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Higgs boson mass bounds in seesaw extended standard model with non-minimal gravitational coupling

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ABSTRACT

In the presence of non-minimal gravitational coupling $\xi H^\dagger H \mathcal{R}$ between the standard model (SM) Higgs doublet H and the curvature scalar $\mathcal{R}$, the effective ultraviolet cutoff scale is given by $\Lambda \approx m_P/\xi$, where m_P is the reduced Planck mass, and $\xi \gtrsim 1$ is a dimensionless coupling constant. In type I and type III seesaw extended SM, which can naturally explain the observed solar and atmospheric neutrino oscillations, we investigate the implications of this non-minimal gravitational coupling for the SM Higgs boson mass bounds based on vacuum stability and perturbativity arguments. A lower bound on the Higgs boson mass close to 120 GeV is realized with type III seesaw and $\xi \sim 10$–10^3.

The search for the SM Higgs boson is arguably the single most important mission of the LHC. According to precision electroweak data and the direct lower mass bound from LEP II, a Higgs boson mass in the range of 114.4 GeV $\lesssim m_H \lesssim$ 180 GeV [1] is favored. If one takes the reduced Planck mass $m_P = 2.4 \times 10^{18}$ GeV as a natural cutoff scale of the SM, theoretical considerations based on vacuum stability and perturbativity arguments narrow the SM Higgs boson mass bounds somewhat, namely 128 GeV $\lesssim m_H \lesssim$ 175 GeV [2,3]. Very recently, it has been reported [4] that the SM Higgs boson mass in the mass range 158 GeV $\lesssim m_H \lesssim$ 175 GeV is excluded at 95% C.L. by the direct searches at the Tevatron.

Clearly, if there exists some new physics beyond the SM between the electroweak scale and the reduced Planck scale, it can affect these theoretical Higgs boson mass bounds. The seesaw mechanism is a simple and promising extension of the SM to incorporate the neutrino masses and mixings observed in solar and atmospheric neutrino oscillations. There are three main seesaw extensions of the SM, type I [5], type II [6], and type III [7], in which new particles, singlet right-handed neutrinos, SU(2) triplet scalar, and SU(2) triplet right-handed neutrinos, respectively, are introduced. These new particles contribute to the renormalization group equations (RGEs) at energies higher than the seesaw scale and as a result, the Higgs boson mass bounds can be significantly altered. The important implications of seesaw models on the Higgs boson mass bounds have been investigated with the reduced Planck mass cutoff for the various seesaw models, type I [8,9], type II [10] and type III [9].

In general, the non-minimal gravitational coupling between the SM Higgs doublet and the curvature scalar,

$$\xi H^\dagger H \mathcal{R}, \tag{1}$$

can be introduced in the SM. This coupling opens up a very intriguing scenario for inflationary cosmology, namely, the possibility that the SM Higgs field may play the role of inflation field, and this has been investigated in several recent papers [11–17]. As pointed out in [18], in the presence of the non-minimal gravitational coupling, it is natural to identify the effective ultraviolet cutoff scale as

$$\Lambda \approx \frac{m_P}{\xi}, \tag{2}$$

for $\xi \gtrsim 1$, rather than m_P. Note that the cutoff may depend on the background field value which in our case is of order the electroweak scale (see last refs. in [11] and [17]).

In this Letter, we extend previous work on the Higgs boson mass bounds in type I and III seesaw extended SM [8,9] to the case with non-minimal gravitational coupling. The ultraviolet cutoff scale is taken to be $\Lambda = m_P/\xi$ in our analysis. We will show that the gravitational coupling as well as type I and III seesaw effects can dramatically alter the vacuum stability and perturbativity bounds on the SM Higgs boson mass. In particular, the vacuum stability bound on the Higgs boson mass can be lowered to 120 GeV or so, significantly below the usual lower bound of about 128 GeV found in the absence of seesaw and with $\xi = 0$.

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In type I seesaw, three generations of SM-singlet right-handed neutrinos ψ_i $(i = 1, 2, 3)$ are introduced. The relevant terms in the Lagrangian are given by

$$\mathcal{L} \supset -y_{ij}\overline{\ell_i}\psi_j H - M_R\overline{\psi_i^c}\psi_i, \tag{3}$$

where ℓ_i is the i-th generation SM lepton doublet. For simplicity, we assume in this Letter that the three right-handed neutrinos are degenerate in mass (M_R). At energies below M_R, the heavy right-handed neutrinos are integrated out and the effective dimension five operator is generated by the seesaw mechanism. After electroweak symmetry breaking, the light neutrino mass matrix is obtained as

$$\mathbf{M}_\nu = \frac{v^2}{2M_R}\mathbf{Y}_\nu^T\mathbf{Y}_\nu, \tag{4}$$

where $v = 246$ GeV is the VEV of the Higgs doublet, and $\mathbf{Y}_\nu = y_{ij}$ is a 3×3 Yukawa matrix.

The basic structure of type III seesaw is similar to type I seesaw, except that instead of the singlet right-handed neutrinos, three generations of fermions which transforms as $(\mathbf{3}, 0)$ under the electroweak gauge group $SU(2)_L \times U(1)_Y$ are introduced:

$$\psi_i = \sum_a \frac{\sigma^a}{2}\psi_i^a = \frac{1}{2}\begin{pmatrix} \psi_i^0 & \sqrt{2}\psi_i^+ \\ \sqrt{2}\psi_i^- & -\psi_i^0 \end{pmatrix}. \tag{5}$$

With canonically normalized kinetic terms for the triplet fermions, we replace the SM-singlet right-handed neutrinos of type I seesaw in Eq. (3) by these SU(2) triplet fermions. Assuming degenerate masses (M_R) for the three triplet fermions, the light neutrino mass matrix via type III seesaw mechanism is obtained as

$$\mathbf{M}_\nu = \frac{v^2}{8M_R}\mathbf{Y}_\nu^T\mathbf{Y}_\nu. \tag{6}$$

For a renormalization scale $\mu < M_R$, the heavy fermions are decoupled, and there is no effect on the RGEs for the SM couplings. However, in the presence of the non-minimal gravitational coupling, a factor $s(\mu)$ defined as

$$s(\mu) = \frac{1 + \frac{\xi\mu^2}{m_P^2}}{1 + (6\xi + 1)\frac{\xi\mu^2}{m_P^2}}, \tag{7}$$

is assigned to each term in the RGEs associated with the physical Higgs boson loop corrections [11,12,15]. In our analysis, we employ 2-loop RGEs for the SM couplings. Since the SM beta functions suitably modified with the s-factor are known only at 1-loop level, we employ the beta functions with the s-factor for 1-loop corrections, while the beta functions for 2-loop corrections are without the s-factor. We have checked that the effects of the s-factor in beta functions for 2-loop corrections are negligible as far as our final results are concerned [13]. In fact, even for 1-loop corrections, the effects of the s-factor are not so important, namely, negligible for the perturbative bound, while the perturbative bound is reduced by, at most, a few GeV.

For the three SM gauge couplings with a renormalization scale $\mu < M_R$, we have

$$\frac{dg_i}{d\ln\mu} = \frac{b_i}{16\pi^2}g_i^3 + \frac{g_i^3}{(16\pi^2)^2}\left(\sum_{j=1}^3 B_{ij}g_j^2 - C_i y_t^2\right), \tag{8}$$

where g_i $(i = 1, 2, 3)$ are the SM gauge couplings,

$$b_i = \left(\frac{81+s}{20}, -\frac{39-s}{12}, -7\right), \qquad B_{ij} = \begin{pmatrix} \frac{199}{50} & \frac{27}{10} & \frac{44}{5} \\ \frac{9}{10} & \frac{35}{6} & 12 \\ \frac{11}{10} & \frac{9}{2} & -26 \end{pmatrix},$$
$$C_i = \left(\frac{17}{10}, \frac{3}{2}, 2\right), \tag{9}$$

and we have included the contribution from the top Yukawa coupling (y_t). We use the top quark pole mass $M_t = 173.1$ GeV and the strong coupling constant at the Z-pole (M_Z) $\alpha_S = 0.1193$ [19]. For the top Yukawa coupling, we have

$$\frac{dy_t}{d\ln\mu} = y_t\left(\frac{1}{16\pi^2}\beta_t^{(1)} + \frac{1}{(16\pi^2)^2}\beta_t^{(2)}\right). \tag{10}$$

Here the one-loop contribution is

$$\beta_t^{(1)} = \left(4 + \frac{s}{2}\right)y_t^2 - \left(\frac{17}{20}g_1^2 + \frac{9}{4}g_2^2 + 8g_3^2\right), \tag{11}$$

while the two-loop contribution is given by [20]

$$\begin{aligned}\beta_t^{(2)} = &-12y_t^4 + \left(\frac{393}{80}g_1^2 + \frac{225}{16}g_2^2 + 36g_3^2\right)y_t^2 \\ &+ \frac{1187}{600}g_1^4 - \frac{9}{20}g_1^2g_2^2 + \frac{19}{15}g_1^2g_3^2 - \frac{23}{4}g_2^4 \\ &+ 9g_2^2g_3^2 - 108g_3^4 + \frac{3}{2}\lambda^2 - 6\lambda y_t^2.\end{aligned} \tag{12}$$

In solving the RGE for the top Yukawa coupling, its value at $\mu = M_t$ is determined from the relation between the pole mass and the running Yukawa coupling [21,22],

$$\begin{aligned}M_t \simeq m_t(M_t)\Bigg(1 &+ \frac{4}{3}\frac{\alpha_3(M_t)}{\pi} \\ &+ 11\left(\frac{\alpha_3(M_t)}{\pi}\right)^2 - \left(\frac{m_t(M_t)}{2\pi v}\right)^2\Bigg),\end{aligned} \tag{13}$$

with $y_t(M_t) = \sqrt{2}m_t(M_t)/v$, where $v = 246$ GeV. Here, the second and third terms in parenthesis correspond to one- and two-loop QCD corrections, respectively, while the fourth term comes from the electroweak corrections at one-loop level.

The RGE for the Higgs quartic coupling is given by [20],

$$\frac{d\lambda}{d\ln\mu} = \frac{1}{16\pi^2}\beta_\lambda^{(1)} + \frac{1}{(16\pi^2)^2}\beta_\lambda^{(2)}, \tag{14}$$

with

$$\begin{aligned}\beta_\lambda^{(1)} = &(3 + 9s^2)\lambda^2 - \left(\frac{9}{5}g_1^2 + 9g_2^2\right)\lambda \\ &+ \frac{9}{4}\left(\frac{3}{25}g_1^4 + \frac{2}{5}g_1^2g_2^2 + g_2^4\right) + 12y_t^2\lambda - 12y_t^4,\end{aligned} \tag{15}$$

and

$$\begin{aligned}\beta_\lambda^{(2)} = &-78\lambda^3 + 18\left(\frac{3}{5}g_1^2 + 3g_2^2\right)\lambda^2 \\ &- \left(\frac{73}{8}g_2^4 - \frac{117}{20}g_1^2g_2^2 - \frac{1887}{200}g_1^4\right)\lambda - 3\lambda y_t^4 \\ &+ \frac{305}{8}g_2^6 - \frac{289}{40}g_1^2g_2^4 - \frac{1677}{200}g_1^4g_2^2 - \frac{3411}{1000}g_1^6 \\ &- 64g_3^2y_t^4 - \frac{16}{5}g_1^2y_t^4 - \frac{9}{2}g_2^4y_t^2 \\ &+ 10\lambda\left(\frac{17}{20}g_1^2 + \frac{9}{4}g_2^2 + 8g_3^2\right)y_t^2 \\ &- \frac{3}{5}g_1^2\left(\frac{57}{10}g_1^2 - 21g_2^2\right)y_t^2 - 72\lambda^2y_t^2 + 60y_t^6.\end{aligned} \tag{16}$$