5435/81

IC/81/222
INTERNAL REPORT
(Limited distribution)

International Atomic Energy Agency
and
United Nations Educational Scientific and Cultural Organization
INTERNATIONAL CENTRE FOR THEORETICAL PHYSICS

PREONS AND SUPERSYMMETRY

(To honour Francis Low's sixtieth birthday)

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MIRAMARE - TRIESTE
November 1981

1. An important aspect of preonic theories is the construction of composite fields and the commutation relations [1] among them, using preonic fields (with their canonical commutation relations) as input. We wish to remark that superfields appear ideally suited for playing the role of preonic fields. The basis for this is the remarkable group property possessed by chiral superfields.

In this note we shall assume that supersymmetry holds for preonic fields and that it is broken just below the ionization energy (possibly 10^5 GeV or higher) for the formation of quarks and leptons as preonic composites [2].

2. The preonic theory we choose to illustrate these remarks with is the theory of three preon types [2]: f (flavons), c (chromons) and the singlet s (henceforth called the "drone"). These are (2,1,1); (1,2,1); (1,1,4) and (1,1,1) representations of $SU_L(2) \times SU_R(2) \times SU_C(4)$. In addition each one of the preons carries $U(1) \times U(1)$ quantum numbers which permit of their binding into quarks and leptons into composites. The quarks and leptons themselves are neutral relative to these U(1)'s. The difficult problem in supersymmetry theories always is the breaking of this symmetry. By using these U(1)'s we shall show that it is possible to break supersymmetry as well as SU(4) (and some of the U(1)'s) simultaneously. Thus if the scale of $SU_C(4)$ spontaneous breaking is of the order [3] of 10^4-10^5 GeV, this could also be the scale of spontaneous breaking of supersymmetry, in contrast to [4] other recent attempts which break supersymmetry either at the $SU(3) \times SU(2) \times U(1)$ level of 300-1000 GeV, or assume that it is broken only at Planck energies.

3. Let Φ_- and Φ_+ represent left- [5] and right-handed chiral superfields, each describing particles of spin $\frac{1}{2}$ and zero. The (Majorana) θ expansion of these fields is:

$$\Phi_\mp(x,\theta) = \frac{1}{4} \pm \frac{1}{4} (\bar{\theta}\not{\partial}\gamma_5\theta)\left[A_\mp(x) + \bar{\theta}\psi_\mp(x) + \frac{1}{2}\bar{\theta}\left(\frac{1 \mp i\gamma_5}{2}\right) F_\mp(x)\right] .$$

The chiral fields are annihilated by covariant operators $D_\mp$, i.e. $D_-\Phi_+ = D_+\Phi_- = 0$, where $D_\pm = \frac{1 \pm i\gamma_5}{2}\left(\frac{\partial}{\partial\bar{\theta}} - \frac{i}{2}\not{\partial}\theta\right)$. Note that Φ_-^* behaves like Φ_+ so far as chirality is concerned. In particular the operation D_-D_- on Φ_- gives rise to a field of the plus type Φ_+'.

The crucial property for our purposes is the group property

$$\Phi_-(x,\theta)\,\Phi_-'(x,\theta) = \Phi_-''(x,\theta)$$

and likewise for the fields $\Phi_+(x,\theta)$. Thus any product-field created by multiplying any number of − (or +) type of preonic chiral fields leads to a − (or +) type of composite. Chiral spin $\frac{1}{2}$ and spin zero preons composed supersymmetrically in this manner do not give rise to any spin-one composites.

Spin-one composites can of course be constructed by multiplying preonic superfields of opposite chiralities. Thus a product field of two preonic fields Φ_+ and Φ_- gives rise to a general superfield $\Phi(x,\theta)$ describing spin zero, spin $\frac{1}{2}$ as well as spin-one objects. Such general superfields can be decomposed through chiral plus "transverse vector" projections as follows:

Define

$$E_+ = -\frac{1}{4\partial^2}(D_-D_-)(D_+D_+)$$

$$E_- = -\frac{1}{4\partial^2}(D_+D_+)(D_-D_-)$$

$$E_1 = 1 - E_+ - E_-$$

$$(E_\pm^2 = E_\pm)$$

then

$$\Phi(x,\theta) = (E_- + E_+ + E_1)\,\Phi(x,\theta)$$

$$= \Phi_- + \Phi_+ + \Phi_1 ,$$

where

$$\Phi_1(x,\theta) = A_1(x) + \bar{\theta}\psi_1(x) + \frac{1}{4}\bar{\theta}\, i\gamma_\nu\gamma_5\theta\, A_{\nu 1}(x)$$

$$+ \frac{1}{4}\bar{\theta}\theta\,\bar{\theta}\,(i\not{\partial}\psi_1) + \frac{1}{32}(\bar{\theta}\theta)^2\,\partial^2 A_1 .$$

The field Φ_1 contains spin zero $(A_1(x))$, spin $\frac{1}{2}$ $(\psi_1(x)$, Majorana, $\psi = C^{-1}\bar{\psi}^T)$, and spin-one $A_{\nu 1}(x)$ pieces.

4. When dealing with symmetry groups of $SU_L(n) \times SU_R(n)$ variety, cancellation of anomalies often requires the introduction of additional mirror fields [6]. These are fields with the same transformation character under SU(n) as the original fields but carrying opposite chirality. Thus given a $\Phi_-(x)$ preonic field, a composite mirror field Φ_+^M can be constructed by using the singlet drone field s_- in the following manner:

$$\Phi_+^M = D_- D_- (\Phi_- s_-) \quad .$$

This construction is of course not unique; in fact all composites $D_- D_- (\Phi_- s_-^r)$ are possible mirror fields, as indeed are the fields $E_-(\Phi_- s_-^{*r})$. If the singlets s_- carry U(1) quantum numbers for binding to Φ_- - as in practice they will (see later) - the composite mirror fields will be distinguished from each other through their U(1) labels.

Another distinguishing label is provided by f, the $\mathbb{F}$-number [7] (associated with the operator $i\bar{\theta}\gamma_5 \frac{\partial}{\partial\theta}$) which may be defined for a general superfield as a combination of an intrinsic gauge transformation $e^{if\alpha}$ combined with the transformation $\theta \to e^{-\alpha\gamma_5}\theta$. Thus let,

$$\Phi(x,\theta) \to \Phi'(x,\theta) = e^{if\alpha}\, \Phi(x, e^{\alpha\gamma_5}\theta) \quad .$$

Expanding $\Phi(x,\theta)$ into components

$$\Phi(x,\theta) = A(x) + \bar{\theta}\psi(x) - \frac{1}{4}\bar{\theta}\theta\, F(x) + \frac{1}{4}\bar{\theta}\gamma_5\theta\, G(x)$$

$$+ \frac{1}{4}\bar{\theta}i\gamma_\nu\gamma_5\theta\, V_\nu(x) + \frac{1}{4}\bar{\theta}\theta\,\bar{\theta}\chi(x) + \frac{1}{32}(\bar{\theta}\theta)^2\, D(x)$$

we can read off $\mathbb{F}$-numbers associated with the components:

$$\begin{aligned} \mathbb{F} &= f && \text{for } A,\ V_\nu,\ D \\ &= f+1 && \text{for } \psi_-,\ \chi_+ \\ &= f-1 && \text{for } \psi_+,\ \chi_- \\ &= f+2 && \text{for } F+iG \\ &= f-2 && \text{for } F-iG \quad . \end{aligned}$$

(Here $\psi_\mp = \frac{1 \mp i\gamma_5}{2}\psi$ and likewise for $\chi_\mp$.) From this we read that for a chiral field $\Phi_-(x)$, carrying intrinsic $\mathbb{F}$-number f, the $\mathbb{F}$ numbers associated with the components A_-, ψ_- and F_- are f, f+1, f+2. For $\Phi_+(x)$ with intrinsic $\mathbb{F}$-number f, the components A_+, ψ_+, F_+ carry f, f-1 and f-2, respectively, while the covariant operators D_- and D_+ add $\mathbb{F}$-numbers -1 and +1 to the fields they act upon.

To take an example, assume Φ_- carries intrinsic $\mathbb{F}$-number f, while the singlet s_- carries f_s. Then the mirror composite defined as $D_- D_-(\Phi_- s_-)$ has intrinsic $\mathbb{F}$-number $f + f_s - 2$ and the alternative mirror $E_-(\Phi_- s_-^*)$ carries $f - f_s$. The F-number assignments differentiate the two types of fields.

All gauge Lagrangians conserve $\mathbb{F}$-number provided the gauge field carries intrinsic $\mathbb{F} = 0$, with the component assignments, $A_\nu, D = 0$. The projection χ_- of the Majorana gaugino must be assigned F-number -1 for the conservation to hold. The renormalizable matter Lagrangians $(\Phi_- \Phi_-')_F = D_- D_-(\Phi_- \Phi_-')$ and $D_- D_-(\Phi_- \Phi_-' \Phi_-'')$ conserve $\mathbb{F}$-number provided (in an obvious notation) $f_- + f_-' - 2 = 0$ in the first and $f_- + f_-' + f_-'' - 2 = 0$ for the second term. Likewise for matter Lagrangians $D_+ D_+(\Phi_+ \Phi_+')$ and $D_+ D_+(\Phi_+ \Phi_+' \Phi_+'')$, the corresponding requirements are $f_+ + f_+' + 2 = 0$ and $f_+ + f_+' + f_+'' + 2 = 0$.

To summarize [8], if gauge particles A_ν, D carry $\mathbb{F} = 0$, while the gauginos $\chi_\mp$ carry $\mathbb{F} = \mp 1$ ($\chi_- = C\bar{\chi}_+^T$, Majorana condition) $\mathbb{F}$-number is conserved for gauge Lagrangians. For matter fields Φ_- and Φ_+, assign intrinsic $\mathbb{F}$ numbers $f_\mp$. The components $A_\mp$, $\psi_\mp$, $F_\mp$ then carry $f_\mp$, $f_\mp \pm 1$, $f_\mp \pm 2$. For conservation of $\mathbb{F}$ number in pure matter renormalizable interactions, e.g. $(\Phi_- \Phi_-' \Phi_-'')$ we need to satisfy conditions like $f_- + f_-' + f_-'' - 2 = 0$.

5. To make quarks and leptons $[(2,1,4)_L$ and $(1,2,\bar{4})_L]$ out of preons: [9]

$$\begin{aligned} f_- &= (2,1,1) \\ f_-' &= (1,2,1) \\ c_- &= (1,1,4) \\ c_-' &= (1,1,\bar{4}) \\ s_- &= (1,1,1) \end{aligned}$$

we need for example the composites $(f_- c_- s_-^r)$ and $(f_-' c_-' s_-^{r'})$ where r and r' are arbitrary positive numbers, so far as $SU_L(2) \times SU_R(2) \times SU_C(4)$ transformations are concerned. We should additionally assign $U(1) \times U(1) \times \cdots$ quantum numbers to the f's, the c's and the s's such that the attractive U(1) forces bind them into quarks and leptons $(2,1,4)$ and $(1,2,\bar{4})$ with the further proviso that the composite quarks and leptons are also U(1) neutral [10]. These particular requirements are easily met: what is difficult to achieve is the orderly (spontaneous) breaking of supersymmetry as well as of internal symmetries.

To show that this may be done in principle, we demonstrate that it is possible to assign the U(1)'s in such a manner that the four-colour baryon-lepton symmetry $SU_C(4)$ and supersymmetry are broken simultaneously by the same set of Higgs, possibly at energies around 10^4-10^5 GeV. To show this we follow a procedure [10] devised by Weinberg to break $SU(3) \times SU(2) \times U(1)$ symmetry and supersymmetry together. Weinberg employs the Fayet-Iliopoulos mechanism [11], with an extra $\widetilde{U}(1)$ together with two super-Higgs multiplets. The quarks and leptons remain massless, even though other particles acquire masses.

6. From henceforth we ignore $SU_L(2) \times SU_R(2)$ and the flavons. The chromons are:

$$c_- \sim 4_{0,x}$$
$$c_-' \sim \bar{4}_{0,x} \quad ,$$

where $(0,x)$ specify $U_Y(1) \times U_{Y'}(1)$ quantum numbers. We take two Higgs multiplets with quantum numbers

$$H_- \sim 4_{1,1}$$
$$H_-' \sim 4_{-1,1}$$

and a singlet

$$s_- \sim 1_{0,-2} \quad .$$

In this model electric charge is $Q = T_4^4 + Y$ which (we show) will remain unbroken. For $U_{Y'}(1)$ anomaly cancellations we need mirrors, but these will be ignored. The gauge Lagrangian plus the matter terms $h(H_-' H_- S_-)_F$ give rise to the following potential for the scalar components of the appropriate multiplets:

$$V = \frac{G^2}{2} \left[(\bar{H}H)^2 - 2\,|H'H|^2 + |H'\bar{H}'|^2 - \frac{1}{4}(\bar{H}H - H'\bar{H}')^2\right]$$
$$+ \frac{g^2}{2}\left[\bar{H}H - H'\bar{H}' - \xi'\right]^2$$
$$+ \frac{g'^2}{2}\left[\bar{H}H + H'\bar{H}' - 2\,|S|^2 - \xi''\right]^2$$
$$+ h^2\left[(\bar{H}H + H'\bar{H}')\,|S|^2 + |H'H|^2\right] \quad .$$

Here G, g and g' are $SU_C(4)$, $U_Y(1)$, $U_{Y'}(1)$ gauge coupling parameters. After some work, the potential minimizes, with $SU(4) \times U_Y(1) \times U_{Y'}(Y)$ breaking down to $SU_C(3) \times U_Q(1)$ with

$$\langle H\rangle = \begin{pmatrix} 0 \\ 0 \\ 0 \\ a \end{pmatrix} \qquad \langle H'\rangle = \begin{pmatrix} 0 \\ 0 \\ 0 \\ a' \end{pmatrix} \quad \langle S\rangle = 0 \quad ,$$

where

$$0 = 2g'^2\,(|a|^2 + |a'|^2 - \xi'') + h^2\,(|a|^2 + |a'|^2)$$
$$0 = \frac{3}{4}G^2(|a|^2 - |a'|^2) + g^2\,(|a|^2 - |a'|^2 - \xi')$$
$$+ g'^2\,(|a|^2 + |a'|^2 - \xi'') + h^2\,|a'|^2 \quad .$$

The quadratic terms in V giving masses for the 3 and $\bar{3}$ scalar components of H and H' are:

$$(G^2 - h^2)\,(a'\underset{\sim}{\bar{H}} - a\underset{\sim}{H}')\cdot(a'\underset{\sim}{H} - a\underset{\sim}{\bar{H}}') \quad .$$

These terms are positive definite provided we impose the requirement $G^2 > h^2$. The quadratic (mass) terms in V which mix Re H_4, Re H_4', S_1 = Re S, S_2 = Re S read:

$$(\tfrac{3}{4}G^2 + 2g^2 + 2g'^2)\,(a\,(\mathrm{Re}\,H)^2 + a'(\mathrm{Re}\,H')^2)$$

$$-(\tfrac{3}{2}G^2 + 2g^2 - 2g'^2 - 2h^2)\,2aa'\,(\mathrm{Re}\,H)\,(\mathrm{Re}\,H') + 2h^2\,(a^2 + a'^2)\,(S_1^2 + S_2^2)$$

One can check that positivity is assured since

$$(2g'^2 + h^2) > 0 \quad .$$

The Goldstone fields Im H and Im H' are massless.

Finally the inclusion of the chromons $c_{0,x}$ and $c'_{0,x}$ gives for mass terms of their scalar components:

$$\begin{aligned} = \ & |c_4|^2 \left(\frac{3}{4}(a^2 - a'^2) + x(a^2 + a'^2 - \xi'')\right) \\ & + |\bar{c}_4|^2 \left(-\frac{3}{4}(a^2 - a'^2) + x(a^2 + a'^2 - \xi'')\right) \\ & + |\underset{\sim}{c}|^2 \left(-\frac{1}{4}(a^2 - a'^2) + x(a^2 + a'^2 - \xi'')\right) \\ & + |\underset{\sim}{\bar{c}}'|^2 \left(\frac{1}{4}(a^2 - a'^2) + x(a^2 + a'^2 - \xi'')\right) \quad . \end{aligned}$$

Clearly x can be chosen such that the positivity of these terms is ensured. The spin $\frac{1}{2}$ components of the chromons are desirably massless at this stage. There are Goldstinos but we do not discuss them in this note.

To summarize we have demonstrated that it is possible to break supersymmetry as well as an internal symmetry (like $SU_C(4) \times U_Y(1) \times U_{Y'}(1)$ down to $SU_C(3) \times U_Q(1)$) with the same set of Higgs, while the spin $\frac{1}{2}$ chromons remain massless. The lesson of the calculation above lies for us in its stressing yet again the role of the U(1)'s. These appear as a necessary feature for simultaneous supersymmetry and internal symmetry breaking, in addition to being needed for providing forces to bind preons together. In a recent set of papers [12], it was suggested that such U(1)'s may possibly be associated with (analogue) electric and magnetic monopole charges for preons. In a further paper we propose to consider the role of supersymmetry in the context of such a preonic theory.

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In this note we work with general values for intrinsic F-number, $f_\mp$ for chiral fields $\Phi_\pm$. Thus F-number need not be tied to the fermion number of the spin $\frac{1}{2}$ components of Φ_- and Φ_+.

8) The conventional parity operation does not commute with the F operation as defined. For Abelian or non-Abelian gauge Lagrangians containing matter multiplets of Φ_- as well as Φ_+ type (and with F number assignments for matter and gauge fields as in the text) a parity operation can be defined for mixtures of scalar and spinor components, in a manner so as to conserve both parity and the F number. In Ref.5 and Ref.7 however, we preserved the commutativity of F-number and the parity operation by working with N = 2 extended supersymmetry, where the gauge field V is supplemented with a chiral multiplet s_+ (in the adjoint representation of the internal symmetry) such that the gauginos are 4-component Dirac - rather than Majorana - particles. In Ref.7 N = 2 extended supersymmetry was called by us complex supersymmetry.

9) If one is economy-minded, the preons c'_- could themselves be considered composites: for example $E_-(c^*_- s^r_-)$ so far as the $SU(2) \times SU(2) \times SU(4)$ quantum numbers are concerned. In the appropriate range of energies where both c_- and c'_- are considered elementary and structureless, we shall of course have to ensure that the U(1) and other quantum numbers of c_- and c'_- match; see Sec.6. For example, the f-label of c'_- constructed as above (with r = 1) is $f_s - f_-$. This would equal the f-label of c_- only if $f_- = \frac{1}{2} f_s$.

10) S. Weinberg, HUTP-81/A047 (1981), has proposed to relax this restriction: quarks and leptons may carry a further $\tilde{U}(1)$ quantum number beyond the quantum numbers of $SU_c(3) \times SU(2) \times U(1)$.

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Note added to Int.Rep. IC/81/222

One may inquire how far back the notion of pre-pre-...preons may be carried? Can we eventually arrive at a single "monotheistic" chiral supermultiplet s_- carrying none other than U(1) charges? In such an approach gauge groups like $SU_L(2) \times SU_R(2) \times SU_C(4)$ (as well as the vector mesons associated with them) are considered as arising at successive composite levels, the renormalizable interaction at each level being the leading part of an effective interaction based on an expansion in powers of the radii of the relevant composites.

Clearly a U(1) gauging of a single chiral multiplet s_- will give rise to non-renormalizable anomalies; thus s_- would need to be supplemented with a mirror supermultiplet s_+. Alternatively, one may conceive of a non-gauge Yukawa (renormalizable) coupling s_-^3 and attribute the next level of composites to binding by this force. This is not an attractive suggestion; however, if the possibility of binding exists, there may arise at the second level, the mirror composites needed for a further U(1) gauging. These are

$$\phi_\mp^{(p)} = E_\mp(s_-^p(\bar{s}_-s_-)^r) \approx E_\mp(s_-^{p+r}\bar{s}_-^r),\; r > 0,\; p \geq 0 .$$ 13)

Disregarding any problems connected with the non-locality of the chiral projection operators $E_\mp$, $\phi_\mp^{(p)}$ provides us with a pair of mirror fields (for each value of the U(1) charge p). An anomaly free gauge theory with a U(1) composite gauge vector multiplet may now be motivated at this level.

At the next level, starting with one pair $\phi_-^{(p)}$ and $\phi_+^{(p)}$ new bound state composites comprising three pairs of mirror fields (each carrying U(1) charges of magnitude pp') may arise. These are

$$(\phi_-^{(p)})^{p'+r'}\;(\overline{\phi_+^{(p)}})^{r'} \text{ and its mirror } (\phi_+^{(p)})^{p'+r'}\;(\overline{\phi_-^{(p)}})^{r'} \qquad \text{(A)}$$

and

$$E_\pm[(\phi_-^{(p)})^{p'+r'}\;(\bar{\phi}_-^{(p)})^{r'}],\; E_\pm[(\phi_+^{(p)})^{p'+r'}\;(\bar{\phi}_+^{(p)})^{r'}] \qquad . \qquad \text{(B)}$$

Assuming that there is a mass degeneracy among the four particles comprised in set (B), there is the possibility that in addition to the U(1) gauge, of the level before, there also exist the composite gauges $SU_L(2) \times SU_R(2)$, operative below the dissociation energy of the composites of set (B).

In fact, if supersymmetry were broken at this stage, such that the fermions and the bosons, contained in these four supermultiplets of set (B) were not supersymmetrically degenerate, the non-supersymmetric composite gauge group could be as large as $SU_L(2) \times SU_R(2) \times SU_C(4)$. The gauge particles associated with spin-zero bosons - which we have called $SU_C(4)$ - would necessarily be pure vectors (rather than axial vectors). In this approach the distinction between colour and flavour quantum numbers would be a consequence of supersymmetry breaking. Quarks and leptons would now form as non-supersymmetric U(1)-neutral composites of these preons at the next level. Such a model is of course different from the model considered in the earlier sections of this note, in that the level at which supersymmetry breaks is even prior to the emergence of $SU_L(2) \times SU_R(2) \times SU_C(4)$.

The labels r and r' may be construed as generation labels, as has been suggested by a number of authors. To illustrate, assume that we start with $p = 1$ ϕ_-, ϕ_+ pre-preons. The composites $E_\pm(\phi_-^{p'+r'}\;\phi_-^{r'})$ and $E_\pm(\phi_+^{p'+r'}\;\phi_+^{r'})$ are formed as a consequence of the interplay of $(p' + r')r'$ mutually attractive and $\frac{(p' + r')(p' + r' - 1)}{2} + \frac{r'(r' - 1)}{2}$ repulsive U(1) forces among the constituents of the composites. Very naively these forces balance when $r' = \frac{p'(p' - 1)}{2}$. Thus given p', the number of distinct generations r' may be limited by the relations $r' \leq \frac{p'(p' - 1)}{2}$.

One may ask the question: why supersymmetry in the first place for the "monotheistic" supermultiplet s_-? Our motivation for this - or rather for the stage when one works with s_- and its mirror s_+ - has been in the context of the (analogue) dual Abelian electric and magnetic $U_E(1) \times U_M(1)$ theory of Ref.12. Disregarding supersymmetry, in such a theory, the "natural" pair of pre-preons would appear to be an electrically charged object (e,0) and a dual magnetically charged object (0,g), at the first level. The preons (from which quarks and leptons are made) could then be the composites (e,g) of these, created through the electromagnetic forces $U_E(1) \times U_M(1)$. Assuming that the pre-preons (e,0), (0,g) are spin-zero objects, the preons (e,g) would carry the field spin $\left|\frac{eg}{4\pi}\right| = \frac{N}{2}$ (N integer).

Now whatever the value of N, the neutral preon-anti-preon composites (which make up quarks and leptons) cannot carry any except integer (or zero) spins. This is because a neutral composite made from (e,g) and (-e,-g) can have no field spin. In order that quarks and leptons do manifest half-integral spin, the set of preons or pre-preons must contain objects carrying both integer (or zero) as well as half-integer spins. This appears to motivate supersymmetry at the basic preonic or pre-preonic level. The implementation of this idea will need a super-symmetrization of the dual electric and magnetic theory.

Footnote 13

One must remember that there is a limiting relationship implied in the construction of composite fields; for example, $\phi^{(1)}(x)$ is defined through the spacelike limit $x_1 \to x_2 \to x_3 \to x$ $(s(x_1)\ (\bar{s}(x_2)\ s(x_3)))$. There is an arbitrariness in the taking of this limit mathematically reflected in the order in which the x_1, x_2, x_3 approach each other and physically representing the distinction of whether or not the $(\bar{s}(x_2), s(x_3))$ composite forms first. In this note these possibilities have not been exploited.