PRESENT STATUS OF SUPERSYMMETRY

L. O'Raifeartaigh

School of Theoretical Physics

Dublin Institute for Advanced Studies

Dublin, Ireland.

The principles and historical development of the recently developed theory of supersymmetry are reviewed. The present status of the theory (its remarkably elegant properties, experimental problems and outlook for the future) are then discussed in a qualitative manner.

The last two or three years have seen the emergence of a remarkable new symmetry called supersymmetry. This symmetry has the property of allowing fields of different intrinsic spin (and, in particular, fermions and bosons) to appear in the same irreducible multiplets. For this reason a more accurate name might be spin-mixing symmetry.

The basic idea of supersymmetry was first put forward by Ramond[1] in 1971 in the context of the string theory of the dual model for strong interactions. The idea was then developed, and superfields introduced, in the context of the string theory[2] by Iwasaki and Kikkawa, Gervais and Sahita and Neveu and Schwartz all independently. It was then realized by Wess and Zumino[3] in 1973 that the algebra of supersymmetry could be taken from the string theory context, where it operated in 1 + 1 dimensions, and set in the context of conventional field theory in 3 + 1 dimensions. Finally in 1974 the Wess-Zumino theory was formulated in terms of superfields in 3 + 1 dimensions by Salam and Strathdee.

Although supersymmetry appeared in this way only in recent years, it could, in principle, have appeared at any time since 1928. This is because the basic idea of supersymmetry consists in carrying Dirac's factorization of the d'Alembertian operator one step further, and factorizing the Dirac operator itself. Thus, where Dirac succeeded in constructing operators $\gamma^\mu p_\mu = \not{p}$ which satisfied the algebra

$$\{\not{p}, \not{p}\} = 2p^2 , \qquad [p^2, \not{p}] = 0 , \tag{1}$$

supersymmetry succeeds in constructing operators G_α which satisfy the algebra

$$\{G_\alpha, G_\beta\} = (C\not{p})_{\alpha\beta} , \qquad [P_\mu, G_\alpha] = 0 , \tag{2}$$

where $\alpha, \beta = 1 \ldots 4$ are Dirac indices. More precisely, supersymmetry shows that there exist representations of the algebra (2). Furthermore, one finds that for at least two of the representations, renormalizable supersymmetric invariant Lagrangians (Lagrangians invariant with respect to the algebra (2)) can be constructed, and these Lagrangians have very remarkable properties.

Representations of Supersymmetry.

The representations of the algebra (2) are most conveniently defined in terms of superfields. The simplest kind of superfield, corresponding to one of the two types of representation for which a renormalizable supersymmetric Lagrangian has been constructed, is a field $\Phi(x\,\theta)$, which depends not only on the space-time coordinates x, but also on a very special kind of internal variable θ. This variable transforms with respect to the Lorentz group like a real Dirac (Majorana) 4-spinor, and satisfies the anti-commutation relations

$$\{\theta_\alpha, \theta_\beta\} = 0 \tag{3}$$

Note the similarity between (3) and (2). However, the right hand side of (3) is zero, so that the variables θ are independent of P_μ, and are nilpotent. As a consequence of the nilpotency, the expansion of $\Phi(x\,\theta)$ in powers of θ terminates and one obtains

$$\Phi(x\theta) = A(x) + \bar{\theta}\psi(x) + \tfrac{1}{4}\bar{\theta}\left[F + \gamma_5 G + i\gamma_\mu\gamma_5 V_\mu\right]\theta + \tfrac{1}{4}(\bar{\theta}\theta)\,\bar{\theta}\lambda(x) + \tfrac{1}{32}(\bar{\theta}\theta)^2 D(x)\,, \tag{4}$$

where $A(x)$, $\psi(x)$ etc. are conventional fields. For simplicity we shall assume that $\Phi(x\,\theta)$ is real and is a Lorentz scalar

$$\Phi^*(x\theta) = \Phi(x\theta)\,, \qquad \left(U(\Lambda)\Phi\right)(x\theta) = \Phi(\Lambda' x, S'\theta)\,, \tag{5}$$

where S is the real Dirac (Majorana) four-dimensional representation of $SL(2,C)$. The conventional fields then have the appropriate Lorentz transformation properties; $A(x)$, $F(x)$, $D(x)$ scalars, $G(x)$ pseudo-scalar, $\psi(x)$, $\lambda(x)$ Majorana spinors and $V_\mu(x)$ vector.

To obtain the representation of the algebra (2) on $\Phi(x\theta)$ one now defines the supersymmetric transformations

$$\left(U(\varepsilon)\Phi\right)(x_\mu, \theta) = \Phi\left(x_\mu + \tfrac{i}{2}\bar{\varepsilon}\gamma_\mu\theta,\ \theta + \varepsilon\right) \tag{6}$$

where the ε are variables similar to the θ and anti-commuting with them. Then if we let G_α denote the infinitesimal generators of the transformations (6) we find that

$$G_\alpha = \frac{\partial}{\partial \theta_\alpha} + \frac{i}{2}\left(\bar{\theta}\gamma^\mu \partial_\mu\right)_\alpha \tag{7}$$

and it is trivial to verify that these G_α satisfy the algebra (2). Note that the supersymmetric transformations (6) are direct translations in θ-space, and induce the formal translations $x \to x + \frac{i}{2}\bar{\varepsilon}\gamma\theta$ in x-space. The quantities $x + \frac{i}{2}\bar{\varepsilon}\gamma\theta$ are to be understood in the sense of Taylor expansions of the conventional fields

$$A\left(x + \frac{i}{2}\bar{\varepsilon}\gamma\theta\right) = A(x) + \frac{i}{2}\bar{\varepsilon}\gamma\theta\, A(x)_{,\mu} + \dots \tag{8}$$

where the expansion terminates on account of the nilpotency of θ and ε. From (7), (5) and (8) one can obtain the direct action of the G_α on the conventional fields and it is

$$\begin{aligned}
\delta A &= \psi \cdot \delta\theta \\
\delta\psi &= \tfrac{1}{2}(F + \gamma_5 G + i\gamma_5 \not{V})\delta\theta && - \tfrac{i}{2}(\not{\partial} A)\delta\theta \\
\delta F &= \tfrac{1}{2}\bar{\chi}\cdot\delta\theta && - \tfrac{i}{2}\delta\bar{\theta}\not{\partial}\psi \\
\delta G &= \tfrac{1}{2}\bar{\chi}\gamma_5\delta\theta && - \tfrac{i}{2}\delta\bar{\theta}\gamma_5\not{\partial}\psi \\
\delta V_\mu &= \tfrac{i}{2}\bar{\chi}\gamma_\mu\gamma_5\delta\theta && - \tfrac{i}{2}\delta\bar{\theta}\, i\not{\partial}\gamma_5\gamma_\mu\psi \\
\delta\chi &= \tfrac{1}{2}D\,\delta\theta && - \tfrac{i}{2}(\not{\partial}F - \gamma_5\not{\partial}G + i\gamma_\nu\gamma_5 A_\nu) \\
\delta D &= && - i\,\delta\bar{\theta}\not{\partial}\chi
\end{aligned} \tag{9}$$

where

$$\delta X \equiv (G_\alpha X)\,\delta\theta_\alpha \quad \text{and} \quad \not{\partial} = \gamma^\mu\partial_\mu, \quad \bar{\psi} = \psi^\dagger\gamma_0 = \tilde{\psi}\gamma_0 .$$

Note that the $\frac{\partial}{\partial\theta_\alpha}$ part of G_α acts as a raising operator and the $(\bar\theta\not\partial)_\alpha$ part as a lowering operator, and that the lowering operation is always accompanied by a divergence.

The real scalar superfield just described is one of the representations of the supersymmetric algebra (2) for which a Lagrangian has been constructed. The second kind of superfield for which a Lagrangian has been constructed is obtained by making a "chiral-reduction" of $\Phi(x\theta)$. To make a chiral reduction one notes that

$$\{G_\alpha, \tilde{G}_\beta\} = 0, \qquad \tilde{G}_\alpha = \frac{\partial}{\partial\theta_\alpha} - \frac{i}{2}(\bar\theta\not\partial)_\alpha \tag{10}$$

that is, that the quantities $\tilde{G}_\alpha$, which are similar to the G_α but have a minus sign between the lowering and raising parts, commute with the G_α . This property of the $\tilde{G}_\alpha$ guarantees that the fields $\Psi_\pm(x\theta)$ which satisfy the conditions

$$\left[\left(\frac{1\pm\gamma_5}{2}\right)\tilde{G}\ \right]_\alpha \Psi_\pm(x\theta) = 0, \tag{11}$$

are again superfields. (The chiral projections $(1\pm\gamma_5)/2$ are used because the condition $\tilde{G}\Psi = 0$ would be too strong and would kill the superfield.) Fields $\Psi_\pm(x\theta)$ satisfying (11) are called chiral scalar superfields. The parity-invariant combination $\Psi(x\theta) = \Psi_+(x\theta) + \Psi_-(x\theta)$ is then a smaller field than $\Phi(x\theta)$ and is thus a reduction of it. It has the same number of formal components as $\Phi(x\theta)$, but the fields are inter-related. For example, for $\Phi(x\theta)$ the spinor fields $\psi(x)$ and $\lambda(x)$ are independent, but for $\Psi(x\theta)$, $\lambda(x) = \not\partial\psi(x)$.

Before going on to discuss the supersymmetric Lagrangians for the scalar and chiral-scalar superfields, there is one point that should be mentioned. That is, in the above we have assumed that θ is independent of x , and one might ask what happens if we let θ be x -dependent. The answer is that if we let θ be linear in x , $\theta_\alpha(x) = \overset{\circ}{\theta}_\alpha + \theta_\alpha^\mu x_\mu$, then instead of obtaining the generators of the translation group on the right hand side of (2) one obtains the whole conformal group. For arbitrary χ -dependence, one obtains the infinite-dimensional Einstein group. The conformal case was actually the case first considered by Wess and Zumino[3]. An interesting feature in the conformal case is that the unrestricted linear transformations generated by $\varepsilon_\alpha(x) = \mathring{\varepsilon}_\alpha + \varepsilon_\alpha^\mu x_\mu$ in θ -space, generate the restricted bilinear conformal transformations $x_\mu \to x_\mu + \frac{1}{2}\bar{\varepsilon}(x)\gamma_\mu \theta(x)$ in x -space. Thus θ -space appears to be a more basic space for the conformal group.

Supersymmetric Lagrangians.

The construction of renormalizable supersymmetric invariant Lagrangians for is based on two observations:

(1) The product of two scalar superfields is again a scalar superfield;

(2) The supersymmetric variation of the coefficient of $(\bar{\theta}\theta)^2$ in a superfield is a pure divergence.

The first observation follows from the nilpotency of θ , which makes the expansion terminate at $(\theta\bar{\theta})^2$ no matter how many products of superfields are taken, and the fact that the Lorentz and supersymmetric transformations (5) and (6) are carried by the arguments of the fields and hence are the same for products of superfields. The second observation follows from the fact that the lowering operation $((\theta\not{\partial})_\alpha$ part of $G_\alpha)$ is always a pure divergence, while for the coefficient of $(\theta\bar{\theta})^2$ the raising operation $((\partial/\partial\theta)_\alpha$ part of $G_\alpha)$ vanishes.

The procedure for forming Lagrangians is then simple in principle (in practice it can become quite complicated): To construct a kinetic or mass-term for a Lagrangian one takes a product of two superfields, expands it in powers of θ and takes the coefficient of $(\bar{\theta}\theta)^2$ as the Lagrangian density. To construct an interaction one takes a product of three superfields and repeats the process. (A product of more than three will also lead to a Lagrangian but it will not be renormalizable.) Since the supersymmetric variation of the Lagrangian densities constructed in the above way will then be a pure divergence, the Lagrangian itself will be supersymmetric invariant. Similar procedures hold for chiral scalar superfields.

Let us now simply describe what emerges from the procedures just described for the superfields $\Phi(x\theta)$ and $\Psi_{\pm}(x\theta)$. Initially $\Psi_{\pm}(x\theta)$ has the expansion

$$\Psi_{\pm}(x\theta) = A_{\pm}(x) + \bar{\theta}\cdot\psi_{\pm}(x) + \tfrac{i}{4}\bar{\theta}\gamma_{\mu}\gamma_5\theta A_{\pm}(x)_{,\mu} + \ldots \tag{12}$$

where the terms omitted are, like the third term, dependent on $A_{\pm}$ and $\psi_{\pm}$, where A_{+} is a complex scalar and ψ_{+} is the positive chiral projection of a Majorana spinor. The $\Psi_{\pm}(x\theta)$ together where $\Psi_{+}^{*} = \Psi_{-}$, have essentially the content

$$\Psi_{\pm}(x\theta) = \{A(x), B(x), \psi(x)\} = \{0^{+}, 0^{-}, \tfrac{1}{2}^{i}\}, \tag{13}$$

where $A_{\pm} = A \pm iB$, and the bracket on the right denotes the spin and parity. (The parity of a Majorana spinor is necessarily pure imaginary.) The Lagrangian which emerges from the supersymmetric procedure for $\Psi_{\pm}(x\theta)$ is[3]

$$\mathcal{L} = \mathcal{L}_{k.e.} + \frac{m}{2}\bar{\psi}\psi + g\bar{\psi}(A+i\gamma_5 B)\psi + \frac{1}{2}F^{*}F, \tag{14}$$

where $\mathcal{L}_{k.e.}$ is a conventional kinetic energy term for $\{A, B, \psi\}$ and

$$F = mA_{+} + gA_{+}^{2}. \tag{15}$$

Note that the potential energy $F^{*}F/2$ is positive indefinite and zero at $A_{+} = 0$. What is so special about this Lagrangian? At first sight it appears to be a conventional Yukawa-$P(\phi^4)$ interaction for $\{A, B, \psi\}$. What is special about it is that, whereas the most general parity-invariant Yukawa-$P(\phi^4)$ Lagrangian would have eight independent parameters (three masses, three boson coupling constants and three fermion coupling constants) the Lagrangian (14) has only two independent parameters, m and g. Thus there is a huge reduction in the number of independent parameters. But why go through all this machinery in order to reduce the number of parameters? Why not arbitrarily set parameters equal in the conventional eight-parameter Lagrangian? The answer is that arbitrary relations among the parameters will not be stable, i.e. will not be maintained after renormalization. Because the relations in the supersymmetric case are derived from a symmetry which is respected by the interaction, the relationships implied by the symmetry are expected to be stable. A calculation, first of the one-loop corrections, and then to all orders, shows that these expectations are indeed realized for the Lagrangian (14).

For the scalar superfield $\Phi(x,\theta)$ the spin-parity content is at first sight larger, namely,

$$\Phi(x,\theta) = \{ A, \psi, F, G, V_\mu, \lambda, D \} \tag{16}$$

However, it turns out that the Lagrangian is such that D is a function of the other fields, and if we have gauge as well as supersymmetric invariance (supergauge invariance) then a supergauge can be chosen so that the fields A, ψ, F, G vanish[5]. Thus in the supergauge invariant case the essential content of is

$$\Phi(x\,\theta) \simeq \{ V_\mu, \lambda \} = \{ 1^-, \tfrac{1}{2}^{i} \} . \tag{17}$$

Thus a super-gauge field $V_\mu(x)$ comes accompanied by a Majorana spinor field $\lambda(x)$. Correspondingly the Lagrangian for $\Phi(x\,\theta)$ is relatively simple, namely,[5]

$$\mathcal{L} = -\frac{1}{4}\,\mathrm{Tr}\,F_{\mu\nu}F^{\mu\nu} - \frac{i}{2}\,\mathrm{Tr}\,\bar{\lambda}\,[\not{D},\lambda] , \tag{18}$$

which is just an ordinary Yang-Mills Lagrangian for $V_\mu(x)$ and $\lambda(x)$. Even in the special supergauge $A = \psi = F = G = 0$ the Lagrangian (18) still retains, of course, the conventional gauge invariance under $V_\mu \to S^{-1} V_\mu S + S^{-1}\partial_\mu S$. Finally one can let the Yang-Mills superfield $\Phi(x,\theta) \simeq \{V_\mu, \lambda\}$ interact with the chiral scalar superfield. One then obtains a Lagrangian of the form[5]

$$\mathcal{L} = -\frac{1}{4} \operatorname{tr} F_{\mu\nu} F^{\mu\nu} - \frac{i}{2} \operatorname{tr} \bar{\lambda} [\not{D}, \lambda]$$

$$- \frac{1}{2} \bar{\psi} \not{\partial} \psi - \frac{1}{2} (\mathcal{D}_\mu A_+)^\dagger (\mathcal{D}_\mu A_+) \tag{19}$$

$$+ e \{ i \bar{\lambda}_\alpha (A^\dagger + i\gamma_5 B^\dagger) \tau_\alpha \psi \} + h.c. \} + \frac{e^2}{2} \{ A^\dagger \tau_\alpha B + h.c. \}^2$$

$$+ \frac{m}{2} \bar{\psi} \psi + g \bar{\psi} (A + i\gamma_5 B) \psi + \frac{1}{2} F^\dagger F \ .$$

where $\Psi_\pm(x,\theta) \simeq \{A, B, \psi\}$ is a multiplet of chiral scalar superfields, $A_+ = A + iB$, $\mathcal{D}_\mu = \partial_\mu + ieV_\mu$, and τ_α are the group generators.

Balance-Sheet.

Having described the basic principles, we must now consider the advantages and disadvantages of supersymmetry. First let us consider the credit side.

Apart from the general elegance and beauty of this symmetry, there are five specific areas in which it is advantageous, or at least interesting. The areas can be specified as follows:

(i) Renormalization

(ii) Spontaneous Symmetry Breaking

(iii) Yang-Mills Theory (Unified Gauge Theory)

(iv) Asymptotic Freedom

(v) Mixing of Lorentz and Internal Symmetry.

We discuss these briefly in turn.

(i) The renormalization properties constitute perhaps the most striking feature of supersymmetry. Cancellations occur at almost every step so that the actual renormalization is reduced to a minimum. For example, for the chiral scalar field Lagrangian (14) there are no quadratic and linear divergences, and there is only <u>one</u> logarithmic divergence[3][6]. Furthermore, this logarithmic divergence serves as the renormalization constant for all the masses and all the wave-functions, both fermions and bosons. (Thus Kallen's prediction that at least one of the renormalization constants must be infinite is just barely fulfilled!) Another result is that for supersymmetry the vacuum expectation value of any n-point function is zero[7]. This means in particular that the Lagrangian need not be normal-ordered.

(ii) There are two kinds of spontaneous symmetry breaking that one can consider, namely the spontaneous breaking of supersymmetry itself, and the spontaneous breaking of internal symmetry by supersymmetry. One finds that the spontaneous breakdown of supersymmetry is a relatively rare occurrence[8][9], whereas the spontaneous breakdown of an internal symmetry, triggered by supersymmetry, happens frequently [4][9]. Thus with respect to spontaneous symmetry breaking, supersymmetry resembles a diamond, which itself is hard to cut, but which serves to cut glass.

(iii) There is a natural way to combine Yang-Mills gauge invariance and supersymmetry into a supergauge invariant theory, and indeed the required supergauge Lagrangian is just that given above in (18) and (19). Furthermore, if the internal symmetry of this Lagrangian is spontaneously broken (which frequently happens as discussed in (ii)) we may, by proper choice of group and representation[4][9], pick up masses for all but an abelian set of the Yang-Mills fields. In that case we obtain an infra-red convergent supergauge theory, or in other words, a supersymmetric unified gauge theory. This result is not trivial because, unlike the conventional Yang-Mills theory, supergauge theory completely determines the <u>form</u> of the potential (see equation (14)) leaving only the choice of group and representation free.

(iv) It is now well-known that, in contrast to abelian fields, a Yang-Mills field, either in self-interaction or in interaction with fermions, is asymptotically free[10]. That is, the renormalized coupling constant tends to zero as the scale parameter tends to infinity. However, it is also known that, in general, a Yang-Mills field in interaction with boson fields is not, in general, asymptotically free[10]. The reason is that renormalizability requires the introduction of a second (quartic) coupling constant for the boson field, and the latter constant is not, in general, asymptotically free. It turns out that the renormalization properties of supersymmetry are such that this problem goes away, and a supersymmetric Yang-Mills field in interaction with either one or two super matter fields (which necessarily include bosons as well as fermions) is asymptotically free[11].

v) Finally, supersymmetry can be used to obtain a nontrivial mixing of Lorentz and internal symmetry[4][9]. The idea is very simple. Given a superfield $\Phi(x\,\theta)$ there are two ways to introduce an internal symmetry as follows:

Trivial (direct product) way: $\Phi(x\,\theta) \to \Phi^a(x\theta)$

Non-trivial way: $\Phi(x\,\theta) \to \Phi(x\,\theta^a)$

where a is the internal symmetry index. In the second case the expansion of $\Phi(x\,\theta)$ becomes

$$\Phi(x\,\theta^a) = A(x) + \bar{\theta}^a\,\Psi_a(x) + \ldots \tag{20}$$

so that the fields of different spin have different internal spin also. Furthermore, the relativistic and supersymmetric transformation laws remain compatible,

$$\left(U(\Lambda)\Phi\right)(x\,\theta_a) = \Phi\left(\Lambda^{-1}x,\, S^{-1}\theta_a\right), \tag{21}$$

$$\left(U(\varepsilon_a)\,\Phi\right)(x_\mu\,\theta_a) = \Phi\left(x_\mu + \tfrac{i}{2}\,\bar{\varepsilon}_a\,\gamma_\mu\,\theta_a\,,\ \theta_a + \varepsilon_a\right). \tag{22}$$

This result establishes that, in principle at any rate, Lorentz and internal symmetry can be successfully combined. However, for the moment, the success remains at the level of principle, because it has been shown[12] that under rather general and plausible assumptions (20) is actually the most general combination of Lorentz and internal symmetry that can be constructed, and it so happens that the spin-isospin correlations obtained from (20) are not found experimentally. Note that (20) combines only the spin and the internal symmetry. There is no mass-breaking either for the supersymmetry or the internal symmetry.

Finally, we must come to the debit side of the balance sheet. The sad fact is that in spite of its intrinsic beauty, supersymmetry has not yet found any useful application. There are three basic reasons for this. The first is the one mentioned above, that the spin-isospin correlation for supersymmetry is far from experiment, but there are two more profound difficulties. These difficulties stem from the fact that supersymmetry forces fermions and bosons to behave in a similar manner. The first difficulty then is that the fermions and bosons in a supersymmetric multiplet have the same mass. Apart from the case of the neutrino and photon, which both have zero mass, this result is in manifest disagreement with experiment. There would be no great problem if we could have a spontaneous breakdown of supersymmetry, since the breakdown would allow different fermion and boson masses to emerge. But as we have mentioned above, a spontaneous breakdown of supersymmetry is a relatively rare event. The second difficulty is that either the fermion number is not conserved (or, more exactly, is conserved only modulo two) or the boson number is conserved. Neither of these alternatives agrees, of course, with experiment.

I should hasten to add that these difficulties are not completely insurmountable, in the sense that they have already been overcome in particular models[(4)(8)(9)(13)]. Indeed one might reverse the argument, as Iliopoulos has suggested, and use mass and particle-number breaking as criteria for selecting models. However, so far no systematic way of overcoming these two difficulties has emerged, and until it does, or until some of the particular models which overcome these difficulties become realistic in other respects, the experimental identification of superfields will present a serious difficulty. Perhaps the immediate future of supersymmetry is to go on ice, as Yang-Mills theory did from 1954 to 1967, until one knows how and where to use it. Perhaps it will never become useful experimentally, but serve as a simple model, analogous to the Lee model[14], on which ideas can be tested. Even in that limited context, I think that the intrinsic beauty and elegance of supersymmetry will serve to keep interest in it alive for some time to come.

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