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Article

126 GeV Higgs Boson Associated with D-term Triggered Dynamical Supersymmetry Breaking

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Abstract: Continuing with our previous work on D-term triggered dynamical supersymmetry breaking, we consider a system in which our generic $\mathcal{N} = 1$ action is minimally extended to include the pair of Higgs doublet superfields charged under the overall $U(1)$ together with μ and $B\mu$ terms. The gauge group is taken to be $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)$. We point out, among other things, that the Higgs mass less than the Z-boson mass at tree level can be pushed up to be around 126 GeV by D-term contributions of the overall $U(1)$. This is readily realized by taking a $U(1)$ gauge coupling to be $\mathcal{O}(1)$.

Keywords: supersymmetry breaking; D-term; Higgs boson

The ATLAS [1] and CMS [2] collaborations have recently announced that a Higgs boson was discovered at the Large Hadron Collider (LHC), and its mass is found to be around 126 GeV. Although the observed data for a variety of Higgs boson decay modes are found to be consistent with the Standard Model (SM) expectations, the new physics beyond the SM is indispensable for explanations of various unsolved issues in the SM, such as the origin of dark matter and dark energy.

Supersymmetry (SUSY) is one of the promising candidates of new physics beyond the SM. In the minimal SUSY standard model (MSSM), it is well known that the lightest Higgs boson mass at the tree level is smaller than the Z-boson mass and can be enhanced up to 130 GeV through the quantum corrections by top and stop [3]. The observed Higgs boson mass around 126 GeV requires a heavy stop mass or a large A-term for stop, which leads to some amount of fine-tuning of the parameters, *i.e.*, the "little hierarchy problem". There are two typical extensions of the MSSM to overcome this issue. One is the next-to-MSSM (NMSSM) [4], and the other is the $U(1)$ extension of the SM gauge group [5]. In the former case, a gauge singlet chiral multiplet is introduced and coupled to the Higgs doublets in the superpotential. F-term contributions can enhance the Higgs mass after developing the vacuum expectation value (VEV) of the singlet. In the latter case, D-term contributions can enhance the Higgs mass if the Higgs doublets are charged under the extended $U(1)$ gauge group. We adopt here the latter approach, since it can be naturally incorporated in our recently proposed mechanism of D-term-triggered dynamical SUSY breaking (DDSB) [6–8], where a nonvanishing D-term VEV of the overall $U(1)$ gauge group is obtained in Hartree–Fock analysis, as is in the Bardeen–Cooper–Shrieffer (BCS) model and the Nambu–Jona–Lasino (NJL) theory [9,10]. Furthermore, our mechanism is a generalization of the Dirac gaugino scenario [11–14], which has recently been given attention as a solution to the natural SUSY breaking spectrum with a 126-GeV Higgs mass, and many pieces of work from various viewpoints have been done till now [15–41].

In this letter, we investigate the implications of the mechanism of DDSB uncovered in [6–8], coupling the system to the MSSM Higgs sector, which includes the μ and $B\mu$ terms. The pair of Higgs doublet superfields H_u, H_d is taken to be charged under the overall $U(1)$:

$$\begin{aligned}\mathcal{L}_{\text{Higgs}} &= \int d^4\theta \left[H_u^\dagger e^{-g_Y V_1 - g_2 V_2 - 2q_u g V_0} H_u + H_d^\dagger e^{g_Y V_1 - g_2 V_2 - 2q_d g V_0} H_d\right] \\ &\quad + \left[\left(\int d^2\theta \mu H_u \cdot H_d\right) - B\mu H_u \cdot H_d + \text{h.c.}\right] \qquad (1)\end{aligned}$$

We have adopted notation $X \cdot Y \equiv \epsilon_{AB} X^A Y^B = X^A Y_A = -Y \cdot X$, $\epsilon_{12} = -\epsilon_{21} = \epsilon^{21} = -\epsilon^{12} = 1$. $V_{1,2,0}$ are the vector superfields of the SM gauge group and that of the overall $U(1)$, respectively, and the corresponding gauge couplings are denoted by $g_{Y,2}$ and g, respectively. Unlike the MSSM case, the soft scalar Higgs masses $m_{H_u}^2|H_u|^2, m_{H_d}^2|H_d|^2$ are not introduced, since they are induced by D-term contributions in our framework.

1. Mechanism of D-Term-Triggered Dynamical Supersymmetry Breaking

Before going into the analysis on the objective of this letter, we summarize here the basic qualitative features of the mechanism of D-term-triggered dynamical supersymmetry breaking proposed in [6–8].

Our underlying theory before the Higgs sector is coupled is given below by Equation (2). By matching the tree parts against the one-loop part in the effective potential (which is the Hartree–Fock approximation) upon extremization with respect to the order parameter $\langle D^0 \rangle$, we obtain the gap equation (Equation (21) of [6]; Equations (4.26) and (4.28) of [8]). The gap equation is a self-consistency condition of the Hartree–Fock approximation, and finding the explicit numerical solutions to the gap equation in [6,8] demonstrates the self-consistency of the framework. Once D-term vev is generated, the equation of motion for the D-term tells us that the non-vanishing D-term vev implies the formation of the Dirac condensate, the reasoning of which is parallel to NJL theory in the auxiliary field formalism. See Equation (2.10) of [8] and the equation below to Equation (4) of [6].

There are two fundamental scales in our original theory. One is set by the mass parameter M_{prep} contained in the prepotential function $\mathcal{F}$. The other is the mass parameter M_{sup} contained in the superpotential W. The SUSY breaking scale, namely the order parameter $\langle D^0 \rangle$, is found to be given by their geometric mean; $\langle D^0 \rangle \sim M_{prep} M_{sup}$ (see Equation (3.13) of [8] for the derivation). Therefore, the SUSY breaking scale can be arbitrarily large, depending on how large these two parameters are. All of the adjoint multiplets of the standard model group appearing in our theory receive mass of order M_{sup} (in this paper, we will not really consider including the MSSM matter sector belonging to the fundamental and anti-fundamental representations of the SM group: we only include the MSSM Higgs sector.)

While the Hartree–Fock approximation exploits the one-loop effective potential in the auxiliary field formalism, the one-loop effective potential is matched to be in the same order as the tree-level potential, and this leads to the gap equation. The solution is transcendental in the Planck constant and, in this sense, is non-perturbative. (It is well-known in the NJL-type models that once the auxiliary fields are eliminated, this approximation is equivalent to the bubble summation of the fermion loops. See, for instance, [42].) Since we ignore the two-loop and higher in the auxiliary field formalism, quantum fluctuations are still assumed to be small on this new vacuum. The solution to the gap equation itself (which was obtained from the D variation of tree and one-loop effective action) demonstrates the generation of the SUSY breaking term. All of these discussions are, of course, consistent with perturbative SUSY non-renormalization theorem, which applies to the F-term alone.

As we have already mentioned, there are two scales in our theory. By making $m \ll M_{prep} = \text{cutoff scale}$, we obtain $\langle D^0 \rangle \ll (\text{cutoff scale})^2$. Our gap parameter Δ is a dimensionless parameter obtained from $\langle D^0 \rangle$ and is determined to be $\mathcal{O}(1)$ by the gap equation. We work, therefore, consistently in the weak field regime, where our effective description is valid. The SUSY breaking scale is much larger than the electroweak scale, but still much smaller than the cutoff.

2. Lagrangian and Effective Potential Extended

2.1. Lagrangian

Continuing [8], we work with the general $\mathcal{N} = 1$ supersymmetric action consisting of chiral superfield Φ^a in the adjoint representation and the vector superfield V^a that has been shown to break supersymmetry dynamically by the nonvanishing D^0-term:

$$\mathcal{L}_{\text{DDSB}} = \int d^4\theta K(\Phi^a, \bar{\Phi}^a) + (gauging) + \int d^2\theta \text{Im} \frac{1}{2} \tau_{ab}(\Phi^a) \mathcal{W}^{\alpha a} \mathcal{W}^b_\alpha + \left(\int d^2\theta W(\Phi^a) + c.c. \right) \quad (2)$$

There are three input functions: the Kähler potential $K(\Phi^a, \bar{\Phi}^a)$ with its gauging, the gauge kinetic superfields $\tau_{ab}(\Phi^a)$ that are the second derivatives of a holomorphic function $\mathcal{F}(\Phi^a)$ and a superpotential $W(\Phi^a)$.

As in [8], we make the following assumptions:

(1) Third derivatives of $\mathcal{F}(\Phi^a)$ at the scalar VEV's are non-vanishing.
(2) The superpotential at the tree-level preserves $\mathcal{N} = 1$ supersymmetry.
(3) The vacuum is taken to be in the unbroken phase of the gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)$.

The last assumption has been made for a technical reason and is not essential to the mechanism of dynamical supersymmetry breaking.

Some comments are in order for phenomenological applications of the theory based on this action. First of all, the term we have proposed $\tau_{0a}\mathcal{W}^0\mathcal{W}^a$ in our starting action Equation (2) is very similar to the super-soft operator $A\mathcal{W}'\mathcal{W}^a$, which appears in the more phenomenological operator analysis. (see, for instance, [14]). The presence of the operator of this type alone has been known to be phenomenologically dangerous, as it leads to a massless particle in the imaginary part of the complex adjoint scalars. The point we make here is, however, that our starting action Equation (2) consists not only of such a part, which contains this operator, but also of the superpotential part, which cannot be obtained just by an operator analysis. The superpotential contains the scale M_{sup}, different both from the cutoff scale M_{prep} and from the electroweak scale contained in the Higgs sector, and the resulting potential has a positive curvature everywhere near the extremum. The ordinary $U(1)$ invariance of the complex scalar field is kept intact, and unbroken $\mathcal{N} = 1$ supersymmetry of this term ensures that the tree spectrum obtained here is a massive $\mathcal{N} = 1$ supermultiplet consisting of two real spinless particles and two polarization states of spin 1/2 particles. See, for instance, [43] for illustrative explicit computation of the mass spectrum in the model of [44]. Therefore, all of the adjoint scalars in the standard model group receive masses of order M_{sup}, and the D-term supersymmetry breaking mechanism gives a boson-fermion splitting. There is no such light scalar in our theory to begin with, in contrast with the operator analysis of [14].

It is also known that the operator of the type $\tau^{ab}\tau_{ab}\mathcal{W}^0\mathcal{W}^0$ causes negative mass squared to the imaginary part of the adjoint scalars. It is clear that our action does not contain such a dangerous operator, and our theory is free from such difficulty.

To simplify the analysis in what follows, while keeping the essence, we adopt the simplest prepotential and superpotential exploited in [8] of a 5×5 complex matrix scalar superfield φ:

$$\mathcal{F} = \frac{c}{2N}\mathrm{tr}\varphi^2 + \frac{1}{3!MN}\mathrm{tr}\varphi^3 \ , \quad W = \frac{m^2}{N}\mathrm{tr}\varphi + \frac{d}{3!N}\mathrm{tr}\varphi^3 \tag{3}$$

where c is a pure imaginary number (as discussed in [8]), and m, M are mass parameters. Here, $N = 5$ and M (real number) sets the scale in the prepotential, which is the cutoff scale.

We embed the generators of the gauge group into the bases, which expand φ:

$$\varphi \equiv \begin{pmatrix} T_8 & 0 \\ 0 & T_3 \end{pmatrix} + \sqrt{\frac{3}{5}}Y \begin{pmatrix} -\frac{1}{3}\mathbf{1}_3 & 0 \\ 0 & \frac{1}{2}\mathbf{1}_2 \end{pmatrix} + \frac{\mathbf{1}_5}{\sqrt{10}}S \ , \quad T_3 = \sum_{a=1}^{3} T^a \left(\frac{\sigma^a}{2}\right) \tag{4}$$

We have represented the overall $U(1)$ and $U(1)_Y$ generators to be proportional to the unit matrix and the traceless diagonal generator, respectively. We analyze the case in which only S receives its VEV, namely the unbroken $U(5)$ vacuum of the superpotential. We will make a comment for those cases in which these do not hold, which lead to the kinetic mixing. We drop octet T_8, as it is irrelevant to the analysis below.

After a simple calculation, we obtain the non-vanishing prepotential derivatives:

$$
\begin{aligned}
\mathcal{F}_{aa} &= \frac{c}{10} + \frac{3}{3!5\sqrt{10}M}\left(\sqrt{\frac{3}{2}}Y + S\right), \ \mathcal{F}_{00} = \frac{c}{10} + \frac{3S}{3!5\sqrt{10}M}, \\
\mathcal{F}_{YY} &= \frac{c}{10} + \frac{3}{3!5\sqrt{10}M}\left(\sqrt{\frac{1}{6}}Y + S\right), \ \mathcal{F}_{a0} = \frac{3}{3!5\sqrt{10}M}T^a, \\
\mathcal{F}_{aY} &= \frac{3}{3!5\sqrt{10}M}\sqrt{\frac{3}{2}}T^a, \ \mathcal{F}_{0Y} = \frac{3}{3!5\sqrt{10}M}Y
\end{aligned} \tag{5}
$$

their VEVs:

$$
\begin{aligned}
\langle \mathcal{F}_{aa} \rangle &= \langle \mathcal{F}_{YY} \rangle = \langle \mathcal{F}_{00} \rangle = \frac{c}{10} + \frac{3}{3!5\sqrt{10}M}\langle S \rangle, \\
\langle \mathcal{F}_{a0} \rangle &= \langle \mathcal{F}_{aY} \rangle = \langle \mathcal{F}_{0Y} \rangle = 0
\end{aligned} \tag{6}
$$

and the derivatives of the superpotential:

$$
\begin{aligned}
\partial_a W &= \frac{3d}{3!5\sqrt{10}}T^a\left(\sqrt{\frac{3}{2}}Y + S\right), \\
\partial_0 W &= \frac{m^2}{\sqrt{10}} + \frac{3d}{3!10\sqrt{10}}\left(\sum_a T^aT^a + Y^2 + S^2\right), \\
\partial_Y W &= \frac{3d}{3!5\sqrt{10}}\left(\frac{3}{4}\sum_a T^aT^a + \frac{1}{4}Y^2 + \sqrt{\frac{3}{2}}SY\right)
\end{aligned} \tag{7}
$$

We choose $c = 10i$, but $\langle S \rangle$ is complex, not necessarily real.

In this letter, we add Equation (1) to Equation (2) and consider a part of $\mathcal{L}_{\text{DDSB}} + \mathcal{L}_{\text{Higgs}}$ relevant to 126 GeV Higgs:

$$
\begin{aligned}
\mathcal{L} &= \mathcal{L}_{\text{Higgs}} + \int d^2\theta \mathrm{Im}\frac{1}{2}\mathcal{F}_{ab}(\Phi^a)\mathcal{W}^{\alpha a}\mathcal{W}^b_\alpha \\
&= \mathcal{L}_{\text{Higgs}} + \frac{1}{4}\left[\int d^2\theta(\mathcal{W}^a\mathcal{W}^a + \mathcal{W}^Y\mathcal{W}^Y + \mathcal{W}^0\mathcal{W}^0) + \text{h.c.}\right] \\
&\quad + \frac{1}{4}\left[\int d^2\theta(\mathcal{F}_{aaY}Y\mathcal{W}^a\mathcal{W}^a + \mathcal{F}_{aa0}S\mathcal{W}^a\mathcal{W}^a + \mathcal{F}_{YYY}Y\mathcal{W}^Y\mathcal{W}^Y + \mathcal{F}_{YY0}S\mathcal{W}^Y\mathcal{W}^Y\right. \\
&\quad \left. + \mathcal{F}_{000}S\mathcal{W}^0\mathcal{W}^0 + \mathcal{F}_{a0a}T^a\mathcal{W}^a\mathcal{W}^0 + \mathcal{F}_{aYa}T^a\mathcal{W}^a\mathcal{W}^Y + \mathcal{F}_{0YY}Y\mathcal{W}^0\mathcal{W}^Y) + \text{h.c.}\right]
\end{aligned} \tag{8}
$$

The third prepotential derivatives, which are now real numbers, can be read off from Equation (6).

In our analysis, we take that the value of D^0 VEV is determined essentially by our Hartree–Fock approximation in [8]. This source of supersymmetry breaking is then fed to the Higgs sector, and its effects are given by a tree-level analysis. We will argue the validity of this procedure below.

Let us make the comment that the $B\mu$ term in Equation (1) is generated once we include an operator $(\mathcal{W}^0/M_{prep})^2 H_u H_d$ in the superpotential as an interaction term of our starting action. Here, M_{prep} is the scale that we have introduced in the prepotential function $\mathcal{F}$ and is regarded as a cutoff scale. It is easy to see that the $(D/M_{prep})^2 H_u H_d$ term is generated after $d^2\theta$ Grassmann integrations. Picking up the vev of D, we conclude that $m^2 H_u H_d$ with $m \sim \langle D^0 \rangle / M_{prep}$ is generated in the potential.

2.2. Higgs Potential and Variations

Let us extract the part relevant to the Higgs potential in Equation (8).

$$\begin{aligned}\mathcal{L}_{\rm pot} &= |F_{H_u}|^2 + \left(-\frac{g_Y}{2}D^Y - q_u g D^0\right)|H_u|^2 - g_2 H_u^\dagger \sum_a D^a \frac{\sigma^a}{2} H_u \\ &+|F_{H_d}|^2 + \left(\frac{g_Y}{2}D^Y - q_d g D^0\right)|H_d|^2 - g_2 H_d^\dagger \sum_a D^a \frac{\sigma^a}{2} H_d \\ &-(\mu H_u\cdot F_{H_d} + \mu F_{H_u}\cdot H_d + B\mu H_u\cdot H_d + \text{h.c.}) + \frac{1}{2}\left(\sum_a D^a D^a + (D^Y)^2 + (D^0)^2\right) \\ &+\frac{1}{2}\sum_{A,B,C=a,Y,0} \text{Im}(\mathcal{F}_{ABC}\varphi^C)D^A D^B + \Gamma^{1-\text{loop}}(D^0) \end{aligned} \tag{9}$$

where $\varphi^C = (T^a, Y, S)$. The one-loop part of the effective potential in [6,8] is denoted by $\Gamma^{1-\text{loop}}(D^0)$. Fermionic backgrounds are not needed in the potential analysis of Higgs and are not included in Equation (9).

Let us vary $\mathcal{L}_{\rm pot}$ with respect to the auxiliary fields, replacing φ^C by their VEV $\langle\varphi^C\rangle = (0, 0, \langle S\rangle)$.

$$\delta D^a: \quad 0 = (1 + \text{Im}\mathcal{F}'''\langle S\rangle)D^a - g_2 H_u^\dagger \frac{\sigma^a}{2} H_u - g_2 H_d^\dagger \frac{\sigma^a}{2} H_d \tag{10}$$

$$\delta D^Y: \quad 0 = (1 + \text{Im}\mathcal{F}'''\langle S\rangle)D^Y - \frac{g_Y}{2}|H_u|^2 + \frac{g_Y}{2}|H_d|^2 \tag{11}$$

$$\delta D^0: \quad 0 = (1 + \text{Im}\mathcal{F}'''\langle S\rangle)D^0 - q_u g|H_u|^2 - q_d g|H_d|^2 + \frac{\partial\Gamma^{1-\text{loop}}(D^0)}{\partial D^0} \tag{12}$$

Note that $\mathcal{F}_{aa0} = \mathcal{F}_{YY0} = \mathcal{F}_{000} \equiv \mathcal{F}'''$ and that Equation (12) with $q_u = q_d = 0$ is in fact the gap equation of [6,8]. Eliminating the auxiliary fields (approximately), we obtain the Higgs potential:

$$\begin{aligned} V_{\rm Higgs} &= \frac{g_2^2}{2(1+\text{Im}\mathcal{F}'''\langle S\rangle)}\left(H_u^\dagger\frac{\sigma^a}{2}H_u + H_d^\dagger\frac{\sigma^a}{2}H_d\right)^2 + \frac{g_Y^2}{8(1+\text{Im}\mathcal{F}'''\langle S\rangle)}\left(|H_u|^2 - |H_d|^2\right)^2 \\ &+\frac{1}{2(1+\text{Im}\mathcal{F}'''\langle S\rangle)}\left(q_u g|H_u|^2 + q_d g|H_d|^2 - \left.\frac{\partial\Gamma^{1-\text{loop}}(D^0)}{\partial D^0}\right|_{D^0=D^{0*}}\right)^2 \\ &+|\mu|^2(|H_u|^2 + |H_d|^2) + (B\mu H_u\cdot H_d + \text{h.c.}) \end{aligned} \tag{13}$$

Here, we have denoted by D^{0*} the solution to Equation (12), the improved gap equation. The deviation δD^{0*} of the value from D^{0*} in [8] is, in fact, small by the ratio of the electroweak scale and SUSY breaking scale. Therefore, we approximate the solution to the improved gap equation by the value of D^{0*} in [8], denoted as $\langle D^0\rangle$. Taking into account the fact that $\text{Im}\mathcal{F}'''\langle S\rangle \sim \langle S\rangle/M \ll 1$, we neglect the term $\text{Im}\mathcal{F}'''\langle S\rangle$ at the leading order. The resulting Higgs potential at the leading order is given by:

$$\begin{aligned} V_{\rm Higgs} &\simeq \frac{g_2^2}{2}\left(H_u^\dagger\frac{\sigma^a}{2}H_u + H_d^\dagger\frac{\sigma^a}{2}H_d\right)^2 + \frac{g_Y^2}{8}\left(|H_u|^2 - |H_d|^2\right)^2 \\ &+\frac{1}{2}\left(q_u g|H_u|^2 + q_d g|H_d|^2 - \langle D^0\rangle\right)^2 + |\mu|^2(|H_u|^2 + |H_d|^2) + (B\mu H_u\cdot H_d + \text{h.c.}) \\ &= \frac{g_2^2 + g_Y^2}{8}\left[|H_u^0|^2 - |H_d^0|^2\right]^2 + \frac{1}{2}\left(q_u g|H_u^0|^2 + q_d g|H_d^0|^2 - \langle D^0\rangle\right)^2 \\ &+|\mu|^2\left(|H_u^0|^2 + |H_d^0|^2\right) - \left(B\mu H_u^0 H_d^0 + \text{h.c.}\right) \\ &= \frac{g_2^2 + g_Y^2}{32}v^4 c_{2\beta}^2 + \frac{v^2}{2}\left[\mu^2 - B\mu s_{2\beta}\right] + \frac{1}{8}\left((q_u s_\beta^2 + q_d c_\beta^2)gv^2 - 2\langle D^0\rangle\right)^2 \end{aligned} \tag{14}$$

where we have restricted the potential to the CP-even neutral sector of Higgs doublets $H_u = (H_u^+, H_u^0)^T$, $H_d = (H_d^0, H_d^-)^T$ in the second line, since we are interested in the Higgs mass. In the last line, the neutral components of Higgs fields are defined as:

$$H_u^0 = \frac{1}{\sqrt{2}}\left[s_\beta(v+h) + c_\beta H + i(c_\beta A - s_\beta G^0)\right] \tag{15}$$

$$H_d^0 = \frac{1}{\sqrt{2}}\left[c_\beta(v+h) - s_\beta H + i(s_\beta A + c_\beta G^0)\right] \tag{16}$$

and we use the shorthand notations:

$$s_\beta \equiv \sin\beta, \quad c_\beta \equiv \cos\beta, \quad t_\beta \equiv \tan\beta, \quad s_{2\beta} \equiv \sin 2\beta, \quad c_{2\beta} \equiv \cos 2\beta \tag{17}$$

G^0 is the would-be Nambu–Goldstone boson eaten as the longitudinal component of the Z-boson. The VEV of Higgs field is $v \simeq 246$ GeV and $\frac{g_Y^2+g_2^2}{4}v^2 = M_Z^2$ in this convention.

3. Estimate of the Higgs Mass

We are now ready to calculate the Higgs mass. As in the MSSM, the minimization of the scalar potential $\partial V_{\text{Higgs}}/\partial v^2 = \partial V_{\text{Higgs}}/\partial\beta = 0$ allows us to express μ and $B\mu$ in terms of other parameters.

$$\mu^2 + \frac{M_Z^2}{2} = \frac{1}{2c_{2\beta}}\left((q_u s_\beta^2 + q_d c_\beta^2)gv^2 - 2\langle D^0\rangle\right)\left(q_u s_\beta^2 - q_d c_\beta^2\right) g \tag{18}$$

$$\begin{aligned} M_A^2 \equiv \frac{2B\mu}{s_{2\beta}} &= 2\mu^2 + \frac{q_u+q_d}{2}g\left((q_u s_\beta^2 + q_d c_\beta^2)qv^2 - 2\langle D^0\rangle\right) \\ &= -M_Z^2 + \frac{q_u - q_d}{2c_{2\beta}}g\left((q_u s_\beta^2 + q_d c_\beta^2)gv^2 - 2\langle D^0\rangle\right) \end{aligned} \tag{19}$$

It is straightforward to obtain the mass matrix for CP-even Higgs from the second derivative of the potential,

$$\mathcal{M}^2 = \begin{pmatrix} m_{hh}^2 & m_{hH}^2 \\ m_{hH}^2 & m_{HH}^2 \end{pmatrix} \tag{20}$$

where each component is given by:

$$m_{hh}^2 = M_Z^2 c_{2\beta}^2 + g^2 v^2\left(q_u s_\beta^2 + q_d c_\beta^2\right)^2 \tag{21}$$

$$m_{HH}^2 = M_A^2 + M_Z^2 s_{2\beta}^2 + g^2 v^2 s_{2\beta}^2\left(\frac{q_u - q_d}{2}\right)^2 \tag{22}$$

$$m_{hH}^2 = -M_Z^2 s_{2\beta}c_{2\beta} + g^2 v^2 s_{2\beta}\left(q_u s_\beta^2 + q_d c_\beta^2\right)\left(\frac{q_u - q_d}{2}\right) \tag{23}$$

The eigenvalues of this mass matrix are found as:

$$\frac{1}{2}\left[m_{hh}^2 + m_{HH}^2 \pm \sqrt{(m_{hh}^2 - m_{HH}^2)^2 + 4m_{hH}^4}\right] \tag{24}$$

and the lightest CP-even Higgs mass is:

$$m_{\text{Higgs}}^2 = \frac{1}{2}\left[m_{hh}^2 + m_{HH}^2 - \sqrt{(m_{hh}^2 - m_{HH}^2)^2 + 4m_{hH}^4}\right] \tag{25}$$

In order for the μ-term to be allowed in the superpotential, we must have a condition $e_u + e_d = 0$, which is also required from an anomaly cancellation condition for the overall $U(1)$. Then, the Higgs mass can be expressed as:

$$\begin{aligned} m^2_{\text{Higgs}} &= \frac{1}{2}\left[M_Z^2 + M_A^2 + q_u^2 g^2 v^2 - \sqrt{\left(M_A^2 - M_Z^2 c_{4\beta} - c_{4\beta} q_u^2 g^2 v^2\right)^2 + s_{4\beta}^2\left(M_Z^2 + q_u^2 g^2 v^2\right)^2}\right] \\ &= \frac{1}{2}\left[\tilde{M}_Z^2 + M_A^2 - \sqrt{\left(\tilde{M}_Z^2 + M_A^2\right)^2 - 4\tilde{M}_Z^2 M_A^2 c_{2\beta}^2}\right] \end{aligned} \tag{26}$$

where $\tilde{M}_Z^2 \equiv M_Z^2 + q_u^2 g^2 v^2$. It is interesting to see the correspondence between our expression of Higgs mass Equation (26) and that in the MSSM,

$$m^2_{\text{MSSM Higgs}} = \frac{1}{2}\left[M_Z^2 + M_A^2 - \sqrt{\left(M_Z^2 + M_A^2\right)^2 - 4M_Z^2 M_A^2 c_{2\beta}^2}\right] \tag{27}$$

As in the case of MSSM, the upper bound of Higgs mass can be obtained by taking a decoupling limit $M_A^2 \to \infty$,

$$m^2_{\text{Higgs}} \to \tilde{M}_Z^2 c_{2\beta}^2 \tag{28}$$

$\tilde{M}_Z$ can be large enough by taking $\mathcal{O}(1)$ charge and coupling $q_u g$:

$$\tilde{M}_Z \sim \sqrt{(90\text{ GeV})^2 + (246\text{ GeV})^2} \sim 262\text{ GeV} \tag{29}$$

Let us go back to the minimization conditions of Higgs potential with $q_u + q_d = 0$,

$$\mu^2 + \frac{M_Z^2}{2} = \frac{q_u g}{2c_{2\beta}}\left(-c_{2\beta} q_u g v^2 - 2\langle D^0\rangle\right) \tag{30}$$

$$M_A^2 = 2\mu^2 = -M_Z^2 - \frac{q_u g}{c_{2\beta}}\left(q_u c_{2\beta} g v^2 - 2\langle D^0\rangle\right) \tag{31}$$

which leads to:

$$M_Z^2 + M_A^2 = -\frac{q_u g}{c_{2\beta}}\left(c_{2\beta} q_u g v^2 + 2\langle D^0\rangle\right) \tag{32}$$

In order to satisfy this condition, the dominant part in the right-hand side of Equation (32) $q_u g\langle D^0\rangle/c_{2\beta}$ is required to be negative.

Using these conditions, we can eliminate M_A^2 in Higgs mass Equation (26).

$$\begin{aligned} m^2_{\text{Higgs}} &= \frac{1}{2}\left[-\frac{2q_u g}{c_{2\beta}}\langle D^0\rangle - \sqrt{\left(-\frac{2q_u g}{c_{2\beta}}\langle D^0\rangle\right)^2 + 8c_{2\beta} q_u g \tilde{M}_Z^2\langle D^0\rangle + 4c_{2\beta}^2 \tilde{M}_Z^4}\right] \\ &\simeq \tilde{M}_Z^2 c_{2\beta}^2\left(1 + \frac{c_{2\beta}\tilde{M}_Z^2}{2q_u g\langle D^0\rangle} s_{2\beta}^2\right) \end{aligned} \tag{33}$$

where the approximation $\langle D^0\rangle \gg \tilde{M}_Z^2$ is applied in the second line.

A plot for the 126-GeV Higgs mass as a function of $\tan\beta$ and $q_u g$ is shown below (Figure 1). Here, we have taken $q_u g > 0$ and $\cos 2\beta < 0$ to satisfy the condition $q_u g\langle D^0\rangle/c_{2\beta} < 0$. We can immediately see that the 126-GeV Higgs mass is realized by $\mathcal{O}(1)$ charge and coupling $q_u g$, namely without fine-tuning of parameters. Furthermore, we found that the result is insensitive to the values of D-term VEV with the assumption $\langle D^0 \rangle \gg \tilde{M}_Z^2$. This fact is naturally expected from the non-decoupling nature of Higgs mass.

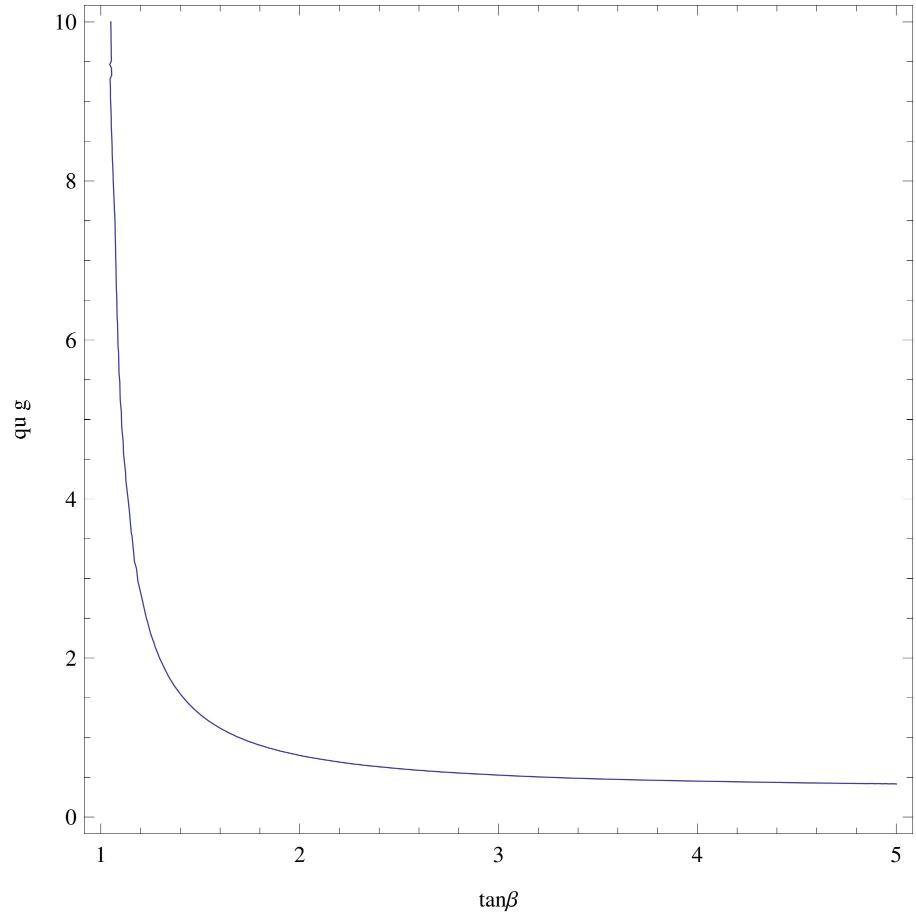

Figure 1. A plot for the 126-GeV Higgs mass as a function of $\tan\beta$ and $q_u g$. The result is insensitive to the values of the D-term VEV with the assumption $\langle D^0 \rangle \gg \tilde{M}_Z^2$.

4. Summary

In this letter, we have examined Higgs mass in the theory of D-term-triggered dynamical SUSY breaking minimally extended to couple the Higgs sector of the MSSM. Since the Higgs doublets are charged under the overall $U(1)$ in our framework, the soft scalar Higgs masses are induced by the overall $U(1)$ D-term contributions after SUSY breaking, unlike the MSSM case. These D-term contributions can enhance the Higgs mass, which is less than Z-boson mass at the tree-level in the MSSM. We have shown that the 126-GeV Higgs mass is naturally realized by taking an overall $U(1)$ gauge coupling to be $\mathcal{O}(1)$.

We give a comment on anomaly cancellation of the overall $U(1)$ gauge symmetry, which arises when the MSSM matter sector is added to our current construction. If we consider that the SM matter and Higgs are charged under the overall $U(1)$ and the superpotential allows the SM Yukawa coupling and μ-term, then the overall $U(1)$ will be found to be anomalous. To resolve this problem, we need to introduce some additional fields charged under the overall $U(1)$. Actually, we can easily confirm that all of the anomalies for $U(1)(SU(3)_C)^2, U(1)(SU(2)_L)^2$, $U(1)(U(1)_Y)^2, U(1), (U(1))^3$ are completely canceled by introducing two SM singlets with appropriate $U(1)$ charges, for instance. Yukawa coupling and the μ-term are allowed in the superpotential under the conditions:

$$q_u + e_q + e_{\bar{u}} = 0 \tag{34}$$

$$q_d + e_q + e_{\bar{d}} = 0 \tag{35}$$

$$q_d + e_l + e_{\bar{d}} = 0 \tag{36}$$

$$q_u + q_d = 0 \tag{37}$$

where $e_{q,l}$ is a $U(1)$ charge of the $SU(2)_L$ doublet quark, lepton superfields Q, L, $e_{\bar{u},\bar{d}}$ are those of the $SU(2)_L$ singlet quark superfields $\bar{U}, \bar{D}$ and $e_{u,d}$ are those of Higgs doublet superfields $H_{u,d}$. The anomaly cancellation conditions are given as:

$$U(1)(SU(3)_C)^2: \quad 2e_q + e_{\bar{u}} + e_{\bar{d}} = 0 \tag{38}$$

$$U(1)(SU(2)_L)^2: \quad 3e_q + e_l = 0 \to e_l = -3e_q \tag{39}$$

$$U(1)(U(1)_Y)^2: \quad \frac{1}{36}e_q + \frac{4}{9}e_{\bar{u}} + \frac{1}{9}e_{\bar{d}} + \frac{1}{4}e_l + e_{\bar{e}} = 0 \tag{40}$$

$$U(1): \quad 3\left[6e_q + 3(e_{\bar{u}} + e_{\bar{d}}) + 2e_l + e_{\bar{e}}\right] + 2(q_u + q_d) + \sum_i q_i = 0 \tag{41}$$

$$(U(1))^3: \quad 3\left[6e_q^3 + 3(e_{\bar{u}}^3 + e_{\bar{d}}^3) + 2e_l^3 + e_{\bar{e}}^3\right] + 2(q_u^3 + q_d^3) + \sum_i q_i^3 = 0 \tag{42}$$

where we have introduced additional singlet fields under the Standard Model gauge group and their $U(1)$ charge q_i. For instance, if we introduce two singlets with $U(1)$ charges $q_{1,2}$ satisfying $q_1 q_2 = \frac{2}{3}\left(\frac{67}{31}\right)^2 q_u^2$, we find that all of the anomalies are in fact canceled.

Some comments on an overall $U(1)$ gauge group are given. In order to realize the 126-GeV Higgs boson mass, the value of the $U(1)$ gauge coupling is required to be order one. Such a large coupling causes a Landau pole of order 1000 TeV, if $U(1)$ charged fields are the MSSM Higgs doublet superfields $H_{u,d}$, and two singlet superfields to cancel anomalies $U(1), U(1)^3$. Identifying the landau pole with the M_{prep}, we find $\langle D^0 \rangle \sim (100\ \text{TeV})^2$. The $U(1)$ gauge symmetry is broken at the weak scale by the VEV of Higgs fields, and the $U(1)$ gauge boson acquires a mass as gv. This implies an $\mathcal{O}(1)$ mass mixing with the Z boson and is not phenomenologically viable as it stands. Therefore, we need an extension of our model to avoid this issue.

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Author Contributions

Hiroshi Itoyama and Nobuhito Maru conceived and discussed ideas of this work, did calculations, and wrote the paper.

Conflicts of Interest

The authors declare no conflict of interest.

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