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Grand Unified Theories Based on Local Supersymmetry[†)]

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General features are discussed of grand unified theories with local supersymmetry broken at high energies by the super-Higgs effect. The low-energy effective Lagrangian is a globally supersymmetric one with all the explicit soft breakings. It is argued that the energy splittings among the vacua due to gravitational effects must be small in order to be able to pick out the correct vacuum. Also discussed are the mechanisms of breaking $SU(2)$ by the gravitational effects and of suppressing monopole production in the early universe.

§ 1. Introduction

It is now believed that global supersymmetry (SUSY) may naturally resolve the hierarchy problem in grand unified theories.[1),2)] Unfortunately such attempts seem to have the following shortcomings:

(i) It is difficult to break SUSY spontaneously as it is known to remain unbroken in perturbation theory if it is so at the tree level. We must introduce, at least, three chiral fields for the sole purpose of breaking SUSY.[3)]

(ii) Theories with spontaneously broken SUSY is characterized by the unrealistic mass formula[4)]

$$\sum_J(-1)^{2J}(2J+1)m_J{}^2=0\,, \tag{1}$$

and by a positive vacuum energy (cosmological constant), which must be very close to zero on observational grounds.

(iii) In order not to contradict experimental and astrophysical constraints,[5)] the resulting Goldstino, distinct from the observed neutrino, must couple weakly to matter, or it must be eliminated.

(iv) Globally supersymmetric theories often possess degenerate minima so that we cannot determine which vacuum should be realized.

(v) Since the unification scale is near the Planck mass, it is not completely obvious whether one can still neglect the effects of gravity.

Although these problems except for the last may be overcome by considering radiative corrections and other mechanisms,[***)] global symmetry does not seem to be the most fruitful approach. We are then naturally led to loçal SUSY.[7)] The purpose of this paper is to discuss such theories and show not only that the above-mentioned problems may be evaded but also that phenomenologically interesting models can be constructed. For example, SUSY is easily broken by the super-Higgs effect, which at the same time

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***) A possible way out of the difficulties (ii) and (iii) is the supersymmetric dipole mechanism.[6)] eliminates the Goldstino and cosmological constant. Furthermore, when n chiral multiplets are coupled to supergravity, Eq. (1) is modified to [8)]

$$\sum_J(-1)^{2J}(2J+1)m_J{}^2=2(n-1)m_{3/2}^2\,, \tag{2}$$

where $m_{3/2}$ is the gravitino mass. We shall exhibit some general aspects and a viable scenario for such theories, instead of trying to build specific models.[*)]

§ 2. Super-Higgs effect

We first note that, in constructing realistic models, it is enough to consider only $N=1$ SUSY since the fermions in $N\geqq 2$ SUSY transform like real representations of the gauge group.[1)] It is true that $N\geqq 2$ SUSY can be broken down to $N=1$ theory in supergravity[10)] (but not in global SUSY), but then we would be left effectively with models of the type to be discussed here. We consider supergravity as such an effective theory and do not worry about its nonrenormalizability.

Let us next consider the coupling of chiral fields to $N=1$ supergravity. The scalar potential in such a case is known to take the form[8),9),11)]

$$V=e^{\kappa^2|z_a|^2}\left\{\left|\frac{\partial W}{\partial z_a}+\kappa^2 z_a{}^\dagger W\right|^2-3\kappa^2|W|^2\right\}+\frac{1}{2}|D_A|^2\,, \tag{3}$$

where $G=\kappa^2/8\pi$ is the gravitational constant, W the superpotential for chiral fields z_a, and $D_A=z_a{}^\dagger(t_A)^{ab}z_b$ with t_A being the representation of the A-th generator for these scalars, including a coupling-constant factor. The last term in Eq. (3) is positive definite and must be small. This automatically leads to a zero vacuum expectation value(VEV) for the scalar partners of quarks and leptons. Then this term vanishes in all the examples considered here and consequently will be dropped hereafter.

In order to implement the supe.-Higgs effect, we take the superpotential[12)]

$$W=m^2(z+B_0), \tag{4}$$

where z is a gauge singlet chiral field, and B_0 and m are constants. By imposing the conditions $\partial V/\partial z=0$ and $V=0$, we find[9)]

$$\langle z\rangle=\frac{a+\sqrt{3}b}{\kappa}\,,\qquad B_0=-\frac{\sqrt{2}a+\sqrt{3}b}{\kappa}\,, \tag{5}$$

where $a,b=\pm1$, and the gravitino acquires the mass $m_{3/2}=\kappa m^2\cdot e^{2+\sqrt{3}ab}$ by the super-Higgs effect. It turns out that we must choose $ab=-1$ since the scalar A and pseudoscalar components B of $z[z-\langle z\rangle=(A-iB)/\sqrt{2}]$ have the masses $-2\sqrt{3}abm_{3/2}^2$ and $2(2+\sqrt{3}ab)\times m_{3/2}^2$, respectively. (If $ab=+1$, then there are solutions for $\partial W/\partial z+\kappa^2 z^\dagger W=0$, leading to unbroken SUSY.) We put $b=-a$ but leave a as a free parameter taking values ±1. On the other hand, the cosmological constraints[13)] imply that $m_{3/2}\gtrsim 10^{3-4}$ GeV or $\lesssim 1$ keV. The latter possibility reduces Eq. (2) practically back to Eq. (1) and hence is not attractive.[**)] We are then forced to choose $m\gtrsim 10^{10}$ GeV, which suggests that $\kappa m^2\gtrsim 10^3$ GeV is

*) Some specific models have recently been considered by several authors.[9)]

**) Since SUSY is broken, radiative corrections can also modify Eq.(1).

the electroweak mass scale. We shall see that this possibility in fact leads to many interesting consequences.

§ 3. $SU(5)$ models

3.1. *Low-energy effective Lagrangian*

We now proceed to examination of some $SU(5)$ models based on this scheme. We do this by expanding the potential in $\varkappa$ with the superpotential

$$W=g+m^2(z+B_0), \tag{6}$$

where g is that for the scalars z_i in appropriate representations of $SU(5)$. A straightforward calculation yields

$$V=e^{\varkappa^2\{|z_i|^2+|z|^2-\langle z\rangle^2\}}\{V_0+V_1+V_{2m}+V_{2r}+O(\varkappa^3)\} \tag{7}$$

with

$$V_0=\left|\frac{\partial\tilde{g}}{\partial z_i}\right|^2, \tag{8}$$

$$V_1=am_{3/2}\left(-z_i\frac{\partial\tilde{g}}{\partial z_i}+\sqrt{3}\,\tilde{g}+\text{h.c.}\right), \tag{9}$$

$$V_{2m}=m_{3/2}^2\{\sqrt{3}A^2+(2-\sqrt{3})B^2+|z_i|^2\}, \tag{10}$$

$$V_{2r}=\varkappa^2\Big[z_i\frac{\partial\tilde{g}}{\partial z_i}\Big\{\tilde{g}^\dagger+\frac{\tilde{m}^2}{\sqrt{2}}(A+iB)\Big\}+\text{h.c.}$$
$$-(2\sqrt{3}-1)|\tilde{g}|^2-(\sqrt{6}-\sqrt{2})i\tilde{m}^2B(\tilde{g}-\tilde{g}^\dagger)\Big], \tag{11}$$

where $\tilde{g}=e^{2-\sqrt{3}}g$ and $\tilde{m}^2=e^{2-\sqrt{3}}m^2$. In the limit $\varkappa\to 0$ with $\varkappa m^2(\propto m_{3/2})$ held fixed, only Eqs. (8)~(10) survive.[14] Notice that (8) is a globally supersymmetric term*) accompanied by the explicit breakings**) (9)~(11) due to the gravitational effects. The requirement of naturalness[2] is therefore satisfied if we take the above limit as we do in ordinary field theories. However, since we intend to discuss the case when g involves the grand unified mass M assumed to be of order 10^{16} GeV and larger than m, other terms cannot be neglected either. Thus we should also keep terms which do not vanish in the limit $\varkappa\to 0$ with $\varkappa M^2$ *and* $\varkappa m^2$ held fixed. Since such SUSY-breaking new terms can arise only from terms of order $\varkappa^2M^4$ in Eq. (7), they have dimension less than three and hence are soft, maintaining the naturalness of the theory.

3.2. *$SU(2)$ breaking at the tree level*

It should also be noted that *all* the scalar fields automatically possess a common mass $m_{3/2}$ because of the terms in Eq. (10).[16] These terms tend to prohibit scalars from developing nonvanishing VEVs. As we shall see shortly, these often make the electroweak gauge group $SU(2)$ unbroken.

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*) The structure of the fermion part is the same as in global SUSY with the superpotential $\tilde{g}$.

**) In particular, (9) and (10) are soft in the sense that no quadratic divergences are generated.[15]

However, this is not always the case. As an example, consider the superpotential[9)]

$$g=\lambda_1\left(\frac{1}{3}\mathrm{Tr}\,\Sigma^3+\frac{M}{2}\mathrm{Tr}\,\Sigma^2\right)+\lambda_2 H'(\Sigma+3M')H+\lambda_3 SH'H\,, \tag{12}$$

where Σ, H, H' and S are scalar fields in **24**, **5**, $\bar{\mathbf{5}}$ and **1** representations of $SU(5)$, respectively, and $\lambda_i(i=1,2,3)$, M and M' are parameters. The minima of the potential are determined by Eq. (8) in the lowest order of x. In this case, we have three minima at

$$\langle H\rangle=\langle H'\rangle=0,\quad \langle S\rangle=\text{undetermined},$$

$$\text{(i)}\ \ \langle\Sigma\rangle=0\,,\qquad \text{(ii)}\ \ \langle\Sigma\rangle=\frac{R}{3}\mathrm{diag}(1,1,1,1,-4),$$

$$\text{(iii)}\ \ \langle\Sigma\rangle=M\,\mathrm{diag}(2,2,2,-3,-3), \tag{13}$$

corresponding to $SU(5)$, $SU(4)\otimes U(1)$ and $SU(3)\otimes SU(2)\otimes U(1)$ phases, respectively. One may hope that the singlet field S in Eq. (12) may serve to make the $SU(2)$ doublets in H and H' light naturally. In global SUSY, the minimization condition with respect to $\langle S\rangle$ ensures that these doublets are massless at the stage of $SU(5)$ breaking — the so-called "sliding singlet" mechanism.[3)] Unfortunately, in our case the additional term (10) makes large VEV for S energetically unfavorable and in general tends to prevent this mechanism from working. Moreover, in the model (12), a closer examination of the full potential shows that, in order to break $SU(2)$ at the energy scale of xm^2 ($\gtrsim 10^3$ GeV), we *must* set $M'=M$ and $\langle S\rangle=0$ (at order x^0). In fact, we can then find a solution in which $SU(3)\otimes SU(2)\otimes U(1)$ is broken to $SU(3)\otimes U(1)$ at the energy scale of xm^2.[9)] Thus the gravitational effects can break $SU(2)$ group. (Another mechanism will be discussed later.)

3.3. *Avoidance of anti-de Sitter vacua and nearly degenerate vacua*

This solution suffers from the presence of a large cosmological constant of order x^2M^6, but it can be chosen to vanish by adding a constant to g. Once this is done, the other two solutions in Eq. (13) correspond to large negative-vacuum-energy ($\sim -x^2M^6$) states, which we call anti-de Sitter vacua hereafter. Nevertheless, we may choose any state as our vacuum since all these vacua are stable.[17)] This is frustrating, however, since we are still ignorant about the correct choice of the vacuum despite the fact that the degeneracy has been lifted. This ambiguity cannot be resolved even when we trace back the history of the universe because of the negligible thermal effects near the expected phase transition points compared with the large gravitational splittings.[17)]

Fortunately, such anti-de Sitter vacua can be avoided in the presence of the super-Higgs term (4).[9)] The question then arises of what are the necessary conditions for this evasion of anti-de Sitter vacua. To investigate this problem, suppose $g\sim\Delta$ at the stationary points which would be near the points $\partial g/\partial z_i=0$. From Eqs. (3) and (6), it is easy to see that the vacuum energy is then of order $x^2m^4M^2-xm^2\Delta-x^2\Delta^2$, which can be positive only when $\Delta\lesssim m^2M\ll M^3$. We have investigated some models and found that, if we try to make Δ smaller than M^3, we are inevitably forced to have vanishing Δ in the limit $x\to 0$. For example, consider

$$g=\lambda_1 X(\mathrm{Tr}\,\Sigma^2-M^2)+\lambda_2 H'(\Sigma+M')H+\lambda_3 SH'H\,, \tag{14}$$

where X is a singlet of $SU(5)$ and other fields are the same as in Eq.(12). This superpotential vanishes at the following minima in the lowest order of x:

$$\langle H\rangle=\langle H'\rangle=\langle X\rangle=0\,,\qquad \langle S\rangle=\text{undetermined}\,,$$

$$\text{(i)}\ \ \langle\Sigma\rangle=\frac{M}{\sqrt{30}}\,\mathrm{diag}(2,2,2,-3,-3),\quad \text{(ii)}\ \ \langle\Sigma\rangle=\frac{M}{\sqrt{20}}\,\mathrm{diag}(1,1,1,1,-4). \qquad (15)$$

Examination of the minimum points in the next order of x shows that they are modified to

$$\langle H\rangle=\langle H'\rangle=\langle S\rangle=0\,,\qquad \langle X\rangle=\frac{xm^2a}{2\lambda_1}\,,$$

$$\text{(i)}\ \ \langle\Sigma\rangle=\frac{M}{\sqrt{30}}(1+\varepsilon)\mathrm{diag}(2,2,2,-3,-3),\quad \text{(ii)}\ \ \langle\Sigma\rangle=\frac{M}{\sqrt{20}}(1+\varepsilon)\mathrm{diag}(1,1,1,1,-4) \qquad (16)$$

with common $\varepsilon=-(\sqrt{3}-1)x^2m^4/4\lambda_1{}^2M^2$, and these states are degenerate in energy. This implies that we have no anti-de Sitter vacuum. Arnowitt et al.[9] considered a more complicated potential with additional **24** and **1** fields and found that the $SU(4)\otimes U(1)$ state is higher than the $SU(3)\otimes SU(2)\otimes U(1)$ only by energy of the order of TeV. Although the precise amount of the energy splitting depends on the details of the models considered, the general picture emerging from this approach is that the degeneracy of vacua is in general removed by the gravitational effects *only slightly* ($\lesssim$ TeV) if we are to avoid anti-de Sitter solutions like Weinberg's.

This is welcome from the cosmological point of view. In fact, we can determine in this case which vacuum will actually be realized by tracing back the history of the universe: At a very high temperature, the potential has additional contributions of the form $-(\pi^2/90)(N_B+\frac{7}{8}N_F)T^4$ with N_B and N_F being the number of helicity states of light bosons and fermions, respectively, and thermal effects always favor the most symmetric phase. As we cool down in this phase, the growth of the coupling constant with decreasing temperature would reduce the coefficients of T^4 in the free energy density owing to the "confinement of massless particles" in the strong coupling region, causing a rapid transition to broken symmetry.[18] If the transition from $SU(5)$ takes place towards the $SU(4)\otimes U(1)$ phase (this is expected at $T\sim10^9$GeV), similar strong coupling pheno mena of $SU(4)$ would occur (at $T\sim10^6$GeV) and finally we are left with the $SU(3)\otimes SU(2)\otimes U(1)$ phase. This is quite plausible because the $SU(3)\otimes SU(2)\otimes U(1)$ phase is the only vacuum that does not possess strong coupling down to the temperature $T\sim1$ GeV.

There is an argument that, in order for this mechanism to work, the barrier between the different vacua must be small; the barrier height must be smaller than 10^{10} GeV.[19] However, in the models like (14), it is possible to satisfy this condition by choosing small λ_1.

In the example (14), only the transition from $SU(4)\otimes U(1)$ to $SU(3)\otimes SU(2)\otimes U(1)$ is expected since $SU(5)$ is not included among the almost degenerate vacua. We also note that in this scheme we need not make the energy of $SU(4)\otimes U(1)$ phase higher than that of $SU(3)\otimes SU(2)\otimes U(1)$, because we enter the strong coupling region, where the above transition is expected to occur, at a temperature much larger than that the energy splittings among the vacua; the energy splittings are negligible compared with the thermal energy at the expected transition point. It seems that this scenario is inevitable in order to avoid Weinberg's situation where one could not determine the correct vacuum.

3.4. $SU(2)$ *breaking by radiative corrections*

Another apparent difficulty with the solution (i) in Eq. (16) is that $SU(2)$ remains unbroken. This arises precisely because gravity gives positive contributions to the masses of the Higgs doublets [Eq. (10)]. At first sight, the gravitational effects in general seem to make unbroken $SU(2)$ favorable, as we have already remarked. Actually this difficulty can be resolved by the gravitational effects themselves in the following way. Now that the gravitino is massive, the four-fermion coupling of the gaugino λ to the gravitino characteristic of supergravity

$$\kappa^2 \bar{\lambda}\lambda\bar{\psi}_\mu\psi^\mu , \tag{17}$$

will induce gaugino mass terms of order $m_{3/2}$ (at the one loop level) if we take $1/\kappa$ as the cutoff for the gravitino loop.[11),20)] There will also appear radiative gaugino masses coming from radiative graphs involving the gauge multiplets since SUSY is broken.[3)] Finally there may be direct gaugino masses if the coupling of $N=1$ supergravity to the Yang-Mills strength is non-minimal.[8)] This is indeed the case for $N \geqq 4$ supergravity. In any case it is quite plausible that the gauginos acquire nonvanishing masses of order $m_{3/2}$ owing to the gravitational effects. These in turn induce positive masses for scalar partners of quarks, giving rise to negative contributions to the masses of the Higgs doublets.[3)]

Suppose these additional contributions are summed up to yield

$$-\alpha^2 m_{3/2}^2(|H|^2+|H'|^2). \tag{18}$$

We find that if $M'=\frac{3}{\sqrt{30}}M$*) and

$$\alpha^2 > \sqrt{3}-\frac{3}{4} , \tag{19}$$

then the model (14) in fact possesses the following solution with $SU(2)$ broken at the scale of κm^2:

$$\langle S\rangle=\frac{\kappa m^2 a}{\lambda_3}y^2 , \quad \langle X\rangle=\frac{\kappa m^2 a}{2\lambda_1} , \quad \langle H_i\rangle=\langle H_i'\rangle=\frac{\kappa m^2}{\lambda_3}y\delta_{i5} ,$$

$$\langle\Sigma\rangle=\frac{M}{\sqrt{30}}\,\mathrm{diag}(2, 2, 2, -3+\varepsilon', -3-\varepsilon'), \tag{20}$$

where

$$y^2=\frac{2\lambda_3{}^2}{\lambda_2{}^2}\left(\frac{3}{4}-\sqrt{3}+\alpha^2\right), \quad \varepsilon'=\frac{3\sqrt{5}\kappa m^2 a}{\sqrt{2}\lambda_2 M}\left\{1-\frac{2}{\sqrt{3}}(1-y^2)\right\}. \tag{21}$$

This mechanism would work favorably provided that there is some lower bound of order 10 GeV for the mass of the top quark, as evaluated in Ref. 3). It should also be noted that

*) This condition is necessary to keep the Higgs doublets light at $SU(5)$ breaking stage. The "sliding singlet" does not work here either.

$\sqrt{3}-\frac{3}{4}$ is less than unity. This implies that the mass terms of the Higgs doublets arising from Eqs. (10) and (18) need not be negative but must be small. We emphasize that this mechanism is quite general in this type of unified models, in contrast to that discussed in the model (12).

3.5. *Monopole suppression*

One more problem we should discuss in this model is related to the monopole production in the early universe. Since $SU(5)$ is broken to $SU(4)\otimes U(1)$ at the temperature near the unification scale, too many monopoles would be produced through thermal fluctuations. To avoid this, it is enough to have $SU(5)$ phase as well. This can be achieved, for example, by taking the superpotential

$$g=\lambda_1\{\mathrm{Tr}(\Sigma^2\Pi)+M\ \mathrm{Tr}(\Sigma\Pi)\}+\lambda_2 H'(\Sigma+M')H+\lambda_3 SH'H\ , \tag{22}$$

where Π is an additional adjoint scalar field. This potential gives $SU(5)$, $SU(4)\otimes U(1)$ and $SU(3)\otimes SU(2)\otimes U(1)$ minima almost depenerate ($g=0$ at all the minima). Monopoles may be suppressed sufficiently according to the development of the universe described before, since the transition from $SU(5)$ to $SU(4)\otimes U(1)$ is delayed until the $SU(5)$ coupling grows large ($T_c\sim 10^9$ GeV$\ll M_{\mathrm{monopole}}$).

The number N_m of monopoles produced in the early universe can be evaluated in the same manner as in Ref. 18) and turns out to be

$$N_m\sim cT^3h^{3/4}\ ,$$

where $c=(3/\pi)^{1/4}\Gamma(5/4)\simeq 0.90$ and h is a dimensionless constant. The temperature T^* at which the transition from $SU(5)$ to $SU(4)\otimes U(1)$ is completed is found to be[18)]

$$1/T^*\simeq 1/T_c+O(10)/h^{1/4}M_{\mathrm{P}}\ .$$

Since $T_c\simeq T^*\simeq 10^9$ GeV, this implies $h^{1/4}\sim 10^{-8\sim -9}$, leading to

$$N_m/T^3\simeq 10^{-27}\ .$$

This is well below the bound $\lesssim 10^{-24}$, as we expected.

3.6. *Flavor-changing-neutral currents*

Finally we note that this breaking mechanism of SUSY due to gravitation does not discriminate flavors. This leads to the natural suppression of flavor-changing-neutral currents(FCNC) according to Inami and Lim.[9),21)]

§ 4. Concluding remarks

We have thus shown that $N=1$ supergravity seems to provide phenomenologically interesting models without difficulties inherent in global SUSY and predict new physics near TeV region. To summarize the main points of the scenario outlined here, the gravitational effects break SUSY in such a way that (i) the unrealistic mass formula (1) is no longer true even before the introduction of radiative corrections, (ii) cosmological constant can be made to vanish, (iii) Goldstino is absent, (iv) the degenerate vacua in global SUSY are slightly shifted and as a result the correct vacuum may be chosen by tracing back the history of the universe, (v) $SU(2)$ is also broken by gravitation, (vi) monopoles may be suppressed, and (vii) FCNC is suppressed naturally. The resulting low-energy effective theory is given by a globally supersymmetric theory with all the explicit soft breakings[15] originating from (9), (10) and (17).

There are also remaining problems. It seems that an adjustment of parameters is necessary to make the Higgs doublets light. The "sliding singlet" does not work, at least, in its simplest form. This may be overcome by the mechanism of "missing partner", e.g., by using **50** representation of $SU(5)$.[22] The question of how the cosmological constant can be made to vanish naturally and that of to what extent we can narrow down the choice of the superpotential are also left unsolved. Perhaps, these problems may be resolved when we succeed in reducing theories based on $N \geqq 2$ supergravity to the models discussed here.

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Note added: The soft breaking terms (9) transform like the A-terms of chiral multiplets and in general have coefficients of intermediate mass scale ($\sim m_{3/2}M \sim 10^{11}\text{GeV}^2$) while the terms (10) transform like the C-terms of vector multiplets with coefficient of low energy mass scale ($\sim m_{3/2}^2$). Our scheme is different in this special breaking pattern of SUSY and other points discussed in the text (e.g., § 3. 6) from models with SUSY broken softly by hand. A similar model with intermediate breakings introduced by hand is discussed in H. Nishino and S. Watamura, Univ. of Tokyo Preprint UT-Komaba 82-14 (1982).