Origin and Phenomenology of Weak-Doublet Spin-1 Bosons

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We study phenomenological consequences of the Standard Model extension by the new spin-1 fields with the internal quantum numbers of the electroweak Higgs doublets. We show, that there are at least three different classes of theories, all motivated by the hierarchy problem, which predict appearance of such vector weak-doublets not far from the weak scale. The common feature for all the models is the existence of an $SU(3)_W$ gauge extension of the weak $SU(2)_W$ group, which is broken down to the latter at some energy scale around TeV. The Higgs doublet then emerges as either a pseudo-Nambu-Goldstone boson of a global remnant of $SU(3)_W$, or as a symmetry partner of the true eaten-up Goldstone boson. In the third class, the Higgs is a scalar component of a high-dimensional $SU(3)_W$ gauge field. The common phenomenological feature of these theories is the existence of the electroweak doublet vectors ($Z^*, W^*$), which in contrast to well-known $Z'$ and $W'$ bosons posses only anomalous (magnetic moment type) couplings with ordinary light fermions. This fact leads to some unique signatures for their detection at the hadron colliders.

PACS numbers: 12.60.Cn, 14.70.Pw, 12.10.Dm, 13.85.Fb

I. INTRODUCTION

The main theoretical motivation for beyond the standard model physics around TeV energies is provided by the Hierarchy Problem, inexpressible quantum stability of the weak interaction scale with respect to the ultraviolet cutoff. This problem suggests the existence of some new regulating physics not far above the weak scale. Needless to say, understanding experimental consequences of the latter is of fundamental importance.

Recently [1], it was pointed out that possible existence of massive vector fields ($V_μ ≡ (Z_μ^*, W_μ^*)$) with the internal quantum numbers identical to the Standard Model Higgs (or the lepton) doublet, can result in some interesting phenomenological consequences. Due to their quantum numbers, to the leading order such vectors can only have magnetic type interactions with the Standard Model fermions,

\[
\frac{1}{M} D_{[μ} V_{μ]}^c \left( g_{LR} Q_L \sigma^{μν} d_R + g_{LR} \bar{L} σ^{μν} e_R \right) + \frac{g_μ^R}{M} D_{[μ} V_{μ]}^c \bar{Q} L σ^{μν} u_R + \text{h.c.},
\]

where $V_μ^c ≡ (-W_μ^*, Z_μ^*)$ is the charge-conjugated doublet; $Q_L ≡ (u_L, d_L)$ and $L ≡ (ν_L, e_L)$ are the left-handed quark and lepton doublets respectively. $D_{μ}$ are the usual $SU(2)_W × U(1)_Y$ covariant derivatives, and the obvious group and family indexes are suppressed. $M$ is the scale of the new physics and $g_{LR}^{u,d,e}$ are dimensionless constants.

Up until now, no theoretical motivation for the existence of such states was given. It is the purpose of this paper to provide such a motivation from the Hierarchy Problem point of view. We shall show that such states are predicted by three different classes of theories that represent different approaches for explaining the relative lightness of the Higgs doublets.

The crucial common feature of all three approaches is, that they are based on the existence of $U(3)_W ≡ SU(3)_W × U(1)_Y$ gauge extension of the $SU(2)_W × U(1)_Y$ electroweak group, which is spontaneously broken down to the latter at scale $M$. The weak doublet vectors $V_μ$ are then identified with the $SU(2)_W$ doublet components of the 8-dimensional gauge multiplet of the $SU(3)_W$ group, which under $SU(2)_W$ subgroup decomposes as

\[
s = 3 + 2 + 2 + 1, \quad (2)
\]

where numbers refer to the dimensionality of the corresponding $SU(2)_W$ representations. The vector fields $V_μ$ ($V_μ^*$) obviously belong to fragments $2$ ($2^*$), and become massive during the spontaneous symmetry breaking $U(3)_W → SU(2)_W × U(1)_Y$. The lightness of the Higgs doublets is guaranteed, because they are related to $V_μ$ vectors by symmetry. This relation in three different approaches is established as follows.

A. Pseudo-Goldstone Higgs

In the first approach the lightness of the Standard Model Higgs doublet is achieved because it is a pseudo-Goldstone boson of a spontaneously broken $G_{global} ≡ G × G$ symmetry of the scalar potential, whereas only the diagonal $G_{local} ≡ G$ part of it is gauged. This idea was originally proposed in the context of $G = SU(6)$ grand unification [2], but our current focus will be the $G = U(3)_W$ realization of it [3]. In the latter realization $SU(2)_W × U(1)_Y$ is embedded into $G_{local} = U(3)_W$ group as a maximal subgroup. As explained above, the 8-dimensional gauge multiplet of $SU(3)_W$, on top of the 3 electroweak gauge bosons and an extra singlet, contains...
a complex doublet $V_\mu$ with the quantum numbers of the Standard Model Higgs.

The spontaneous breaking of the local symmetry $U(3)_{\text{local}} \rightarrow SU(2)_W \times U(1)_Y$ is triggered by the two independent Higgs triplets, $3_H$ and $3'_{H'}$, which under the $SU(2)_W$ subgroup decompose as $3_H = 1_H + 2_H$ and $3'_{H'} = 1'_{H'} + 2'_{H'}$ respectively. The fragments $2_H$ and $2'_{H'}$ are doublets of $SU(2)_W$ and have the quantum numbers of the electroweak doublets. The non-zero vacuum expectation values (VEVs) are developed by the singlet component, $\langle 1'_H \rangle \neq 0$ and $\langle 1_H \rangle \neq 0$. As a result of this breaking, gauge bosons $V_\mu$ become massive. The following combination,

$$2_{Gold} \equiv \frac{\langle 1_H \rangle 2_H + \langle 1'_H \rangle 2'_{H'}}{\sqrt{\langle 1_H \rangle^2 + \langle 1'_H \rangle^2}}$$

(3)

is eaten-up and becomes a longitudinal components of $V_\mu$. Whereas the orthogonal state,

$$2_{Higgs} \equiv \frac{\langle 1'_H \rangle 2_H - \langle 1_H \rangle 2'_{H'}}{\sqrt{\langle 1_H \rangle^2 + \langle 1'_H \rangle^2}}$$

(4)

is a pseudo-Nambu-Goldstone, which is massless at the tree-level and gets the suppressed mass only at the loop level, and therefore remains lighter than the symmetry breaking scale. This pseudo-Goldstone plays the role of the Standard Model Higgs doublet.

**B. Goldstone Sister Higgs**

In the second approach [4], the Higgs mass is protected because it is related by symmetry to an exact Goldstone boson that becomes a longitudinal component of $V_\mu$. We shall refer to this scenario as “Goldstone-Sister Higgs”. The gauge symmetry structure of the simplest model is identical to the previous case. There is an exact $U(3)_W$ gauge symmetry that incorporates the Standard Model group as its maximal subgroup. Again, the spontaneous breaking is triggered by two Higgs triplets, $3_H$ and $3'_{H'}$. However, no approximate global symmetry is required. Instead, the two Higgs triplets are related by an exact custodial symmetry, such as, the permutation or an $SU(2)_{\text{cust}}$ symmetry that transforms the triplets into each other $3_H \equiv 3'_{H'}$.

The lightness of the Higgs doublet is then guaranteed by the following effect. Breaking of $U(3)$ symmetry is triggered by the VEV of the singlet component of the $3'_{H'}$-triplet. During this breaking, $V_\mu$ becomes massive and eats up the doublet $2'_{H'}$. Thus $2'_{H'}$ becomes a longitudinal polarization of a massive gauge field

$$V_\mu \rightarrow V_\mu + \frac{1}{M_V} \partial_\mu (2'_{H'}) ,$$

(5)

where $M_V$ is the mass of $V_\mu$. Since $2'_{H'}$ is a true eaten-up Goldstone, it cannot have any contribution to its mass from the scalar potential, but only from the kinetic mixing with the gauge field. But, since the un-eaten doublet $2_H$ is related to $2'_{H'}$ by the custodial symmetry, the former also stays massless at the three-level. In this way the mass of the physical Higgs doublet is protected by its sister doublet becoming a Goldstone particle. We shall consider this scenario in more details below.

**C. Higgs as Extra Dimensional Gauge Field**

Finally, the third class of theories in which appearance of the doublet gauge fields is the must, is the one in which the Standard Model Higgs doublet $H$ is an extra dimensional component of a high-dimensional gauge field $E \otimes F$. For understanding the key idea of this approach, it suffices to consider a simplest case of a vector field in five dimensional Minkowski space, $V_A$, where $A = \mu, 5$ is the five dimensional Lorentz index. In the approach of [3, 6], the Higgs is identified with the fifth component of the gauge field, which is a four-dimensional Lorentz-scalar, $H \equiv V_5$. Obviously, since $H$ and $V_\mu$ are the components of the same high-dimensional gauge field, their internal quantum numbers must be identical.

Thus, the existence of the weak-doublet vector particles is reinforced by the high-dimensional gauge symmetry. This symmetry is spontaneously broken by compactification. In this way the mass of the Higgs doublet is controlled by the compactification scale, as opposed to the high-dimensional cutoff of the theory. The realistic model building in this class of theories is much more involved than in the previous two cases. For us the only important aspect is the model-independent property of the existence of the massive vector doublet $V_\mu$. This property is guaranteed by the symmetry and is insensitive to the concrete model building.

Having specified the class of the theories of our interest, let us turn to the interaction between $V_\mu$ and the Standard Model fermions. In all three classes of theories, the coupling (11) even if not present at the tree-level can (and in general will) be generated by the loop corrections, as it is permitted by all the symmetries of the low energy theory. We shall first demonstrate how this generation happens by considering a toy model reduced to its bare essentials, and later illustrate it on a detailed example of Goldstone-Sister Higgs [4].

**II. TOY MODEL**

In this section we shall discuss a generation of coupling (11) in a simple toy model. The latter consists of two sectors: Hypothetical heavy particles and chiral massless fermions of ordinary matter. The corresponding La-
The Lagrangian reads,
\[ \mathcal{L} = -D^\mu V^\dagger_{\bar{v}} D^{\mu} V_{\nu} + M^2 V^\dagger_{\bar{v}} V^\mu + \partial_\mu \psi^* \partial^\mu \psi - M^2 \phi^2 \phi + \sum_{k=1}^{2} \bar{\psi}_k (i \not{D} - m) \psi_k + g \bar{\psi}_2 \gamma^\mu \psi_1 \not{V}_{\mu} + g V^\dagger_{\bar{v}} \psi_1 \gamma^\mu \psi_2 \]
\[ + \sum_{k=1}^{2} \frac{h}{2} \bar{\psi}_2 (1 + \gamma^5) \phi^2 \psi_2 + \frac{h}{2} \bar{\psi}_1 (1 + \gamma^5) \phi^2 \psi_1, \tag{6} \]
where the first line represents the bilinear Lagrangian of the $SU(2)_W$-doublet vector fields, $V_\mu (V^\dagger_{\bar{v}})$, and of a complex singlet scalar field, $\phi$. The second line describes the kinetic and mass terms of a singlet ($\psi_1^\dagger$) and a doublet ($\psi_2^\dagger$) heavy fermions and their interactions with the doublet vector fields. The primed fermions are non-chiral, meaning that each of them comes in both left and right chiralities. The last two lines include the kinetic terms of the ordinary chiral massless fermions, the left-handed doublets, $\psi_2$, and the right-handed singlets, $\psi_1$, and their interactions with the heavy fields.

The one-loop diagram in Fig. 1 leads to new anomalous coupling of the vector fields with the ordinary fermions.

\[ \Delta \mathcal{L} = \frac{ig^2 h^2}{4} \bar{\psi}_2 \int \frac{d^4 k}{(2\pi)^4} \left( 1 + \gamma^5 \right) \frac{\not{p} - \not{k} + m}{(\not{p} - \not{k})^2 - m^2} \gamma^\mu \]
\[ \times \frac{\not{p} - \not{k} + m}{(\not{p} - \not{k})^2 - m^2} \left( 1 + \gamma^5 \right) \frac{1}{k^2 - M^2} \psi_1 V_{\mu}, \]
\[ = \frac{gh^2}{32\pi^2 m} \mathcal{I} (q^2, m^2, m^2) \bar{\psi}_2 \not{\sigma}^{\mu\nu} \psi_1 \partial_\mu V_\nu, \tag{7} \]
where
\[ \mathcal{I} = \int_0^1 x^2 dx \int_0^1 \frac{y dy}{x + \frac{m^2}{x} (1 - x) - \frac{m^2}{x} x^2 y (1 - y)} \tag{8} \]
is a slow varied function at $q^2 \ll m^2 \sim M^2$.

### III. A MODEL

In this section we shall discuss an explicit example of the realistic model. As such we shall choose a model based on the idea of Goldstone-Sister Higgs [4]. As discussed above, in this theory the electroweak $SU(2)_W \times U(1)_Y$-symmetry is enhanced to $U(3)_W$. Obviously the gauge sector contains an additional electromagnetodoublet vector field ($V_\mu$), which after spontaneous breaking $U(3)_W \rightarrow SU(2)_W \times U(1)_Y$ becomes massive. This breaking is realized by the VEV of a Higgs triplet $3_H$, the weak-doublet part of which ($2_H^t$) becomes a longitudinal component of $V_\mu$. In order to trigger the second Standard Model stage of symmetry breaking $SU(2)_W \times U(1)_Y \rightarrow U(1)_{EM}$, the theory contains a second Higgs triplet $3_H^t$. The key point is that the two triplets are related by an additional custodial $SU(2)_{cust}$ symmetry. The motivation for such symmetry is, that the physical Higgs doublet ($2_H^t$) is a partner of the Goldstone boson $2_H^t$ that is eaten up via the Higgs effect of $U(3)_W \rightarrow SU(2)_W \times U(1)_Y$ breaking. Due to this symmetry relation, the SM Higgs doublet remains naturally light.

It is the simplest to discuss this mechanism directly in supersymmetric Grand Unified Theory (GUT) context, in which the $U(3)_W \equiv SU(3)_W \times U(1)_Y$ group together with the color $SU(3)_C$ is embedded into the $SU(6)$ group as a maximal subgroup, $G_{CW} \equiv SU(3)_C \times SU(3)_W \times U(1)_W \subset SU(6)$. Notice that this embedding allows to treat color and weak $SU(3)$-groups in a completely democratic way.

Notice, that for our present purposes having the full $SU(6)$-symmetry group is completely unessential. The latter is anyway broken down to its $G_{CW}$ subgroup at scales much above the energies of our present interest. So we could have equally well limit ourselves by $G_{CW}$ symmetry. However, the analysis is much more convenient in terms of $SU(6)$ representations, rather than in terms of its subgroups. It also allows us to understand the particle content in term of representations of a more familiar $SU(5)$ subgroup. Due to this advantages, we shall use $SU(6)$-classification for the particles. If needed, the reader can easily perform decompositions into the $G_{CW}$-reduced representations instead.

The full symmetry of the model is thus $SU(6) \times SU(2)_{cust}$, where $SU(2)_{cust}$ is an additional custodial symmetry that relates the SM Higgs doublet to an eaten-up Goldstone boson.

The chiral superfield content is:

1. Higgs sector:
   - $35$-plet $\Sigma_i^j \ (i,j = 1, \ldots, 6)$ and $\{6, \bar{2}\} \equiv H_{A1}$, $\{6, \bar{2}\} \equiv H_{A2}$, where $i,j$ are $SU(6)$ and $A = 1, 2$ are $SU(2)_{cust}$ indexes respectively.

2. The SM fermions are embedded in the following anomaly-free set (per generation):
   - $\{15\}_{ij}$, $\{6, \bar{2}\} \equiv \bar{6}^{A3}$ and a singlet $\mathbf{1}$.

We shall denote the superfields by the same symbols as their components. In each case it will be clear from the context which component we are referring to.

The symmetry breaking is achieved by the Higgs part of the superpotential, which has the following form,

\[ W_{Higgs} = \frac{\lambda}{3} \text{Tr} \Sigma^3 + \lambda' H^A \Sigma H_A + M' \bar{H}^A H_A. \tag{9} \]
The vacuum of the theory is:

\[ \Sigma = \begin{pmatrix} 1 & 1 & M' \\ 1 & -1 & 0 \\ -1 & 0 & \mu \end{pmatrix} \times \begin{pmatrix} \lambda \\ \lambda \\ \lambda \end{pmatrix}, \]

\[ \bar{H}^1 = H_2 = \begin{pmatrix} 0 \\ 0 \\ \mu \end{pmatrix}, \]

which leads to the symmetry breaking \( SU(6) \rightarrow SU(3)_C \times SU(2)_W \times U(1)_Y \). The value of \( \mu \) is undetermined in SUSY limit.

We are interested in the situation when the scale \( \mu \) is not very high, around TeV or so, whereas the scale \( M'/\lambda' \) is much larger. In this way, below the latter scale the low energy symmetry group is \( G_{CW} \). The further breaking of this symmetry to the Standard Model group, is triggered at much lower scale \( \mu \) by the VEV of the \( \text{6-plets} \).

As a result of this breaking, the Higgs doublets in \( \bar{H}^1 \) and \( H_2 \) are eaten up by the \( SU(2)_W \)-doublet gauge fields \( V_\mu(V_\mu^\dagger) \) that reside in the adjoint representation of \( SU(3)_W \). The doublets in \( \bar{H}^2 \) and \( H_1 \) are physical and at the tree level (in SUSY limit) have very small masses \( M \sim \lambda' \mu^2/M' \). An additional contribution comes from loop corrections after supersymmetry breaking (see [4] for details).

Notice, that as a result of the symmetry structure of the theory, the color triplet partners of the Higgs doublets are automatically super-heavy, with the masses \( 2M' \), and decouple from the low energy spectrum. This solves the doublet-triplet splitting problem in SUSY GUTs.

The fermion masses are generated from the following interactions. The up-type quark masses come from

\[ \varepsilon^{AB} H_{Ai} H_{Bj} 15_{[km]} 15_{[nl]} \varepsilon^{ijklmn}, \]

and the masses of down-type quarks, charged leptons and heavy states come from

\[ \varepsilon_{AB} \tilde{H}^{A1} 15_{[ij]} 15_{[kl]} + H_{Ai} \tilde{6}^{A1}. \]

From the above couplings the following interactions between the doublet vector and light fermions are generated. The superdiagram containing external light up-type quarks is given in Fig. 2. The corresponding effective operator of interest is,

\[ \tilde{H}^{B1}(F_{\mu\nu})_a^\dagger H_{Ai} 15_{[km]} \sigma_{\mu\nu} 15_{[ln]} \varepsilon_{BA} \delta^{ijklmn}, \]

where \((F_{\mu\nu})_a^\dagger\) is the \( SU(6) \) gauge field strength, the fermionic components are taken from 15-plets and bosonic components from the rest. After substituting the VEVs of 6-plet Higgses, this operator reduces to the magnetic coupling between \( V_\mu \) and the up-type quarks given by the last term in (11). Coupling to down-type quarks and leptons is generated through the diagram of Fig. 3, which leads to the following effective operator,

\[ H^{*BA}(F_{\mu\nu})_a^\dagger 15_{[ij]} \sigma_{\mu\nu} \tilde{6}^{A1} \varepsilon_{BA}. \]

After substitution of the VEVs, the above interaction reduces to the magnetic coupling of \( V_\mu \) with \( d \)-quark and charged leptons, given by the first two terms in (11).

Notice, that due to supersymmetry violating insertions in the vertexes, the flavor structure of the operators (11) and (12) is not necessarily aligned with the flavor structure of quark and lepton mass matrices. This means, that the exchange by \( V_\mu \) could potentially contribute into new flavor and CP violating interactions.

Notice also, that the same operator (12) could potentially give a contribution to \((g-2)\) of leptons. Indeed, if we instead of inserting the \( SU(2)_W \)-doublet VEV of \( H_2 \), insert the VEV of a physical Higgs doublet living in \( H_1 \), the operator (12) will reduce to the magnetic moment coupling of photon. However, this magnetic coupling will mix the light left-handed lepton residing in 15-plet with the heavy right-handed lepton from \( \tilde{6}^1 \), as opposed to the Standard Model right-handed lepton residing in \( \tilde{6}^1 \), thus, giving no contribution to \((g-2)\). However, a non-zero contribution could arise in case of a small mixing between the doublets from \( H^1 \) and \( H^2 \). Such mixing could arise from the calculable radiative corrections, and thus, could relate the phenomenological signatures of \( V_\mu \), with the value of \((g-2)\).

This study will not be attempted in the present work. Instead, we shall focus on characteristic signatures of resonance production and decay of \( V_\mu \)-bosons.
IV. CONSEQUENCES FOR COLLIDERS

In paper [7] it has been shown that tensor current interaction leads to a new angular distribution in comparison with well-known vector interactions. It was realized later [1] that this property ensures distinctive signature for their detection at the hadron colliders.

The hadron colliders, due to their biggest center-of-momentum (CM) energy $\sqrt{s} \sim$ several TeVs and their relatively compact sizes, still remain the main tools for discoveries of very heavy particles. The presence of partons with a broad range of different momenta allows to flush the entire energetically accessible region, roughly, up to $\sqrt{s}/6$. The production mechanism for new heavy bosons at a hadron collider is the $q\bar{q}$ resonance fusion.

In this paper we will consider the resonance production and decay of the above-introduced heavy spin-1 gauge bosons into the light lepton pairs, electrons or muons. For such high energies it is convenient to use the helicity formalism, since the helicity is a good quantum number for massless particles.

This, in a way, fixes the dominant production and decay mechanisms. For example, the decay angular distribution in the CM frame of a particle with spin-$s$ and helicity $\lambda$ with $-s \leq \lambda \leq s$ decaying into two massless particles with helicities $\lambda_1$ and $\lambda_2$ can be written as [8]

$$\frac{d\Gamma_s}{d\cos \theta \ d\phi} = \frac{1}{64\pi^2 M} |M_{\lambda_1\lambda_2}^{s}(\theta, \phi)|^2,$$

where the helicity amplitude

$$M_{\lambda_1\lambda_2}^{s}(\theta, \phi) = \sqrt{\frac{2s+1}{4\pi}} e^{i(\lambda-s)\phi} d_\lambda^s(\theta) M_{\lambda_1\lambda_2}$$

is expressed through the difference $\delta \equiv \lambda_1 - \lambda_2$ and the reduced decay amplitude $M_\lambda^{s}$, which is a function of $s$ and the outgoing helicities, but is independent of the polar ($\theta$) and the azimuthal ($\phi$) angles.

Up until now only resonance production and decay of spin-1 bosons with maximal helicities $\lambda = \pm 1$ have been considered. They are associated with additional $U(1)'$ gauge symmetries and are usually called $Z'$ particles. The Lorentz structure of its couplings to each fermion flavor is characterized by two generally independent constants $g^f_{LL}$ and $g^f_{RR}$.

$$\mathcal{L}_{Z'} = \sum_f \left( g^f_{LL} \bar{\psi}_L^{j\mu} \psi_R^{j\mu} + g^f_{RR} \bar{\psi}_R^{j\mu} \psi_L^{j\mu} \right) Z'_\mu.$$

Experimental determination of these coupling constants or disentangling among the different models is a rather hard task and cannot be fulfilled, for example, with the first LHC data. In the best case only the specific symmetric angular distribution over the polar angle $\theta$,

$$\frac{d\Gamma_1(q\bar{q} \rightarrow Z' \rightarrow \ell\ell)}{d\cos \theta} \propto |d_0^1|^2 + |d_{-1}^1|^2 \sim 1 + \cos^2 \theta,$$

allows to distinguish its production from distributions of spin-0,

$$\frac{d\Gamma_0(q\bar{q} \rightarrow h \rightarrow \ell\ell)}{d\cos \theta} \propto |d_0^0|^2 \sim 1$$

and spin-2

$$\frac{d\Gamma_2(q\bar{q} \rightarrow G^* \rightarrow \ell\ell)}{d\cos \theta} \propto |d_{11}^2|^2 + |d_{-11}^2|^2 \sim 1 - 3 \cos^2 \theta + 4 \cos^4 \theta$$

$$\frac{d\Gamma_2(q\bar{q} \rightarrow G^* \rightarrow \ell\ell)}{d\cos \theta} \propto |d_{21}^2|^2 + |d_{-21}^2|^2 \sim 1 - \cos^4 \theta$$

resonances.

Another possibility is the resonance production and decay of longitudinal spin-1 bosons with $\lambda = 0$, but this possibility is not widely discussed.

While the $Z'$ bosons with helicity $\lambda = \pm 1$ are produced in left(right)-handed quark and right(left)-handed antiquark fusion, the longitudinal $Z^*$ bosons can be produced through the new chiral couplings [1].

$$\mathcal{L}_{Z^*} = \sum_{f=d,e} \left( g^f_{LL} \bar{\psi}_L^\gamma \psi_R^\gamma \partial_\mu Z^*_\mu + g^f_{RR} \bar{\psi}_R^\gamma \psi_L^\gamma \partial_\mu Z^*_\mu \right),$$

with the complex constants $g^f_{LL} = (g^f_{LL})^*$. In left-handed or right-handed quark-antiquark fusion [9].

The new couplings lead to a different angular distribution

$$\frac{d\Gamma_1(q\bar{q} \rightarrow Z^* \rightarrow \ell\ell)}{d\cos \theta} \propto |d_{01}^1|^2 \sim \cos^2 \theta,$$

than the previously considered ones. At first sight, the small difference between the distributions [15] and [25] seems unimportant. However, the absence of the constant term in the latter case results at least in two potential experimental signatures.

First of all, the known angular distributions for scalar [19], vector [15] and spin-2 [20,21] bosons include a nonzero constant term, which leads to the kinematic singularity in $p_T$ distribution of the final lepton

$$\frac{1}{\cos \theta} \propto \frac{1}{\sqrt{(M/2)^2 - p_T^2}}$$

in the narrow width approximation $\Gamma \ll M$

$$\frac{1}{(s-M^2)^2 + M^4 \Gamma^2} \approx \frac{\pi}{MT} \delta(s-M^2).$$

This singularity is transformed into a well known Jacobian peak due to a finite width of the resonance. In contrast to this, the pole in the decay distribution of the $Z^*$ bosons is canceled out and the lepton $p_T$ distribution even reaches zero at the kinematical endpoint $p_T = M/2$. Therefore, the $Z^*$ boson decay distribution has a broad...
The lepton $p_T$ distributions from the $Z^*$ (solid) and $Z'$ (dotted) bosons decays.

smooth hump with the maximum below the kinematical endpoint, instead of a sharp Jacobian peak (Fig. 4).

Another striking feature of the distribution (23) is the forbidden decay direction perpendicular to the boost of the excited boson in the rest frame of the latter (the Collins–Soper frame [10]). It leads to a peculiar “swallow-tail” shape of the angular distribution with a profound dip at $\cos \theta_{\text{CS}}^* = 0$ in the Collins–Soper frame (Fig. 5) [9].

In conclusion, we would like to emphasize the difference between group properties of the gauge spin-1 $Z'$ and $Z^*$ bosons. While the $Z'$ bosons are described by real representations transforming as triplets or singlets of $SU(2)_W \times U(1)_Y$ group, the $Z^*$ bosons, like the Higgs fields, are assigned to the complex representation and transform as doublets. This results in another important experimental consequence.

Together with the neutral $Z^*$ bosons the doublets always contain also the charged bosons, $W^\pm$, which decay into a charged lepton and an undetected neutrino. Therefore, the angular distribution in Fig. 5 is experimentally unaccessible for them in the lepton channel. Only the $p_T$-distribution of the charged lepton can be measured. However, the distribution of the $W^*$ bosons differs drastically from the distribution of the $W'$ bosons (Fig. 4). Hence, even relatively small decay width of the $W^*$ bosons leads to a wide hump without the Jacobian peak, that obscures their identification as resonances at the hadron colliders.

The only way to access the angular distribution for $W^*$ bosons like in Fig. 5 is kinematical reconstruction of their decays into heavy quarks, $t\bar{b}$, which, in spite of the strong QCD background, can be identified via $b$-tagging. The presence of the heaviest $t$ quark in the final state will lead to an additional contribution,

\[ |d_{01}|^2 \sim 1 - \cos^2 \theta, \]

(26)
to the angular distribution proportional to the ratio $m_t^2/M^2$ due to helicity flip of the $t$ quark.

Acknowledgements

We are grateful to C. Wetterich for discussions. The work of MC was partially supported by Grant-in-Aid for Scientific Research 104/15.05.2009 from the Sofia University. The research of GD is supported in part by European Commission under the ERC advanced grant 226371, by David and Lucile Packard Foundation Fellowship for Science and Engineering and by the NSF grant PHY-0758032.

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