under T-parity.

In the LHT-I, the Higgs couplings with the heavy quarks are given by

$$L_\kappa \simeq -\frac{\lambda_1}{\sqrt{2}} f \bar{u}_L u'_R - m_q \bar{q}_L q_R - m_\chi \bar{\chi}_L \chi_R + \text{h.c.},$$

where $c_\xi \equiv \cos\frac{\sqrt{2}(v+h)}{f}$ and $s_\xi \equiv \sin\frac{\sqrt{2}(v+h)}{f}$. After diagonalization of the mass matrix in Eq. (9), we can obtain the T-odd mass eigenstates u_-, q_- and χ_-. In fact, there are three generations of T-odd particles, and we assume they are degenerate. The mass eigenstates t and T can be obtained by mixing the interaction eigenstates in Eq. (10), and their Higgs couplings are the same as those of LH with $x = 0$.

In the LHT-II, the Higgs couplings with the first two generations of heavy quarks are given by

$$L_q \simeq -\frac{\lambda_1}{\sqrt{2}} f \bar{u}_L u'_R - \bar{u}_L q_R + \bar{u}_L \chi_R - m_q \bar{q}_L q_R - m_\chi \bar{\chi}_L \chi_R + \text{h.c.}.$$

The mass eigenstates of u_-, q_- and χ_- and their Higgs couplings can be obtained by diagonalizing the mass matrix in Eq. (11).

The Higgs couplings with the third generation of heavy quarks are given by

$$L_q \simeq -\frac{\lambda_1}{\sqrt{2}} f \bar{u}_L u'_R - \bar{u}_L q_R - \bar{U}_L q_R - \bar{U}_L u'_R + \bar{u}_L \chi_R - m_q \bar{q}_L q_R - s_\Sigma \bar{u}_L u_R + c_\Sigma \bar{U}_L U_R + \text{h.c.},$$

where c_t is taken as $1/\sqrt{2}$. After diagonalization of the mass matrix in Eq. (12), we can obtain the mass eigenstates t, T, q_- and χ_- as well as their Higgs couplings.

For the SM down-type quarks (leptons), the Higgs couplings of LHT-I and LHT-II have two different cases [15]:

$$g_{h\bar{d}d} \simeq \frac{m_d}{v}\left(1 - \frac{v^2}{f^2}\right) \quad \text{for Case A,}$$

$$g_{h\bar{d}d} \simeq \frac{m_d}{v}\left(1 + \frac{v^2}{f^2}\right) \quad \text{for Case B.}$$

The relation of down-type quark couplings also applies to the lepton couplings.

C. Simplest little Higgs model (SLH)

The SLH model [4] is based on $[SU(3) \times U(1)_X]^2$ global symmetry. The gauge symmetry $SU(3) \times U(1)_X$ is broken down to the SM electroweak gauge group by two copies of scalar fields Φ_1 and Φ_2, which are triplets under $SU(3)$ with aligned VEVs f_1 and f_2.

The gauged $SU(3)$ symmetry promotes the SM fermion doublets into $SU(3)$ triplets. The Higgs couplings with the quarks are given by

$$L_t \simeq -\lambda c_{\beta t} c'_1 (s_1 t'_L + c_1 T'_L) + s_{\beta t} c'_2 (s_2 t'_L + c_2 T'_L) + \text{h.c.},$$

$$L_d \simeq -\lambda c_{\beta d} c'_1 (s_1 d'_L + c_1 D'_L) + s_{\beta d} c'_2 (s_2 d'_L + c_2 D'_L) + \text{h.c.},$$

$$L_s \simeq -\lambda c_{\beta s} c'_1 (s_1 s'_L + c_1 S'_L) + s_{\beta s} c'_2 (s_2 s'_L + c_2 S'_L) + \text{h.c.},$$

where $f = \sqrt{f_1^2 + f_2^2}$, $\tan\beta = f_2/f_1$, $c_\beta = f_1/f$, $s_\beta = f_2/f$, and

$$c_1 \equiv \cos\frac{h+v}{\sqrt{2} t_\beta f}, \quad s_1 \equiv \sin\frac{h+v}{\sqrt{2} t_\beta f},$$

$$c_2 \equiv \cos\frac{(h+v)(t_\beta^2+1)}{\sqrt{2}t_\beta f}, \quad s_2 \equiv \sin\frac{(h+v)(t_\beta^2+1)}{\sqrt{2}t_\beta f}.$$

After diagonalization of the mass matrices in Eqs. (13), (14) and (15), we can obtain the mass eigenstates (t, T), (d, D) and (s, S), which was performed numerically in our analysis, and the relevant couplings with the Higgs boson can be obtained.

The Higgs coupling with the charged bosons is given by [19],

$$y_W W^+ W^- h + 2y_{W'} W'^+ W'^- h,$$

where

$$y_W \simeq v_{SM}\left(1 - \frac{v_{SM}^2}{12f^2}\frac{\beta+1}{\beta}\right),$$

$$y_{W'} \simeq \frac{v_{SM}^3}{36f^4}\frac{(\beta+1)^2}{\beta}.$$

The Yukawa and gauge interactions break the global symmetry and then provide a potential for the Higgs boson. However, the Coleman-Weinberg potential alone is not sufficient since the generated Higgs mass is too heavy. Therefore, one can introduce a tree-level term which can partially cancel the Higgs mass,

$$\mu^2(\Phi_1^\dagger\Phi_2 + \text{h.c.}) = 2\mu^2 f^2 s_\beta c_\beta \cos\frac{\sqrt{2}s_\beta c_\beta h}{f}.$$

The Coleman-Weinberg potential involves the following parameters: f, t_β, λ, μ, m_h, v. Due to the modification of the observed W-boson mass, v is defined as [19]

$$v = v_{SM}\left[1 - \frac{v_{SM}^2}{12f^2}\frac{\beta+1}{\beta} + \frac{v_{SM}^4}{180f^4}\left(\frac{\beta^4+1}{\beta^2} + \frac{t_\beta^4+1}{t_\beta^2}\right)\right].$$

Assuming that there are no large direct contributions to the potential from physics at the cutoff, we can determine other parameters in Eq. (20) from f, t_β and m_h with the definition of v in Eq. (21).

III. THE DI-PHOTON $pp \to \gamma\gamma$ SIGNAL AT LHC

A. Calculations

At the LHC the cross section of single Higgs production via gluon-gluon fusion can be given by

$$\hat{\sigma}(gg \to h) = \frac{\pi^2}{8m_h^3}\Gamma(h \to gg)\tau_0\delta(1-\tau_0),$$

where $\tau_0 = m_h^2/s$ with $\sqrt{s}$ being the center-of-mass energy of the LHC. The total cross section is

$$\sigma(gg \to h) = \int_{\tau_0}^{1} dx_1 \int_{\tau_0/x_1}^{1} dx_2\, f_g(x_1, \mu_F^2) f_g(x_2, \mu_F^2) \hat{\sigma}(gg \to h),$$

where $f_g(x, \mu_F^2)$ is the parton distribution function of the gluon. Equation (22) shows that $\sigma(gg \to h)$ has a strong correlation with the decay width $\Gamma(h \to gg)$.

Now we discuss the Higgs decays in little Higgs models. For the tree-level decays $h \to XX$, where XX denotes WW, ZZ or SM fermion pairs, the little Higgs models give corrections via the corresponding modified couplings:

$$\Gamma(h \to XX) = \Gamma(h \to XX)_{SM} \left(\frac{g_{hXX}}{g_{hXX}^{SM}} \right)^2,$$

where $\Gamma(h \to XX)_{SM}$ is the SM decay width, and g_{hXX} and g_{hXX}^{SM} are the couplings of hXX in the little Higgs models and SM, respectively.

For a low Higgs mass, the loop-induced decay $h \to gg$ will be important. The general expression for the effective coupling hgg is shown in Appendix A. In the SM, the main contributions are from the top quark loop, and the little Higgs models give corrections via the modified couplings $ht\bar{t}$. In addition, the decay width of $h \to gg$ can also be corrected by the loops of heavy partner quark T in LH (T, D and S in SLH) (new T-even and T-odd quarks in LHT-I and LHT-II).

The general expression for the effective coupling $h\gamma\gamma$ is shown in Appendix A. In the SM, the top quark loop and W-boson loop give the main contributions to the decay $h \to \gamma\gamma$. The little Higgs models give corrections via the modified couplings $ht\bar{t}$ and hWW. In addition to the loops of the heavy quark mentioned in the decay $h \to gg$, the decay width of $h \to \gamma\gamma$ can also be corrected by the loops of W_H, Φ^+, Φ^{++} in LH, LHT-I and LHT-II (W' in SLH). Note that in the lepton sector, LHT-I, LHT-II and SLH also predict some neutral heavy neutrinos, which do not contribute to the couplings of $h\gamma\gamma$ at the one-loop level. Although charged heavy leptons are predicted in LHT-I and LHT-II, they do not have direct couplings with the Higgs boson.

In addition to the SM decay modes, the Higgs boson in LHT-I, LHT-II and SLH has some new important decay modes which are kinematically allowed in the parameter space.

In LHT-I the breaking scale f may be as low as 500 GeV [20], and the constraint in LHT-II is expected to be even weaker [9]. For a lower value of f, the lightest T-odd particle A_H may have a light mass, so that the decay $h \to A_H A_H$ can be open, with partial width

$$\Gamma(h \to A_H A_H) = \frac{g_{hA_HA_H}^2}{32\pi m_h} \sqrt{1 - x_{A_H}} \left(1 - x_{A_H} + \frac{3}{4} x_{A_H}^2 \right),$$

where $x_{A_H} = 4m_{A_H}^2/m_h^2$, and $g_{hA_HA_H}$ is the coupling constant of hA_HA_H. However, in LH the electroweak precision data requires f larger than a few TeV [6] and thus the decay $h \to A_HA_H$ is kinematically forbidden.

In SLH, the new decay modes are $h \to \eta\eta$ and $h \to Z\eta$, with partial widths given by

$$\Gamma(h \to \eta\eta) = \frac{g_{h\eta\eta}^2}{32\pi m_h}\sqrt{1-x_\eta},$$

$$\Gamma(h \to Z\eta) = \frac{g_{hZ\eta}^2}{16\pi m_h}\lambda^{1/2}(1, x_\eta, x_Z)\left(1 - \frac{x_\eta}{2} - \frac{x_Z}{2} - \frac{x_\eta x_Z}{2} + \frac{x_\eta^2}{2} + \frac{x_Z^2}{2}\right),$$

where $x_\eta = 4m_\eta^2/m_h^2$ and $\lambda(1, x, y) = (1-x-y)^2 - 4xy$.

B. Numerical results and discussions

In our calculations the SM input parameters are taken from [21]. For the SM decay channels, the relevant higher-order QCD and electroweak corrections are considered using the code Hdecay [22]. We focus on a light SM-like Higgs boson, with mass in the range 110–140 GeV.

In the LH model the new free parameters are f, c, c', c_t and x, where $0 < c < 1$, $0 < c' < 1$, $0 < c_t < 1$, $0 < x < 1$.

Taking $f = 1$ TeV, $f = 2$ TeV and $f = 4$ TeV, we scan over these parameters in the above ranges and show scatter plots. The parameter c_t can control the Higgs couplings with t, T and m_T, which is involved in the calculation of $\Gamma(h \to t\bar{t})$, $\Gamma(h \to gg)$ and $\Gamma(h \to \gamma\gamma)$. For a light Higgs boson, the decay mode $h \to t\bar{t}$ is kinematically forbidden. For $\Gamma(h \to gg)$ and $\Gamma(h \to \gamma\gamma)$, the c_t dependence of the top-quark loop and T-quark loop can cancel each other to a large extent (see Eq. (3)). Therefore, the rate $\sigma(gg \to h) \times \mathrm{BR}(h \to \gamma\gamma)$ is not sensitive to c_t for a light Higgs boson.

The rate $\sigma(pp \to \gamma\gamma)$ for the LH model is shown in Fig. 1 [Figure 1: see original paper] normalized to the SM prediction. We can see that the LH model always suppresses the rate $\sigma(pp \to \gamma\gamma)$, but the suppression can only reach about 10% for small values of f. As f increases, the suppression becomes smaller, and the rate is not sensitive to the parameters c, c', c_t and x. For example, for $f = 4$ TeV, the scatter plots are shown in line with the rate being around 99.6 percent of the SM prediction.

In LHT-I and LHT-II, the parameters c, c' and x are fixed as $c = c' = 1/\sqrt{2}$ and $x = 0$. Similar to the LH model, the result is not sensitive to c_t in LHT-I and LHT-II. Taking $c_t = 1/\sqrt{2}$ can simplify the top quark Yukawa sector in LHT-II [9, 16], and this choice is also favored by electroweak precision data [20]. The new heavy quarks can contribute to the decay widths of $h \to gg$ and $h \to \gamma\gamma$ via loops, which are not sensitive to the actual values of their masses as long as they are much larger than half of the Higgs boson mass [14].

The rate $\sigma(pp \to \gamma\gamma)$ for LHT-I and LHT-II is shown in Fig. 2 [Figure 2: see original paper] normalized to the SM prediction. We can see that LHT-I and LHT-II always suppress the rate, and the suppression is much more sizable than that of LH. For each model the rate in Case A is smaller than the rate in Case B because the coupling $hb\bar{b}$ in Case A is less suppressed than in Case B. Besides, we see that for $f = 500$ GeV and m_h in the range 130-140 GeV, the rate in both models drops drastically. The reason for such severe suppression is that the new decay mode $h \to A_H A_H$ is open and dominant in this parameter space, making the total decay width of the Higgs boson much larger than the SM value.

In SLH the new free parameters are f, t_β, x_t, λ (m_T), x_d (m_D) and x_s (m_S). As shown above, λ, μ (m_η) can be determined by f, t_β, m_h and v assuming that physics at the cutoff does not give large direct contributions to the potential. Ref. [4] shows that LEP-II data requires $f > 2$ TeV, and ref. [23] gives a lower bound of $f > 4.5$ TeV from the oblique parameter S. Recent studies of Z leptonic decay and $\tau^+\tau^-\gamma$ process at the Z-pole show that the scale f should be larger than 5.6 TeV and 5.4 TeV, respectively [24]. Here we assume the new flavor mixing matrices in lepton and quark sectors are diagonal [5, 25], so that f and t_β are free from experimental constraints of lepton and quark flavor violating processes.

For the perturbation expansion to be valid, t_β cannot be too large for a fixed f. If we require $(v^2/f^2) < 0.1$ in the expansion of v, t_β should be below 10, 20, and 28 for $f = 2$ TeV, 4 TeV, and 5.6 TeV, respectively. The small masses of the d quark and s quark require that x_d, λ and x_s are very small, respectively, so there is almost no mixing between the SM down-type quarks and their heavy partners, and the results are not sensitive to them. We take x_d, λ and x_s such that the masses of D and S are in the range 0.5–2 TeV with other parameters fixed as in our calculation.

The rate $\sigma(pp \to \gamma\gamma)$ for SLH at the LHC is shown in Fig. 3 [Figure 3: see original paper]. We can see that SLH always suppresses the rate, and the suppression is more sizable for large t_β. When t_β is large enough, such as $t_\beta = 10$ for $f = 2$ TeV ($t_\beta = 20$ for $f = 4$ TeV or $t_\beta = 25$ for $f = 5.6$ TeV), the new mode $h \to \eta\eta$ is open and dominant, which can further suppress the rate (the suppression can be up to 90%).

If the ultraviolet completion of the theory can give sizable contributions to the Coleman-Weinberg potential, the correlation of the parameters x_t, λ, μ (m_η), f, t_β, m_h and v can be loosened greatly. In Fig. 4 [Figure 4: see original paper], we scan the following parameter space:

$$0.5 \text{ TeV} < m_D(m_S) < 2 \text{ TeV}, \quad 1 < t_\beta < 30,$$

where the parameter x_t, λ is replaced with m_T and the bound $(v^4/f^4)/(v^2/f^2) < 0.1$ is still valid. To avoid suppression of the rate by the new decay modes $h \to \eta\eta$ and $h \to \eta Z$, we take $m_\eta > 2m_h$, so that the result is independent of the parameter m_η (μ). Fig. 4 shows that, compared to the SM prediction, SLH still suppresses the rate $\sigma(pp \to \gamma\gamma)$ in more general parameter space.

From our results we see that compared to the SM prediction the rate $\sigma(pp \to \gamma\gamma)$ is always suppressed in these typical little Higgs models. Now we analyze this suppression in detail. Equations (A4) and (A8) show that the effective coupling hgg is not sensitive to heavy quark masses as long as they are much larger than half of the Higgs boson mass. Therefore, according to Eqs. (A9) and (A10), the effective coupling hgg is approximately proportional to $\sum_i y_{f_i}$, where y_{f_i} is defined in Eq. (A2).

Table 1 shows the value of y_{f_i} corresponding to the quark f_i and BR$(h \to \gamma\gamma)$ normalized to SM prediction in these little Higgs models. The parameters are fixed as $c = c' = c_t = 1/\sqrt{2}$ and $x = 0$ in LH, and $t_\beta = 10$, $m_T = 450$ (798) GeV, $m_D = 548$ (1039) GeV, $m_S = 597$ (1132) GeV, $m_\eta = 42.7$ (179) GeV corresponding to $f = 2$ (4) TeV in SLH.

Because of the sizable suppression of the coupling $hb\bar{b}$, BR$(h \to \gamma\gamma)$ is generally not suppressed, unless the new decay mode is open and dominant, as shown for SLH with $f = 2$ TeV in Table 1 (the new mode is $h \to \eta\eta$). In these little Higgs models, $\sum_i y_{f_i}$ can be respectively less than 1, which shows that $\sigma(gg \to h)$ is suppressed compared to the SM prediction. There are common reasons for these models: (i) All models are based on non-linear sigma models, and the Yukawa coupling $ht\bar{t}$ is suppressed with the expansion of the non-linear sigma fields. (ii) The top quark partner cancels the quadratic divergence of Higgs mass contributed by the top quark, which induces opposite signs for the Yukawa couplings of the top quark and its partner.

The forthcoming measurement of the di-photon signal at the LHC will allow probing of these little Higgs models. For example, if the signal rate is found to be above the SM prediction, these little Higgs models will be immediately disfavored. If the signal rate is found to be much lower than the SM prediction, then SLH and LHT will be favored. However, due to the free parameters involved in the signal rate for each model, it will be hard for the LHC to clearly discriminate between these different little Higgs models. For precision tests of different models, the ILC collider is necessary [27].

IV. CONCLUSION

We performed a comparative study for the LHC di-photon signal by considering four different little Higgs models, namely the LH, LHT-I, LHT-II and SLH. We obtained the following observations: (i) Compared with the SM prediction, the di-photon signal rate is always suppressed in these models; (ii) The suppression extent is different in different models, which is below 10% in the LH but can reach 90% in the LHT-I, LHT-II and SLH, especially in the parameter space where new decay modes ($h \to \eta\eta$ for SLH and $h \to A_H A_H$ for LHT-I and LHT-II) are open and dominant. Therefore, discovering the light Higgs predicted by these little Higgs models through the di-photon channel at the LHC will be more difficult than discovering the SM Higgs boson.

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Appendix A: The effective couplings of Higgs-photon-photon and Higgs-gluon-gluon

The effective Higgs-photon-photon coupling can be written as [14, 28]

$$\mathcal{L}_{h\gamma\gamma} = I F_{\mu\nu} F^{\mu\nu} h,$$

where $F^{\mu\nu}$ is the electromagnetic field strength tensor. With the Higgs boson couplings to the charged fermion f_i, vector boson V_i and scalar S_i given by

$$y_{f_i} \bar{f}_i f_i h + y_{V_i} V_i V_i h + y_{S_i} S_i S_i h,$$

the factor I in Eq. (A1) can be written as

$$I = \frac{1}{16\pi^2 v}\left[\sum_{f_i} N_c^{f_i} Q_{f_i}^2 y_{f_i} F_{1/2}(\tau_{f_i}) + \sum_{V_i} Q_{V_i}^2 y_{V_i} F_1(\tau_{V_i}) + \sum_{S_i} Q_{S_i}^2 y_{S_i} F_0(\tau_{S_i})\right],$$

where Q_X (X denotes f_i, V_i and S_i) is the electric charge for a particle X running in the loop, and $N_c^{f_i}$ is the color factor for f_i. The dimensionless loop factors are

$$F_{1/2}(\tau) = 2\tau[1 + (1-\tau) f(\tau)],$$
$$F_1(\tau) = -2 - 3\tau - 3\tau(2-\tau) f(\tau),$$
$$F_0(\tau) = \tau[1 - \tau f(\tau)],$$

where $\tau_X = 4m_X^2/m_h^2$ and

$$f(\tau_X) = \begin{cases} [\sin^{-1}(1/\sqrt{\tau_X})]^2, & \tau_X \geq 1 \\ -\frac{1}{4}\left[\ln\frac{\eta_+}{\eta_-} - i\pi\right]^2, & \tau_X < 1 \end{cases}$$

with $\eta_\pm = 1 \pm \sqrt{1-\tau_X}$. When the masses of particles in the loops are much larger than half of the Higgs boson mass, we can obtain

$$F_{1/2}(\tau_{f_i}) \to \frac{4}{3}, \quad F_1(\tau_{V_i}) \to 7, \quad F_0(\tau_{S_i}) \to 1.$$

The effective Higgs-gluon-gluon coupling can be written as [14, 28]

$$\mathcal{L}_{hgg} = I_{hgg} G_{\mu\nu}^a G_a^{\mu\nu} h,$$

where $G^a_{\mu\nu} = \partial_\mu g^a_\nu - \partial_\nu g^a_\mu + g_s f^{abc} g_{\mu b} g_{\nu c}$ and the factor I_{hgg} from the contributions of quarks running in the loops is given by

$$I_{hgg} = \frac{\alpha_s}{8\pi v} \sum_q y_q F_{1/2}(\tau_q),$$

with $\tau_q = 4m_q^2/m_h^2$. Once the interactions in Eq. (A2) are given, we can obtain the effective $h\gamma\gamma$ and hgg couplings from the above formulas. The relevant Higgs interactions in LH, LHT-I, LHT-II and SLH are listed in Sec. II. Here the Higgs interactions with light fermions are not given since their contributions can be ignored.

References

[1] N. Arkani-Hamed, A. G. Cohen and H. Georgi, Phys. Lett. B 513, 232 (2001); N. Arkani-Hamed, et al., JHEP 0208, 021 (2002).
[2] D. E. Kaplan and M. Schmaltz, JHEP 0310, 039 (2003); I. Low, W. Skiba, and D. Smith, Phys. Rev. D 66, 072001 (2002); S. Chang and J. G. Wacker, Phys. Rev. D 69, 035002 (2004); T. Gregoire, D. R. Smith, and J. G. Wacker, Phys. Rev. D 69, 115008 (2004); W. Skiba and J. Terning, Phys. Rev. D 68, 075001 (2003); S. Chang, JHEP 0312, 057 (2003); H. Cai, H.-C. Cheng, and J. Terning, JHEP 0905, 045 (2009); A. Freitas, P. Schwaller, and D. Wyler, JHEP 0912, 027 (2009).
[3] N. Arkani-Hamed, A. G. Cohen, E. Katz, A. E. Nelson, JHEP 0207, 034 (2002).
[4] M. Schmaltz, JHEP 0408, 056 (2004).
[5] T. Han, H. E. Logan and L. T. Wang, JHEP 0601, 099 (2006).
[6] C. Csaki, et al., Phys. Rev. D 67, 115002 (2003); Phys. Rev. D 68, 035009 (2003); J. L. Hewett, F. J. Petriello, T. G. Rizzo, JHEP 0310, 062 (2003); M. C. Chen, S. Dawson, Phys. Rev. D 70, 015003 (2004); M. C. Chen, et al., Mod. Phys. Lett. A 21, 621 (2006).
[7] H. C. Cheng and I. Low, JHEP 0309, 051 (2003).
[8] H. C. Cheng and I. Low, JHEP 0408, 061 (2004); I. Low, JHEP 0410, 067 (2004).
[9] H. C. Cheng, I. Low and L. T. Wang, Phys. Rev. D 74, 055001 (2006).
[10] A. Martin, hep-ph/0602206.
[11] See, e.g., S. Mukhopadhyay, B. Mukhopadhyaya, A. Nyffeler, JHEP 1005, 001 (2010); R. S. Hundi, B. Mukhopadhyaya, A. Nyffeler, Phys. Lett. B 649, 280 (2007); S. Yang, Phys. Lett. B 675, 352 (2009); C.-X. Yue, et al., Europhys. Lett. 77, 51003 (2007); J. J. Liu et al., Phys. Rev. D 70, 015001 (2004); A. Datta, et al., Phys. Lett. B 659, 308 (2008); L. Wang, et al., Phys. Rev. D 76, 017702 (2007); Phys. Rev. D 75, 074006 (2007); Phys. Rev. D 82, 095009 (2010); P. Kai, et al., Phys. Rev. D 76, 015012 (2007); X. F. Han, L. Wang, J. M. Yang, Nucl. Phys. B 825, 222 (2010); K. Cheung, J. Song, P. Tseng, and Q.-S. Yan, Phys. Rev. D 78, 055015 (2008); W. Kilian, D. Rainwater, and J. Reuter, Phys. Rev. D 71, 015008 (2005).

[12] ATLAS Collaboration, ATLAS-CONF-2011-085 (2011).
[13] K. Hsieh and C. P. Yuan, Phys. Rev. D 78, 053006 (2008); J. Cao, et al., arXiv:1103.0631; I. Low and S. Shalgar, JHEP 0904, 091 (2009); S. Moretti and S. Munir, Eur. Phys. J. C 47, 791 (2006); U. Ellwanger, arXiv:1012.1201; M. Almarashi and S. Moretti, arXiv:1011.6547; Y.-B. Liu, X. L. Wang, Phys. Lett. B 694, 417 (2011).
[14] T. Han, H. E. Logan, B. McElrath, L.-T. Wang, Phys. Lett. B 563, 191 (2003).
[15] C. R. Chen, K. Tobe, C. P. Yuan, Phys. Lett. B 640, 263 (2006).
[16] L. Wang, J. M. Yang, Phys. Rev. D 77, 015020 (2008); Phys. Rev. D 79, 055013 (2009).
[17] T. Han, H. E. Logan, B. McElrath, L.-T. Wang, Phys. Rev. D 67, 095004 (2003).
[18] J. Hubisz and P. Meade, Phys. Rev. D 71, 035016 (2005).
[19] K. Cheung and J. Song, Phys. Rev. D 76, 035007 (2007).
[20] J. Hubisz, P. Meade, A. Noble, M. Perelstein, JHEP 0601, 135 (2006).
[21] C. Amsler, et al., Phys. Lett. B 667, 1 (2008).
[22] A. Djouadi, J. Kalinowski and M. Spira, Comput. Phys. Commun. 108, 56 (2006).
[23] G. Marandella, C. Schappacher and A. Strumia, Phys. Rev. D 72, 035014 (2005).
[24] A. G. Dias, et al., Phys. Rev. D 77, 055001 (2008); A. Gutierrez-Rodriguez, Mod. Phys. Lett. A 25, 703-713 (2010).
[25] J. I. Illana, M. D. Jenkins, Acta Phys. Polon. B 40, 3143 (2009); F. d. Aguila, J. I. Illana, M. D. Jenkins, Nucl. Phys. Proc. Suppl. 205, 158 (2010).
[26] F. d. Aguila, J. I. Illana, M. D. Jenkins, JHEP 1103, 080 (2011).
[27] L. Wang, F. Xu and J. M. Yang, JHEP 1001, 107 (2010).
[28] J. F. Gunion, H.E. Haber, G.L. Kane, and S. Dawson, The Higgs Hunter's Guide, Addison-Wesley, Reading, MA, U.S.A. (1990).

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