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ON THE SPECTRUM OF COMPOSITE RESONANCES

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Abstract

I discuss the infrared mass spectrum of strongly-coupled gauge theories, that induce the Higgs as a composite pseudo-Nambu-Goldstone boson. The set of composite states accompanying the Higgs is determined by the symmetries of the theory. Here we estimate their mass spectrum by non-perturbative techniques inspired by QCD, as well as by exploiting gauge-gravity duality.

1 Composite Higgs: motivations and relevant energy scales

As the Large Hadron Collider (LHC) did not find new states significantly coupled to the Standard Model (SM) below the TeV scale, any SM extension by such heavy states suffers from a little hierarchy problem, as the mass of the scalar Higgs boson lies close to the 100 GeV scale. Still, some SM extensions have the potential to address the big hierarchy between the TeV scale and the Planck scale. One possibility is to avoid elementary scalar fields, and assume the observed Higgs is a composite object, with a compositeness scale $f \gtrsim 1$ TeV. This scenario requires a strongly-coupled sector, whose spectrum generically includes several additional composite states besides the Higgs. The mass of the lowest-lying states cannot exceed $\sim 4\pi f$, and some could be significantly lighter and within the LHC reach. Definite predictions for the mass spectrum require to specify the strongly-coupled theory in the ultraviolet (UV). Here we will assume it is a gauge theory of fermions, that confines in the infrared. We will estimate its mass spectrum in some well-defined approximations, by employing non-perturbative techniques inspired by QCD [1), as well as gauge-gravity duality techniques [2).

In models where the Higgs is a pseudo-Nambu-Goldstone boson (pNGB) the electroweak scale, $v \simeq 246$ GeV, is induced in two steps. The theory has a global (flavour) symmtry G_F, that is spontaneously broken to a subgroup H_F at the scale f. The electroweak symmetry $SU(2)_L \times U(1)_Y$ is embedded in

H_F, and the set of NGB includes the SM Higgs doublet. Weak sources of explicit symmetry breaking – typically loops involving the top-quark Yukawa coupling – misalign the vacuum, inducing an effective potential for h, whose minimum determines v. The electroweak precision parameters as well as the Higgs couplings receive corrections of order v^2/f^2, and present data already imply $f \gtrsim 1$ TeV.

If the scale f is induced by strong dynamics, it is protected from large radiative corrections from UV physics, and the pNGB Higgs is a composite object. [3] The spectrum of composite resonances has typical mass gap $m_* \sim g_* f$, where $1 \lesssim g_* \lesssim 4\pi$ is the generic inter-resonance coupling. Since only resonances significantly lighter than $\sim 4\pi f$ have chances to be discovered at the LHC, our aim is to investigate the strong dynamics in order to find a rationale for the lightness of some composite states, besides the pNGB Higgs. In some instances light states are also welcome to minimise the fine-tuning in the Higgs potential.

2 UV-complete composite-Higgs models

A prototypical strongly-coupled sector is provided by an asymptotically-free gauge theory, with a hypercolour gauge group G_C and fermion matter fields only (no scalars). We will assume that the theory enters a strongly-coupled, walking (approximately scale-invariant) regime at some UV scale Λ_{UV}, and eventually develops a mass gap at some IR scale m_*. A large walking region, that is, a hierarchy $m_* \ll \Lambda_{UV}$, is required to induce the SM Yukawa couplings and to suppress flavour violation at the same time.

The choice of the appropriate gauge theory of fermions requires some exercise in group theory. In order to correctly describe electroweak symmetry breaking and preserve the SM custodial symmetry to a good approximation, the flavour-symmetry-breaking pattern should satisfy $G_F \to H_F \supset SU(2)_L \times SU(2)_R$, and the associated set of NGBs should include the Higgs transforming as $h \sim (2_L, 2_R)$. A generic gauge theory of fermions has flavour symmetry $G_F = SU(N_1) \times \cdots \times SU(N_k) \times U(1)^{k-1}$, where N_i is the number of Weyl fermions in the representation R_i of the gauge group G_C. The minimal possibility satisfying the above requirements is provided by $G_F = SU(4) \to H_F = Sp(4)$, that corresponds to four Weyl fermions in a pseudoreal representation of G_C. The simplest pseudoreal representation is the fundamental of a group $Sp(2n)$. Thus, we are led to choose as hypercolour group $G_C = Sp(2N_C)$, with Weyl fermions $\psi^a \sim \Box_{Sp(2N_C)}$, where $a = 1, 2, 3, 4$ is the flavour index.

Once the hypercolour theory confines, the constituent degrees of freedom, ψ^a and the hypergluons, are replaced by composite, hypercolour-singlet states. They are associated to operators constructed out of the constituent fields, in given Lorentz and flavour representations. Let us limit ourselves to fermion-bilinear operators, which excite several spin-0 and spin-1 composite states, including the NGB Higgs, as illustrated in table 1. Scalars organise into a flavour-singlet $\sigma \sim 1_{Sp(4)}$ and a flavour-multiplet $S^{\hat{A}} \sim 5_{Sp(4)}$, where $\hat{A}$ runs over the five broken generators. Pseudoscalars sit in the same representations, $\eta' \sim 1_{Sp(4)}$ and $G^{\hat{A}} \sim 5_{Sp(4)}$. The latter is the NGB multiplet, that is massless in the chiral limit, $G = \{h, \eta\} \sim \{(2_L, 2_R), (1_L, 1_R)\}$: note that the Higgs doublet is accompanied by an electroweak singlet state. On the other hand η' is expected to be massive, because the associated flavour symmetry, an axial $U(1)_\psi$, is anomalous with respect to G_C, in analogy with the axial $U(1)$ in QCD. Coming to spin-one states, vectors organise in a multiplet $V^A_\mu \sim 10_{Sp(4)}$, where A runs over the ten unbroken generators. Axial vectors transform as $a_\mu \sim 1_{Sp(4)}$ and $A^{\hat{A}}_\mu \sim 5_{Sp(4)}$. It is also possible to establish spectral sum rules [1], that relate the masses and decay constants of the various states.

Before discussing the dynamics, let us generalise the model to the case of a large number of flavours N_F. In fact, in a realistic model the group G_F should contain several other symmetries besides $SU(2)_L \times SU(2)_R$. Firstly, in order to induce the SM Yukawa couplings, one needs to mix the SM fermions with composite operators. The latter should have the same colour and electroweak charges as the various SM

	Lorentz	$Sp(2N_C)$	$SU(4)$	$Sp(4)$
ψ_i^a	$(1/2,0)$	$\square_i$	4^a	4
$\overline{\psi}_{ai} \equiv \psi_{aj}^\dagger \Omega_{ji}$	$(0,1/2)$	$\square_i$	$\overline{4}_a$	4^*
$M^{ab} \sim (\psi^a \psi^b)$	$(0,0)$	1	6^{ab}	$5+1$
$\overline{M}_{ab} \sim (\overline{\psi}_a \overline{\psi}_b)$	$(0,0)$	1	$\overline{6}_{ab}$	$5+1$
$a^\mu \sim (\overline{\psi}_a \overline{\sigma}^\mu \psi^a)$	$(1/2,1/2)$	1	1	1
$(V^\mu, A^\mu)_a^b \sim (\overline{\psi}_a \overline{\sigma}^\mu \psi^b)$	$(1/2,1/2)$	1	15_b^a	$10+5$

Table 1: The transformation properties of the constituent fermions, and of the spin-0 and spin-1 fermion bilinears, in the $N_F = 2$ model. The hypercolour $Sp(2N_C)$ indexes $i, j, \ldots$ are contracted by the antisymmetric invariant tensor Ω_{ij}, and brackets stand for hypercolour-invariant contractions. Spinor indexes are understood, and $a, b, \ldots$ are flavour $SU(4)$ indexes.

fermions, therefore the whole $SU(3)_c \times SU(2)_L \times U(1)_Y$ needs to be embedded within G_F. Secondly, the SM global symmetries, such as baryon and lepton number, or custodial, should be also included in G_F, to avoid that hypercolour dynamics violates these symmetries too strongly. Thus, one is led to introduce additional constituent fermions, $\psi^a \sim \square_{Sp(2N_C)}$, with $a = 1, \ldots, 2N_F$, corresponding to the flavour-symmetry-breaking pattern $G_F = SU(2N_F) \to H_F = Sp(2N_F)$, with $N_F \gtrsim 5$ depending on the model details. [2] One also needs [4, 1] to introduce constituents fermions X in larger representation of $Sp(2N_C)$, in order to build hypercolour-singlet trilinear operators such as $(\psi\psi X)$, that interpolate fermionic composite states, such as top-quark partners. We argue that, to preserve asymptotic freedom, it is preferable to minimise the number of X flavours, and rather assign the required SM charges to the $2N_F$ copies of ψ. Here we will neglect the X sector, and discuss only the spectrum of ψ-bilinear operators.

To go beyond the symmetry considerations above, and derive a quantitative estimate of the mass spectrum, one needs to model the hypercolour dynamics, either numerically on the lattice, or by some analytical approximations in the large-N_C limit. We will show that the latter provide relatively rapid and general estimates for the spectrum, complementarily to lattice computations, which are currently limited to $Sp(2N_C)$ theories with $N_C = 1, 2$. [5] In order to determine the spectrum of composite states associated to a given operator, one has to determine the poles of the associated two-point correlation function. Let us consider, for illustration, the case of vector currents,

$$i \int d^4x \, e^{iq\cdot x} \langle \text{vac}| T\{\mathcal{J}_\mu^A(x) \mathcal{J}_\nu^B(0)\} |\text{vac}\rangle = \Pi_V(q^2) \delta^{AB} (q_\mu q_\nu - \eta_{\mu\nu} q^2) \quad (1)$$

where $\mathcal{J}_\mu^A = \overline{\psi} \overline{\sigma}_\mu T^A \psi$. In the large-$N_C$ limit one expects the form factor to behave as a sum over narrow resonances, $\Pi_V(q^2) \simeq \sum_n f_{Vn}^2 (q^2 - m_{Vn}^2)^{-1}$. Our aim is to estimate the position of the poles, m_{Vn}^2, and similarly for other two-point correlators. We will discuss two methods that provide an analytic approximation for such correlators.

3 Spectrum of mesons à la Nambu-Jona Lasinio

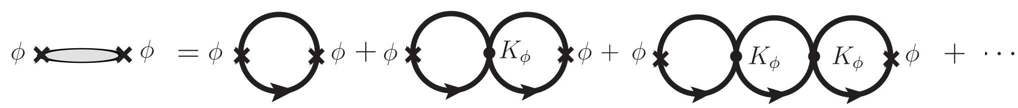
The Nambu-Jona Lasinio (NJL) model approximates strong dynamics by effective four-fermion interactions. This corresponds to give a dynamical mass to the gauge bosons and decouple them, writing an effective Lagrangian for the constituent fermions only. For the $Sp(2N_C)$ hypercolour theory, the

$SU(2N_F)$-invariant Lagrangian for $N_F = 2$ reads [4]

$$\mathcal{L}_{scalar} = \frac{\kappa_A}{2N_C}(\psi^a\psi^b)(\overline{\psi}_a\ \overline{\psi}_b) - \frac{\kappa_B}{8N_C}\left[\epsilon_{abcd}(\psi^a\psi^b)(\psi^c\psi^d) + h.c.\right] , \qquad (2)$$

where for simplicity we included only scalar-scalar operators. The κ_A operator is induced by a tree-level hypergluon exchange, while the κ_B operator accounts for the anomaly of the axial $U(1)_\psi$ symmetry. One can show [4, 1] that this NJL Lagrangian can describe spontaneous breaking $SU(4) \to Sp(4)$, by inducing a non-zero mass gap, $N_C M_\psi = (\kappa_A + \kappa_B)\langle\psi\psi\rangle \neq 0$, where M_ψ is the dynamical mass for the fermions.

To estimate two-point correlators, one can resum massive fermion loops, at leading order in $1/N_C$:



Here ϕ is the meson associated with a given fermion bilinear, K_ϕ is the corresponding four-fermion coupling, and the resummation describes the composite meson propagator,

$$\overline{\Pi}_\phi(q^2) \equiv \frac{\tilde{\Pi}_\phi(q^2)}{1 - 2K_\phi\tilde{\Pi}_\phi(q^2)} , \qquad (3)$$

where $\tilde{\Pi}_\phi$ is the one-loop function. The resummation of the geometric series induces a pole in the composite propagator $\overline{\Pi}_\phi$, for some specific value of q^2, that defines the meson mass in the NJL approximation.

In fig.1 we show our results for the pole of each meson correlator, as a function of the dimensionless four-fermion coupling $\xi \equiv (\kappa_A+\kappa_B)\Lambda^2/(4\pi^2)$, where Λ is the cutoff of the fermion loops. One can check [1] that $\xi \geq 1$ is needed to induce a non-zero M_ψ and global symmetry breaking, while $\xi \leq (1 - \ln 2)^{-1}$ is needed for the mass gap not to exceed the cutoff, $M_\psi \leq \Lambda$. The NGB decay constant f can also be computed [1] as a function of ξ: in fig.1 the meson masses are given in units of f. Note the pole positions do not scale with N_C, however $f \sim N_C^{1/2}$, therefore the physical masses decrease with the number of colours if $f \simeq$ TeV is kept fixed. The only exception is the η' pole, that scales as $N_C^{-1/2}$, as the axial anomaly vanishes in the large-N_C limit.

Assuming the dynamics is dominated by a current-current operator (corresponding to a single hypergluon exchange), one can relate the scalar-scalar and vector-vector operators by using Fierz identities: this fixes the relative size of spin-zero and spin-one meson masses. The latter are always heavy, $\gtrsim 5f$, while spin-zero mesons can become light in several cases. First, NGBs are massless, $M_h = M_\eta = 0$, as we neglected possible sources of $SU(4)$ explicit breaking. Second, the singlet pseudoscalar η' also becomes light in the large-N_C limit. Third, the singlet scalar σ becomes light as the four-fermion coupling approaches the critical value $\xi = 1$. This lightness indicates that the four-fermion operator becomes marginal as $\xi \to 1$, that is, the explicit breaking of scale invariance vanishes, and σ can be interpreted as an approximate dilaton.

4 Spectrum of mesons via gauge-gravity duality

If the hypercolour sector is close to a fixed point, it behaves as an approximately Conformal Field Theory (CFT). The CFT in the limit of large number of colours, N_C, and large 't Hooft coupling, $\lambda \equiv g_C^2 N_C$, has a holographic description in terms of a five-dimensional (5d) theory of gravity in the classical and weakly-coupled regime, with Anti-de Sitter (AdS) metric, $\mathrm{d}s^2 = \mathrm{d}r^2 + e^{2A(r)}\mathrm{d}x_{1,3}^2$ with warp factor $A(r) = r$. [6] Holography implications for composite Higgs scenarios are reviewed in Ref. [7].

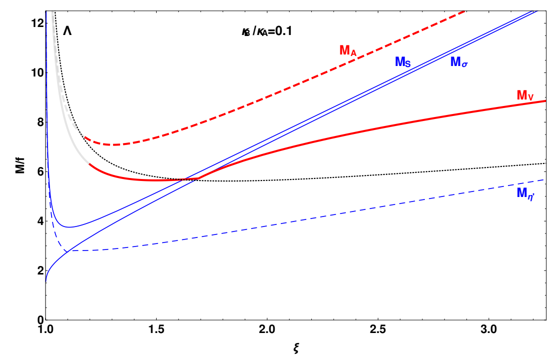


Figure 1: The masses of the spin-zero (blue) and spin-one (red) mesons in units of the Goldstone decay constant f, as a function of the dimensionless four-fermion coupling $\xi \equiv (\kappa_A + \kappa_B)\Lambda^2/(4\pi^2)$, for $\kappa_B/\kappa_A = 0.1$ and $N_C = 4$. The Goldstone multiplet (not shown) is massless, $M_G = 0$, and the two axial-vector multiplets are degenerate, $m_a = M_A$. See Ref. [1] for more details.

According to the holographic dictionary, the CFT global symmetry corresponds to a 5d gauge symmetry G_F, and CFT operators O_Φ are associated to 5d fields Φ, in the same G_F representation and with the same spin. Moreover, CFT correlators correspond to 5d correlators built from the bulk action on-shell, $S_{bulk}^{on-shell}$, in particular they scale in the same way with N_C and N_F. [2] For example, the glue-glue correlator $\langle G_{ij} G_{ij} \rangle \sim N_C^2$ can be extracted from a 5d gravity action, $S_{bulk}[R] \propto N_C^2$, with R the Ricci scalar. On the other hand, the fermion-fermion correlator $\langle \psi_i^a \psi_j^a \Omega_{ij} \rangle \sim N_C N_F$ is associated to a 5d scalar action, $S_{bulk}[\text{Tr } \Phi^{ab}] \propto N_C N_F$, with a 5d scalar Φ^{ab} dual to the operator $(\psi^a \psi^b)$.

The CFT departure from scale invariance in the IR with a mass gap m_* can be described by adding a 5d scalar field with non-flat profile, $\sigma(r)$. The latter back-reacts on the metric, inducing a warp factor $A(r) \neq r$, that is, a departure from AdS. Let us consider the case $\sigma(r) \equiv \text{Tr } [\Phi^{ab}(r)]/N_F$, that is, the scalar associated to flavour-symmetry breaking. The gravity-scalar interplay is described by

$$S_{bulk} \supset N_C^2 \int \mathrm{d}^5x \sqrt{-g} \left[\frac{R}{4} - \frac{\Lambda_c}{2} - x_F \left(\frac{1}{2} g^{MN} \partial_M \sigma \partial_N \sigma + V(\sigma) \right) \right] \tag{4}$$

where $x_F \equiv N_F/N_C$ and Λ_c a cosmological constant. We are interested in the large N_F case, to accommodate all the required SM global symmetries. Thus, we are led to consider the Veneziano limit, with large N_C and constant $x_F \sim 1$. This implies that the back-reaction of the flavour sector on the 5d geometry is an order-one effect. Indeed, for some appropriate choice of the potential $V(\sigma)$, motivated by string theory compactifications, the equations of motion imply [2] that $A(r)$ and $\sigma(r)$ have a singularity at some finite value $r = r_{IR}$, which corresponds to the dynamical generation of a mass gap, $m_* \neq 0$. This opens the possibility to relate m_* and the G_F spontaneous-breaking scale f, that is associated to $\langle O_\sigma \rangle$.

The nature of scale-invariance breaking is determined by the UV behaviour of $\sigma(r)$,

$$\sigma(r) \underset{r\to\infty}{\simeq} \left(\sigma_- e^{-\Delta_- r} + \sigma_+ e^{-\Delta_+ r} \right) \; , \quad \Delta_\pm = 2 \pm \sqrt{4 + m_\sigma^2} \; , \tag{5}$$

where the bulk mass m_σ as well as the values of $\sigma_\pm$ depend on the choice of $V(\sigma)$. We are interested in the regime $-4 \leq m_\sigma \leq 0$, corresponding to a departure from scale-invariance in the IR. The dual operator $\mathcal{O}_\sigma$ has scaling dimension Δ_+. A non-zero σ_- corresponds to a relevant deformation of the CFT, $\Delta\mathcal{L}_{CFT} = \sigma_- \mathcal{O}_\sigma$. This amounts to explicit breaking of scale invariance and G_F: the CFT couplings acquire non-zero β-functions, and there are no massless dilaton nor exact NGBs. A non-zero σ_+ corresponds to a vacuum

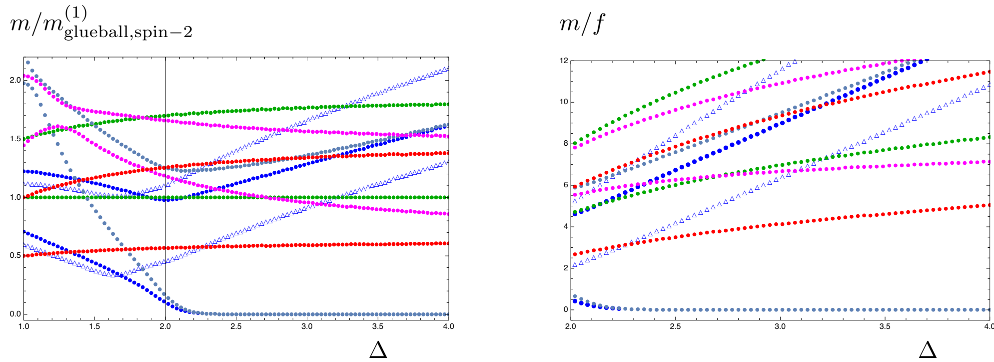


Figure 2: The masses of the flavour-singlet scalars (blue), flavour-multiplet scalars (blue triangles), pseudoscalars (cyan), vectors (red), axial vectors (magenta) and spin-two glueballs (green), in units of the lightest spin-two mass (left panel) and in units of the Goldstone decay constant f (right panel), as a function of the anomalous dimension Δ.

expectation value, $\langle \mathcal{O}_\sigma \rangle \sim \sigma_+$, that controls spontaneous breaking of scale invariance and G_F. In models with $\sigma_- \to 0$ one then finds a massless dilaton as well as massless NGBs.

In order to estimate the composite mass spectrum, let us extract the poles of two-point correlators using the gauge-gravity duality. The solution of the equations of motion for $A(r)$ and $\sigma(r)$ fixes the 5d background. One can expand S_{bulk} around such background, to quadratic order in the field fluctuations, for any 5d field ϕ_i, dual to the CFT operator of interest $\mathcal{O}_i$. Solving the equations of motion linear in the fluctuations, one can compute S_{bulk} on shell. The latter determines the CFT correlators, according to

$$\langle \mathcal{O}_1 \mathcal{O}_2 \rangle = \lim_{r\to\infty} \frac{\delta^2 S^{on-shell}_{bulk}[\phi_i]}{\delta\phi_1 \delta\phi_2} \ . \tag{6}$$

Let us consider the CFT correlator in momentum space, and call the 4d momentum q. The 5d field fluctuations satisfy appropriate boundary conditions [2)] only for discrete values of q^2, that correspond to the mass of the composite states. For example, the axial-vector transverse correlator takes the form

$$\langle J^\mu(q) J^\nu(-q) \rangle = -\lim_{r\to\infty} \frac{\delta^2 S^{on-shell}_{bulk}}{\delta A_\mu(-q,r)\delta A_\nu(q,r)} \sim \lim_{r\to\infty} \left[e^{2A(r)} (\eta^{\mu\nu} - q^\mu q^\nu/q^2) \frac{\partial_r A_\rho(q,r)}{A_\rho(q,r)} \right] \ . \tag{7}$$

In this case the poles are given by the values of q^2 where the 5d gauge field vanishes asymptotically, $\lim_{r\to\infty} A_\rho(q,r) = 0$. Moreover, the value of f^2 is given by the residue of this correlator at $q^2 = 0$.

In the left panel of fig.2 we show our preliminary result for the spectrum of bosonic resonances (spin 0, 1 and 2), as a function of the parameter Δ, defined by $\sigma(r) \sim e^{-\Delta r}$ for $r \to \infty$, in units of the smallest spin-two mass. The dilaton and NGBs remain massless for $2 < \Delta < 4$ (spontaneous symmetry breaking), while they are lifted for $0 < \Delta < 2$ (explicit symmetry breaking). With the given choice of model parameters the dilaton mass grows faster than the pNGB one. Scale invariance may be broken explicitly by additional flavour-singlet sources, that raise the mass of the dilaton only. [2)] From the phenomenological perspective, one should include radiative corrections from SM couplings, in particular the top quark Yukawa, before comparing the pNGB mass with the 125 GeV scalar observed at the LHC. An even lighter dilaton cannot be excluded, as its couplings to the SM are suppressed. [8)] In the spontaneous breaking regime, one can estimate f and display the spectrum in units of f, see the right panel of fig.2. We remark that the exact value of f is determined by the overall normalisation of the gauged-G_F kinetic term. 2) While the f scaling with N_C and N_F is generic, the order-one normalisation can be predicted only in a complete top-down model coming e.g. from a specific string theory.

5 Perspective

We showed that, in UV-complete models for the Higgs compositeness, a plethora of composite states are expected to accompany the Higgs boson. We focused on the composite bosons of the hypercolour theory, associated to fermion-bilinear operators. Similarly, one can consider composite fermions, associated to fermion-trilinear operators, relevant to induce Yukawa couplings, especially the large top-quark one. In order to study the strongly-coupled sector, one needs to model non-perturbative effects, making radical assumptions to simplify the dynamics. We considered the NJL model and a gauge-gravity duality model, showing that they catch several essential features of the mass spectrum, and thus provide an important guidance for future searches, even though quantitative estimates are model-dependent. The crucial phenomenological question is whether some new states could be significantly lighter than the compositeness scale, $m_* \sim 10$ TeV. The rationale for light spin-zero states is to realise a hierarchically large (small) spontaneous (explicit) breaking scale for global symmetries. Light spin-one-half states may occur in the case of approximate chiral symmetries of the hypercolour theory. Our models provide exploratory tools to characterise such composite physics at the high energy frontier.

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