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ABSTRACT

In locally supersymmetric grand unified models we show rather a model independent upper bound 3×10^{11} GeV for the scale of supersymmetry breaking, which is derived by considering $SU(2) \times U(1)$ breaking at electro-weak mass scale. This bound necessarily implies the existence of new particles (superpartners) below 10^4 GeV.

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If supersymmetry (SUSY) is relevant in particle physics, it is known to be broken at the scale $m_S \gtrsim 10^{11}$ GeV (the Weinberg bound [1]) or at $m_S \lesssim 10^6$ GeV (the Pagels-Primack bound [2]). In this paper, we would like to point out that there should be an upper bound of 3×10^{11} GeV order on the Weinberg region in general local SUSY gauge theories by considering $SU(2) \times U(1)$ breaking. This will lead to a phenomenologically important conclusion that superpartners of gauge bosons, quarks and leptons are lurking around TeV region, but are not much heavier than that.

SUSY has been considered in last several years [3] for a solution to the infamous gauge hierarchy problem.[4] Even though there does not exist a generally accepted explanation of this problem at present, we take the following assumptions in this paper toward a solution of the gauge hierarchy problem:

(A) Low energy electro-weak gauge group $SU(2) \times U(1)$ is broken by two elementary Higgs doublets;

(B) N=1 supergravity [5] is valid below $1/\kappa = 2.4 \times 10^{18}$ GeV.

In view of Assumption (A), we note that the single most plausible alternative scenario, the hypercolour idea [6], is in conflict with the feeble flavour-changing neutral current.[7] Two Higgs doublets seem to be phenomenologically required in SUSY GUTs, since four or more Higgs doublets in general lead to too large values for $\sin^2\theta_W$. Assumption (B) is required for consistency. If there exists an elementary scalar field, the natural mass scale cannot be much larger than the mass of the scalar field.[8] The only known way out of this problem is local SUSY. Even if one ever solves the gauge hierarchy problem without SUSY, there always exist dangerous graviton loop diagrams for the light scalar mass corrections. To remove them by cancellations, we need at least local SUSY, with which the graviton loop can be cancelled by the gravitino loops.

Let us express SUSY breaking mass scale by m_S. We have two more mass scales, $1/\kappa$ and the grand unification scale M_{GUT}. Since we are interested in obtaining an upper bound on m_S, we consider the region $m_S \gtrsim 10^{11}$ GeV given by Weinberg.[1)] The SUSY breaking effects of light fields may then be of order

$$m_S^2, \quad (\kappa m_S^2)^2, \quad (\alpha m_S^2/M_{GUT})^2, \ \ldots \ .$$

However, the light Higgs fields should not feel SUSY breaking of order m_S^2, when we accept Assumption (A). Since the second and the third terms are of the same order, we will treat the SUSY breaking effect for light fields as κm_S^2.

When the SUSY breaking is greater than 10^{11} GeV, it is no longer justified to treat the gravity separately from other interactions because its effect κm_S^2 is at least of order TeV. The scalar potential due to the A-components of chiral superfields z_a with a given superpotential W is

$$V = E\tilde{V} + \frac{1}{2} D^{\alpha} D_{\alpha}, \tag{1}$$

where the last term is the usual D-term and

$$E = \exp(\kappa^2 \sum_a |z_a|^2) , \tag{2}$$

$$\tilde{V} = \sum_a \left| \frac{\partial W}{\partial z_a} + \kappa^2 z_a^* W \right|^2 - 3\kappa^2 |W|^2 . \tag{3}$$

Note that we obtain the usual global SUSY potential in the limit $\kappa \to 0$. As in the global SUSY case, nonzero $\partial W/\partial z_a$ characterizes the scale of SUSY breaking $\approx O(m_S^2)$. Then it is easy to observe that all light scalar partners acquire a common mass

$$\kappa^2 \sum_a \left| \frac{\partial W}{\partial z_a} \right|^2 \approx (\kappa m_S^2)^2, \tag{4}$$

which is of great importance, since it is a SUSY version of GIM mechanism for the case $\kappa m_S^2 > m_W^2$.

SUSY breakings which light fields feel through supergravity are quadratic and cubic "soft terms" of the scalar potential.[11)] Usually these soft terms are parametrized by [10)]

$$\tilde{m}^2 \sum_a |y_a|^2 + \left[\left\{ (A-3)\tilde{m}W + \tilde{m} \sum_a y_a \frac{\partial W}{\partial y_a} \right\} + \text{h.c.} \right] , \tag{5}$$

where y_a denotes light scalar fields and A is a numerical constant depending upon specific models. In many cases $\tilde{m}$ coincides with the gravitino mass. The reason that we can get an upper bound on SUSY breaking is that any scalar field cannot escape the fate of feeling the supergravity multiplet through the first term of Eq.(5).

The super Higgs mechanism [9)] occurs when SUSY is broken. The super Higgs sector relevant to this super Higgs mechanism may be hidden [12)] as in the Polony super Higgs sector [13)], or may not be hidden as in many global SUSY models. If the super Higgs sector is not hidden as in geometrical hierarchy models [14)], there exist three mass scales κ^{-1}, M_{dec} and m_S, where M_{dec} is the heavy decoupling mass protecting light fields from strong effects of SUSY breaking [14)]. For $M_{dec} \gg \kappa^{-1}$ it is sufficient to consider the supergravity effects only for the SU(2) X U(1) breaking. For $M_{dec} \ll \kappa^{-1}$ the scale m_S^2/M_{dec} cannot be the scale of electroweak symmetry breaking. If it does not break SU(2) X U(1), we have a large positive squared mass $(m_S^2/M_{dec})^2$ for the Higgs doublets, causing another hierarchy problem. This argument implies $M_{dec} \gtrsim \kappa^{-1}$, or a hidden super Higgs sector for $m_S > 10^{11}$ GeV. Since we are interested in finding out an upper bound in the Weinberg region, it is sufficient to consider a hidden Higgs sector. One example of such super Higgs sector is the Polony form $W_{SH} = m_S^2(z + B)$. However our argument does not depend on the specific form.

To get an upper bound on m_S^2, we must renormalize the soft SUSY breaking terms to low energy. Conveniently enough, Inoue et al. [15)] gave a complete analysis on renormalizations of these parameters in the context of global SUSY. We use their results for arguing the existence of an upper bound on m_S. Since supergravity transmit the SUSY breaking effect to low energy through soft terms, we will consider only gaugino masses M_3, M_2, M_1, and scalar boson masses m_1^2, m_2^2, m_3^2, and trilinear scalar coupling masses m_6, m_8, m_{10}. Other mass parameters of Ref.15 do not appear at the order κm_S^2. The scalar partners of the third generation ℓ_3, τ, q_3, b, t and two Higgs doublets $H_1 = (1, 2, -1/2)$ and $H_2 = (1, 2, 1/2)$ form the soft terms in the Higgs potential,

$$V_{soft} = m_1^2 H_1^\dagger H_1 + m_2^2 H_2^\dagger H_2 - m_3^2 \varepsilon_{ij}(H_{1i}H_{2j} + h.c.)$$
$$- (f m_6 \varepsilon_{ij} H_{1i} \tilde{\ell}_{3j} \tilde{\tau}^c + h m_8 \varepsilon_{ij} H_{1i} \tilde{q}_{3j} \tilde{b}^c$$
$$+ \tilde{h} m_{10} H_{2j} \tilde{q}_{3j} \tilde{t}^c + h.c.) \quad , \qquad (6)$$

where $\sim$ represents a squark or slepton. The condition for SU(2) X U(1) breaking is

$$m_1^2 + m_2^2 > 2|m_3^2| , \qquad m_3^4 > m_1^2 m_2^2 \qquad (7)$$

and m_Z^2 is given by

$$m_Z^2 = - \left[m_1^2 - m_2^2 + (m_1^2 + m_2^2)\cos 2\theta \right] / \cos 2\theta \ , \qquad (8)$$

where

$$\sin 2\theta = 2m_3^2/(m_1^2 + m_2^2). \qquad (9)$$

The m_3^2 term in Eq.(6) needs more explanation in effective theories of local SUSY GUTs. If $m_3^2 = 0$ as in models without a sliding singlet, tree-level breaking of SU(2) X U(1) does not occur in view of Eq.(7). However, SU(2) X U(1) can be broken radiatively, if $m_1^2 + m_2^2$ is renormalized to zero around m_W. This possibility is considered by Ellis et al..[16] They obtained

$m_1^2 \lesssim m_Z^2$ even at 300 GeV, implying $m_S \lesssim 1.5 \times 10^{10}$ GeV. This value is smaller than the upper bound we would like to derive.

When a light singlet is present, the m_3^2 term will be nonzero. For a light singlet X, its couplings are given by

$$W_X = \lambda X H_1 H_2 + \lambda' X^3 \ , \qquad (10)$$

from which we can write the potential energy in terms of v.e.v.'s:

$$V_{v_1 v_2 v'} = \frac{1}{32}(g^2 + g'^2)(v_2^2 - v_1^2)^2 + \frac{1}{2}m_1^2 v_1^2 + \frac{1}{2}m_2^2 v_2^2$$
$$+ m'^2 v'^2 + \lambda A m'' v' v_1 v_2 + \frac{2}{3}\lambda' A m''' v'^3$$
$$+ \frac{1}{4}\lambda_1^2 v_1^2 v_2^2 + \frac{1}{2}\lambda_2^2 v'^2 (v_1^2 + v_2^2) + \lambda_3^2 v'^4 + \lambda_4^2 v'^2 v_1 v_2, \qquad (11)$$

where coupling parameters of each term represents its value at m_W. The v.e.v.'s are

$$\langle H_1^T \rangle = (0, v_1/\sqrt{2}), \ \langle H_2^T \rangle = (v_2/\sqrt{2}, 0), \ \langle X \rangle = v'. \qquad (12)$$

The $(\lambda A m'' v' + \lambda_4^2 v'^2)$ term in Eq.(11) can be interpreted as $-m_3^2$ term of Inoue et al.'s. However, we cannot use their results directly, since there is another quartic term of λ_1^2, which was absent in their analysis. After giving X a v.e.v., we obtain the following conditions for the SU(2) X U(1) breaking,

$$\left| m_1^2 + m_2^2 + \lambda^2(v_1^2 + v_2^2) + 2\lambda_2^2 v'^2 \right| \geq 2 \left| \lambda_4^2 v'^2 + \lambda A m'' v' \right| , \qquad (13)$$

$$(\lambda A m'' v' + \lambda_4^2 v'^2)^2 > (m_1^2 + \lambda_2^2 v'^2)(m_2^2 + \lambda_2^2 v'^2) \ . \qquad (14)$$

If these conditions are satisfied, SU(2) X U(1) breaking occurs, and v_1^2 and v_2^2 are of order m_1^2, m_2^2 and m'^2. Let us take m_1 as the largest mass parameter.

The mass of the charged Higgs boson is

$$m^2(H^{\pm}) = m_1^2 + m_2^2 + m_W^2, \tag{15}$$

which is bounded $< (1\text{TeV})^2$ for a perturbation theory.[17] Considering some possible enhancements in combinations of dimensionless coupling constants, we can safely take $m^2(H^{\pm}) \lesssim (10\text{TeV})^2$, which amounts to $m_1^2 \lesssim (10\text{TeV})^2$.

Using the renormalization group equations, we can obtain the bound on $m_1^2(Q=m_S)$, the value at the scale m_S. Relevant renormalization group equations for our discussion are

$$\frac{dh}{dt} = \frac{h}{16\pi^2}\left(-\frac{16}{3}g_c^2 - 3g^2 - \frac{7}{9}g'^2 + ff^* + 6hh^* + \tilde{h}\tilde{h}^*\right), \tag{16}$$

$$\begin{aligned}\frac{dm_1^2}{dt} = \frac{1}{16\pi^2}\Big[&-6g^2M_2^2 - 2g'^2M_1^2 + 2ff^*\left\{m^2(\ell_3) + m''(e_3) + m_1^2 + m_6^2\right\} \\ &+ 6hh^*\left\{m^2(q_3) + m^2(n_3) + m_1^2 + m_8^2\right\} \\ &+ 2\lambda\lambda^*\left\{m^2(X) + m_2^2\right\}\Big],\end{aligned} \tag{17}$$

where $t = \ln(Q/m_W)$. For $|\tilde{h}| \lesssim O(1)$, $|h|^2$ is asymptotically free. It is reasonable to take $|f|$, $|h| \lesssim 0.1$, since the b-quark mass is much smaller than the v.e.v. v_1. In fact, if $h = 0.1$, we have to take $v_1 \approx 50$ GeV, which is unnatural since $v_1^2/v^2 \equiv v_1^2/(v_1^2 + v_2^2) = (50/250)^2 = 1/25 \ll 1$. Therefore, we assume $|f|^2 = 0.01$, $|h|^2 = 0.01$ and $|\lambda|^2 = 0.01$ for obtaining a possible upper bound. We have to solve the coupled renormalization group equations for m_1^2, m_2^2, $m^2(\ell_3)$, $m^2(e_3)$, m_6^2, etc.. We found it reasonable to obtain an upper bound by approximating $m_1^2(t) = m_2^2(t) = m^2(\ell_3) = m^2(e_3) = m_6^2(t)/A$ in Eq.(17). For $A \leq 10$, we obtain a bound

$$\frac{m_1^2(m_S)}{m_1^2(m_W)} \lesssim 10, \tag{18}$$

which gives the bound of SUSY breaking scale

$$m_S \lesssim 3 \times 10^{11}\ \text{GeV} \tag{19}$$

through the relation $m_1^2(m_S) = \kappa m_S^2$. If the fourth generation is present, $|f|$ and $|h|$ can be large, and our bound is no longer valid. A detailed analysis of SU(2) X U(1) breaking with a sliding light singlet and numerical evaluations of coupled renormalization group equations for m_1^2, m_2^2, ..., m_6^2 will be presented elsewhere.

The scale of SUSY breaking we have obtained coincides with the scale of axion cosmology.[18] It will therefore be very interesting, if one can construct a reasonable invisible axion model[19] with a common scale for SUSY and Peccei-Quinn symmetry breaking.

It is to be emphasized that this upper bound is common to all natural GUT models based on Assumptions (A) and (B) together with small Yukawa couplings.

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