Effects of dynamical breaking in SUSY SU(5) to two standard models

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Supersymmetric (SUSY) models and dynamical breaking of symmetries have been used to explain hierarchies of mass scales. We find that a chiral representation, $\overline{10} \otimes \begin{pmatrix} 5 \\ \phi \end{pmatrix} \otimes 2 \cdot 5$ in SUSY SU(5) in the hidden sector, breaks global SUSY dynamically, by producing a composite field $\phi$ below the SU(5) confinement scale. Starting with this dynamically generated scale, we attempt to explain two vastly different mass scales in two standard models of physics, one in particle physics and the other in cosmology. Gravitational effects transmit this dynamical breaking to the SM superpartners and the quintessential vacuum energy. The SM superpartners feel the effects just by the magnitude of the gravitino mass while the smallness of the quintessential vacuum energy is due to the composite nature of $\phi$. The composite $\phi$ has a global symmetry which is hardly broken in SUSY and hence its phase can be used toward a quintessential axion for dark energy of the Universe.

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I. INTRODUCTION

Mass scales are the most important physical parameters in all physics disciplines. In any theory, the definition of mass scale is given by the Planck mass $M_P$. The next important mass scale determines a physics discipline. It is a mystery how this next scale has arisen from the fundamental Planck mass. In the standard model (SM) of particle physics, the next scale is the electroweak scale $v_{ew} = 246$ GeV. In the standard Big Bang cosmology (SBBC), the next scale is the current value of the Hubble parameter $H_0$. In the current dark energy (DE) dominated Universe, DE can be considered as the next scale, which we follow here. In nuclear physics, atomic physics, and condensed matter physics, the next scales are 7 MeV, 1 eV, $10^{-3}$ eV, respectively, which can be derived in principle from the electroweak scale. Cosmological dark matter has its root in particle species [1] and DE looks for solutions in the framework of general relativity [2]. Therefore, understanding mass parameters in two standard models, in the SM and in the SBBC, is the key in understanding all physical mass parameters.

Gauge forces in the SM can be unified into a grand unified theory (GUT) at the scale $10^{15-17}$ GeV [3]. Understanding the ratio of the electroweak scale and the GUT scale is known as the gauge hierarchy problem [4]. If this hierarchy of masses is not fine-tuned, the best ideas so far suggested are behind two pillars: dynamical symmetry breaking (DSB) [5] and breaking of $N=1$ supersymmetry (SUSY) [6]. In particular, the DSB idea is most natural in the sense that an exponentially small mass parameter can be obtained by the evolution of gauge couplings of the confining force. If the confinement is working for generating small mass scales, then it can generate even very small mass parameter such as the DE scale of $10^{-3}$ eV because a composite particle with a large engineering dimension may be the source of this DE scale.

The $N=1$ SUSY idea accompanies superpartners of the SM particles around the TeV scale. Once the small electroweak scale is introduced as a parameter at the GUT scale then it can be used down at the electroweak scale via the non-renormalization theorem but it lacks in explaining the exponentially small scale itself. Currently, TeV scale SUSY particle effects have not shown up at the Large Hadron Collider, and we need to raise the superparticle masses above a few TeV [7][8]. By raising the superpartner masses above the TeV scale, there must be a kind of small fine-tuning of parameters, the so-called the little hierarchy problem (a problem on the SUSY particle masses and the electroweak scale), which can be understood now in various methods [9].

Since SUSY must be broken, generating a SUSY breaking scale is the key in SUSY solutions of gauge hierarchy problem. In this regard, a DSB in SUSY models is promising toward understanding hierarchical mass scales. Phenomenologically, the SUSY idea must relate the electroweak scale to the SUSY breaking scale, which was popular in the early 1980’s under the name of gravity mediation [10][11]. Supersymmetry breaking by gluino condensation needs the gluino condensation scale of order $10^{13}$ GeV [12] to obtain TeV scale gravitino mass. Since it is required to raise the gravitino mass above several TeV [7][8], the gluino condensation idea may not work if the hidden sector and the visible sector gauge couplings are unified at the GUT scale. On the other hand, if an F-term of chiral field is the source for SUSY breaking, then the intermediate mass scale about $5 \times 10^{10}$ GeV is needed for gravitino mass of order the TeV scale [13]. In addition, this intermediate scale has been useful for the “invisible” axion and the $\mu$-term in supergravity [14][15].
We find that a chiral representation, $\mathbf{10} \oplus \mathbf{5} \oplus 2 \cdot \mathbf{5}$ in SUSY SU(5) in the hidden sector, breaks global SUSY dynamically, by producing a composite field $\phi$ below the SU(5) confinement scale $\Lambda$. The composite field $\phi$ respects a global symmetry $U(1)$, which is shown by matching anomalies above and below the confinement scale. The F-term scale of $\phi$ can be larger than $5 \times 10^{10}$ GeV, raising the masses of SM superpartners much above TeV. Since $\phi$ is composite, breaking the resulting global $U(1)$ of $\phi$ can be very tiny, which can be the source of DE scale in the Universe [1].

II. DYNAMICAL BREAKING IN A SUSY GUT

Early ideas on supersymmetric QCD(SQCD) in SU($N_c$) gauge group [16–18] are DSB models, which focussed on the global symmetry structure below the confinement scale. Here, $N_f$ number of vector-like quarks, i.e. $N_f$ L-chirality representations of $N_c$ and $N_f$ L-chirality representations of $N_c$ are introduced, and the global symmetry can include a big size non-Abelian family group such as SU($N_f$)$_L$ × SU($N_f$)$_R$. The results are summarized in [16–20]. Our model here does not belong to SQCD but to a kind of a chiral-family model in SUSY grand unification (SGUT). The first study on DSB in SGUT was presented in [19], where one chiral family in the SU(5) SGUT was suggested to induce SUSY breaking based on the flat direction argument. In the one-family SGUT SU(5), however, a superpotential cannot be written and only an argument toward SUSY breaking has been presented.

In this paper, we find a chiral representation, $\mathbf{10} \oplus \mathbf{5} \oplus 2 \cdot \mathbf{5}$, in SU(5) SGUT with superpotentials given above and below the confinement scale and attribute the superpotential source of SUSY breaking in supersymmetric SMs (SSMs) to the confinement scale $\Lambda$ of this hidden sector SU(5). For a non-Abelian gauge symmetry SU(5) and a non-Abelian global symmetry SU(2), the representations are written as, (SU(5)$_{\text{gauge}}$, SU(2)$_{\text{global}}$), and we introduce

$$\bar{\Psi} = (\mathbf{10}, 1), \bar{\psi}_1 = (\mathbf{5}, 1), \psi_2 = (\mathbf{5}, 2), D_1 \sim (1, 2).$$ (1)

We are guided to study Eq. (1) from our earlier work [21]. These fields, being complex, can have phases. So, the full global symmetry we consider is SU(2)$_L$ × U(1)$_\psi$ × U(1)$_{\bar{\psi}}$ × U(1)$_{\psi_2}$ × U(1)$_{D_1}$. Note that our global symmetry is complete.

There are three SU(2) invariant terms in the superpotential,

$$W_0 \equiv \frac{1}{4} \bar{\psi}^{\alpha \beta} \psi^{\dagger \beta} \psi^{\dagger \alpha} \epsilon_{ij}, \quad \bar{\psi}^{\alpha \beta} \psi^{\dagger \alpha} \psi^{\dagger \beta}, \quad \frac{1}{3!} \bar{\Psi}^{\alpha \beta \gamma \delta \epsilon} \psi^{\dagger \alpha} \psi^{\dagger \beta} \psi^{\dagger \gamma} \psi^{\dagger \delta} \epsilon_{\alpha \beta \gamma \delta \epsilon},$$ (2)

which act as three constraints on the U(1) phases,

$$U(1)_\psi + 2U(1)_{\psi_2} = 0, \quad U(1)_{\bar{\psi}_1} + U(1)_{\psi_2} + U(1)_{D_1} = 0, \quad 2U(1)_{\bar{\psi}_1} + U(1)_{\psi_1} = 0.$$ (3)

Therefore, there remains only one U(1) global symmetry, say U(1)$_\psi$, and the other phases can be written as, $U(1)_\psi = -2U(1)_{\bar{\psi}_1} + U(1)_{\psi_2} = -\frac{1}{2} U(1)_\psi$, and $U(1)_{D_1} = \frac{3}{2} U(1)_\psi$. Global symmetries survive below the confinement scale, where only SU(5) singlets will be considered. If the confinement preserves SUSY and the confining SU(5)$_{\text{gauge}}$ is not broken, we can consider the following SU(5)$_{\text{gauge}}$–singlet chiral fields,

$$\phi = \frac{1}{3!} \bar{\Phi}^{\alpha \beta \gamma \delta \epsilon} \Phi^{\dagger \alpha} \Phi^{\dagger \beta} \Phi^{\dagger \gamma} \Phi^{\dagger \delta} \epsilon_{\alpha \beta \gamma \delta \epsilon}, \quad \Phi_i = \bar{\psi}^{\dagger \alpha} \psi_{2 \alpha}.$$ (4)

In the confinement process, the anomaly matching of global symmetries is required [22]. Since we have the global symmetry SU(2) × U(1)$_\psi$, it is a question how an anomalous U(1)$_\psi$ survives the confinement process. Not to worry about the perturbative effects of fermion representations, let us gauge the SU(2) part. For the global U(1), consider the chiral transformations on fermions by an angle $\theta$ of U(1). If there is anomalies above the confinement scale, then due to the non-perturbative effects through instantons there is an effective $\theta$ term for the non-Abelian SU(2). Use this chiral transformation of U(1) such that quark phases become 0. Below the confinement scale, the phase of composite fermion due to perturbative effects is then 0. But the anomalous $\theta$ term, generated non-perturbatively, survives the confinement process as discussed in axion physics. We will consider a global symmetry in the end, which will corresponds to an

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1 Two family SU(5) Georgi-Glashow model [8], 2(10 $\oplus$ 5) in Ref. [20], is different from Eq. (1).

2 Here, we express charges as U(1)’s.
above and below the confinement scale. But we need not consider $U(1)$ because there is no gauge symmetry we can consider below the confinement scale.

considered only matter fields. Considering all fermions, the anomaly free combination $U(1)_{AF}$ is also anomaly-free if we consider only matter fields. Considering all fermions, the anomaly free combination $U(1)_{AF}$ is $U_{AF} = U(1)_R + U(1)_A$, which is listed in Table I $U(1)_{AF}$ is a kind of R symmetry. For composites $\phi, \Phi_i$ and $S$, $U(1)_{AF}$ are also listed in Table I $U(1)_{AF}$, being $SU(5)_{gauge}$-anomaly free, there is no constraint for the sum of $U(1)_{AF}$ of composites to satisfy because there is no gauge symmetry we can consider below the confinement scale.

Let us check the matching of anomalies $U(1)_{global} \times SU(2)_{global} \times SU(2)_{global}$ above and below the confinement scale. For this, the $U(1)_A$ column of Table I is relevant. Above the confinement scale, we consider doublets of global $SU(2)$ only in $(5, 2)$ from which we obtain $U(1)_A - SU(2)_{global} - SU(2)_{global}$ anomaly of $(+1) \times 5 = 5$ units. Below the confinement scale, we consider only $\Phi_i$ which carries $U(1)_A - SU(2)_{global} - SU(2)_{global}$ anomaly of 5 units. So, they perfectly match. For the $SU(2)$ discrete anomaly, we have six (equal number) copies (together with $D_1$) of doublets above the confinement scale and two (equal number) copies (together with $D_1$) of doublets below the confinement scale. Therefore, the Witten anomaly is also matched.

Therefore, we can consider the following superpotential below the confinement scale, consistently with our global

\begin{table}[h]
\centering
\begin{tabular}{|c|c|c|c|c|c|c|c|c|c|c|}
\hline
 & $2(R)$ & $SU(2)$ & $U(1)_\psi$ & $U(1)_{\bar{\psi}}$ & $U(1)_{D_1}$ & $U(1)_A$ & $U(1)_R$ & $U(1)_{AF}$ & Dimension \\
\hline
$\vartheta$ & 0 & 0 & 0 & 0 & 0 & 0 & +1 & +1 & $\frac{1}{2}$ \\
$\bar{\psi} \sim (\Omega, 1)$ & -3 & 1 & +1 & 0 & 0 & 0 & 0 & 0 & 0 & - \\
fermion & -1 & 1 & 0 & +1 & 0 & 0 & +2 & -1 & +1 & - \\
$\bar{\psi}_i \sim (5, 1)$ & - & 2 & 0 & 0 & +1 & 0 & +1 & +1 & +1 & 1 \\
fermion & +1 & 2 & 0 & 0 & +1 & 0 & +1 & +1 & +1 & - \\
$\psi_2 \sim (5, 2)$ & -2 & 0 & 0 & +1 & 0 & +1 & +1 & +1 & +1 & - \\
fermion & +1 & 2 & 0 & 0 & +1 & 0 & +1 & +1 & +1 & - \\
$D \sim (1, 2)$ & -2 & 0 & 0 & +1 & 0 & +1 & +1 & +1 & +1 & - \\
fermion & +1 & 2 & 0 & 0 & +1 & 0 & +1 & +1 & +1 & - \\
$W^n \sim \chi^n$ & $2N_c$ & - & 0 & 0 & 0 & 0 & 0 & +1 & +1 & $\frac{3}{2}$ \\
$\Lambda^b$ & - & - & - & - & - & - & - & $\frac{2b}{3}$ & - & b \\
$\phi$ & - & 1 & - & - & - & - & +2 & +2 & +4 & 1 \\
fermion & - & 1 & - & - & - & - & +2 & +1 & +3 & - \\
$\Phi_i$ & - & 2 & - & - & - & - & +5 & $\frac{1}{2}$ & $\frac{11}{2}$ & 1 \\
fermion & - & 2 & - & - & - & - & +5 & $\frac{1}{2}$ & $\frac{11}{2}$ & 1 \\
$S$ & - & 1 & 0 & 0 & 0 & 0 & 0 & +2 & +2 & 1 \\
fermion & - & 1 & 0 & 0 & 0 & 0 & 0 & +1 & +1 & - \\
$D' \sim (1, 2)$ & -2 & - & - & - & +1 & - & +1 & +1 & +1 & - \\
fermion & - & 2 & - & - & - & +1 & - & +1 & +1 & - \\
\hline
\end{tabular}
\caption{Global symmetries with $N_c = 5$.}
\end{table}
In this case, the vacuum energy
\[ W = M^2 \phi + \frac{N_c(N_c^2 - 1)}{32\pi^2} \mu_0^2 S \left( 1 - a \log \frac{\Lambda^3}{S\mu_0} \right) + \cdots \]
(7)

where \( \Lambda \) is the definition for the gaugino condensation scale leading to the Veneziano-Yankielowicz composite \( S \) [24]. [Note that ours as an effective one still respects the R symmetry in contrast to the phenomenological Polonyi potential \( W = m^2 z + \text{constant} \) [25].] Considering the R charges of Table III we note that \( M^4 \Phi_i D_i \), allowed by SU(2) x U(1), is forbidden, \textit{i.e.} it is not allowed by SUSY. So, the SUSY conditions for the composites are

\[ \frac{\partial W}{\partial \phi} = 0 : M^2 = 0, \]
\[ \frac{\partial W}{\partial S} = 0 : \mu_0^2 \left( 1 + a - a \log \frac{\Lambda^3}{S\mu_0} \right) = 0, \]
(8)

where the first equation cannot be satisfied. The \( D_{11} \) related ones are irrelevant in the study of confinement. One can add one more pair of doublets \( D_2 \) and \( D_3 \) to break the global symmetry SU(2) x U(1) to a global U(1). At this stage, it is difficult to break the remaining U(1) by another term in the superpotential. Let the remaining global symmetry be SU(2) x U(1) and \( S \). The VEV \( \langle S \rangle \) is determined from the SUSY condition at \( \langle \Lambda^3 / \mu_0^2 \rangle e^{-(1+a-1)} \). Equations in [8] show that the SUSY breaking is not by \( \langle S \rangle \) but through the \( M^2 \) term due to the O’Raifeartaigh mechanism [20].

We view the failure \( M^2 = 0 \) in Eq. [8] as an invalid assumption in our discussion, \textit{i.e.} either on SUSY or on SU(5)\textsubscript{gauge}. We can argue that the SUSY breaking is favored because of the following reason. Condensate \( \Phi_i \), transforming nontrivially under SU(5)\textsubscript{gauge}, will be quickly dissolved via gluon scattering. On the other hand, the SU(5)\textsubscript{gauge}–singlet condensate \( \phi \) is continuously formed until the confinement process ends. To break SU(5)\textsubscript{gauge}, we must consider the VEV of the condensate \( \Phi_i \) which does not stay long when SU(5) gluons are active. Hence, instead of considering \( \langle \Phi_i \rangle \) which breaks SU(5)\textsubscript{gauge}, we better consider the superpotential [7]. So, we conclude that SUSY is broken by the O’Raifeartaigh mechanism.

One of the popular scenarios for SUSY breaking is the gaugino condensation triggering the \( F \)-term(s) of singlet chiral field(s) \( z \). The gauge kinetic function \( f(z) \) in \( f d^2 \theta f(z) W^\alpha W_\alpha \) can take a form \( f(z) = 1 + (z/M) \) in this case [12] and \( z \) must carry \( R \) charge 0. From matter fields of \( E_8 \times E_8 \), we encountered that there does not appear a singlet field through compactification with all U(1) charges being zero [21, 27]. Therefore, a candidate can be only from the antisymmetric tensor field \( B_{MN} \), which was suggested in [28, 29]. It is commented that the only possible bilinear of Majorana-Weyl gluinos in 10 dimensions is \( \text{Tr} \chi M_{NP} \chi (M, N, P = 1, \cdots, 10) \) and hence the coupling can be \( B^{MNP} \text{Tr} \chi M_{NP} \chi [29] \). One needs a large value, \( \gtrsim 10^{13} \) GeV, for gluino condensation mass of gluino of dimension \( \frac{3}{2} \). With our \( \phi \) of dimension 1, the same amount of SM SUSY breaking is obtained by our \( M \) at \( 5 \times 10^{10} \) GeV. Therefore, if SUSY breaking between \( 10^{13} \) GeV and \( 5 \times 10^{10} \) GeV is needed then our mechanism is useful.

Since \( V \) is bounded by the SUGRA correction, \( \phi \) get a VEV of order \( M_P \). So the SUGRA corrections are not negligible. So the Polonyi problem (=moduli problem) is there as is well known.

### III. Effects of Supergravity

The SM particles acquire the SUSY breaking effects from the hidden sector by supergravity effects, proportional to the gravitino mass [30]. The scalar partners of quarks and leptons, and the fermionic partner of gauge bosons acquire masses at the order the gravitino mass \( m_3^2 \approx |F|^2 / 3M_P^2 \) where \( F \) is the SUSY breaking \( F \)-term [13], and for the minimal Kahler potential \( K = |\phi|^2 + \cdots \)

\[ F = DW = (\partial_\phi W) + (\partial_\phi K)W/M_P^2 = M^2 + \phi^* M^2 (\phi + N_c \Lambda^3 e^{i\alpha}/M^2)/M_P. \]
(9)

In this case, the vacuum energy \( V \) is [30],

\[ \frac{V}{\exp(|\phi|^2/M_P^2)} = |F|^2 - 3|W|^2/M_P^4 \]
\[ = M^4 \left( 1 + \frac{2|\phi|^2}{M_P^2} + \frac{|\phi|^4}{M_P^4} \right) + \frac{\Lambda^6 N_c^2 |\phi|^2}{M_P^2} + \frac{M^2 \Lambda^3 N_c (\phi e^{-i\alpha} + \phi^* e^{i\alpha})}{M_P^2} - \frac{3}{M_P^2} |M^2 \phi + N_c \Lambda^3 e^{i\alpha}|^2, \]
(10)
where we considered only vacuum values of $M^2$ and $S$ terms in $W$ and $\alpha = 0$ or $\pi$ for the gluino condensation to break $U(1)_R$ to $\Z_2$. The dynamically generated superpotential of gluino condensation corresponds to the $N_f = 0$ case in SQCD, i.e. the Affleck-Dine-Seiberg superpotential $W = N_c \Lambda^3$ [13] where $\Lambda^3$ carries two units of $R$ charge. Let us define the $R = 2$ chiral field $S$ is normalized to the gluino condensation as

$$\langle \lambda^a \lambda^a \rangle = \frac{-32 \pi^2 \Lambda^3}{\mu^2_0} \equiv -(N_c^2 - 1) \mu_0^2 S,$$

$$|\Lambda| = \left(\frac{(N_c^2 - 1)}{32 \pi^2}\right)^{1/3} \mu_0^{2/3} |S|^{1/3},$$

$$W = N_c \Lambda^3 = \frac{N_c (N_c^2 - 1)}{32 \pi^2} \mu_0^2 S.$$

(11)

Let us determine $\langle \phi \rangle \equiv \nu e^{i \delta}$ such that $V = 0$ for two vacua of gluino condensates. To illustrate solutions, let us keep the leading terms in the $(1/M^2)$ expansion to obtain the minimization condition

$$M^4 v^3 - 4M^2 \Lambda^3 N_c \cos(\delta - \alpha) v^2 - 3N_c^2 \Lambda^6 v - 2M_P^2 M^2 \Lambda^3 N_c \cos(\delta - \alpha) = 0$$

(12)

which has solutions. Depending on the range of $M^2$, it has no, one, two or three solutions. The magnitude of $M^2$ arises from the original superpotential [3],

$$\frac{\lambda_0}{3!} \bar{\Psi}^\alpha \bar{\Psi}^\gamma \phi \bar{\Psi}^\beta \epsilon_{\alpha \beta \gamma \delta} \rightarrow \lambda_0 \mu_0^2 \phi,$$

(13)

leading to $M^2 = \lambda_0 \mu_0^2$. Even though we keep all terms in the $1/M^2$ expansion, we expect that there exist solutions for $V = 0$ for a reasonably large value of the original $\lambda_0$. In string compactification, $\lambda_0$ is calculable [31]. In Fig. 1, we sketch a rough idea of the vacuum energy dependence on the composite $\phi$. The curvature at the $v$ direction is known as the Polonyi problem [32]. In our case the maximum point in the center is about $\Lambda^4$ and $f_{\text{quintessence}}$ is about
TABLE II: U(1) charges of SU(5)' fields. Here, U(1)$_R$ is defined by $Q_R = \frac{1}{2}(Q_1 + Q_2 + Q_3) + \frac{1}{3}Q_6 + 2Q_{20}$. 

\[ \lesssim M_P, \text{ and a rough estimate of the mass of real } \phi \text{ is } \gtrsim (10^{12} \text{ GeV})^2/M_P \approx 10^4 \text{ GeV}. \text{ So, the Polonyi problem can be evaded for a little bit larger value of } A \text{ compared to } \mu_0. \]

The complex $\phi$ accompanies a global U(1) symmetry, as commented below Eq. (5), which is very difficult to be broken because of $\phi$ being composite. Determination of $\langle \phi \rangle$ from Eqs. (10) and (12) is a fine-tuning.

### IV. QUINTESSENTIAL AXION AS DARK ENERGY

In addition, we need a mechanism to break the global U(1) explicitly because gravity does not allow an exact global symmetry. If the magnitude of breaking is tiny, $\sim 10^{-3} \text{ eV}$, and the radius of the circle in the inset of Fig. 1 is large, $\lesssim M_P$, it is the quintessential axion. Since our global symmetry has originated from the anomaly, one can consider the QCD anomaly first. But it is too big for the cosmological constant. Anyway, it must be reserved for the invisible axion. The next obvious contribution is the SU(2)$_W$ anomaly which is sufficiently small,

$$\sim e^{-2\pi/\alpha_2} v_{ew}^4 \sim 10^{-81.6} v_{ew}^4 \sim 10^{-5} (10^{-3} \text{ eV})^4.$$  

But, it is too small for dark energy of a quinessential axion. We note, however that, there can be a larger contribution as discussed in Ref. [37]. In Fig. 1 we depict the vacuum energy as a function of $\phi$. The $Z_2$ vacua from the gaugino condensation are denoted as the black and green bullets, which are determined by fine-tuning as commented below Eq. (12). In the inset, the quintessential field is along the red circle with radius $f_{\text{quintessence}}$ and the star describes the current magnitude of dark energy. The height at the star is small because of the composite nature of our $\phi$ and $f_{\text{quintessence}}$ can be large, depending on models. Thus, determination of $\langle \phi \rangle$ depicted in the inset with the observed dark energy needs a knowledge of more couplings in the superpotential, whose discussion is outside the scope of this paper.

This interesting spectrum of the hidden sector particles was obtained before from string compactification [21], where the visible sector flipped SU(5) model with three families are realized. Recently, a successful $Z_{4\text{LR}}$ symmetries needed for proton longevity [40] has been assigned successfully in this model [42].

Finally, we list the SGUT non-trivial spectra in Table II obtained from $Z_{12-I}$ orbifold compactification [21], where many vector-like pairs of $(5 \oplus \overline{5})$ of [21] are not listed. From this Table, we can remove the vector-like representation $F_3' \oplus F_4'$ just below the string scale by a superpotential term $\sim F_3' F_4'$. Then, we obtain the spectra presented in Eq. (1).

### V. CONCLUSION

We showed that a chiral representation, $\overline{10} \oplus \overline{5} \oplus 2 \cdot \overline{5}$ in supersymmetric SU(5), breaks global supersymmetry dynamically. This mechanism of supersymmetry breaking can generate smaller mass scales in physics if appropriate...
discrete symmetries are provided [51]. We obtained this needed spectra from the hidden sector of heterotic string as shown in a $\mathbb{Z}_{12−1}$ orbifold compactification. There can be many useful applications of this dynamical breaking of supersymmetry. Firstly, the little hierarchy of factor $r$ in particle physics, i.e. supersymmetry appearing above $r$ TeV, is obtained if SU(5) confines at $\gtrsim \sqrt{7} \cdot 5 \cdot 10^{10}$ GeV. Second, dark energy in the Universe can be attributed to the composite nature of $\phi$ arising from the SU(5) confinement.

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