Pramāṇa – J. Phys., Vol. 25, No. 4, October 1985, pp. 353–361. © Printed in India.

Gravity-induced weak symmetry breaking and supergravity*

S P MISRA

Institute of Physics, Bhubaneswar 751 005, India

Abstract. We give here a review of the recent developments of grand unified theories based on $N = 1$ supergravity. We start with a brief introduction of supersymmetry and supergravity multiplets, and then discuss the construction of an invariant Lagrangian. The phenomena of gravity-induced weak symmetry breaking via the super Higgs effect at the tree level, corresponding to the conventional SU(5) gauge group, are then considered. We then extend this idea to the larger group SO(10), showing two possible breaking chains given as (i) $SO(10) \times \text{SUSY} \to SU(2)_L \times U(1)_R \times U(1)_{B-L} \times SU(3)_C$ ($\equiv G_{2113}$) $\times \text{SUSY} \to U(1)_{em} \times SU(3)_C$ (G_{LE}) predicting a second Z-boson having mass lower than 1 TeV, and (ii) $SO(10) \times \text{SUSY} \to SU(2)_L \times SU(2)_R \times SU(4) \to (\equiv G_{224}) \times \text{SUSY} \to SU(2)_L \times U(1)_Y \times SU(3)_C$ ($\equiv G_{213}$) $\times \text{SUSY} \to U(1)_{em} \times SU(3)_C$. We also consider the radiative breaking of weak symmetry via renormalisation group effects, which predicts the top quark mass. Some experimental signatures of the supersymmetric particles are investigated and possible future outlook is discussed.

Keywords. Supergravity; grand unification; weak symmetric breaking

On the occasion of the Symposium celebrating Bose Statistics, it may be appropriate to consider internal symmetries which connect bosons and fermions. Supersymmetry, as such a boson-fermion symmetry, has been an extremely significant symmetry principle developed since early 1970's (Wess and Zumino 1974; Salam and Strathdee (1974). This could circumvent many no-go theorems (Coleman and Mandula 1967) regarding intermingling of internal symmetry with space-time symmetries, since here one dealt with graded Lie algebra, as opposed to Lie algebra. Further, theoretical developments with local form of supersymmetry included within its description the spin-2 graviton, with a hypothetical spin-3/2 particle, called gravitino (Nilles 1984). This gave rise to the hope that perhaps finally gravitation can also be absorbed into these schemes with a complete unification of forces.

However, in spite of such aesthetically appealing features, the scheme could not be implemented with any phenomenologically realistic model until late 1982. It was then realised simultaneously in many almost parallel attempts, that in fact local supersymmetry can trigger the breaking of Glashow-Salam-Weinberg symmetry at the weak scale (Chamseddine *et al* 1982). This phenomenon had many interesting aspects. It gave rise to a belief that gravitational coupling constant, so far regarded as irrelevant for high-energy physics, can play a vital role for weak-symmetry breaking. Further, it appears to indicate that phase transition at Planck scale (Polonyi 1977) could have observable effect at scales as low as 100 GeV! Some aspects of this highly interesting developments of past two years will be the subject matter of the present talk.

The basic philosophy of this approach is of two specific types. Sometimes, a grand unification group, SU(5) or SO(10) or one of a more esoteric variety, is regarded as the

* Invited talk presented at the International Symposium on Theoretical Physics, Bangalore, November 1984.

final symmetry group below Planck scale. Then, two or three stages of symmetry breaking can take place at successive mass scales. These are, (i) grand-unification mass-scale; (ii) any other intermediate mass scale, and, (iii) the weak scale. The last scale is taken to be induced by supergravity. One advantage of this approach is that it ensures the stability of the gauge-hierarchy of the lower scales due to non-renormalisation theorems of supersymmetry. The smallness of the weak-scale compared with earlier scales here arises from the large value of Planck mass-scale compared to hypothetical supersymmetry breaking scale.

First we recall the basic features of supersymmetric models. Here, quarks and leptons belong to what are known as chiral supermultiplets. Such multiplets contain scalars, as superpartners of the quark and leptons, usually called squarks and sleptons. Since they are not observed, they should be reasonably massive (Haber and Kane 1984). Usually models predict their masses to be of the same order as the weak scale. The Higgs particles also belong to chiral supermultiplets, naturally accompanied by corresponding spin-$\frac{1}{2}$ superpartners known as Higgsinos. Further, the gauge bosons, belong to the vector supermultiplets which have also spin-$\frac{1}{2}$ superpartners, called gauginos. We give in table 1, the signatures of the particles which would be observed in the weak-symmetry breaking scale, i.e. say, around 50–300 GeV. Local supersymmetry or supergravity also includes spin-3/2 gravitino in addition to graviton in a supermultiplet, as given in the same table.

Local supersymmetry transformations are best understood through the superspace algebra (Wess and Bagger 1981), given as

$$\begin{aligned} x^m &\to x^m + i(\theta\sigma^m\bar{\xi}(x) - \xi(x)\sigma^m\bar{\theta}), \\ \theta^\mu &\to \theta^\mu + \xi^\mu(x), \\ \bar{\theta}_{\dot{\mu}} &\to \bar{\theta}_{\dot{\mu}} + \bar{\xi}(x). \end{aligned} \tag{1}$$

In the above θ and $\bar{\theta}$ are two-component Weyl spinors and $\xi(x)$, $\bar{\xi}(x)$ yield translations in superspace co-ordinates x^m, θ and $\bar{\theta}$. Obviously, there is a mixing of space-time co-ordinates x^m and θ, $\bar{\theta}$. θ and $\bar{\theta}$ are anticommuting C-numbers, or Grassman variables. Superfields are operator functions of super-coordinate and its components are ordinary field operators with specific Lorentz transformation properties.

Table 1. Superpartners of ordinary particles.

	Spin-1	Spin-$\frac{1}{2}$	Spin-0
Scalar multiplets or Matter multiplets		quarks q_L, q_R leptons l_L, l_R *Higgsinos*	*squarks* $\tilde{q}_L, \tilde{q}_R$ *sleptons* $\tilde{l}_L, \tilde{l}_R$ Higgs
Vector Multiplets	Photon γ	*Photino* $\tilde{\gamma}$	
	Gluon g	*Gluino* $\tilde{g}$	
	W-boson W	*Wino* $\tilde{W}$	
	Z-boson Z	Zino $\tilde{Z}$	
Gravity multiplet	Graviton, spin-2		
	Gravitino, spin-3/2		

Supergravity multiplets are best introduced through differential forms (Wess and Bagger 1981). With $\{Z^M\} = \{x^m, \theta^\mu, \bar{\theta}_{\dot{\mu}}\}$ as co-ordinates, the vielbein is a 1-form given as,

$$E^A = \mathrm{d}Z^M E^A_M. \tag{2}$$

The connection 1-form

$$\phi^B_A = \mathrm{d}Z^M \phi^B_{MA} \tag{3}$$

is the basic geometrical object in superspace, where, A, B are group indices. The connection 1-form always absorbs the Lorentz group structure, and shall, in addition absorb also any internal group structure. Suitable constraints are imposed on torsion 2-form which is given as

$$T^A = DE^A = \mathrm{d}E^A + E^B \phi^A_B \tag{4}$$

and curvature 2-form which is given as

$$F = \mathrm{d}\phi + \phi\phi. \tag{5}$$

In the above, the 'wedge' is omitted in writing the product of differential forms. One then defines the supergravity transformations for chiral multiplets, vector multiplets as well as the gravity multiplet. The algebra is complicated and is being omitted in the present talk.

The phenomenology of supergravity can be predicted with the construction of an invariant Lagrangian (Cremmer *et al* 1983). This consists of the field operators with the chiral supermultiplet $S \equiv (Z_i, \chi_{Li}, h_i)$, vector supermultiplet $V \equiv (B^\alpha_\mu, \lambda^\alpha_R, D^\alpha)$ and the gravity supermultiplet given as $(e^m_\mu, \psi_{\mu L}, u, A_\mu)$. In the above, Z_i, χ_{Li} are the scalar and spin-$\frac{1}{2}$ fields, B^α_μ and λ^α_R are the spin-1 gauge fields and spin-$\frac{1}{2}$ gaugino fields, and e^m_μ, $\psi_{\mu L}$ are spin-2 graviton and spin-3/2 gravitino fields. The remaining are auxillary fields. i and α are indices for gauge symmetry, with α standing for the adjoint representation of the gauge group. In the above framework, we have used the Wess-Zumino gauge (Wess and Zumino 1974) to describe the vector multiplet. Otherwise, this supermultiplet will have a larger number of fields, which are gauge-related.

The action integral in superspace is given as (Cremmer *et al* 1983),

$$I = \int \mathrm{d}^4x\, \mathrm{d}^4\theta\, E\left[\phi(S, \bar{S}e^{2v}) + \mathrm{Re}\left(\frac{1}{R}g(s)\right) + \mathrm{Re}\left(\frac{1}{R}f_{\alpha\beta}(S)\, W^\alpha_a \varepsilon^{ab} W^\beta_b\right)\right] \tag{6}$$

In the above, E is the determinant of vielbein, and R is the scalar curvature, both being in superspace. Further ϕ, g and f are arbitrary functions of chiral multiplets S, which are separately gauge-invariant. Now, in order to consider the dynamics, one first eliminates the auxillary fields from the Lagrangian obtained from (6) through equations of motion. One nice feature of supergravity is that the Lagrangian for the gravity multiplet then arises automatically from the first term of the right side of (6) after a so-called Weyl rescaling. Further, one may combine the two functions ϕ and g into a single function $\mathscr{G}$ with suitable transformations in superspace. We note that,

$$\mathscr{G} = J - \log\left(\frac{1}{4}|g|^2\right) \tag{7}$$

where, for a flat Kahler metric,

$$J = \frac{1}{2} Z_i^* Z_i. \tag{8}$$

In the above, gravitational coupling constant is hidden since we have choosen $k = (8\pi G)^{1/2} = 1$. Explicitly retaining k, the Lagrangian finally yields the potential given as

$$V = \frac{1}{2} \exp\left(\frac{k^2}{2} |\phi|^2\right) \left[|g_\phi + \frac{1}{2} k^2 \phi^+ g|^2 - \frac{3}{2} k^2 |g|^2 \right] + \frac{1}{2} D_\alpha^2. \tag{9}$$

In the above, ϕ stands for an arbitrary scalar field, and D_α is the gauge auxillary field, and g is the superpotential of equation (7). ϕ summation is understood.

We shall now qualitatively discuss how from the potential (9), supergravity can lead to gravity-induced symmetry breaking at the weak scale. We first note that in (9) if we take k as zero, then we get the potential for global supersymmetry. From this potential one can easily construct grand-unified models which maintains supersymmetry, but lead to unbroken Salam-Weinberg symmetry below the grand-unification mass scale. We now switch on the gravitational coupling with $k = (8\pi G)^{1/2} \neq 0$ in (9) as the next lowest order approximation. The parameter that enters for such a correction, is given as $km \simeq 10^{-8}$ where m is a mass scale associated with the supersymmetric potential. This gravity-induced correction to the potential causes the breaking of Salam-Weinberg symmetry.

We shall first review as an illustration, the tree-level breaking pattern corresponding to the conventional SU(5) gauge group (Chamseddine *et al* 1982). In this framework one may start with superpotential $g = g_1(\phi) + g_2(Z)$ where g_1 contains the gauge super-multiplets, and g_2 is a function of Polonyi singlet Z. One may take for example,

$$g_1(\phi) = \lambda_1 \left[\frac{1}{3} \mathrm{Tr}\Sigma^3 + \frac{1}{2} M \mathrm{Tr}\Sigma^2 \right] + \lambda_2 H'_x [\Sigma_y^x + 3M'\delta_y^x] H^y$$
$$+ \lambda_3 U H'_x H_x + \varepsilon_{uvwxy} H^u M^{vw} m_1 M^{xy} + H'_x M^{xy} m_2 M'_y \tag{10}$$

where Σ_y^x is in the adjoint representation {24} of SU(5), H^x and H'_x are in the {5} and {5*} representation, M_j^{xy} and M'_{jx} are the matter multiplets in the {10}, {5*} representations, m_1 and m_2 are matrices in the generation space. The Polonyi superfield Z may be e.g. taken in the form

$$g_2(Z) = m^2(Z + B_0), \tag{11}$$

where B_0, m are constants of dimension of mass. In the expression of g_1, the field Σ_y^x breaks SU(5) gauge group to the standard Salam-Weinberg symmetry $SU(2)_L \otimes U(1)_Y \otimes SU(3)_C (\equiv G_{213})$ while preserving global supersymmetry at this initial stage. Then the fields H, H' and U break the Salam-Weinberg symmetry. Supersymmetry in the visible sector is also broken at the same weak scale via the super Higgs mechanism with $g_2(Z)$. This breaking of supersymmetry comes about by minimising the potential of (9). Here it may be noted that the singlet Z takes a vacuum expectation value of the order of Planck scale, which ultimately causes the supersymmetry breaking as well as the weak symmetry breaking at the weak scale. This theory predicts the masses of the gravitino and the scalar partners of the quarks and leptons to be of the order of weak scale. The exact values of masses are very much model-dependent.

We may also consider the grand-unification gauge group SO(10) which has a richer structure in the sense of various possible intermediate symmetries (Maalampi and Pulido 1983). Before we do so, however we consider a single stage breaking of SO(10) to the gauge group (Mahapatra *et al* 1984) $SU(2)_L \otimes U(1)_R \otimes U(1)_{B-L} \otimes SU(3)_C$ ($\equiv G_{2113}$) as the low-energy gauge group as opposed to the standard model. The interesting feature here is that the right symmetry is broken in the charged sector at the grand-unification scale and the residual symmetry G_{2113} is broken via supergravity corrections. The neutral sector is different from that of the standard model in having two neutral Z-bosons. The surprising feature is the consistency of this model with the observed Z-boson mass along with a second Z-boson, whose mass may be lower than 1 TeV. Because of its peculiarity, and as an illustration of supergravity, we now discuss it in some detail.

In this framework one takes the chiral Higgs multiplets $T\{210\}$, $\psi_1\{16_1\}$, $\psi_2\{16_2\}$, $\bar{\psi}_1\{\overline{16}_1\}$, $\bar{\psi}_2\{\overline{16}_2\}$ and the gauge singlet $U\{1\}$ which constitute the observable sector g_1 of the superpotential g, given as

$$g_1 = \lambda_1\left[\frac{M}{2}TT - \frac{1}{3}TTT\right] + \lambda_2\bar{\psi}_1(T - aM)\psi_1 + \lambda_3\bar{\psi}_2(T - a'M)\psi_2 + \lambda_4 U(\psi_1\bar{\psi}_1 + \psi_2\bar{\psi}_2). \tag{12}$$

As usual one may take the Polonyi part of (11) to break the supersymmetry. By Higgs mechanism one then breaks SO(10) symmetry to G_{2113} by giving non-zero vacuum expectation value (VEV) to $T\{210\}$ with the condition that global supersymmetry (SUSY) is good at the initial stage. Then by super Higgs mechanism, through the non-zero vevs of ψ_1, ψ_2, U and Z, supersymmetry is broken simultaneously with the present weak symmetry G_{2113}.

Here the charged current sector is the same as in the standard model. The neutral current sector including the electromagnetic current is given as

$$H^{NC}_{\text{int}} = g_L W_{3L} J_{3L} + g_{3R} W_{3R} J_{3R} + g_C C J_{(B-L)/2}. \tag{13}$$

This yields the effective neutral current-current interaction as,

$$J^{\text{eff}}_{NC} = \frac{4G_F}{\sqrt{2}}\,[J^2_{ZL} + \eta^2 J^2_{ZR}], \tag{14}$$

where,

$$J_{ZL} = J_{3L} - \frac{e^2}{g_L^2} J_{\text{em}} \quad \text{and} \quad J_{ZR} = J_{3R} - \frac{e^2}{g_R^2} J_{\text{em}}.$$

Further, $\eta = (v_1/v_2)$ where v_1 is the symmetry breaking scale for the left-handed weak doublet and v_2 is that of the right-handed weak doublet respectively of ψ_1 and ψ_2. Here it is to be noted that in the standard model $\eta = 0$. In table 2 we give respectively the masses of two neutral Z-bosons for various values of η. It appears that we may even expect a second neutral Z-boson with mass around 500 GeV. The renormalisation group calculations (Mahapatra *et al* 1984) show that here the grand unification mass scale is of the order of 4×10^{15} GeV, and is consistent with the present bounds on proton life-time.

Now we may consider another example of SO(10) symmetry breaking with

Table 2. Masses of two neutral Z-boson (Z_1 and Z_2).

Inputs	$\eta = v_1/v_2$	M_{z1} (GeV)	M_{z2} (GeV)
$v_1 = 175\,\mathrm{GeV}$	0·10	93	889
$M_W = 82\,\mathrm{GeV}$	0·175	92·5	509
	0·25	92	358
	0·35	91	259

intermediate symmetries. In particular, we consider the breaking pattern (Mahapatra and Misra 1984), given as,

$$\begin{aligned} \mathrm{SO(10)} \otimes \mathrm{SUSY} &\to \mathrm{SU(2)}_L \otimes \mathrm{SU(2)}_R \times \mathrm{SU(4)}_C \otimes \mathrm{SUSY} \\ &\to \mathrm{SU(2)}_L \otimes \mathrm{U(1)}_Y \otimes \mathrm{SU(3)}_C \otimes \mathrm{SUSY} \\ &\to \mathrm{U(1)}_Q \otimes \mathrm{SU(3)}_c \end{aligned} \tag{15}$$

To achieve the above mentioned breaking pattern one may take for example, the Higgs multiplets $S\{54\}, \Gamma\{126\}\bar{\Gamma}\{\overline{126}\}, H\{10\}, H'\{10'\}$ and the singlet $U\{1\}$. As usual like the previous cases, through a minimisation of the scalar potential one can obtain the specific pattern (15). Here also supergravity correction becomes responsible for the weak symmetry breaking, and the presence of supersymmetry maintains the stability of gauge hierarchy including that of the intermediate scale.

In all the above cases, we note that the small parameter, which relates the supergravity with weak symmetry breaking is given by $\xi'^2 = k^2m^2 \simeq 10^{-16}$. Here the mass scale m is associated with the Polonyi singlet Z as mentioned in (11). In fact the supergravity correction constitutes an expansion in this small parameter $km = \xi' \simeq 10^{-8}$. Hence it is thought desirable to examine whether the corrections of different orders can be used for the symmetry breaking at the intermediate scale as well as for the weak symmetry breaking. This aspect has recently been analysed (Mishra and Misra 1984) and it is observed that supergravity correction can in fact reproduce multiple energy scales in a natural way.

Another pertinent attempt for gravity-induced weak symmetry breaking consists of writing the low-energy potential with explicit soft-supersymmetry breaking terms which can arise through supergravity (Barbieri *et al* 1982; Ibanez 1982, Hall *et al* 1983). This method usually does not assume any specific grand-unification symmetry. The scalar potential in such a model is given as, with M as the Planck mass,

$$\begin{aligned} V = \exp\left(\frac{|Z_i|^2 + |Y_a|^2}{M^2}\right)\Bigg[&|h_i + (Z_i^*/M^2)\tilde{g}|^2 + \left|g_a + \frac{Y_a^*}{M^2}\tilde{g}\right|^2 \\ &- \frac{3}{M^2}|\tilde{g}|^2 \Bigg] + \frac{1}{2}D_\alpha^2, \end{aligned} \tag{16}$$

which is the equivalent of expression of (9). We shall assume that the scales in $g(Y_a)$ ($\equiv g_1$, in (10)) are small compared to M. Hence we have to take the approximation (in the so-called flat limit) as $M \to \infty$ keeping $m_{3/2}$ (mass of the gravitino) fixed. Low-energy effective potential is then obtained by replacing Y_a''s by their VEVs, and taking the limit $M \to \infty$. One then obtains the scalar potential with explicit soft super-symmetry breaking terms, given as (Nilles 1984),

$$V = |\hat{g}_a|^2 + m_{3/2}^2 |Y_a|^2 + m_{3/2}[Y_a \hat{g}_a + (A-3)\hat{g} + \text{h.c.}], \tag{17}$$

where g is the low-energy superpotential. This potential is also related to the interaction at higher scales through renormalisation effects (Ibanez and Lopez 1984; Inoue *et al* 1982). This analysis also enables us to obtain weak symmetry breaking in supergravity models, through radiative corrections (Pran Nath 1983), introducing a negative mass square for the Higgs field. This feature will be illustrated below.

In the tree-level breaking one needs the gauge singlet field U to break the $SU(2) \otimes U(1)$ invariance. But now we shall show that in the absence of this singlet, weak symmetry can be broken radiatively via renormalisation effects. We shall demonstrate this feature in a simple toy model involving only two Higgs doublets and the usual matter fields. In this model one takes the potential as (Pran Nath 1983).

$$U = U_M + U_H \tag{18}$$

where U_H is the Higgs sector potential, given as

$$\begin{aligned} U_H = {} & m_H^2 \bar{H}^\alpha H^\alpha + m_{H'}^2 \bar{H}'_\alpha H'_\alpha + m_3 m_g (H'_\alpha H^\alpha + \text{h.c.}) \\ & + \frac{1}{8} g_1^2 (\bar{H}^\alpha H^\alpha - \bar{H}'_\alpha H'_\alpha)^2 + \frac{1}{8} g_2^2 (\bar{H}\tau_i H - \bar{H}'\tau_i H')^2 \end{aligned} \tag{19a}$$

and

$$\begin{aligned} U_M = {} & h_t^2 |\tilde{t}_R \tilde{t}_L|^2 + h_t^2 |\tilde{t}_R H|^2 + m_g A_t h_t (\tilde{t}_R H \tilde{t}_L + \text{h.c.}) \\ & + m_{\tilde{t}_R}^2 |\tilde{t}_R|^2 + m_{\tilde{t}_L}^2 |\tilde{t}_L|^2. \end{aligned} \tag{19b}$$

Here in the matter sector, we have retained only the top-quark Yukawa coupling. This is because the top-quark mass is much higher as compared to the masses of the other quarks and hence the corresponding Yukawa coupling h_t dominates in U_M. One then obtains the renormalisation group equations (RNGE) for the Higgs masses, and the right-handed and left-handed stop masses, as,

$$\begin{aligned} \frac{dm_H^2}{dt} &= \left(\frac{3\alpha_t}{2\pi}\right)\left[m_H^2 + m_{\tilde{t}_R}^2 + m_{\tilde{t}_R}^2 + A_t^2 m_g^2\right] \\ \frac{dm_{\tilde{t}_R}^2}{dt} &= \left(\frac{2\alpha_t}{2\pi}\right)\left[m_H^2 + m_{\tilde{t}_R}^2 + m_{\tilde{t}_L}^2 + A_t^2 m_g^2\right] \\ \frac{dm_{\tilde{t}_L}^2}{dt} &= \left(\frac{\alpha_t}{2\pi}\right)\left[m_H^2 + m_{\tilde{t}_R}^2 + m_{\tilde{t}_L}^2 + A_t^2 m_g^2\right]. \end{aligned} \tag{20}$$

In the above, $\alpha_t = (h_t^2/4\pi)$; $t = \ln(\mu/M)$ with $m_H \leqslant \mu \leqslant M$. We also have the RNGE for the coupling constants given as

$$\frac{d\alpha_t}{dt} = \left(\frac{3\alpha_t^2}{\pi}\right) - \left(\frac{\alpha_t}{\pi}\right)\left[\frac{8}{3}\alpha_3 + \frac{3}{2}\alpha_2 + \frac{13}{18}\alpha_1\right], \tag{21a}$$

$$\frac{dA_t}{dt} = \left(\frac{3\alpha_t}{\pi}\right)A_t. \tag{21b}$$

At $t = 0$, the appropriate boundary conditions are given as, $\alpha_t = \alpha_0$, $A_t = A_0$ and $m_H^2 = m_{\tilde{t}_R}^2 = m_{\tilde{t}_L}^2 = m_g^2$. Ignoring the gauge coupling constants, one then obtains the solutions for scalar masses as,

$$m_H^2 = \frac{3m_g^2}{2(1-\xi)}\left[1 - \frac{A_0^2\xi}{3(1-\xi)}\right] - \frac{1}{2}m_g^2, \tag{22a}$$

$$m_{\tilde{t}_R}^2 = \frac{2m_g^2}{2(1-\xi)}\left[1 - \frac{A_0^2\xi}{3(1-\xi)}\right], \tag{22b}$$

$$m_{\tilde{t}_L}^2 = \frac{m_g^2}{2(1-\xi)}\left[1 - \frac{A_0^2\xi}{3(1-\xi)}\right] + \frac{1}{2}m_g^2, \tag{22c}$$

where $\xi = (3\alpha_0 t/\pi)$. Thus for a large enough $-\xi > 0$, m_H^2 is found to be negative, implying a spontaneous gauge symmetry breaking. In the above toy model, using the fact that the symmetry breaking here takes place at the weak scale, one can obtain the top-quark mass $m_t = h_t v$ as around 100 GeV. There are alternative and refined (Nilles 1984) versions of a similar nature which will permit the top-quark mass to be as small as we like, even of the order of 25 GeV.

Another recent analysis has been made by Roy and Majumdar (1984) which links the mass of the Higgs scalar with the top quark mass. In fact they obtained that under reasonable assumptions, the mass of top-quark lying between 25 and 55 GeV may imply Higgs mass to be probably half of the Z-boson mass or less.

We shall next consider possible experimental signatures of the supersymmetric particles (Ellis and Sher 1984; Reya and Roy 1984). These may be produced in e^+e^- annihilation, or lepton-hadron deep inelastic collisions or in $p\bar{p}$ collider. The production cross-sections of these superpartners are roughly known, since they have the same coupling strengths as ordinary particles. These cross-sections do obviously also depend on the masses of the superpartners which are unknown. Hence, to discover that such particles are produced, we must have some qualitative guidelines regarding their possible decay modes. This is supplied by the existence of R-parity, which remains conserved. Ordinary particles have even R-parity and the superpartners have *odd* R-parity. Thus when a superpartner of relatively high mass is produced, it should finally yield a low mass stable superpartner. One expects this to be either the photino or scalar neutrino. Thus to recognise the production of these particles in any reaction we might expect missing energy that cannot be explained from normal sources like neutrino production. For example, one could expect a scalar quark decaying into a photino and a quark jet. One will then have a monojet structure with missing energy. One can also see gluino to decay to a photino and a gluon jet, where the gluon jet may be found to be rather broad.

Ellis and Sher (1984) have analysed the monojet events in the $p\bar{p}$ collider where they observed that if these events are caused by squarks, then the photino mass is expected to be around 8 GeV, the gluino mass is around 50 GeV, and what is most interesting, the *selectron* may have a mass around 20 GeV. The last mentioned particle may be observed in high energy e^+e^- collider experiments. The above results are based on the assumption that the top-quark mass is around 40 GeV.

In the above we have considered the phenomenological success of supergravity theories, where, gravitational effects are retained in the lowest order. We shall now consider certain demerits of this framework and some parallel alternative attempts which appear to be very attractive. The major difficulty of supergravity is that it gives rise to a theory which is not renormalizable. The Lagrangian in fact is not even a polynomial. Within the last one or two years another type of local supersymmetry theories of extended objects known as superstrings (Green and Schwartz 1983) have been considered. These theories lead to at best cubic and quartic couplings of field operators which are always renormalisable, and, often finite. In fact the theory appears to include an infinite number of supermultiplets having the mass gap of the order of Planck scale. The lowest lying levels are identified with the observed particles. The theory appears to be highly restrictive in the sense that there is very little arbitrariness in the construction of the interaction or the gauge group. However, since it is in higher dimensions it involves Kaluza-Klein compactification (Weinberg 1984) of a 10-dimensional space as well as spontaneous symmetry breaking of probably SO(32) or $E_8 \times E_8$ as the gauge group (Green and Schwartz 1984). The subject matter of local supersymmetry investigated in this manner is still under active consideration as an incomplete theory, and, may as well become the successful theory of tomorrow.

Acknowledgement

I would like to thank Snigdha Mishra for discussions while preparing the manuscript.

References

Barbieri R, Ferrara S and Savoy C A 1982 *Phys. Lett.* **B119** 343
Chamseddine A, Arnowitt R and Pran Nath 1982 *Phys. Rev. Lett.* **49** 970
Coleman S and Mandula J 1967 *Phys. Rev.* **159** 1251
Cremmer E *et al* 1983 *Nucl. Phys.* **B212** 413
Ellis J and Sher M 1984 CERN-TH. 3968
Green M and Schwartz J 1983 *Nucl. Phys.* **B218** 43
Green M and Schwartz J 1984 *Phys. Lett. B* (to be published)
Haber H E and Kane G L 1984 *Phys. Rep.* (to be published)
Hall L, Lykken J and Weinberg S 1983 *Phys. Rev.* **D27** 2359
Ibanez L E 1982 *Phys. Lett.* **B118** 73
Ibanez L E and Lopez C 1984 *Nucl. Phys.* **B233** 511
Inoue K *et al* 1982 *Prog. Theor. Phys.* **67** 1889
Maalampi J and Pulido J 1983 *Nucl. Phys.* **B228** 242
Mahapatra S, Mishra S and Misra S P 1984 *Phys. Rev.* D30 2019
Mahapatra S, Misra S P 1984 *Phys. Rev. D* (to be published)
Mishra S and Misra S P 1984 (to be submitted)
Nilles H P 1984 *Phys. Rep.* **110** 1
Polonyi J 1977 University of Budapest Rep. No. KFKI-1977-93 (unpublished)
Pran Nath 1983 *Lectures at summer workshop on particle physics, Trieste*
Reya E and Roy D P 1984 DO-TH. 84/11
Roy P and Majumdar P 1984 *Phys. Rev.* **D30** 2432
Salam A and Strathdee J 1974 *Nucl. Phys.* **B76** 477
Weinberg S 1984 UTTG-12-84
Wess J and Bagger J 1981 Princeton lectures (Parts I & II)
Wess J and Zumino B 1974a *Nucl. Phys.* **B70** 39
Wess J and Zumino B 1974b *Nucl. Phys.* **B78** 1