Probing Extra Matter in Gauge Mediation
Through the Lightest Higgs Boson Mass

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Abstract

We discuss the implications of the excesses in LHC Higgs boson searches on the
gauge mediated supersymmetric standard model, for the mass range 120 – 140 GeV.
We find that a relatively heavy lightest Higgs boson mass in this range can be
reconciled with light SUSY particles, $m_{\text{gluino}} < 2 \text{ TeV}$, if there are additional fields
which couple directly to the Higgs boson. We also find that the mass of this extra
matter can be predicted rather precisely in gauge mediation for a given Higgs boson
and gluino mass.
1 Introduction

Models with gauge mediated supersymmetry (SUSY) breaking \cite{1,2,3,4,5,6,7} are one of the most attractive and phenomenologically viable realizations of the minimal supersymmetric standard model (MSSM). One widely acknowledged prediction of gauge mediation is that the lightest Higgs boson mass is near the edge of the presently allowed mass region. This small Higgs mass is due to a nearly vanishing $A$ terms. For example, the lightest Higgs boson mass is predicted to be smaller than 120 GeV for $m_{\text{gluino}} \lesssim 2$ TeV.

The ATLAS and CMS collaborations at the LHC experiment, on the other hand, have recently reported excesses in the Higgs boson searches in the mass range $m_{h^0} = 120 - 140$ GeV with a confidence level close to $3\sigma$ \cite{9}. This appears to be troubling news for gauge mediation since a Higgs mass this size would require very heavy SUSY particles. If the lightest Higgs boson is $\sim 140$ GeV, the gluino mass of gauge mediation would be far beyond the reach of the LHC experiments, i.e. $m_{\text{gluino}} \gg 2$ TeV.

In this paper, we point out that such seeming tension between the Higgs search and the SUSY search in gauge mediation follows from the assumption that there is no additional matter beyond the MSSM. In contrast to such views, we claim that such a large Higgs boson mass suggests the existence of additional fields which couple to the Higgs boson. In fact in this scenario we will see that the lightest Higgs boson can be heavier than 140 GeV even for $m_{\text{gluino}} \lesssim 2$ TeV\footnote{The enhancement of the lightest Higgs boson mass was discussed in Refs. \cite{11,12,13}}.

The organization of the paper is as follows. In section 2, we discuss the current constraints on the mass and the coupling constants of the additional matter fields. In section 3, we discuss the effects of this additional matter on the lightest Higgs boson mass. The final section is devoted to discussions and conclusions.

2 Models with extra matter

In most gauge mediation models, the messengers are assumed to be a fundamental and anti-fundamental of the minimal grand unified gauge group, $SU(5)$. In this way, gauge mediation is consistent with grand unified theories (GUT), which seem very likely given that the gauge couplings unify in the MSSM. It is important to note that there is room
for extra low energy matter multiplets beyond the above mentioned messenger. That is, keeping perturbative unification at the GUT scale, we can still add

- up to one pair of $\bar{5} + 10$
- up to three pairs of $5 + \bar{5}$
- up to one pair of $10 + \bar{10}$

along with an arbitrary number of gauge singlets.

In the above lists, the first set of fields corresponds to a fourth-generation of matter. The CMS collaboration, however, has ruled out a fourth generation for a Higgs boson mass in the range $120 - 600$ GeV at a 95% C.L. due to the Higgs production cross section being too large. Since we are interested in the Higgs boson mass in the range of $120 - 140$ GeV, we do not pursue this possibility, but rather concentrate on the other two possibilities with vector-like multiplets (see discussion on the Higgs production cross section in the case of vector-like extra matter in the appendix A). We assume that the masses of these vector like fields are similar to the Higgsino masses, i.e. the so-called $\mu$-term, otherwise their masses should naturally be at the GUT scale.

The addition of a $10 + \bar{10}$ is particularly interesting, since the doublet Higgs $H_u$ may couple to this extra matter via $W = H_u 10 \bar{10}$. From this additional Higgs interaction, the lightest Higgs boson mass can be increased due to radiative corrections. The $5 + \bar{5}$ mentioned above can also couple to $H_u$ if we also introduce a singlet $1$, $W = H_u 5 1$. In the following, however, we mainly concentrate on the model with an extra $10 + \bar{10}$. This matter content is more advantageous since the gauge mediated contribution to the soft masses of the $10 + \bar{10}$ are much larger.

### 2.1 Constraints on extra matter

Here we concentrate on the constraints on the additional matter,

$$10 = (Q_E, \bar{U}_E, \bar{E}_E), \quad \bar{10} = (\bar{Q}_E, U_E, E_E),$$

where their Standard Model charge assignments should be clear. In terms of these fields, the Yukawa interaction $W = H_u 10 \bar{10}$ and invariant mass term $W = M_T 10 \bar{10}$ can be

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2 In this paper, we assume that the extra matter forms a complete multiplet.

3 We may allow a pair of $\bar{5} + 10$ and $5 + \bar{10}$ keeping the perturbative unification when the messenger scale is very high. This set of extra matter corresponds to a vector-like fourth generation of matter.
rewritten as,

\[ W = M_T \bar{Q}EQ_E + M_T \bar{U}EU_E + M_T \bar{E}EE_E + \lambda_u H_u Q_E \bar{U}_E , \]  

(2)

where \( \lambda_u \) denotes a coupling constant. As a result of the Higgs coupling, the masses of the additional up-type quarks are split,

\[ M_{\text{up-type}} \simeq M_T \pm \frac{\lambda_u \sin \beta v}{2} , \]  

(3)

with the Higgs vacuum expectation value \( v \simeq 174 \text{ GeV} \). The masses of the other extra matter multiplets remain degenerated and are given by \( M_T \).

As we will show in the next section, the lightest Higgs boson mass is considerably enhanced for \( \lambda_u \simeq 1 \). From this point on, we will assume that \( \tan \beta = \langle H_u \rangle / \langle H_d \rangle = O(10) \), since we are interested in models with a relatively heavy lightest Higgs boson.\(^4\)

In the superpotential in Eq. (2), we have implicitly assumed that the extra matter does not mix with the matter multiplets of the MSSM. These couplings can be forbidden by assuming some type of parity for which the extra matter is odd. If the mixing between the MSSM and the extra matter are completely forbidden, however, some of the extra charged particles can be stable, which conflicts with cosmological constraints. To avoid such difficulties, we assume that the additional matter has a small amount of mixings with the MSSM matter multiplets (\( \bar{5}_{i=1-3}, 10_{i=1-3} \)), such as

\[ W = \epsilon_i H_u 10_i \bar{10} + \epsilon_i H_d \bar{5}_i 10 . \]  

(4)

Even with this small additional mixing the extra matter in \( 10 \) can decay into the MSSM fields. The \( \bar{10} \) field can also decay into the MSSM fields through the mass terms in Eq. (2).\(^5\) Particularly, the lightest additional colored particle, which is an up-type quark, can decay into a \( W \)-boson and a down-type quark.

The most stringent collider constraint on the mass of these heavy up-type quarks is

\[ M_{\text{up-type}} > 358 \text{ GeV} . \]  

(5)

\(^4\)In our discussion, we could also include another Yukawa interaction \( W = \lambda_d H_d 10 \bar{10} \). This interaction, however, does not change our discussion even for \( \lambda_d \simeq \lambda_u \) as long as \( \tan \beta = O(10) \).

\(^5\)The amplitudes of the FCNC process such as \( K^0 - \bar{K}^0 \) mixing caused by these terms are suppressed by \( \epsilon^4 \). Thus, these terms do not cause FCNC problems for \( \epsilon \lesssim 10^{-3} \).
The above constraint is for up-type quarks decaying to a $W$-boson and a down type quark and was obtained by the CDF collaboration at the Tevatron experiments using $5.3 \text{fb}^{-1}$ of data in the channel lepton+jets [14].

Using the mass constraints found in Eq. (3), we find that collider constraints place a rough lower bound on $M_T$,

$$M_T \gtrsim 445 \text{ GeV}.$$  

As for the decays of extra scalars, they place much weaker constraints on $M_T$, since their production cross section is suppressed compared with the extra fermions [7].

The new fields which couple to the Higgs boson also give important contributions to the electroweak precision observables, $S$, $T$ and $U$ [17]. To see such contributions, we compute the $S$, $T$ and $U$ parameters at the one-loop level and find

$$S \approx \frac{2}{5\pi} \frac{(\lambda_u \sin \beta v)^2}{M_T^2},$$

$$T \approx \frac{13}{160\pi s_W^2 c_W^2} \frac{(\lambda_u \sin \beta v)^4}{M_T^2 m_Z^2},$$

$$U \approx \frac{13}{140\pi} \frac{(\lambda_u \sin \beta v)^3}{M_T^3},$$

where $s_W^2 \approx 0.231$ and $c_W^2 = 1 - s_W^2$ denote the weak mixing angle. (See also the appendix [3] and Ref. [13] for detailed analysis.) The above results neglect contributions from the additional scalar fields, since they are expected to be heavier than their fermion counterpart.

In Fig. 1, we show the shift in the $S$ and $T$-parameters due to the additional matter for $\lambda_u = 1$ [8]. In this figure, we also show the electroweak precision constraints on new physics at the 68%, 95% and 99% C.L. for $m_{h^0} = 120 \text{ GeV}$, $m_{t\text{op}} = 173 \text{ GeV}$ and $U = 0$ given in Ref. [18]. The figure shows that contributions to the $S$ and $T$ parameters from

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6 Using 37 pb$^{-1}$ data set, the ATLAS collaboration placed a constraint on a heavy quark decaying to a $W$-boson and quarks giving $M_T > 270 \text{ GeV}$ [15]. The CMS collaboration has also reported more stringent constraints on the up-type fermion masses from examining decays to a $W$-boson and bottom quark using $573 - 821 \text{ pb}^{-1}$, which gives $M_T > 450 \text{ GeV}$.

7 The collider constraint on $M_{\text{up-type}}$ can be weaken by assuming the mixing between the extra matter and the MSSM matter is very small, so that the extra colored matter would have a rather long lifetime. In this case, a smaller $M_T$ is allowed.

8 The coupling $\lambda_u = 1$ can be perturbative up to the GUT scale due to the quasi-fixed point behavior [13].
Figure 1: The contributions to the electroweak precision observables $S$ and $T$ for $\lambda_u = 1$ from the additional $10 + \bar{10}$. In this figure, we have neglected the contributions from the extra scalar fields assuming that they are heavier than about a TeV. The ellipses of 68%, 95% and 99% C.L. in Ref. [18] are also shown.

the extra matter satisfy the experimental constraints for $M_T \gtrsim 270$ GeV at 95% C.L.. Moreover, the contribution from extra matter makes the fit to the $Z$-pole data better than the Standard Model as long as

$$300 \text{ GeV} \lesssim M_T \lesssim 1 \text{ TeV}.$$  \hspace{1cm} (8)

As we will see in the next section, it is remarkable that this mass range is the preferred region for a lightest Higgs boson mass around $m_{h^0} \simeq 130 - 140$ GeV for $m_{\text{gluino}} < 2$ TeV.

3 The lightest Higgs boson mass

The scalar potential of the Higgs boson receives additional radiative corrections from the new interactions in Eq. (2). Most importantly, the effective quartic term of the Higgs boson,

$$V = \frac{\lambda_{\text{eff}}}{2} |H_u|^4 ,$$  \hspace{1cm} (9)

receives a positive contribution from its interactions with this extra matter if the extra scalar fields are heavier than their fermionic counterpart.

Calculating the one-loop four-point diagrams, we obtain the effective quartic coupling,

$$\Delta \lambda_{\text{eff}} \simeq \frac{\lambda_u^4}{16\pi^2(x_Q + 1)^2} (6(x_Q + 1)^2 \log(x_Q + 1) - x_Q(5x_Q + 4)) ,$$
Here, \( x_Q \) is defined by,

\[
x_Q = \frac{m_Q^2}{M_T^2},
\]

where we have assumed degenerate gauge mediated SUSY breaking squared masses for the up-type squarks, \( m_Q^2 \simeq m_U^2 \). We have also neglected \( A \)-term contributions to the effective quartic term, since they are generically small in models with minimal gauge mediation.

With these additional contributions to the effective quartic coupling, the lightest Higgs boson mass is,

\[
m_{h_0}^2 \simeq m_Z^2 \cos^2 2\beta + \frac{3}{4\pi^2} g_b^2 m_t^2 \sin^2 \beta \log \frac{m_t^2}{m_i^2} + \frac{4 \Delta \lambda_{\text{eff}}}{(g^2 + g'^2)} m_Z^2 \sin^2 \beta + \cdots ,
\]

where the second term comes from the one-loop top-stop contributions [19], the third term is the new contribution calculated above, and the ellipses denote the higher-loop contributions.

In Fig. 2 we show the Higgs boson mass for a given value of the gluino pole mass and \( x_Q \). In the figure, we have taken \( \lambda_u = 1 \). To obtain the mGMSB contribution to the lightest Higgs boson mass, for a given gluino mass, we used \textit{SoftSusy} [20]. The upper bounds on \( x_Q \), for each line in Fig. 2, corresponds to choosing \( M_T \simeq 450 \text{ GeV} \), which is rough the experimental lower bound.

The figure shows that the Higgs boson can be much heavier than 140 GeV even for \( m_{\text{gluino}} < 2 \text{ TeV} \). Notice that we have not included two-loop contributions from the extra matter fields in our analysis. In the case of the stop-loop contributions, for example, the two-loop contributions reduce the radiative corrections to the effective Higgs quartic term by about 30% [19]. Although the two-loop corrections to the Higgs boson mass from the extra matter goes beyond the scope of our paper, we show the lightest Higgs boson mass with the additional contribution to one-loop effective quartic term suppressed by 30% for comparison. With this suppression, we still find that the these extra matter fields can push the mass of the lightest Higgs boson above 140 GeV for \( m_{\text{gluino}} < 2 \text{ TeV} \).

\[\text{In our analysis, we have neglected the effects of } \lambda_u \text{ on the running of the soft SUSY breaking mass } m_Q \text{ and taken } m_Q = (1.3, 2.8, 4.5) \text{ TeV for } m_{\text{gluino}} = (1, 2, 3) \text{ TeV, respectively.}\]
In the right panel of Fig. 2, we also plot the required invariant mass of the extra matter as a function of the lightest Higgs boson mass. The figure shows that the mass of the extra matter can be predicted rather precisely for a given Higgs boson mass. For example, the lightest Higgs boson mass of 130 – 140 GeV requires the existence of additional matter with mass less than about a TeV for \( m_{\text{gluino}} \lesssim 2 \) TeV. It is also interesting to note that some of the parameter space has already been excluded by the LHC experiments because the lightest Higgs mass is too large i.e. 155 GeV < \( m_{h^0} \) < 190 GeV.

**4 Discussion and conclusion**

In this paper, we discussed the implications of the recently reported excesses in the Higgs boson searches at the LHC experiments in the mass range \( m_{h^0} = 120 - 140 \) GeV for models with gauge mediation. As we have seen, such a relatively heavy lightest Higgs boson can be reconciled with light SUSY particles, i.e. \( m_{\text{gluino}} < 2 \) TeV, if there is extra matter that couples to the Higgs doublet. Furthermore, due to the robust prediction of the lightest Higgs boson mass in gauge mediation models, we find that the mass of
the extra matter can be predicted rather precisely for a given a Higgs boson and gluino mass. It is also remarkable that the invariant mass of the extra matter preferred by a Higgs boson mass in the range $m_{h^0} \simeq 130 - 140$ GeV gives a better fit to the electroweak precision measurements as seen in Fig. 1.

Finally, let us comment on the effects of the $A$-term. In this paper, we have neglected the $A$-term contributions to the lightest Higgs boson mass, since the $A$-terms are predicted to be quite small in minimal gauge mediation models. As recently discussed in Ref. [21], however, the lightest Higgs boson can be significantly increased even if the extra matter is relatively heavy if large $A$-terms are generated. In fact, it is possible to generate large $A$-terms in models with gauge mediation as is discussed in Ref. [22]. In cases with large $A$-terms, the predicted mass for the extra matter in this paper will change.

Note added

Within a few days of completing this manuscript, Ref. [23] was posted on arXiv, which overlaps with some of our discussion on the lightest Higgs boson mass in gauge mediation.

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A Higgs Effective couplings to gluons

In the limit where the additional fermions are much heavier than the other energy scales in the problem, the Higgs effective couplings to gluons can be derived from the effective operator,

$$\mathcal{L}_{\text{eff}} = \frac{\alpha_s}{8\pi} \frac{1}{3} \frac{\partial}{\partial h} \log \det \mathcal{M}(h) \bigg|_{h=0} \times h G^{a}_{\mu
u} G^{a\mu\nu}. \quad (13)$$

10 To make more precise prediction, it is important to include two-loop extra matter contributions to the lightest Higgs boson mass.

11 The analysis of models with extra matter with large $A$-terms will be given elsewhere.
The effective mass matrix $M(h)$ is given by,

$$
M = \begin{pmatrix}
M_T & \lambda_u \sin \beta v \\
\lambda_d \cos \beta v & M_T
\end{pmatrix}
$$

(14)

by replacing $v \rightarrow v + h/\sqrt{2}$. Here, we have included the effects of the Yukawa coupling,

$$
W = \lambda_d H_d \bar{Q} E U E ,
$$

(15)

which was neglected in the main text. In the limit of $M_T \gg v$, we obtain,

$$
L_{\text{eff}} = \frac{\alpha_s}{8\pi} \frac{1}{3} \frac{1}{M_T^2} \frac{1}{\sqrt{2}v} \frac{h}{c^2_W} G_{\mu\nu}^a G^{a\mu\nu} .
$$

(16)

Therefore, compared with the top loop contribution to the effective gluon coupling,

$$
L_{\text{eff}}^{\text{top}} = \frac{\alpha_s}{4\pi} \frac{1}{3} \frac{1}{\sqrt{2}v} \frac{h}{c^2_W} G_{\mu\nu}^a G^{a\mu\nu} ,
$$

(17)

we find that the contribution to the Higgs effective couplings to gluons is negligible for,

$$
M_T^2 \gg \lambda_u \lambda_d \sin 2\beta v^2 .
$$

(18)

In this discussion, we neglected the extra scalar particles. The scalar contributions are, however, further suppressed by a large SUSY breaking masses, and hence, they can be safely neglected.$^{13}$

**B Precision Electroweak Parameters**

In this appendix, we discuss contributions of the extra matter to the electroweak precision observables, $S$, $T$ and $U$, which are defined by $^{17}$,

$$
\begin{align*}
\alpha S &= 4 s_W c_W \left[ \Pi''_{ZZ}(0) - \frac{c_W^2 - s_W^2}{s_W c_W} \Pi'_{Z\gamma}(0) - \Pi''_{\gamma\gamma}(0) \right], \\
\alpha T &= \frac{\Pi_{WW}(0) - \Pi_{ZZ}(0)}{M_W^2}, \\
\alpha U &= 4 s_W^2 \left[ \Pi'_{WW}(0) - s_W^2 \Pi'_{ZZ}(0) - 2 s_W c_W \Pi'_{Z\gamma}(0) - s_W^2 \Pi''_{\gamma\gamma}(0) \right].
\end{align*}
$$

(19)

$^{12}$We have taken the so called decoupling limit for the heavier Higgs bosons.

$^{13}$The scalar contributions to the Higgs effective coupling to gluons can be obtained by replacing $\det M$ to the mass matrix of scalar component, i.e. $\det M_B^2$. 

10
The dominant contributions to the electroweak precision observables, from our additional vector like matter, comes from the fermionic components of the up-type quarks whose mass matrix is given by,

\[ \mathcal{M} = \begin{pmatrix} M_T & \lambda_u \sin \beta v \\ 0 & M_T \end{pmatrix}, \]

where we have again omitted \( W = H_d \bar{10} \bar{10} \).

By computing the one-loop contributions to the vacuum polarization, we obtain the follow corrections to \( S, T \) and \( U \),

\[ S = \frac{1}{6\pi} \frac{1}{(M_H^2 - M_L^2)^2(M_H + M_L)^2} \left[ 3M_H^8 - 10M_H^7M_L - 24M_H^6M_L^2 + 54M_H^5M_L^3 - 54M_H^4M_L^5 + 24M_H^2M_L^6 + 10M_HM_L^7 - 3M_L^8 + 6M_HM_L(M_H^6 - 3M_H^4M_L^2 + 6M_H^3M_L^3 - 3M_H^2M_L^4 + M_L^6) \log \left( \frac{M_H^2}{M_L^2} \right) \right], \]

\[ T = \frac{3}{16\pi s_W^2 c_W} \frac{1}{m_Z^2(M_H^2 - M_L^2)(M_H + M_L)^2} \left[ M_H^6 - 6M_H^5M_L - 3M_H^4M_L^2 + 3M_H^2M_L^4 + 6M_HM_L^5 - M_L^6 + M_HM_L(M_H^4 + 10M_H^2M_L^2 + M_L^4) \log \left( \frac{M_H^2}{M_L^2} \right) \right], \]

\[ U = \frac{1}{6\pi} \frac{1}{m_Z^2(M_H^2 - M_L^2)(M_H + M_L)^2} \left[ -7M_H^8 - 10M_H^7M_L - 16M_H^6M_L^2 - 90M_H^5M_L^3 + 90M_H^3M_L^5 + 16M_H^2M_L^6 + 10M_HM_L^7 + 7M_L^8 + 6(M_H^8 + 4M_H^6M_L^2 + 20M_H^5M_L^3 + 10M_H^4M_L^4 + 20M_H^3M_L^5 + 4M_H^2M_L^6 + M_L^8) \log \left( \frac{M_H^2}{M_L^2} \right) \right]. \]

Here, the masses \( M_{H,L} \) are given by,

\[ M_{H,L}^2 = M_T^2 + \frac{(\lambda_u \sin \beta v)^2}{2} \pm \frac{\lambda_u \sin \beta v}{2} \sqrt{4M_T^2 + (\lambda_u \sin \beta v)^2}, \]

which satisfy,

\[ M_T^2 = M_HM_L. \]

We have also taken the phase convention \( M_T > 0, \lambda_u \sin \beta v > 0 \). By taking the limit where \((\lambda_u \sin \beta v)/M_T\) is small, we obtain the \( S, T \) and \( U \) parameters in Eqs. (7).
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