Hidden-sector-assisted 125 GeV Higgs boson

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In order to significantly raise the mass of the supersymmetry Higgs boson, we consider a radiative correction to it by heavy (~1 TeV) hidden sector fields, which communicate with the Higgs through relatively heavy “messengers” (300–500 GeV). The messenger fields (S, T) are coupled to the Higgs (“yHSHd, Hd” yH ≲ 0.7) and also to hidden sector fields with a Yukawa coupling of order unity. The hidden sector fields are assumed to be large representations of a hidden gauge group, and so their scalar partner masses can be heavier than other typical soft scalars in the visible sector. Even with a relatively small yH (~0.2) or tanβ ~ 10 but without stop-top’s considerable contributions, the radiative correction by such hidden sector fields can be enhanced enough to yield the 125 GeV Higgs mass.

The solution of the naturalness problem and the gauge coupling unification are indeed the great achievements of the minimal supersymmetric standard model (MSSM). By supersymmetry (SUSY) in the MSSM, the smallness of the Higgs mass can be perturbatively valid up to the fundamental scale, and so the standard model (SM) can be naturally embedded in a unified theory in the Planck or string scale. The gauge coupling unification in the MSSM might be strong evidence for it.

Recently, the ATLAS and CMS collaborations reported the excesses of events for the γγ, ZZ* → 4ℓ, and WW* → 2ℓν channels around 125 GeV invariant mass. They seemingly imply the presence of the Higgs with 125 GeV mass at 3σ confidence level [1]. They might be accepted as the signals of the MSSM, since the 125 GeV Higgs mass is still inside the range that the MSSM admits, just assuming the relatively heavy stop (≳ a few TeV) [2]. Indeed the radiative correction by the top (tL,R) and stop (tL,R), and the maximal mixing in (tL,tR) via the SUSY-breaking “A term” can raise the Higgs mass up to 135 GeV in the decoupling limit of the CP odd Higgs [2]. For 125 GeV Higgs mass, however, the naturalness of the relatively light Higgs mass in the MSSM would now become seriously challenged with such a heavy stop.

The naturalness problem of the Higgs mass may be much alleviated in the next-to-minimal supersymmetric standard model (NMSSM) [3]. In the NMSSM, the “μ term” of the MSSM is promoted to a trilinear superpotential “yHSHd∗,” introducing a singlet superfield S [4]. It provides an additional F-term quartic coupling in the Higgs potential apart from the quartic potential coming from the D-term in the MSSM. Thus, the maximum mass of the lightest Higgs is modified as

\[ m_h^2 \approx M_Z^2 \cos^2 2β + y_H^2 v^2 \sin^2 2β + \Delta m_h^2, \]

where \( v^2 \equiv v_u^2 + v_d^2 \approx (174 \text{ GeV})^2 \) and \( \Delta m_h^2 \) indicates the radiative correction by the (s)top. In the MSSM (yH = 0), for instance, the needed \( \Delta m_h^2 \) is around (85 GeV)² for the 125 GeV Higgs mass. To obtain \( m_h \approx 125 \text{ GeV} \) with smaller contributions of the (s)top, \( \Delta m_h^2 \) should be larger than 0.5 for tanβ > 1. In the NMSSM, however, yH is stringently restricted by the Landau pole constraint, yH ≲ 0.7 [3]. Moreover, only a quite narrow range of tanβ, 1 ≲ tanβ ≲ 3, which gives almost the maximal values of \( \sin^2 2β \), is allowed for O(100) GeV stop mass.

The fourth generation of the SM particles or extra vectorlike matter fields, if they exist at the electroweak scale and couple to the Higgs, can also contribute to \( \Delta m_h^2 \) by adding extra radiative corrections to the Higgs mass [5]. However, the introduction of new colored particles with an order-one Yukawa coupling to raise the Higgs mass would exceedingly affect the production rate of two gluons to Higgs (gg → h0) as well as the decay rate of Higgs to two gammas (h0 → γγ); they might result in immoderate deviations from the SM predictions for their rates, unless some new invisible Higgs decay channels open [6]. Moreover, for leaving intact the gauge coupling unification in the MSSM, the newly introduced extra vectorlike matter fields should compose SU(5) or SO(10) multiplets. If the low-energy theory is not embedded in SU(5) [or SO(10)] grand unifications below the string scale but still keeps the gauge coupling unification as seen in many string models [7], one needs to explore other possibilities for \( m_h \approx 125 \text{ GeV} \).

In this paper, we will propose a singlet extension of the MSSM to achieve \( m_h \approx 125 \text{ GeV} \), in which the radiative correction by hidden sector fields mainly contributes to \( \Delta m_h^2 \). Since the radiative correction can be enhanced just with the MSSM singlets but without the (s)top quark’s significant contributions, the gauge coupling unification is maintained, and also the naturalness of the SUSY Higgs can still be supported by the relatively light stop.

We consider the following singlet-extended superpo-
potential:
\[
W = y_H S H_u H_d + m_S S \bar{S} + y_N S N \bar{N} + m_N N \bar{N} + W_{\text{MSSM}},
\]
where \((S, \bar{S})\) and \((N, \bar{N})\) are the neutral superfields under the SM gauge symmetry. For the simple presentation, in Eq. (2) we do not explicitly write down the relevant superpotential of the MSSM including the \(\mu\) term. Since the \(\mu\) term can actually be generated from the tadpole term of \(S\), the SUSY-breaking \(A\) term corresponding to the first term in Eq. (2), the bare \(\mu\) term would not be essential. However, we will keep it to avoid unwanted Peccei-Quinn (PQ) symmetry breaking at the electroweak energy scale. We suppose that the MSSM singlets \(N, \bar{N}\) are proper vectorlike \(n\)-dimensional representations of a certain hidden gauge group. They could survive down to low energies by the global symmetries discussed later. They can communicate with the Higgs sector through the “messenger” fields, \(S\) and \(\bar{S}\). Note that the messenger in this paper does not mean the conventional messenger mediating SUSY-breaking from the hidden sector. They just play the role of connecting the Higgs and \((N, \bar{N})\) sectors. By redefining the superfields, all the parameters in Eq. (2), i.e. \(y_H, y_N, \) and \(m_S\) can be made real. We will discuss later how the effective superpotential Eq. (2) can be obtained at low energies.

As in the NMSSM, the superpotential Eq. (2) yields a quartic Higgs potential at tree-level, i.e. \(\partial W/\partial S \bar{S}\), which could raise the MSSM Higgs mass if the Yukawa coupling constant \(y_H\) was sizable. As mentioned above, however, the Landau pole constraint restricts the size of \(y_H\) to be smaller than 0.7. In this paper, we are interested in the case that \(y_H\) is small enough, \(0.2 \lesssim y_H \lesssim 0.5\), and also \(3 \lesssim \tan \beta \lesssim 10\). In this case, the tree-level quartic contribution \(\sim y_H^2 v^4 \sin^2 \beta \bar{\beta}\) in Eq. (1), which is the dominant correction in the NMSSM, becomes quite suppressed. On the other hand, \(y_N\) of order unity can avoid the Landau pole constraint, because the relatively strong hidden gauge interaction of \(N, \bar{N}\) can prevent the blowing-up of \(y_N\) at high energies. We assume that \(y_N\) is of order unity at the electroweak energy scale.

The SUSY mass parameter \(m_S\) is assumed to be quite heavier than the Higgs mass, say, \(\gtrsim 300\) GeV. The origin of \(m_S\) (and \(m_N\)) as well as the “\(\mu\)” in the MSSM can be explained e.g. by the Giudice-Masiero mechanism. Since \(y_H\) is relatively small and \(m_S\) is quite heavier than the Higgs mass, the mixing angles between the Higgs and singlet sectors are expected to be negligible. If \(y_H\) is small enough, \(m_S^2\) does not run with energy at one-loop level.

As mentioned above, \(y_N\) is of order unity. Accordingly, the soft mass squared of \(\bar{S}\), \(\tilde{m}_S^2\), can be suppressed or even become negative at low energies by the renormalization group (RG) running effect. We will assume \(|\tilde{m}_S|^2 \lesssim m_S^2\). On the other hand, the soft masses of \(N\) and \(\bar{N}\), \(\tilde{m}_N\) and \(\tilde{m}_{\bar{N}}\) can be quite heavier than other soft masses including \(\tilde{m}_S\) at low energies. Actually, it is also possible by the RG effect for \(\tilde{m}_N, \tilde{m}_{\bar{N}}\), since the superfields \(N, \bar{N}\) are nontrivial representations under a non-Abelian gauge group in a hidden sector. We will set \(\tilde{m}_N = \tilde{m}_{\bar{N}}\), which would be reasonable in large classes of supergravity (SUGRA) models, if there is no other Yukawa interaction. Just for the simplicity of our future calculations, but considering the above discussions, we suppose the following hierarchy among the mass parameters in this model:

\[
\tilde{m}_S^2 \lesssim \mu^2 \lesssim m_{3/2}^2, m_S^2 \lesssim \tilde{m}_N^2, m_{3/2}^2,
\]
where \(m_{3/2}^2\) collectively denotes other typical soft parameters. Note that even with this hierarchy, the scalar component of \(\bar{S}\), i.e., \(\bar{S}\) is still much heavier than the Higgs because of the SUSY mass \(m_S\). Since \(m_S\) and \(m_N\) both are much heavier than the Higgs mass, there is no “singlet-ino” (fermionic component of a singlet superfield) lighter than the Higgs.

The relevant scalar potential is given by
\[
V_f = \left| y_H H_u H_d + m_S \bar{S} \right|^2 + \left| m_S \bar{S} \right|^2 + \left| y_H \bar{S} \right|^2 \left( |H_u|^2 + |H_d|^2 \right),
\]
\[
V_{\text{soft}} = \tilde{m}_S^2 |S|^2 + \tilde{m}_N^2 |\bar{N}|^2
\]
\[
\approx \left( y_H m_{3/2}^2 \bar{S} H_u H_d + m_S m_{3/2}^2 \bar{S} + \text{h.c.} \right),
\]
where \(m_{3/2}^{A, B}\) indicate the soft parameters of \(A\) and \(B\) terms, respectively. Here we ignored the \(\mu\) term due to its relative smallness, and set \(\langle N \rangle = \langle S \rangle = 0\) due to their heavy masses. With the hierarchy in Eq. (3), the equations of motion for \(S, \bar{S}\) yield
\[
\langle \bar{S} \rangle \approx -\frac{y_H (m_{3/2}^A - m_{3/2}^B)}{m_S^2} \langle H_u H_d \rangle,
\]
\[
\langle S \rangle \approx -y_H \langle H_u H_d \rangle.
\]

Since the superfields \(N, \bar{N}\) couple to \(S, \bar{S}\) in Eq. (2) and its vacuum expectation value (VEV) should be related to the Higgs VEVs at the minima of \(\langle S, \bar{S} \rangle\), \(N\) and \(\bar{N}\) can get additional SUSY masses once the Higgs fields develop VEVs. The relevant diagrams are displayed in FIG. 1, where \(N, \bar{N}\) couple to the Higgs \(H_u, H_d\) through the off-shell \(S, \bar{S}\). Thus, the mass squared of \(\langle N, \bar{N} \rangle\) becomes
\[
M_N^2 = \left[ m_N - y_H y_N \langle H_u H_d \rangle \right]^2 m_S^2/\tilde{m}_N^2
\]
\[
\approx m_N^2 - 2 \left( y_H y_N \frac{m_N}{m_S} \right) \langle H_u H_d \rangle.
\]
By rolling up the \(N, \bar{N}\) lines of FIG. 1-(a), the radiative Higgs potential or mass can be generated as seen in FIG. 1-(b).
The scalar components of $N, \overline{N}$ possess, of course, the additional mass terms by SUSY-breaking effects. Due to the mass difference between the fermionic and the bosonic modes of $N, \overline{N}$, the one-loop effective Higgs potential is generated after integrating out the quantum fluctuations of $N, \overline{N}$ [10]:

$$\Delta V_{1-loop} = \frac{n}{16\pi^2} \left( \frac{M_N^2 + \tilde{m}_N^2}{\Lambda^2} \right)^2 \left( \log \left( \frac{M_N^2 + \tilde{m}_N^2}{\Lambda^2} \right) - 3 \right) - \frac{M_N^4}{2} \left( \log \left( \frac{M_N^2}{\Lambda^2} \right) - 3 \right),$$

(7)

where $\Lambda$ denotes a renormalization mass scale. This effective Higgs potential by the hidden sector fields is valid only in the energy scales below the messenger scale $m_S$: above $m_S$ scales the nonrenormalizable operator suppressed with $m_S$ in Eq. (6) cannot be regarded as a local operator any longer. With Eq. (7), the radiative mass correction can be estimated as

$$\Delta m_h^2 = \frac{n}{4\pi^2} \left( \frac{y_t m_N}{m_S} \right)^2 \left( y_t^2 v^2 \sin^2 \beta \right) \log \frac{m_N^2 + \tilde{m}_N^2}{m_N^2}. \tag{8}$$

Compared to the case by the (s)top, which is the dominant correction in the MSSM,

$$\Delta m_{h,\text{top}}^2 = \frac{3}{4\pi^2} \left( y_t m_t \right)^2 \sin^2 \beta \log \frac{m_t^2 + \tilde{m}_t^2}{m_t^2}, \tag{9}$$

where $y_t$ stands for the top quark Yukawa coupling, $m_t = y_t v_\text{ew}$ and the mixing effect by the $A$ term is ignored, we have several advantages to enhance the radiative correction to the Higgs mass in Eq. (8): we can take larger $n$ ($> 3$) and heavier $\tilde{m}_N^2$ ($> \tilde{m}_t^2$), leaving intact the naturalness issue associated with the stop mass. Moreover, we can quite efficiently raise the Higgs mass using the factor $(m_N/m_S)^2$ in Eq. (8). As a result, the Higgs mass can reach 125 GeV for $0.2 \lesssim y_t \lesssim 0.7$ and $3 \lesssim \tan \beta \lesssim 10$ with the relatively light stop ($\tilde{m}_t \approx 500$ GeV). See FIGS. 2 and 3.

In fact, large mixing in $(\tilde{t}_L, \tilde{t}_R)$ through the $A$ term, $X_t/m_t \equiv (A_t - \mu \cot \beta)/m_t \approx 2$, where $m_t \equiv \sqrt{\tilde{m}_t^2 + \tilde{m}_t^2}$, is very helpful for raising the Higgs mass with relatively light stop. Particularly, the maximal mixing $(X_t/m_t \approx \sqrt{3})$ can push the Higgs mass up to 135 GeV without any other help [2]. As the mixing deviates from the maximal mixing, however, the enhancement effect drops rapidly.
Without the mixing effect, thus, there is no way to raise the Higgs mass except for increasing the stop mass: the stop mass needed to achieve 125 GeV Higgs is above a few TeV [2]. However, such a heavy stop mass gives rise to fine-tuning of 0.1 or 0.01 percent level among the soft parameters in order to get the Z boson mass of 91 GeV. In our analysis, we turn off the mixing effect and set $\Delta m_{31}^2$ by the (s)top to be around $$(66 \text{ GeV})^2$$, which corresponds to $m_{\tilde{t}} \approx 500 \text{ GeV}$ when the two-loop corrections also included.

In the parameter ranges we have discussed so far, $\Delta m_{31}^2$ in Eq. (8) dominates over the tree-level addition, $y_H^2 v^2 \sin 2\beta$. Normally a one-loop correction is hard to win a tree-level result without extending a system. In our case, it is possible by introducing the $(N, \overline{N})$ sector, which does not participate in the tree-level addition. In our analysis, we do not include the tree-level addition for increasing the Higgs mass.

Apart from Eq. (8) we have another part containing $\log \Lambda$ in the second derivative of the potential Eq. (7), which renormalizes a soft parameter as in the radiative correction by the (s)top. In this case it does the “Bµ” parameter ($\approx m_3^2$). [If one integrates out $(N, \overline{N})$ first, this part also contributes to Eq. (5) through the induced tadpole potential for $\overline{N}$. Inserting the extremum conditions for $\overline{N}, \overline{N}$ in the potential, however, gives the same result discussed here.] Together with the RG solution of the $m_3^2$ appearing in the tree-level potential, it makes a correction to $m_3^2$:

$$
\delta m_{31}^2 \approx \frac{n}{8\pi^2} y_H y_N \frac{m_N}{m_S} \left[ m_N^2 \log \left( \frac{m_N^2 + \tilde{m}_N^2}{m_N^2} \right) + \tilde{m}_N^2 \left( \log \left( \frac{m_N^2 + \tilde{m}_N^2}{m_S^2} \right) - 1 \right) \right].
$$

(10)

The parameters considered so far let $\delta m_{31}^2$ inside the range of $$(600 \text{ GeV})^2 - (1 \text{ TeV})^2$$.

The effective superpotential Eq. (2) can be derived, e.g., by employing the following superpotential:

$$
W_{UV} = y_H S H_u H_d + y_N \overline{N} \overline{N} \overline{N} \overline{N} + \frac{f_1}{M_P} \Sigma_1^2 H_u H_d + \frac{f_2}{M_P} \Sigma_2^2 N \overline{N} + \frac{f_3}{M_P} \Sigma_3^2 \overline{N} \overline{N} + \frac{g_1}{M_P} \overline{\Sigma}_1 \overline{S} \overline{N} + \frac{g_2}{M_P} \overline{\Sigma}_2 \overline{S} \overline{N} + \frac{g_3}{M_P} \overline{\Sigma}_3 \overline{S} \overline{N}
$$

(11)

where $y_H, y_N, f_i,$ and $g_i$ ($i = 1, 2, 3$) are dimensionless couplings, and $M_P$ denotes the reduced Planck mass ($= 2.4 \times 10^{18} \text{ GeV}$). The form of the superpotential in Eq. (11) can be controlled by the two continuous global symmetries, U(1)$_{R}$ and U(1)$_{PQ}$. The U(1)$_{R}$ and U(1)$_{PQ}$ charge assignments for the superfields appearing in the superpotential Eq. (11) are presented in TABLE I.

The $\Sigma$ terms corresponding to the $g_{1,2,3}$ terms in Eq. (11), the soft mass terms, etc. in the Lagrangian admit the VEVs of $\Sigma_{1,2,3}$ and $\overline{\Sigma}_{1,2,3}$ of order $\sqrt{m_{31}^2/2M_P}$ ($\sim \nu_{10} \text{ GeV}$) [11], assuming the gravity-mediated SUSY-breaking scenario. From the $f_{1,2,3}$ terms in Eq. (11), thus, $\mu$ in the MSSM, and $m_N$ and $m_S$ in Eq. (2) are generated, which are of order $m_{31/2}$. Due to the VEVs of $\Sigma$s and $\overline{\Sigma}$s, the U(1)$_{R}$ and U(1)$_{PQ}$ are completely broken at the intermediate scale.

$\Sigma$s and $\overline{\Sigma}$s in Eq. (11) carry some accidental $Z_2$ charges. As a result, the domain wall problem would arise, if the $Z_2$s are broken after inflation. We assume that such discrete symmetries are already broken before or during inflation. If the reheating temperature is lower than $10^9 \text{ GeV}$, the $Z_2$-breaking vacuum still remains the minimum of the potential also after inflation over $\nu_{10}$.

In conclusion, we have discussed the possibility that the SUSY Higgs mass increases through the radiative correction by 1 TeV scale hidden sector fields, which can communicate with the Higgs via the messenger fields with around 300–500 GeV masses. We pointed out that even for $0.2 \lesssim y_H \lesssim 0.5$ or $3 \lesssim \tan \beta \lesssim 10$, which is the excluded region in the NMSSM, 125 GeV Higgs mass can be naturally explained in a broad parameter space with relatively light stop masses ($\sim 500 \text{ GeV}$) but without their mixing effect. The fine-tuning problem associated with the light Higgs and heavy soft scalars can be avoided, basically because the mass parameters considered for 125 GeV Higgs are all just around a few hundred GeV to 1 TeV. Consequently, the naturalness and gauge coupling unification in SUSY models can still be quite valid.

This research is supported by Basic Science Research Program through the National Research Foundation of Korea (NRF) funded by the Ministry of Education, Science and Technology (Grant No. 2010-0009021), and also by Korea Institute for Advanced Study (KIAS) grant funded by the Korean government (MEST).

TABLE I. R and Pecci-Quinn charges of the superfields. The MSSM matter superfields carry the unit $R$ charges, and also the PQ charges of 1/8. $N$ and $\overline{N}$ are assumed to be proper $n$-dimensional vectorlike representations of a hidden gauge group, under which all the MSSM fields are neutral. $\Sigma$s and $\overline{\Sigma}$s carry some $Z_2$ charges.

| Superfields | $H_u$ | $H_d$ | $N$ | $\overline{N}$ | $S$ | $\overline{S}$ | $\Sigma_1$ | $\Sigma_2$ | $\Sigma_3$ |
|------------|------|------|-----|------------|-----|---------|---------|---------|---------|
| U(1)$_{R}$ | $-\frac{1}{2}$ | $-\frac{1}{2}$ | $\frac{1}{2}$ | $-\frac{1}{2}$ | $-1$ | $-1$ | $1$ | $1$ | $2$ |
| U(1)$_{PQ}$ | $-\frac{1}{2}$ | $-\frac{1}{2}$ | $\frac{1}{2}$ | $-\frac{1}{2}$ | $-1$ | $-1$ | $1$ | $1$ | $2$ |

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