Singlet superfield extension of the minimal supersymmetric standard model with Peccei-Quinn symmetry and a light pseudoscalar Higgs boson at the LHC

John E. Kim\textsuperscript{1,3}, Hans Peter Nilles\textsuperscript{2}, Min-Seok Seo\textsuperscript{3}

\textsuperscript{1}GIST College, Gwangju Institute of Science and Technology, Gwangju 500-712, Korea, and
\textsuperscript{2}Bethe Center for Theoretical Physics and Physikalisches Institut der Universität Bonn, Nussallee 12, 53115 Bonn, Germany, and
\textsuperscript{3}Department of Physics and Astronomy and Center for Theoretical Physics, Seoul National University, Seoul 151-747, Korea

Motivated by the $\mu$-problem and the axion solution to the strong CP-problem, we extend the MSSM with one more chiral singlet field $X_{ew}$. The underlying PQ-symmetry allows only one more renormalizable term $X_{ew} H_u H_d$ in the superpotential. The spectrum of the Higgs system includes a light pseudoscalar $a_X$ (in addition to the standard CP-even Higgs boson), predominantly decaying to two photons: $a_X \to \gamma \gamma$. Both Higgs bosons might be in the range accessible to current LHC experiments.

I. INTRODUCTION

The Large Hadron Collider [LHC] at CERN in Geneva, Switzerland is the world’s highest energy accelerator designed to study the physics relevant for the origin of mass. According to the theory called the Standard Model [SM] which describes, with few exceptions, most of the observable phenomena in the universe, the Higgs field is responsible for the masses of the $W$ and $Z$ bosons, as well as that of quarks and leptons. Experiments, published so far, now only allow for the Higgs mass to be in the range between 115 and 127 GeV \cite{1,2}. The two humongous multi-purpose detectors, CMS and ATLAS, at CERN have recently presented hints towards the discovery of the elusive Higgs boson \cite{3}. Details concerning its exact mass and branching ratios will tell us whether we are dealing with the Higgs boson of the SM, or whether physics beyond the SM is required. At the moment all these possibilities are still open.

Supersymmetry is a mild extension of the SM. In its simplest form, the Minimal Supersymmetric Standard Model [MSSM], it favors a rather light CP-even Higgs boson with mass below 130 GeV \cite{4} and properties very similar to those of the SM Higgs boson. There are however two serious problems of the MSSM. These are known as the $\mu$ and strong CP problems. The $\mu$ term ($\mu H_u H_d$, where $H_u$ and $H_d$ denote the Higgs doublet superfields of the MSSM) gives mass to the Higgs doublets of order $\mu$. The problem is that this term breaks no low energy symmetry and is thus naturally of order the Planck scale, but it needs to be at the weak scale. The strong CP problem is a problem of QCD. One popular solution requires the existence of a light pseudo-scalar particle, known as the axion.

In order to address these problems simple extensions of the MSSM have been proposed. The simplest (singlet) extension of the MSSM is the NMSSM \cite{5}, motivated by questions of electroweak symmetry breakdown, the $\mu$-problem and an increase of the upper limit on mass of the lightest Higgs boson \cite{6}. Properties of the Higgs system might change drastically and could be checked by LHC experiments \cite{7}. A relation between the $\mu$ problem and the (invisible) axion solution \cite{8} to the strong CP-problem have been noticed in a particular singlet extension \cite{9} of the MSSM.

In this letter we shall discuss a simple generalization \cite{10} of this scheme (which we denote by $N_{PQ}$MSSM) with additional light supermultiplets, one of which ($X_{ew}$) is protected by the original Peccei-Quinn (PQ) symmetry \cite{11}. The PQ symmetry forbids the renormalizable terms, $H_u H_d$, $X_{ew} X_{ew}^\dagger$, and $X_{ew}^3$. The coefficients of these can be generated by breaking the PQ symmetry at a high energy scale. In particular, we note that the coefficients of $X_{ew}$ and $X_{ew}^3$ are not the PQ symmetry breaking scale, in contrast to most PQ symmetry models where they are generically of order the PQ symmetry breaking scale. The $\mu$ term can be generated by a nonrenormalizable term, which is the one used in the MSSM. With the singlet $X_{ew}$ surviving down to the TeV scale, there is an important renormalizable term $X_{ew} H_u H_d$ which distinguishes the phenomenology of $N_{PQ}$MSSM from that of the MSSM.

The main result of this letter is the observation that such a model predicts the existence of a light pseudoscalar (CP-odd) Higgs boson that could be within reach of the current LHC experiments. In the limit $\mu \sim 0$, we find that there exists a chiral symmetry which is nontrivial only for $H_u, H_d$ and $X_{ew}$. It may be called ‘Higgsino symmetry’. Because of its pseudoscalar nature, such an (axion-like) particle $a_X$ will predominantly decay to two photons, $a_X \to \gamma \gamma$, and could be easily distinguished from the CP-even Higgs boson.
The D-term potential is given as usual, and the TeV scale

\[ V' = |S_1 S_2 - F_1^2|^2 + |S_1 S_2 - F_2^2|^2 + |\tilde{m} X - \eta S_1^2|^2. \]  

(6)

Let the phases of \( S_1 S_2 \) and \( \bar{X} \) be \( \delta_s \) and \( \delta_\bar{X} \), respectively. Then, \( \delta_s = 0 \) and \( \delta_{\bar{X}} - 2\delta_s = 0 \) determine

\[ s_1 = \sqrt{\frac{m_X}{\eta}}, \quad s_2 = \frac{\eta F^2}{2m_X} \]

(7)

where \( F^2 = F_1^2 + F_2^2, s_{1,2} = |S_{1,2}| \) and \( \bar{x} = |\bar{X}| \). With \( m = O(M_P) \sim O(M_{GUT}) \) and \( \bar{x} = O(\text{TeV}) \), we obtain \( \langle S_{1,2} \rangle \) at the intermediate scale. With \( F_{1,2} \) at the intermediate scale, this scenario is realized. Here, we note that \( X_{\text{ew}} \) does not appear in Eq. (3) and survives to the electroweak scale. Integrating out \( \bar{X} \), we consider the following terms in the superpotential

\[ W_{\text{ew}} = -\mu H_u H_d - f_h H_u H_d X_{\text{ew}} \]

(8)

where

\[ f_h = -\sin \alpha + \xi \cos \alpha, \]

(9)

and soft terms of \( H_u, H_d \) and \( X_{\text{ew}} \) (with a tadpole term possibility). The reason that one PQ charge carrying singlet survives below the axion scale comes from the fact that we have one more field with charge \( Q_{\text{PQ}} = -2 \) than fields with \( Q_{\text{PQ}} = +2 \). This asymmetric appearance of the PQ fields is of general phenomena in string compactifications \[ 12 \].

The same objective can be achieved with less fields but with the nonrenormalizable term,

\[ W = -\frac{S_2^2}{M_P} H_u H_d - f_h H_u H_d X_{\text{ew}} \]

\[ + Z_1 (S_1 S_2 - F_1^2) + Z_2 (S_1 S_2 - F_2^2). \]

While there are many ways to introduce the \( N_{\text{PQ}} \) MSSM at the electroweak scale, one aspect is true for all of them: if the PQ symmetry forbids the \( H_u H_d \) term then the \( \mu H_u H_d \) must appear by breaking the PQ symmetry at a high energy scale. In addition, if a light singlet \( X_{\text{ew}} \) carrying the PQ charge \(-2\) survives down to the electroweak scale, then the only additional superpotential term is \( X_{\text{ew}} H_u H_d \), i.e. the \( X_{\text{ew}} X_{\text{ew}}^2 \) and \( X_{\text{ew}}^3 \) terms are not allowed with the minimal form of the Kähler potential.

If \( X_{\text{ew}} \) is charged under a new U(1)' gauge symmetry which is broken at high energy scale, the D-term potential gives rise to a fine-tuning problem. Therefore, we do not introduce a quartic term from the D-term and stabilize \( X_{\text{ew}} \) by the positive quadratic term \( m_{\text{ew}}^2 |X_{\text{ew}}|^2 \).

If \( f_h \) is large, the soft term of \( X_{\text{ew}} \) will be a subject of renormalization group. For \( f_h \approx 1 \) and \( A \approx 200 \text{ GeV} \), we note that \( m_{\text{ew}}^2 \) stays positive even with its vanishing GUT scale value, \( m_{\text{ew}}^2_{\text{GUT}} = 0 \).

| \( H_u \) | \( H_d \) | \( S_1 \) | \( S_2 \) | \( Z_1 \) | \( Z_2 \) | \( X \) | \( X' \) | \( \bar{X} \) |
|---|---|---|---|---|---|---|---|---|
| +1 | +1 | -1 | +1 | 0 | 0 | -2 | -2 | +2 |

\[ Q_{\text{PQ}} \]

\[ R \]

[+1 | +1 | 0 | 0 | 2 | 2 | 0 | 0 | 2]

\[ Q_{\text{PQ}} \]

\[ R \]

TABLE I: The PQ and \( R \) charges of \( H_{u,d}, S_{1,2} Z_{1,2}, X \) and \( \bar{X} \).
III. RAISING THE HIGGS MASS

With this effective superpotential and soft terms we can consider the following potential for Higgs and \( X_{ew} \):

\[
V = |\mu + f_h X_{ew}|^2(|H_u|^2 + |H_d|^2) + f_k^2 |H_u H_d|^2 - m_0^2 |H_u|^2 + m_d^2 |H_d|^2 - (B \mu H_u H_d + c.c.)
+ m_c^2 |X_{ew}|^2 - (AX_{ew} H_u H_d + c.c.) + \frac{1}{8} (g_2^2 + g_Y^2) (|H_u|^2 - |H_d|^2)^2 + \frac{g_2^2}{2} |H_u H_d|^2.
\]

(10)

Now, we can decompose neutral fields as real and complex components, \( \phi = \frac{1}{\sqrt{2}} (\phi^r + i \phi^i) \) where \( \phi = H_u, H_d, X_{ew} \).

At the vacuum, they take VEVs \( v_u, v_d, x \), respectively, and

\[
V_{\min} = \frac{1}{2} (|\mu + f_h x|^2 - m_u^2 v_u^2) + \frac{1}{2} (|\mu + f_h x|^2 + m_d^2 v_d^2)
+ \frac{f_k^2}{4} v_u^2 v_d^2 + \frac{1}{32} (g_2^2 + g_Y^2) (v_u^2 - v_d^2)^2 - B \mu v_u v_d + \frac{A}{\sqrt{2}} x v_u v_d + \frac{1}{2} m_c^2 x^2.
\]

(11)

Thus, the parameters we introduced for a fixed \( X_{ew} \) quantum number are \( x, v, \tan \beta, f_h, \mu, A, B \), \( m_u^2, m_d^2 \) and \( m_c^2 \). Three minimization conditions of \( V_{\min} \) of Eq. (11) and \( v \approx 246 \) GeV reduce the number of independent parameters to six (seven). The CP odd and even mass matrices are

\[
M_\phi^P = \begin{pmatrix}
\left( \frac{1}{\sqrt{2}} x + B \mu \right) v_u \\
\left( \frac{1}{\sqrt{2}} x + B \mu \right) v_d \\
\frac{A}{\sqrt{2}} v_u \\
\frac{A}{\sqrt{2}} v_d
\end{pmatrix}
\]

(12)

\[
M_\phi^E = \begin{pmatrix}
m_0^2 \cos^2 \theta \\
\frac{1}{2} \sin 2 \beta (f_h^2 v_u^2) \\
\frac{1}{2} \sin 2 \beta (f_h^2 v_d^2) \\
\frac{1}{2} \sin 2 \beta (f_h^2 v_d^2) \\
\frac{1}{2} \sin 2 \beta (f_h^2 v_u^2) \\
\frac{1}{2} \sin 2 \beta (f_h^2 v_d^2) \\
\frac{1}{2} \sin 2 \beta (f_h^2 v_u^2) \\
\frac{1}{2} \sin 2 \beta (f_h^2 v_d^2)
\end{pmatrix}
\]

(13)

where \( M_\phi^P = (\frac{1}{2x} (Av_u v_d - \mu f_h (v_u^2 + v_d^2))) \), \( M_\phi^E = M_\phi^P \), \( m_e^2 = f_h (\sqrt{2} \mu + f_h) \), \( m_c^2 = \frac{A v}{\sqrt{2}} \), and \( m_0^2 = \frac{\sqrt{2} A x + 2 B \mu}{\sin 2 \beta} \).

**Eigenvalues of** \( M_\phi^P \): The smallest eigenvalue of CP even Higgs mass matrix \( M_\phi^P \), Eq. (13), is mostly \( H_u \)-like. Since (33) element is inversely proportional to \( x \), the smallest eigenvalue is close to the smallest eigenvalue of the \( 2 \times 2 \) submatrix composed of (11), (13), (31), and (33) elements for large \( x \) whereas that of the \( 2 \times 2 \) submatrix composed of (11), (12), (21), and (22) elements,

\[
2m_{0x}^2 \simeq (m_0^2 + M_0^2) - (m_0^2 + M_0^2)^2 - 4m_0^2 M_0^2 \cos 2 \beta + f_k^2 v_u^2 (f_k^2 v_d^2 - 2m_0^2 - 2M_0^2) \sin^2 2 \beta)^{1/2}
\]

(14)

for small \( x \).

In Fig. 3 we show the Higgs boson masses in the \( x - f_h \) plane for \( \tan \beta = 3 \) and \( 5 \) and for \( \tan \beta = 3 \) with \( A, B \) and without \( A, B \) the radiative corrections. We used \( M_s = 1 \) TeV, \( A_t = 800 \) GeV, and \( B = 500 \) GeV, \( A/f_h = 200 \) GeV. For the quantum corrections, we consider two more parameters:

the geometric mean of the stop masses \( M_s = \sqrt{m_t, m_{\tilde{t}_2}} \) and the \( A_t \) term from the top Yukawa coupling. For \( A_t = 800 \) GeV and \( M_s = 1 \) TeV, the radiative mass shift
to the tree level mass \( m_0^2 \approx 96 \) GeV is about 30 GeV.

If we require that perturbativity holds up to the PQ scale, we need |\( f_h | \leq 0.8 \).

**Eigenvalues of** \( M_\phi^E \): One eigenvalue of \( M_\phi^E \) is 0, corresponding to the longitudinal component of \( Z \) boson. Among the two remaining eigenvalues, the smaller one is

\[
2m_{0x}^2 = (m_0^2 + M_0^2) - (m_0^2 + M_0^2)^2 - 4\mu \tilde{M}^3 \sin^2 2 \beta
\]

(15)

where \( \tilde{M}^3 = 2BM_0^2 - fhA(v_u^2 + v_d^2) \). From Eq. (15), we note that the \( \alpha_X \) mass is small for a small \( \mu \). However, the \( 1/\sin 2 \beta \) dependence is not singular for \( \sin 2 \beta \rightarrow 0 \) because the numerator cancels this divergence. In Fig. 9, the mass of \( \alpha_X \) is shown in the \( x - f_h \) plane for \( \mu = 150 \) GeV and 200 GeV, respectively, and \( A = f_h \times 200 \) GeV, \( B = 500 \) GeV and \( \tan \beta = 3 \). In this parameter range, the lightest eigenvalue of \( M_\phi^E \) is \( X_{ew} \)-like. As in the axion-photon-photon coupling case,
in general there exists an $a_X$-photon-photon coupling as shown in Fig. 2. Since the diagram occurs through the Higgsino line only, the anomaly coupling estimation is simple to give

$$L_{a_X\gamma\gamma} = \frac{\alpha_{em}}{4\pi} \frac{a_X}{x} F_{em\mu\nu} \tilde{F}_{em\mu\nu}$$

(16)

where $F_{em\mu\nu}$ is the electromagnetic field strength and $\tilde{F}_{em\mu\nu}$ is its dual. We note that the coupling is not suppressed by the axion decay constant $F$ but by the VEV of $X_{ew}$. Similar anomaly couplings to $W_{\mu\nu} W^{\mu\nu}$ and $Z_{\mu\nu} Z^{\mu\nu}$ are present with couplings proportional to $a_2$ and $a_Z$, respectively. But, $a_X hh$ and $a_X H^+ H^-$ are not present. Therefore, the production and decay of $a_X$ occur with the electroweak scale. Through the anomalous coupling (16), the electron-positron collider LEP II (with $\sqrt{s} \geq 130$ GeV) could have produced three photon events (through $e^+ e^- \rightarrow \gamma a_X$) for 125 GeV pseudoscalar at a $10^{-5} - 10^{-6}$ pb level,

$$\sigma(s) = \frac{3\alpha_{em}^3}{256\pi} \frac{1}{x^2} \left(1 - \frac{m_{a_X}^2}{2s}\right)^2 \left(1 - \frac{m_{a_X}^2}{3s} + \frac{4m_X^4}{12s^2}\right).$$

(17)

which gives only $\sim 5 \times 10^{-3}$ event for the integrated luminosity of LEP II [15].

With Eq. (16), the decay width is given by

$$\Gamma(a_X \rightarrow \gamma\gamma) = \frac{\alpha_{em}^2 m_{a_X}^3}{64\pi^2 x^2}.$$  

(18)

Since the LHC lower bound on the Higgsino mass is above 200 GeV [10], the $a_X$ decay to two photons for its mass of order 125 GeV with the insertions of $f_h$ and $1/\mu$ is negligible compared to Eq. (18). For $m_{a_X} < 2M_W$, the decay $a_X \rightarrow W +$ lepton + neutrino introduces a suppression factor $m_{lepton}^2 q_W^2 / M_W^4$. A similar remark applies to the $Z_{\mu\nu} Z^{\mu\nu}$ coupling. On the other hand, some superpartner fermion can be lighter than $a_X$.

For the $X_{ew}$-like $a_X$, therefore, decays to the $\gamma\gamma$ and a pair of lightest neutralino modes account for almost 100%. If so, some two photon events with strong dijets in the forward direction, which is the characteristic of vector boson fusion, may come from $a_X$ production and decay, showing a two-photon resonance peak at the $a_X$ mass different from 125 GeV. Such event rate is too small to be observed at the LHC, since the ratio of the $a_X$ production to that of the MSSM Higgs boson $h$ is naively estimated by the ratio of couplings, $\sim \alpha_{a_X}^2 / \alpha_{h}^3 \sim (0.0336/0.118)^4 \approx 0.65 \times 10^{-2}$ [17].

Even though the lightest eigenvalue of $M_P^2$ is $X_{ew}$-like, a small fraction of the Higgs component makes $a_X$ decay

---

1 The term suppressed by $F$ is the axion-photon-photon coupling. The chiral current of the charged Higgsino, $j_\mu^H = \overline{t}^c \gamma_\mu t H$, has the divergence $\partial_\mu j_\mu^H = (\alpha_{em}/2\pi) F_{em\mu\nu} \tilde{F}_{em}^{\mu\nu} + 2\mu H \gamma^\mu \gamma_5 H$ which gives Eq. (16).
to $b\bar{b}$ at tree level. The ratio of BRs to $a_X \rightarrow \gamma\gamma$ and to $a_X \rightarrow b\bar{b}$ in our example is given by

$$R = \frac{\alpha_{em} m_{a_X}^3}{64 \pi^2 x^2} \frac{8 \pi M_W^2}{3 C^2 m_{a_X}^3 m_{a_X}} = \frac{\alpha_{em} m_{a_X}^2}{24 \pi x^2} \frac{M_W^2}{C^2 m_{b}^2},$$

(19)

where $C = [g_2 \tan \beta \tan \gamma - (M_W/x)] \cos \gamma$ with $\tan \gamma = - (\cos \beta \cot \beta)/(|x/v|+\sqrt{2} B \mu/A v)$ in a large tan $\beta$ limit, and $\mu$ is multiplied for three colors of $b$. \textit{R} of Eq. (19) is about $(m_{a_X}/6\pi C x)^2$. In a large tan $\beta$ and a small $\mu$ limits, $C \approx -g_2 v \cos \beta/x$ which is very small.

Note in addition that the absolute production rate through Eq. (18) is relevant at this stage as far as the $b\bar{b}$ mode is swamped by the SM background. If the production through the initial $q\bar{q}$ as in Eq. (17) is estimated for 10 fb$^{-1}$ luminosity of the 7 TeV LHC, we expect roughly 10$^{-2}$ event.

\section{IV. CONCLUSION}

Experiments at the LHC will soon test the properties of the SM Higgs boson [20]. Masses and decay properties of the Higgs system will be crucial for the analysis of potential physics beyond the standard model. In this work we have considered a specific scheme motivated by supersymmetry and the strong CP-problem that predicts a pseudoscalar particle with decay $a_X \rightarrow \gamma\gamma$, that might well be within reach of current LHC experiments.

Acknowledgments: We are grateful to Bumseok Kyae, Hyun Min Lee, and especially to Stuart Raby for helpful discussions. JEk and MSS are supported in part by the National Research Foundation (NRF) grant funded by the Korean Government (MEST) (No. 2005-0093841), and HPN is partially supported by the SFB-Transregio TR33 “The Dark Universe (Deutsche Forschungsgemeinschaft) and the European Union 7th network program “Unification in the LHC era (PITN-GA-2009-237920).