PHYSICAL REVIEW D **90,** 015033 (2014)

Brout-Englert-Higgs boson as spin-0 partner of the Z in the supersymmetric standard model

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(Received 8 June 2014; published 28 July 2014)

Supersymmetric extensions of the standard model lead to gauge/Brout-Englert-Higgs (BEH) unification by providing spin-0 bosons as *extra states for spin-1 gauge bosons* within massive gauge multiplets. They may be described by the spin-0 components of *massive gauge superfields* (instead of chiral superfields as usual). In particular, the 125 GeV/c^2 boson observed at CERN, considered as a BEH boson associated with electroweak breaking and mass generation, may also be interpreted, up to a mixing angle induced by supersymmetry breaking, as *the spin-0 partner of the Z* under *two* supersymmetry transformations, i.e., as a Z that would be deprived of its spin.

DOI: 10.1103/PhysRevD.90.015033 PACS numbers: 12.60.Jv, 12.60.Fr, 14.80.Ly, 14.80.Da

I. INTRODUCTION

Supersymmetric extensions of the standard model lead to *superpartners* for all particles, squarks and sleptons, gluinos, charginos and neutralinos, etc. [1–5]. They differ from ordinary particles by 1/2 unit of spin and are distinguished by an *R-parity* quantum number related to baryon and lepton numbers, discrete remnant of a continuous $U(1)_R$ symmetry, making the lightest superpartner stable.

While the standard model [6,7] involves a single scalar doublet leading to one Brout-Englert-Higgs (BEH) boson [8–11], spontaneous electroweak breaking is induced here by *two doublets* h_1 and h_2. They are responsible for charged-lepton and down-quark masses, and up-quark masses, respectively, leading to additional *charged and neutral spin-0 BEH bosons*. These theories also provide systematic associations between massive gauge bosons and spin-0 BEH bosons, a very nontrivial feature owing to *their different gauge symmetry properties* [1,2,12].

These relations were proposed in 1974 even before the standard model (SM) was considered as "standard," and are at the basis of its supersymmetric extensions, even if they may often go unnoticed. Weak neutral currents were just recently discovered [13] with their structure unknown, and the $W^\pm$ and Z hypothetical. Little attention was paid to fundamental spin-0 particles, the very possibility of their existence getting questioned and frequently denied for many years later.

Proposing relations between massive spin-1 mediators of weak interactions and spin-0 particles associated with electroweak breaking and mass generation [1,12] then amounted to

relate two classes of hypothetical particles,
using an hypothetical symmetry! (1)

And this was at a time when supersymmetry was viewed as an algebraic structure [14–17] very far from being able to describe Nature, for many reasons including an obvious lack of similarities between known bosons and fermions.

Forty years later, the situation has improved considerably. With the introduction of R-odd superpartners and two spin-0 doublets for the electroweak breaking, supersymmetry could indeed be a symmetry of the fundamental laws of physics [2]. The discoveries in 1983 of the $W^\pm$ and Z mediators of weak interactions [18,19], and in 2012 of a new boson considered as a spin-0 BEH boson [20,21] confirmed to a very large extent the validity of the standard model (SM) theory or of a closely approaching one. This gives additional interest to the relations between spin-1 and spin-0 bosons provided by supersymmetric extensions of the standard model.

These relations may be more concretely discussed now that we know, with the Z and h bosons, at least one representative in each class of formerly hypothetical particles. The supersymmetric Standard Model offers a way to view the 125 GeV/c^2 boson recently observed at CERN as *a spin-0 partner of the Z*, up to a mixing angle induced by supersymmetry breaking.

With the two doublets h_1 and h_2 leading to five spin-0 particles, one neutral and two charged ones may be viewed as extra spin-0 states of the Z and $W^\pm$, two neutral ones [and possibly others beyond the minimal supersymmetric standard model (MSSM)] staying unmatched with massive gauge bosons [22].

II. THE SPIN-0 z PARTNER OF THE Z

Within supersymmetry two spin-0 doublets h_1 and h_2 are needed for the electroweak breaking, at first to avoid a massless chiral chargino, allowing for the construction of *two Dirac winos* associated with the $W^\pm$ within a massive gauge multiplet of supersymmetry [1]. These doublets $h_1 = (h_1^0, h_1^-)$ and $h_2 = (h_2^+, h_2^0)$ have weak hypercharges $Y = -1$ and $+1$. By leading to a negative mass^2 term for

h_2, the term $-\xi D'$ [23] associated with $U(1)_Y$ in the Lagrangian density plays a crucial role in triggering spontaneous electroweak breaking and giving masses to the $W^\pm$ and Z, and to their spin-$\frac{1}{2}$ and spin-0 partners.

The auxiliary components $\boldsymbol{D}$ and D' associated with the $SU(2)\times U(1)_Y$ gauge group are expressed as

$$\boldsymbol{D} = -\frac{g}{2}(h_1^\dagger \boldsymbol{\tau} h_1 + h_2^\dagger \boldsymbol{\tau} h_2) + \cdots,$$
$$D' = \xi + \frac{g'}{2}(h_1^\dagger h_1 - h_2^\dagger h_2) + \cdots. \quad (2)$$

The resulting potential reads [1]

$$V = \frac{1}{2}(\boldsymbol{D}^2 + D'^2) + \cdots = \frac{g^2}{8}(h_1^\dagger \boldsymbol{\tau} h_1 + h_2^\dagger \boldsymbol{\tau} h_2 + \cdots)^2 + \frac{1}{2}\left[\xi + \frac{g'}{2}(h_1^\dagger h_1 - h_2^\dagger h_2) + \cdots\right]^2 + \cdots, \quad (3)$$

ignoring for the moment possible soft supersymmetry-breaking terms, considered at a later stage. Its quartic part, fixed by the electroweak couplings g and g' as

$$V_{\text{quartic}} = \frac{g^2+g'^2}{8}(h_1^\dagger h_1 - h_2^\dagger h_2)^2 + \frac{g^2}{2}|h_1^\dagger h_2|^2, \quad (4)$$

appears within supersymmetry as part of electroweak gauge interactions.

The potential is minimum for $\langle h_1^0\rangle = v_1/\sqrt{2}$, $\langle h_2^0\rangle = v_2/\sqrt{2}$. The correspondence with the notations of [1], using two doublets φ'' and φ' with the same $Y=-1$, is as follows:

$$h_1 = \varphi'' = \begin{pmatrix} h_1^0 = \varphi''^0 \\ h_1^- = \varphi''^- \end{pmatrix}, \qquad h_2 = \begin{pmatrix} h_2^+ = -\varphi'^+ \\ h_2^0 = \varphi'^{0*} \end{pmatrix},$$
$$\tan\beta = v_2/v_1 \equiv \tan\delta = v'/v''. \quad (5)$$

The $\mu H_1 H_2$ superpotential term is first taken to vanish, as initially forbidden by a continuous $U(1)_R$ and/or an extra $U(1)_A$ symmetry [1]. The latter acts according to $\varphi'' \to e^{i\alpha}\varphi''$, $\varphi' \to e^{-i\alpha}\varphi'$ as introduced in a pre-SUSY two-doublet model in [24], i.e.,

$$h_1 \to e^{i\alpha} h_1, \qquad h_2 \to e^{i\alpha} h_2, \quad (6)$$

allowing one to rotate the phases of the two doublets independently. Taking $\mu = 0$ in this first stage allows for gauge symmetry to be spontaneously broken with supersymmetry remaining conserved in the neutral sector, shedding light on the relations between massive gauge bosons and spin-0 BE-Higgs bosons provided by supersymmetry.

The initial $U(1)_R$ symmetry survives the electroweak breaking induced by $\langle h_1\rangle$ and $\langle h_2\rangle$. As long as it is present it allows us to benefit, in the absence of a μ term and of direct gaugino mass terms, from Dirac neutralinos as well as charginos, and more specifically, *two Dirac winos and a Dirac zino*, carrying ± 1 unit of the additive quantum number R.

Some attention may be useful in the presence of an extra $U(1)$ symmetry acting on h_1 and h_2 as in (6) [24], that became known later as a $U(1)_{PQ}$ symmetry. Indeed it might lead to a classically massless pseudoscalar A (and associated scalar s_A), jointly described by [1,2]

$$\varphi''^0 \sin\delta + \varphi'^{0*}\cos\delta = h_1^0 \sin\beta + h_2^0 \cos\beta = \frac{s_A + iA}{\sqrt{2}}. \quad (7)$$

These particles, momentarily appearing as classically massless in the spectrum [2], get a mass as in [1] through an explicit breaking of the $U(1)_A$ symmetry [25].

We see from (3) that the term $-\xi D'$ in $\mathcal{L}$ generates a negative mass^2 for h_2, triggering spontaneous electroweak breaking. The origin is a saddle point of the potential, with $m^2(h_2) = -\frac{\xi g'}{2} < 0$, $m^2(h_1) = \frac{\xi g'}{2} > 0$. The would-be spin-0 Goldstone field [with $\delta = \beta$ as indicated in (5)]

$$z_g = -\sqrt{2}\,\text{Im}(\varphi''^0 \cos\delta + \varphi'^0 \sin\delta) = \sqrt{2}\,\text{Im}(-h_1^0\cos\beta + h_2^0 \sin\beta) \quad (8)$$

is eliminated by the Z. The corresponding real part

$$z = \sqrt{2}\,\text{Re}(-\varphi''^0\cos\delta + \varphi'^0\sin\delta) = \sqrt{2}\,\text{Re}(-h_1^0\cos\beta + h_2^0\sin\beta) \quad (9)$$

describes *a scalar BEH boson associated with the Z under supersymmetry*, with the same mass m_Z as long as supersymmetry is unbroken [1,2,12].

This results in the general association

$$Z \overset{\text{SUSY}}{\longleftrightarrow} 2\,\text{Maj. zinos} \overset{\text{SUSY}}{\longleftrightarrow} \text{spin-0 BEH boson}\, z \quad (10)$$

with in this description

$$z = \sqrt{2}\,\text{Re}(-h_1^0\cos\beta + h_2^0\sin\beta). \quad (11)$$

This is also made possible by the $U(1)_R$ symmetry remaining unbroken at this stage, allowing for the two Majorana zinos to combine into a *Dirac zino* of mass m_Z. It implies the existence of a spin-0 BE-Higgs boson of mass

$$m \simeq 91\ \text{GeV}/c^2 \quad (12)$$

up to susy-breaking effects.

This result, valid *independently of* $\tan\beta$, may now be compared with the recent CERN discovery of a new boson with a mass close to 125 GeV/c^2 [20,21].

The spin-0 field z may also be compared with the SM-like BEH field

$$h_{\rm SM} = \sqrt{2}\,\mathrm{Re}(h_1^0 \cos\beta + h_2^0 \sin\beta). \tag{13}$$

This z has Yukawa couplings "of the wrong sign" to down quarks and charged leptons, acquiring their masses through Yukawa couplings to h_1 [2] [26]. It becomes very close to $h_{\rm SM}$ at large $\tan\beta$, with

$$\langle h_{\rm SM}|z\rangle = -\cos 2\beta. \tag{14}$$

We thus rotate neutral chiral superfields as indicated by (7)–(9), according to

$$\begin{cases} H_z = -H_1^0 c_\beta + H_2^0 s_\beta = (z + i z_g)/\sqrt{2} + \cdots, \\ H_A = H_1^0 s_\beta + H_2^0 c_\beta = (s_A + iA)/\sqrt{2} + \cdots. \end{cases} \tag{15}$$

H_z describes the would-be Goldstone field z_g and spin-0 z associated with the Z as in (8) and (11), while H_A describes the scalar and pseudoscalar

$$\begin{cases} s_A = \sqrt{2}\,\mathrm{Re}(h_1^0 s_\beta + h_2^0 c_\beta), \\ A = \sqrt{2}\,\mathrm{Im}(h_1^0 s_\beta + h_2^0 c_\beta), \end{cases} \tag{16}$$

discussed more later.

III. ELECTROWEAK BREAKING AND Z AND z MASSES

From

$$D_3 = \frac{g}{2}(-|h_1^0|^2 + |h_2^0|^2) + \cdots,$$
$$D' = \xi + \frac{g'}{2}(|h_1^0|^2 - |h_2^0|^2) + \cdots, \tag{17}$$

we get for $D_Z = D_3 c_\theta - D' s_\theta$, $D_\gamma = D_3 s_\theta + D' c_\theta$,

$$\begin{cases} D_Z = -\xi s_\theta + \dfrac{\sqrt{g^2 + g'^2}}{2}(|h_2^0|^2 - |h_1^0|^2) + \cdots, \\ D_\gamma = \xi c_\theta + 0 + \cdots. \end{cases} \tag{18}$$

We express $V = \frac{1}{2}(D_Z^2 + D_\gamma^2) + \cdots$ as a function of h_1^0 and h_2^0 as

$$V = \frac{1}{2}\left(-\xi s_\theta + \frac{\sqrt{g^2 + g'^2}}{2}(|h_2^0|^2 - |h_1^0|^2) + \cdots\right)^2 + \cdots. \tag{19}$$

Minimizing this term fixes only $v_2^2 - v_1^2$, leading to a flat direction associated with s_A. $v_2^2 - v_1^2$ adjusts so that

$$\langle D_Z\rangle = -\xi s_\theta + \frac{\sqrt{g^2 + g'^2}}{4}(v_2^2 - v_1^2) \equiv 0. \tag{20}$$

Expanding D_Z in (18) at first order in h_1^0 and h_2^0 we have

$$D_Z = \frac{1}{2}\sqrt{g^2 + g'^2}(-v_1\sqrt{2}\,\mathrm{Re}h_1^0 + v_2\sqrt{2}\,\mathrm{Re}h_2^0) + \cdots$$
$$= m_Z\sqrt{2}\,\mathrm{Re}(-h_1^0 c_\beta + h_2^0 s_\beta) + \cdots = m_Z z + \cdots, \tag{21}$$

providing from $D_Z^2/2 = \frac{1}{2} m_Z^2 z^2 + \cdots$ the supersymmetric mass term for the spin-0 field z.

The parameter ξ associated with $U(1)_Y$ determines m_Z, given by $m_Z^2 = (g^2 + g'^2)(v_1^2 + v_2^2)/4$. For $v_1 \simeq 0$ we would get for Z and z (described by $\simeq \sqrt{2}\,\mathrm{Re}h_2$)

$$m_Z^2 = m_z^2 \simeq -2m^2(h_2) = \xi g', \tag{22}$$

i.e.,

$$m_Z = m_z \simeq \sqrt{\xi g'}, \tag{23}$$

up to radiative corrections, and supersymmetry-breaking effects for m_z. With $g' = e/\cos\theta \simeq .345$ the Z, $W^\pm$ and spin-0 BEH boson masses get *fixed by the ξ parameter* associated with $U(1)_Y$. This leads to

$$\xi \approx \frac{m_Z^2}{g'} \simeq \frac{m_W m_Z}{e} \simeq 2.4 \times 10^4 \ \mathrm{GeV}^2, \tag{24}$$

or equivalently

$$\sqrt{\xi} \approx \frac{v}{2}\frac{1}{\sin\theta}\sqrt{\frac{e}{\cos\theta}} \approx 150 \ \mathrm{GeV}, \tag{25}$$

up to radiative corrections.

The ξ parameter [23] *determines here the $W^\pm$ and Z masses*, a feature that may further persist when R-odd squarks and sleptons acquire large mass², e.g., from the compactification of extra dimensions. More generally we have

$$(-\cos 2\beta) m_Z^2 = \xi g', \tag{26}$$

reducing to (23) for large $\tan\beta$. $\xi = 0$ would be associated with $\tan\beta = 1$, leaving at this stage m_Z and m_W unfixed, at the classical level [27]. In such a situation the scalar s_A associated with this flat direction would describe a classically massless particle with dilatonlike couplings.

We also have from (18) and (20)

$$\langle D_\gamma\rangle = \xi c_\theta = \frac{g}{4 s_\theta}(v_2^2 - v_1^2) = \frac{m_W^2}{e}(-\cos 2\beta). \tag{27}$$

Having at this stage the photino as the Goldstone fermion implies that charged particles only are sensitive to the spontaneous breaking of the supersymmetry. Neutral ones

TABLE I. Minimal content of the supersymmetric Standard Model. Gauginos λ', λ_3 mix with Higgsinos $\tilde{h}_1^\circ$, $\tilde{h}_2^\circ$ into a photino, two zinos, and a Higgsino, further mixed into four neutralinos. The charged $w^\pm$ associated with $W^\pm$ is usually known as $H^\pm$. The scalars (z, s_A) mix into h and H. The N/nMSSM also involves an extra singlet superfield S with a trilinear superpotential coupling $\lambda H_1 H_2 S$, leading to an additional neutralino (singlino) and two singlet bosons.

gluons photon	gluinos $\tilde{g}$ photino $\tilde{\gamma}$	
$W^\pm$ Z	winos $\tilde{W}_{1,2}^\pm$ zinos $\tilde{Z}_{1,2}$ Higgsino $\tilde{h}_A$	$w^\pm$, z, s_A, A } BE-Higgs bosons
	leptons l quarks q	sleptons $\tilde{l}$ squarks $\tilde{q}$

remain mass degenerate within massive (Z) or massless (γ) multiplets of supersymmetry, before the introduction of extra terms breaking the $U(1)_A$ and $U(1)_R$ symmetries, the latter reduced to R-parity.

We now discuss zinos, winos, and charged spin-0 bosons within massive gauge multiplets, as shown in Table I, before returning to spin-0 bosons, and how they may be described by *massive gauge superfields*, in contrast with the usual formalism.

IV. ZINOS AND OTHER NEUTRALINOS

The massive gauge multiplet of the Z [1] includes a Dirac zino, obtained from chiral gaugino and Higgsino components transforming under $U(1)_R$ according to

$$\text{gaugino}\, \lambda_Z \to e^{\gamma_5 \alpha} \lambda_Z, \qquad \text{Higgsino}\, \tilde{h}_z \to e^{-\gamma_5 \alpha} \tilde{h}_z. \tag{28}$$

It may be expressed as a massive Dirac zino with $R = +1$,

$$\lambda_{ZL} + (-\tilde{h}_z)_R = (\lambda_3 c_\theta - \lambda' s_\theta)_L + (\tilde{h}_1^0 c_\beta - \tilde{h}_2^0 s_\beta)_R. \tag{29}$$

Or equivalently as two Majorana zinos, degenerate as long as $U(1)_R$ is preserved, with a mass matrix given in the corresponding 2×2 gaugino-Higgsino basis by [28]

$$\mathcal{M}_{\text{zinos}} = \begin{pmatrix} 0 & m_Z \\ m_Z & 0 \end{pmatrix}. \tag{30}$$

Supersymmetry remains unbroken in this sector, in the absence of direct gaugino (m_1, m_2) and Higgsino (μ) mass terms.

This 2×2 zino mass matrix may be unpacked again into a 4×4 neutralino matrix expressed in the (λ', λ_3, $\tilde{h}_1^0$, $\tilde{h}_2^0$) basis using (29). Including additional $\Delta R = \pm 2$ supersymmetry-breaking contributions from gaugino (m_1, m_2) and Higgsino (μ) mass terms it reads

$$\mathcal{M}_{\text{inos}} = \begin{pmatrix} m_1 & 0 & -s_\theta c_\beta m_Z & s_\theta s_\beta m_Z \\ 0 & m_2 & c_\theta c_\beta m_Z & -c_\theta s_\beta m_Z \\ -s_\theta c_\beta m_Z & c_\theta c_\beta m_Z & 0 & -\mu \\ s_\theta s_\beta m_Z & -c_\theta s_\beta m_Z & -\mu & 0 \end{pmatrix}. \tag{31}$$

For equal gaugino masses $m_1 = m_2$ the photino $\lambda_\gamma = \lambda' c_\theta + \lambda_3 s_\theta$ is a mass eigenstate. The remaining 3×3 mass matrix is expressed in the (λ_Z, $\tilde{h}_1^0$, $\tilde{h}_2^0$) basis (with $\lambda_Z = -\lambda' s_\theta + \lambda_3 c_\theta$) as

$$\begin{pmatrix} m_2 & c_\beta m_Z & -s_\beta m_Z \\ c_\beta m_Z & 0 & -\mu \\ -s_\beta m_Z & -\mu & 0 \end{pmatrix}, \tag{32}$$

as seen from (29). It further simplifies for $\tan\beta = 1$ into

$$\begin{pmatrix} m_2 & m_Z/\sqrt{2} & -m_Z/\sqrt{2} \\ m_Z/\sqrt{2} & 0 & -\mu \\ -m_Z/\sqrt{2} & -\mu & 0 \end{pmatrix}. \tag{33}$$

Next to a pure Higgsino of mass $|\mu|$ corresponding to $(\gamma_5)(\tilde{h}_1^0 + \tilde{h}_2^0)/\sqrt{2}$, the two zinos constructed from λ_Z and $(\tilde{h}_1^0 - \tilde{h}_2^0)/\sqrt{2}$ have the mass matrix

$$\mathcal{M}_{\text{zinos}} = \begin{pmatrix} m_2 & m_Z \\ m_Z & \mu \end{pmatrix}, \tag{34}$$

as obtained directly from (30).

There may also be additional neutralinos, as described by the extra N/nMSSM singlet S with a $\lambda H_1 H_2 S$ superpotential coupling, leading through $\langle H_i \rangle = \frac{v_i}{\sqrt{2}}$ to a $\frac{\lambda v}{\sqrt{2}} H_A S$ superpotential mass term. Here $H_A = H_1^0 s_\beta + H_2^0 c_\beta$ is the same "left-over" chiral superfield as obtained in (15), now acquiring a mass by combining with S [1].

V. THE SPIN-0 $w^\pm$ PARTNER OF THE $W^\pm$, AND ASSOCIATED WINOS

We have, in a similar way,

$$W^\pm \overset{\text{SUSY}}{\longleftrightarrow} 2\,\text{Dirac winos} \overset{\text{SUSY}}{\longleftrightarrow} \text{spin-0 boson}\, w^\pm, \tag{35}$$

with $m_{w^\pm} = m_{W^\pm}$, also up to supersymmetry-breaking effects. This is why the charged boson now known as $H^\pm$ was called $w^\pm$ in [1].

The two doublets being expressed as $\varphi'' = (h_1^0, h_1^-)$ and $\varphi' = (h_2^{0*}, -h_2^-)$, as in (5) with $\delta = \beta$, the would-be Goldstone field

$$w_g^\pm = \varphi''^\pm \cos\delta + \varphi'^\pm \sin\delta = h_1^\pm \cos\beta - h_2^\pm \sin\beta \tag{36}$$

is eliminated by the $W^\pm$. The orthogonal combination

$$w^\pm = \varphi''^\pm \sin\delta - \varphi'^\pm \cos\delta = h_1^\pm \sin\beta + h_2^\pm \cos\beta \qquad (37)$$

(approaching $h_1^\pm$ at large $\tan\beta$) describes *a charged spin-0 BEH boson associated with the* $W^\pm$ [1,2,12].

With

$$h_1^\dagger h_2 = h_1^{0*} h_2^+ + h_1^+ h_2^0 = \frac{v}{\sqrt{2}}(h_1^+ \sin\beta + h_2^+ \cos\beta) + \cdots$$
$$= \frac{v}{\sqrt{2}} w^+ + \cdots, \qquad (38)$$

the quartic terms (4) in the potential,

$$V = \frac{g^2}{2}(h_1^\dagger h_2)^2 + \cdots = \frac{g^2 v^2}{4}|w^+|^2 + \cdots, \qquad (39)$$

generate a mass $m_w = m_W = gv/2$ for

$$w^\pm \equiv H^\pm = h_1^\pm \sin\beta + h_2^\pm \cos\beta. \qquad (40)$$

It is the same as for the $W^\pm$, to which it is related by *two* infinitesimal supersymmetry transformations.

The mass spectrum is given, at this first stage for which supersymmetry is spontaneously broken with the photino as the Goldstone fermion, by [1]

$$\begin{cases} m^2_{w^\pm} = m^2_{W^\pm} = \dfrac{g^2(v_1^2+v_2^2)}{4}, \\ m^2(\text{winos}_{1,2}) = \dfrac{g^2 v_{1,2}^2}{2} = m_W^2(1 \pm \cos 2\beta) \\ \qquad = m_W^2 \mp e\langle D_\gamma\rangle. \end{cases} \qquad (41)$$

Boson-fermion mass² splittings are given by $\pm e\langle D_\gamma\rangle$ (as in [4] in the absence of other sources of supersymmetry breaking), and fixed by (27).

The two Dirac winos are R eigenstates carrying $R = \pm 1$, with masses $gv_1/\sqrt{2}$ and $gv_2/\sqrt{2}$. The wino mass matrix would be supersymmetric [as for zinos in (30)] for $\xi = 0$ so that $\beta = \pi/4$ and m(winos) $= m_W$, with $\langle D_\gamma\rangle = \langle D_Z\rangle = 0$ from (27).

In the presence of additional $\Delta R = \pm 2$ gaugino and Higgsino mass terms further breaking the supersymmetry as well as $U(1)_R$ (for m_2 and μ) and $U(1)_A$ (for μ), the wino mass matrix obtained from (41) reads

$$\mathcal{M}_{\text{winos}} = \begin{pmatrix} m_2 & \frac{gv_2}{\sqrt{2}} = m_W\sqrt{2}s_\beta \\ \frac{gv_1}{\sqrt{2}} = m_W\sqrt{2}c_\beta & \mu \end{pmatrix}. \qquad (42)$$

m_2 and μ jointly allow for both winos to be heavier than m_W (as experimentally required [29]).

For gaugino and Higgsino mass terms related by $m_1 = m_2 = m_3 = -\mu$ (up to radiative corrections), possibly also equal to the gravitino mass $m_{3/2}$, with $\tan\beta = 1$ [30], we get from (34) and (42) remarkable mass relations like, at the classical level,

$$\begin{cases} m^2(\text{winos}) = m_W^2 + m_{3/2}^2, \\ m^2(\text{zinos}) = m_Z^2 + m_{3/2}^2, \\ m(\text{photino}) = m(\text{gluinos}) = m_{3/2}. \end{cases} \qquad (43)$$

This also paves the way for more general situations involving $N = 2$ extended supersymmetry with grand-unification groups [31,32]. Similar mass relations like

$$\begin{cases} m^2(\text{xinos}) = m_X^2 + m_{3/2}^2, \\ m^2(\text{yinos}) = m_Y^2 + m_{3/2}^2 = m_X^2 + m_W^2 + m_{3/2}^2 \end{cases} \qquad (44)$$

are then obtained for xinos, yinos, etc., with a grand-unification gauge group like $SU(5)$ or $O(10)$, ….

Extra compact dimensions may then be responsible for supersymmetry and grand-unification breakings [32], R-odd supersymmetric particles carrying momenta $\pm m_{3/2}$ along an extra dimension. When R-parity is identified with the action of performing a closed loop along such a compact dimension, $m_{3/2}$ and more generally superpartner masses get quantized in terms of its size, according to (43) and (44) with e.g., in the simplest case

$$m_{3/2} = (2n+1)\frac{\pi\hbar}{Lc} = (2n+1)\frac{\hbar}{2Rc}. \qquad (45)$$

But let us return, in a more conservative way, to four spacetime dimensions.

VI. THE PSEUDOSCALAR A AND SCALAR s_A

The potential (3) admits, at this initial stage excluding a $\mu H_1 H_2$ superpotential term [both $U(1)_R$ and $U(1)_A$ symmetries being present] two classically flat directions corresponding to the scalar s_A and pseudoscalar A in (7), both classically massless [2]. A then appears as an "axion" associated with the extra $U(1)_A$ symmetry acting on h_1 and h_2 as in (6) [24], extended to supersymmetry according to [1]

$$H_1 \to e^{i\alpha} H_1, \qquad H_2 \to e^{i\alpha} H_2. \qquad (46)$$

Its scalar partner s_A is also associated with a flat direction, the minimization of the potential (3) fixing only $v_2^2 - v_1^2$.

This "axion" A (a notion unknown at the time, that appeared in a different context several years later) and associated scalar s_A were given a mass in [1] by breaking explicitly the $U(1)_A$ symmetry (6), (46), now often referred to as $U(1)_{PQ}$. This was done by introducing a singlet S

coupled through a trilinear superpotential $\lambda H_1 H_2 S$, and transforming under $U(1)_A$ according to

$$S \to e^{-2i\alpha} S. \tag{47}$$

Its $f(S)$ superpotential interactions, that may include S, S^2, and S^3 terms as in the N/nMSSM, break explicitly $U(1)_A$, the presence of a quasimassless "axion" being avoided.

Explicitly, the potential includes an extra term V_λ, with a vanishing minimum still preserving the supersymmetry. It reads

$$\begin{aligned} V_\lambda &= \left|\frac{\partial \mathcal{W}}{\partial S}\right|^2 = |\lambda h_1 h_2 + \sigma + \cdots|^2 + \cdots \\ &= \frac{\lambda^2 v^2}{2} |\underbrace{h_1 s_\beta + h_2 c_\beta}_{h_A}|^2 + \cdots = \frac{1}{2}\frac{\lambda^2 v^2}{2}(s_A^2 + A^2) + \cdots . \end{aligned} \tag{48}$$

It provides a mass term $(\lambda v/\sqrt{2})$ for the complex field

$$h_A = \frac{s_A + iA}{\sqrt{2}} = h_1 s_\beta + h_2 c_\beta, \tag{49}$$

the would-be "axion" A (and associated scalar s_A) acquiring a mass $m_A = \lambda v/\sqrt{2}$ [1].

In terms of superfields, the $\lambda H_1 H_2 S$ superpotential coupling of the N/nMSSM generates in [1], from $\langle H_1 \rangle = v_1/\sqrt{2}$, $\langle H_2 \rangle = v_2/\sqrt{2}$,

$$\lambda H_1 H_2 S = \frac{\lambda v}{\sqrt{2}}(H_1 s_\beta + H_2 c_\beta) S + \cdots = \frac{\lambda v}{\sqrt{2}} H_A S + \cdots, \tag{50}$$

a supersymmetric mass term $\lambda v/\sqrt{2}$ for H_A and S, possibly to be combined with a $\frac{1}{2}\mu_S S^2$ singlet mass term, if present.

VII. z YUKAWA COUPLINGS "OF THE WRONG SIGN"

The new boson found at CERN close to 125 GeV$/c^2$ [20,21] is considered as a Brout-Englert-Higgs boson [8–11] associated with the electroweak breaking, as expected in the standard model [6,7] where this breaking involves a single spin-0 doublet. But it may also be interpreted, in general up to a mixing angle, as a spin-0 partner of the Z under *two* infinitesimal supersymmetry transformations. The z field in (11) may be compared with the SM-like scalar, obtained from the real part of the neutral component of the "active" doublet combination

$$\varphi_{\rm sm} = \varphi'' \cos\delta + \varphi' \sin\delta = h_1 \cos\beta + h_2^c \sin\beta, \tag{51}$$

such that $\langle \varphi^0_{\rm sm} \rangle = \frac{v}{\sqrt{2}}$, and

$$h_{\rm SM} = \sqrt{2}\,{\rm Re}(h_1^0 \cos\beta + h_2^0 \sin\beta) \tag{52}$$

as in (13). We have $\langle h_{\rm SM} | z \rangle = -\cos 2\beta$, the two fields getting close for large $\tan\beta$, with the z tending to behave very much as the SM-like $h_{\rm SM}$.

More precisely while $h_{\rm SM}$ has standard Yukawa couplings to quarks and charged leptons $m_{q,l}/v = 2^{1/4} G_F^{1/2}$ $m_{q,l}$, the z has almost-identical couplings

$$\frac{m_{q,l}}{v}\, 2T_{3q,l} = 2^{1/4} G_F^{1/2} m_{q,l}\, 2T_{3q,l}. \tag{53}$$

They simply differ by *a relative change of sign for d quarks and charged leptons* (with $2T_{3d,l} = -1$) acquiring their masses through $\langle h_1^0 \rangle$, as compared to u quarks.

This may also be understood by *deducing the scalar couplings of the spin-0 z from the axial couplings of the spin-1 Z*, as follows:

The Z is coupled, with coupling $\sqrt{g^2 + g'^2}$, to the weak neutral current $J_Z^\mu = J_3^\mu - \sin^2\theta J_{\rm em}^\mu$, with an *axial part* $J^\mu_{3\,\rm ax}$ fixed by $T_{3q,l}/2$. It gets its mass by eliminating the Goldstone field z_g, pseudoscalar partner of the scalar z. As seen from the global limit g, $g' \to 0$ for which the Z would become massless and behave like the spin-0 z_g, this z_g has *pseudoscalar* couplings to quarks and leptons given by

$$\begin{aligned} \sqrt{g^2 + g'^2}\, \frac{T_{3q,l}}{2}\, \frac{2m_{q,l}}{m_Z} &= \\ = \frac{m_{q,l}}{v}\, 2T_{3q,l} &= 2^{1/4} G_F^{1/2} m_{q,l}\, 2T_{3q,l}. \end{aligned} \tag{54}$$

This is the same argument as for relating the axial coupling of a U boson to the pseudoscalar coupling of the equivalent axionlike pseudoscalar A or a, with the U, replaced by the Z, considered in the small mass and small coupling limit [33]. The scalar partner z, described by the same chiral superfield H_z as the would-be Goldstone z_g, has *scalar* couplings to quarks and leptons also given by (54) [and as found in (53) by the conventional method]. This may be remembered as

$$\underbrace{\sqrt{g^2 + g'^2}\ \frac{T_{3\,q,l}}{2}}_{\text{axial coupling of } Z}\ \ \frac{2m_{q,l}}{m_Z}\ \ \underset{\leftrightarrow}{=}\ \ \underbrace{2^{1/4}\, G_F^{1/2}\ m_{q,l}\ 2\,T_{3\,q,l}}_{\text{scalar coupling of } z}\,. \tag{55}$$

This also provides the couplings of the spin-0 $w^\pm$ from the $W^\pm$ ones using (36) and (37), leading to the factors $m_{d,e} \tan\beta$ and $m_u \cot\beta$ in the expressions of these couplings.

We recover as expected spin-0 couplings to quarks and leptons proportional to their masses, in contrast with the couplings of the spin-1 Z and $W^\pm$ in the same multiplets of supersymmetry. This is, however, a rather intriguing feature as z and Z, or $w^\pm$ and $W^\pm$ may also be simultaneously described by the same massive gauge superfields $Z(x,\theta,\bar\theta)$ and $W^\pm(x,\theta,\bar\theta)$. It is discussed and understood in [12], showing how the couplings of the spin-0 z and $w^\pm$ get in this description *resurrected* from the supersymmetric mass terms for quarks and leptons, through nonpolynomial field and superfield redefinitions.

In comparison with a standard model h_{SM} boson the z has *reduced trilinear couplings to the* $W^\pm$ *and* Z by a factor $-\cos 2\beta$ owing to (13) and (14), so that

$$\begin{cases} (z\ VV)\ \text{couplings} = (h_{\mathrm{SM}}\ VV)\ \text{couplings} \times (-\cos 2\beta), \\ (z\ ff)\ \text{couplings} = (h_{\mathrm{SM}}\ ff)\ \text{couplings} \times (2T_{3f} = \pm 1). \end{cases} \tag{56}$$

The expected production of a z in the ZZ^* or WW^* decay channels would then be *decreased by* $\cos^2 2\beta$ as compared to a SM boson, with respect to fermionic quark and lepton channels (the change of sign in d-quark and charged-lepton couplings also affecting the $h \to \gamma\gamma$ decay).

But the z does not necessarily correspond to a mass eigenstate, and further mixing effects induced by supersymmetry breaking must be taken into account, as discussed in Sec. IX. The h field presumably associated with the 125 GeV/c^2 boson observed at CERN may then be expressed (in the absence of further mixing effects that could involve an additional singlet) as

$$h = \sqrt{2}\,\mathrm{Re}(-h_1^0 c_{\beta'} + h_2^0 s_{\beta'}) = \sqrt{2}\,\mathrm{Re}(-h_1^0 s_\alpha + h_2^0 c_\alpha) \tag{57}$$

with $\beta' = \frac{\pi}{2} - \alpha$, and

$$\langle z|h\rangle = \cos(\beta - \beta') = \sin(\beta + \alpha). \tag{58}$$

At the same time

$$\langle h_{\mathrm{SM}}|h\rangle = -\cos(\beta + \beta') = \sin(\beta - \alpha), \tag{59}$$

the factor $\cos^2 2\beta$ affecting the ZZ^* or WW^* decay rates of a z being replaced by $\cos^2(\beta + \beta') = \sin^2(\beta - \alpha)$. The physical mass eigenstate h is very close to the z in (11) for $\beta = \beta'$ i.e., $\beta + \alpha \simeq \frac{\pi}{2}$, then justifying *an almost complete association of this* 125 GeV/c^2 *boson with the spin-1 Z.*

VIII. MASSIVE GAUGE SUPERFIELDS FOR SPIN-0 BOSONS

Supersymmetric theories thus allow for associating spin-1 with spin-0 particles within massive gauge multiplets of supersymmetry, leading to *gauge/BE-Higgs unification*, BEH bosons appearing as *extra spin-0 states of massive spin-1 gauge bosons*. We can even use the superfield formalism [34] to jointly describe these massive spin-1, spin-$\frac{1}{2}$ and now also spin-0 particles with massive gauge superfields [12].

Quite remarkably, this is possible *in spite of their different electroweak properties*, spin-1 fields transforming as a gauge triplet and a singlet with spin-0 BEH fields transforming as electroweak doublets. And although gauge and BE-Higgs bosons have *very different couplings to quarks and leptons*, which may first appear very puzzling but is elucidated in [12], using appropriate changes of field and superfield variables.

To do so we must *change picture* in our representation of such spin-0 bosons. The previous z and $w^\pm$ ($\equiv H^\pm$) cease being described by spin-0 components of the chiral superfields H_1 and H_2, to get described, through a *nonpolynomial change of (super)fields*, by the lowest (C) components of the Z and $W^\pm$ superfields. This association can be realized in a supersymmetric way by *completely* gauging away the three chiral superfields H_1^-, H_2^+, and $H_z = -H_1^0 c_\beta + H_2^0 s_\beta$. These complete superfields are now considered as Goldstone chiral superfields and eliminated by being taken identical to their vev's:

$$\begin{cases} H_1^- \equiv H_2^+ \equiv 0, \\ H_z = -H_1^0 c_\beta + H_2^0 s_\beta \equiv -\dfrac{v}{\sqrt{2}}\cos 2\beta. \end{cases} \tag{60}$$

The field degrees of freedom normally described by them, i.e., the spin-0 BEH fields referred to as z and $w^\pm$ in (11) and (40) and associated Higgsino fields are completely gauged away, and naively seem to be "lost" in this description.

But at the same time the corresponding gauge superfields $Z(x,\theta,\bar\theta)$ and $W^\pm(x,\theta,\bar\theta)$ acquire masses in a supersymmetric way, describing new physical degrees of freedom. These correspond precisely to those just "lost" in the gauging away of H_1^-, H_2^+, and H_z in (60). The chiral superfields, normalized so that $\langle H_i^0\rangle = \langle h_i^0\rangle = v_i/\sqrt{2}$, generate mass terms $\frac{1}{2}m_Z^2\,(Z^2)_D$ and $m_W^2\,|W^+|_D^2$ for $Z(x,\theta,\bar\theta)$ and $W^\pm(x,\theta,\bar\theta)$, the linear term in $Z(x,\theta,\bar\theta)$ vanishing owing to (20).

In the superfield formalism for supersymmetric gauge theories [34–36] the Lagrangian density [1] includes the terms

$$\mathcal{L} = \frac{1}{2}[H_1^\dagger \exp(g\boldsymbol{\tau}.\mathbf{W} - g'B)H_1 + H_2^\dagger \exp(g\boldsymbol{\tau}.\mathbf{W} + g'B)H_2]_D - \xi D' + \cdots. \tag{61}$$

We make the generalized gauge choice (60), so that

$$H_1 = \begin{pmatrix} \frac{v_1}{\sqrt{2}} + \cdots \\ 0 \end{pmatrix}, \qquad H_2 = \begin{pmatrix} 0 \\ \frac{v_2}{\sqrt{2}} + \cdots \end{pmatrix}, \tag{62}$$

the … involving the leftover superfield H_A. A second order expansion of $\mathcal{L}$ along the lines of [12], with

$$-\xi D' = \xi \sin\theta D_Z - \xi\cos\theta D_\gamma, \tag{63}$$

generates superfield mass terms for $W^\pm(x,\theta,\bar\theta)$ and $Z(x,\theta,\bar\theta)$. The term linear in $Z(x,\theta,\bar\theta)$, which appears with the coefficient

$$\frac{\sqrt{g^2+g'^2}}{4}(v_1^2-v_2^2)+\xi\sin\theta=-\langle D_Z\rangle\equiv 0, \quad (64)$$

vanishes identically owing to (20).

We get at second order

$$\mathcal{L}=\frac{1}{2}\left[\frac{(g^2+g'^2)(v_1^2+v_2^2)}{4}(W_3\cos\theta-B\sin\theta)^2 + \frac{g^2(v_1^2+v_2^2)}{4}(W_1^2+W_2^2)\right]_D+\cdots, \quad (65)$$

so that

$$\mathcal{L}=\frac{1}{2}m_Z^2(Z^2)_D+m_W^2|W^+|_D^2+\cdots = \frac{1}{2}m_Z^2(2C_ZD_Z-\partial_\mu C_Z\partial^\mu C_Z-Z_\mu Z^\mu+\cdots)-\frac{1}{4}Z_{\mu\nu}Z^{\mu\nu} +\cdots+\frac{D_Z^2}{2}+\cdots+\cdots. \quad (66)$$

After elimination of auxiliary fields through

$$D_Z=-m_Z^2C_Z+\cdots=m_Zz+\cdots,\quad \text{etc.}, \quad (67)$$

it includes the kinetic and mass terms for the gauge boson Z and associated spin-0 boson z,

$$\mathcal{L}=-\frac{1}{4}Z_{\mu\nu}Z^{\mu\nu}-\frac{m_Z^2}{2}Z_\mu Z^\mu-\frac{1}{2}\partial_\mu z\partial^\mu z-\frac{m_Z^2}{2}z^2+\cdots. \quad (68)$$

And similarly for the $W^\pm$ and spin-0 partner $w^\pm$ ($\equiv H^\pm$), keeping also in mind that supersymmetry is spontaneously broken for this superfield when $\tan\beta\neq 1$.

In this picture these spin-0 bosons get described by the lowest (C) spin-0 components of *massive* Z and $W^\pm$ superfields, expanded as $Z(x,\theta,\bar\theta)=C_Z+\cdots-\theta\sigma_\mu\bar\theta Z^\mu+\cdots$, $W^\pm(x,\theta,\bar\theta)=C_W^\pm+\cdots-\theta\sigma_\mu\bar\theta W^{\mu\pm}+\cdots$. Their C components now describe, through *nonpolynomial field transformations* linearized as $z=-m_ZC_Z+\cdots$, $w^\pm=m_WC_W^\pm+\cdots$, the same spin-0 fields z and $w^\pm$ as in the usual formalism (with signs depending on previous choices for the definitions of z and $w^\pm$). We thus have

$$Z(x,\theta,\bar\theta)=\left(\frac{-z}{m_Z}+\cdots\right)+\cdots-\theta\sigma_\mu\bar\theta Z^\mu+\cdots,$$
$$W^\pm(x,\theta,\bar\theta)=\left(\frac{w^\pm}{m_W}+\cdots\right)+\cdots-\theta\sigma_\mu\bar\theta W^{\mu\pm}+\cdots, \quad (69)$$

massive gauge superfields now describing spin-0 fields usually known as BEH fields. Their subcanonical (χ) spin-$\frac{1}{2}$ components, instead of being gauged away as usual, now also correspond to physical degrees of freedom describing the spin-$\frac{1}{2}$ fields previously known as Higgsinos.

IX. THE BE-HIGGS BOSON AS SPIN-0 PARTNER OF Z, IN THE (N/n) MSSM

A. MSSM

This applies to the spin-0 sector of the MSSM. The scalar potential may be expressed by adding to V obtained from (3) the soft dimension-2 supersymmetry-breaking term

$$-m_A^2|h_1s_\beta-h_2^cc_\beta|^2=-m_A^2|\varphi_{\rm in}|^2 \quad (70)$$

including in particular the μ-term contribution. This term, which vanishes for $\langle h_i^0\rangle=v_i/\sqrt{2}$, is a mass term for the doublet $\varphi_{\rm in}$, which has no vev and thus no direct trilinear couplings to gauge boson pairs (only to quarks and charged leptons). It does not modify the vacuum state considered, initially taken as having a spontaneously broken supersymmetry in the gauge-and-Higgs sector. It breaks explicitly the $U(1)_A$ symmetry (6) and (46), lifting the two previously flat directions associated with s_A and A. With

$$|\varphi_{\rm in}|^2=|h_1\sin\beta-h_2^c\cos\beta|^2 =|H^+|^2+\frac{1}{2}A^2+\frac{1}{2}|\sqrt{2}\,{\rm Re}(h_1^0s_\beta-h_2^0c_\beta)|^2, \quad (71)$$

it provides an extra contribution m_A^2 to $m_{H^\pm}^2$, so that

$$m_{H^\pm}^2=m_W^2+m_A^2. \quad (72)$$

Adding the supersymmetric m_Z^2 contribution associated with the z in (11) and supersymmetry-breaking contribution m_A^2 from (70) we get the scalar mass^2 matrix

$$\mathcal{M}_\circ^2=\begin{pmatrix} c_\beta^2m_Z^2+s_\beta^2m_A^2 & -s_\beta c_\beta(m_Z^2+m_A^2)\\ -s_\beta c_\beta(m_Z^2+m_A^2) & s_\beta^2m_Z^2+c_\beta^2m_A^2\end{pmatrix}, \quad (73)$$

verifying

$${\rm Tr}\mathcal{M}_\circ^2=m_H^2+m_h^2=m_Z^2+m_A^2,$$
$$\det\mathcal{M}_\circ^2=m_H^2m_h^2=m_Z^2m_A^2\cos^22\beta, \quad (74)$$

so that

$$m_{H,h}^2=\frac{m_Z^2+m_A^2}{2}\pm\sqrt{\left(\frac{m_Z^2+m_A^2}{2}\right)-m_Z^2m_A^2\cos^22\beta}. \quad (75)$$

It implies $m_h<m_Z|\cos2\beta|$ at the classical level, up to radiative corrections which must be significant if one is to reach $\simeq 125\ {\rm GeV}/c^2$ from a classical value below m_Z. These mass eigenstates behave for large m_A as

$$\begin{cases} H \to \sqrt{2}\,\mathrm{Re}(h_1^0 s_\beta - h_2^0 c_\beta) & (\text{large } m_H \simeq m_A), \\ h \to h_{\mathrm{SM}} = \sqrt{2}\,\mathrm{Re}(h_1^0 c_\beta + h_2^0 s_\beta) & (\text{SM-like}). \end{cases} \tag{76}$$

The h field, presumably associated with the 125 GeV$/c^2$ boson observed at CERN, is then also very close to the z in (11) for large $\tan\beta$, justifying an almost complete association of this 125 GeV$/c^2$ boson with the spin-1 Z.

B. N/nMSSM

This also applies to extensions of the minimal model, as with an extra N/nMSSM singlet S with a trilinear $\lambda H_1 H_2 S$ coupling, making it easier to get from λ large enough spin-0 masses [37]. In the N/nMSSM, first considered without a μ term, the supersymmetric contributions to spin-0 masses are [1]

$$\begin{cases} m_w = m_W, \qquad m_z = m_Z, \\ m\begin{pmatrix} \text{scalar } s_A, \text{ pseudoscalar A} \\ \text{complex singlet} \end{pmatrix} = m_A = \dfrac{\lambda v}{\sqrt{2}}. \end{cases} \tag{77}$$

They correspond, already in the absence of supersymmetry breaking, to the neutral scalar doublet mass^2 matrix

$$\mathcal{M}_\circ^2 = \begin{pmatrix} c_\beta^2 m_Z^2 + s_\beta^2 m_A^2 & s_\beta c_\beta (m_A^2 - m_Z^2) \\ s_\beta c_\beta (m_A^2 - m_Z^2) & s_\beta^2 m_Z^2 + c_\beta^2 m_A^2 \end{pmatrix}, \tag{78}$$

where $m_A = \lambda v/\sqrt{2}$.

Adding as in (70) the supersymmetry-breaking term $-\delta m_A^2 |h_1 s_\beta - h_2^c c_\beta|^2 = -\delta m_A^2 |\varphi_{\mathrm{in}}|^2$ does not modify the vacuum state, while shifting the A and $w^\pm$ mass^2 by the same amount δm_A^2, so that

$$m_A^2 = \frac{\lambda^2 v^2}{2} + \delta m_A^2,$$
$$m_w^2 = m_W^2 + \delta m_A^2 = m_W^2 + m_A^2 - \frac{\lambda^2 v^2}{2}. \tag{79}$$

It provides as in the MSSM an extra contribution to the neutral scalar doublet mass^2 matrix, shifted by

$$\delta\mathcal{M}_\circ^2 = \begin{pmatrix} s_\beta^2 \delta m_A^2 & -s_\beta c_\beta \delta m_A^2 \\ -s_\beta c_\beta \delta m_A^2 & c_\beta^2 \delta m_A^2 \end{pmatrix}. \tag{80}$$

From this shift $m_A^2 = \frac{\lambda^2 v^2}{2} \to \frac{\lambda^2 v^2}{2} + \delta m_A^2$, the mass^2 matrix for the scalar doublet components reads

$$\mathcal{M}_\circ^2 = \begin{pmatrix} c_\beta^2 m_Z^2 + s_\beta^2 m_A^2 & s_\beta c_\beta (\lambda^2 v^2 - m_A^2 - m_Z^2) \\ s_\beta c_\beta (\lambda^2 v^2 - m_A^2 - m_Z^2) & s_\beta^2 m_Z^2 + c_\beta^2 m_A^2 \end{pmatrix}. \tag{81}$$

For $\lambda \to 0$ S decouples and the spectrum (79) and (81) returns to the usual MSSM one. For $\lambda \neq 0$ further contributions involving also a possible singlet mass term $\frac{1}{2}\mu_S S^2$ lead in general to a mixing between neutral doublet and singlet components, with $\mathcal{M}_\circ^2$ embedded into a 3×3 matrix.

X. CONCLUSIONS

Independently of specific realisations (MSSM, N/nMSSM, USSM, …) supersymmetric theories provide spin-0 bosons as extra states for massive spin-1 gauge bosons, despite different symmetry properties and different couplings to quarks and leptons [1,12]. This further applies to supersymmetric grand-unified theories with extra dimensions [31,32]. By connecting spin-1 *mediators of gauge interactions* with spin-0 *particles* associated with *symmetry breaking* and *mass generation*, supersymmetry provides an intimate connection between the electroweak gauge couplings and the spin-0 couplings associated with symmetry breaking and mass generation.

The 125 GeV$/c^2$ boson recently observed at CERN may also be interpreted, up to a mixing angle induced by supersymmetry breaking, as the spin-0 partner of the Z under *two* supersymmetry transformations,

$$\text{spin-1 } Z \overset{\text{SUSY}}{\longleftrightarrow} \overset{\text{SUSY}}{\longleftrightarrow} \text{spin-0 BEH boson}, \tag{82}$$

i.e., as a Z deprived of its spin.

This provides within a theory of electroweak and strong interactions

the first example of two known fundamental particles of different spins related by supersymmetry, (83)

in spite of different electroweak properties. This is considerable progress as compared to the initial situation in (1), bringing further confidence in the relevance of supersymmetry for the description of fundamental particles and interactions.

Supersymmetry may thus be tested in the gauge-and-BE-Higgs sector at present and future colliders, in particular through the properties of the new boson, even if R-odd superpartners were still to remain out of reach for some time.

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[27] In a supersymmetric grand-unified theory with a semisimple gauge group ξ must vanish at the GUT scale, requiring for two spin-0 doublets $v_1 = v_2$ i.e., $\tan\beta = 1$ at this scale.

[28] If $\psi = a_L + b_R$ is a Dirac spinor constructed from two Majorana ones a and b, its mass term may be expressed through a nondiagonal matrix, as $-im\bar{\psi}\psi = -im(\bar{b}\frac{1-i\gamma_5}{2}a + \bar{a}\frac{1+i\gamma_5}{2}b) = -i\frac{m}{2}(\bar{a}b + \bar{b}a)$.

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