Review Article

Yukawa Bound States and Their LHC Phenomenology

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We present the current status on the possible bound states of extra generation quarks. These include phenomenology and search strategy at the LHC. If chiral fourth-generation quarks do exist their strong Yukawa couplings, implied by current experimental lower bound on their masses, may lead to formation of bound states. Due to nearly degenerate 4G masses suggested by Precision Electroweak Test one can employ “heavy isospin” symmetry to classify possible spectrum. Among these states, the color-octet isosinglet vector ω_8 is the easiest to be produced at the LHC. The discovery potential and corresponding decay channels are covered in this paper. With possible light Higgs at ~ 125 GeV two-Higgs doublet version is briefly discussed.

1. Introduction

The observed baryon asymmetry of the universe (BAU) cannot be accounted for within the current particle physics theory, the Standard Model (SM). An additional heavy chiral fourth generation (4G) could provide much needed enhancement of CP-violation for BAU [1, 2], while potentially providing or being part of the mechanism for electroweak symmetry breaking [3–7]. The latter phenomenon is thought to happen through the condensation of $Q\bar{Q}$ state due to strong Yukawa attractive potential. Ever since Nambu’s insight to symmetry breaking and subsequent works with Jona-Lasinio [8] this idea has been with us. Although it is quite old, in advent of the Large Hadron Collider (LHC), this is the first time we may be able to experimentally see whether this happens in Nature. Here we aim to review current status and near future prospect for phenomenological study of Yukawa bound states of the heavy 4G quarks.

Current lower bound from direct searches for the 4G chiral quarks put their Yukawa couplings well beyond perturbative limit. Often cited upper limit from the perturbative unitarity violation (UV) [9–12] for 4G has now been reached at the LHC. Therefore if they exist, their Yukawa couplings are necessarily in nonperturbative regime. Phenomenologically most interesting consequence is the existence of meson-like $Q\bar{Q}$ bound states of these heavy quarks around TeV mass region. The principal binding force is that of the

Yukawa potential. This fact essentially fixes what type of states one would expect once proper quantum numbers, such as heavy “isospin,” which we define shortly, and spin are given. Over the years there have been many works on collider phenomenology of extra generations, for example, by Holdom [13, 14], to which the bound state study covered here is complimentary and many more important works dealing with the different aspects of the subject are not mentioned due to lack of space.

An important information which helps to classify the spectrum comes from the precision electroweak test [15, 16]. The constraint favors a small mass splitting between t' and b' , the 4G up- and down-type quarks, respectively. This near degeneracy in the heavy mass limit, combined with the fact that their mixings to lighter generations are constrained to be small by various flavor constraints, leads to an approximate “isospin” symmetry [17–19]. This approximate symmetry helps one to classify the bound states and identify leading production and decay processes. For example, taken the exact limit, one can see that isotriplets must be produced in pairs; therefore in a realistic situation, the single production is suppressed by small mass splitting between t' and b' and/or small mixing to a lighter generation. From this one can identify most interesting modes as far as the 7 and 8 TeV LHC run are concerned. The Higgs also has to be heavy for the precision test and now it is ruled out between 120 and 600 GeV. Recently announced signals at 5σ from both CMS

and ATLAS [20, 21] point toward more complex situation than just the simple extension of the SM by an extra generation. In particular, the signals are thus far consistent with the SM Higgs and, if confirmed, the simple extension by 4G is ruled out. On the other hand, the question regarding BAU within the SM will remain unresolved with the confirmation of the existence of Higgs. In light of this fact the possible solutions can still include extra generation and/or multi-Higgs scenarios. A recent study of 4G SM with two-Higgs doublet remains consistent with the Precision Electroweak Constraint [22] with 125 GeV Higgs.

In this report, we give a brief review of the status of the LHC phenomenology of Yukawa bound states mainly along the line of [19] by Enkhbat et al. We update these analyses with the available results from the LHC 2012 run. The structure of the review is as follows. In Section 2, we review descriptions of the Yukawa bound states. In Section 3, the expected signal types and their rates are given. Section 4 is devoted for further phenomenological discussions. In Section 5, a two-Higgs doublet model extension is discussed. The conclusion is given in Section 6.

2. Yukawa Bound State Descriptions

The current direct search of sequential fourth generation quarks, that is, one additional generation of left-handed weak doublet and right-handed weak singlet, has now reached the lower limit of $m_Q = 611$ GeV at 95% C.L. [23]. Their existence means that there is a new nonperturbative regime of Yukawa interaction. Such a strong regime will lead to a bound state of these quarks. Whether they are part of electroweak symmetry breaking or not, a phenomenological study of the bound states due to the strong Yukawa potential is not only interesting but very crucial for heavy quark search program.

Emergence of a strong dynamics could lead to a rich spectrum and phenomenology as witnessed in the case of QCD. Although the tree-level UV is related to more asymptotically higher energy unlike confinement in the infrared regime there is a lesson from the QCD and chiral perturbation theory. Traditionally the tree-level unitarity constraint has been considered as some upper limit for the masses of possible extra heavy fermions above which some new degree of freedom or new dynamics should appear to cure the apparent unitarity violation. In reality it is a perturbative statement and should be considered with a grain of salt. One even could argue that the effective field theory should be able to cure UV if a correct frame work is supplemented. In particular, the chiral perturbation theory with only light pions has UV at $\sqrt{s} \simeq 450$ MeV in the scalar-isoscalar channel, which can be treated by inverse amplitude method [24]. The justification of this approach is provided by existence of σ -like object $f_0(600)$. Regarding the UV, there is another dynamical argument based on large- N expansion, which points to much higher bound [25] around ~ 3 TeV for 4G quark masses.

To examine what bound states are phenomenologically more interesting one needs a tool to identify the quantum numbers of these states. An important constraint on the 4G spectrum is the precision electroweak test, which requires the

t' and b' masses to be nearly equal. This “new isospin” makes it easier to list the lowest lying states. Unlike technicolor-like theories, one expects colored objects as they are far above the QCD scale. So for every color singlet meson-like states there is a color-octet counterpart.

Here we use the same QCD nomenclature for these states and add subscripts 1 and 8 as the singlet and octet color representations, respectively. Thus we denote isovector states $\bar{b}'t', \bar{t}'b', (\bar{t}'t' - \bar{b}'b')/\sqrt{2}$ as $\pi_{1,8}, \rho_{1,8}$ and isoscalars $(\bar{t}'t' + \bar{b}'b')/\sqrt{2}$ as $\eta_{1,8}, \omega_{1,8}$.

At this scale the QCD coupling is subleading compared to the Yukawa couplings. The potential from the Yukawa couplings consists of contributions from the Higgs and pseudoscalar Goldstones. While the Higgs part is universally attractive the Goldstone potential crucially depends on the isospin and spin of the state. Since we are considering a heavy Higgs scenario for the electroweak precision constraints, the main binding force is that of the Goldstone exchange. One can quickly establish the following facts by examining the Goldstone exchange diagrams: while isotriplet scalars $\pi_{1,8}$ and isosinglet vector $\omega_{1,8}$ states have attractive Goldstone potential, isosinglet scalar $\eta_{1,8}$ and isotriplet vector $\rho_{1,8}$ have repulsive one. Therefore η - and ρ -like states are most likely unbound while the lower bound states are π 's and ω 's. This is quite distinct compared to the well-studied scenario of technicolor, and its colored versions where one has ρ -like resonance as the defining prediction. Also the colored meson ω_8 may be accessible since it can be produced singly at the LHC.

More quantitative treatments for a strongly bound relativistic bound state remain challenging. The most reliable one is the lattice approach which is still being pursued. Thus far the best estimates have been provided by Bethe-Salpeter approach, which has many shortcomings. There is a crucial qualitative difference between the Bethe-Salpeter and simpler approaches of relativistic expansion based on variational method. The former takes into account the contribution of mixing between positive and negative frequency solutions. This has been shown to have a dramatic effect in [17] as illustrated in the case of pions. In the following we review essential results of this study and one based on relativistic expansion [18].

2.1. Relativistic Bethe-Salpeter Approach. Here we explain the spectrum emerging from the study of relativistic bound states based on Bethe-Salpeter equation. In the early nineties, in a series of papers Jain et al. [17] studied strong Yukawa bound states. They have shown using Bethe-Salpeter equations numerically that in the strong Yukawa coupling limit highly relativistic deeply bound states form. These are coupled set of integral equations applied for the bispinor wave function of the $Q\bar{Q}$ bound states. The bispinor amplitude $\chi_{\pm}^{\pi}(q)$ for exchange momentum q between Q and $\bar{Q}$ has two components corresponding to the negative and positive frequencies. The mixing between these states, which is not taken into account in a simple relativistic expansion approach [18], plays a crucial role. In the following we give the integral equation, without the explicit form of its kernel $V(q, q')$ (Bethe-Salpeter

FIGURE 1: The V_+^π (solid) and V_-^π (dashed) potentials in (1) for the π_1 state, for $m_Q = 400, 500, 600, 700$ GeV, and $m_H = 700$ GeV, plotted along $q' = q$. Both vanish for $q = 0$, or zero relative motion, but go to rather large values for large q , with V_-^π turning on rather sharply for $q \gtrsim 100$ GeV.

potential) which can be found in [17]. Instead we plot it for the lowest lying state π_1 along the steepest direction $q = q'$, which suffices to show the importance of this approach. The Bethe-Salpeter equations for $\chi_\pm^\pi(q)$ are given by

$$(M \mp 2\omega) \chi_\pm^\pi(q) = \pm \int dq' \frac{q'}{q} [V_+^\pi(q, q') \chi_+^\pi(q') + V_-^\pi(q, q') \chi_-^\pi(q')]. \quad (1)$$

In the heavy isospin limit the possible states are classified by their definite isospins $I = 0, 1$.

In Figure 1, the Bethe-Salpeter potential for π_1 along the steepest direction $q = q'$ is plotted for various choices of the heavy quark mass m_Q . $V_\pm$ are the diagonal and mixing contributions, respectively, for χ_+ and χ_- amplitudes. As we can see the mixing potential V_- from the negative frequency solution is dramatically enhanced as we increase the mass. This contribution causes a “collapse” of the state; namely the bound state mass becomes zero, at a much lower value of m_Q compared to the case when the mixing is ignored [17]. Similar feature can be seen for $\omega_{1,8}$ states. On the other hand this dominant contribution is repulsive for ρ - and η -like states. For this reason the lowest π and ω states are the focus of phenomenological study for the LHC. While a small amount of mass splitting is required between b' and t' for electroweak precision constraint, here the degenerate case is taken for simplicity. In [17], the authors did not discuss color-octet mesons which are the more interesting objects for a hadron collider like the LHC. Nevertheless, they have calculated the isosinglet and isotriplet meson masses in different formalisms. Among these, Bethe-Salpeter formalism with both positive and negative frequency solutions leads to strongest binding and one can see that it even gives a collapse at ~ 500 GeV. Their numerical solution as binding energy versus the 4G mass is shown in Figure 2.

FIGURE 2: The masses of color singlet isotriplet π and isosinglet states versus the 4G quark mass from [17] in various formalisms. Curves S1 (iso-singlet ω_1) and T1 (isotriplet π_1) are the results from covariant gauge formalism. S2 (isosinglet) and T2 (isotriplet) are the results from Salpeter’s equation using only positive frequency components. S3 (isosinglet) and T3 (isotriplet) are the results using both positive and negative frequency components. For more details of these formalisms see [17].

FIGURE 3: Variational binding energy of color singlet channels as the function of the heavy quark mass $m_q \equiv m_Q$.

Ishiwata and Wise studied the Yukawa bound states by variational method in relativistic expansion approach in [18]. The essence of their result is captured in the binding energy obtained by minimization of the $Q\bar{Q}$ potential. We show their plot here in Figure 3, where they chose $m_h = 130$ GeV and get similar result for $m_h = m_Q$. Around $m_Q \approx 400$ GeV Yukawa contribution becomes dominant where binding energy changes dramatically. This “kink” behavior is due to the constraint $a \geq \sqrt{3}/m_Q$ they put to have valid

FIGURE 4: Production mechanism for the color-octet, isoscalar-vector meson ω_8 at hadron colliders.

relativistic expansion. Here a is the length characterizing the size of the bound system. From this one can conclude that the relativistic expansion, although limited by the constraint on a , also indicates that the Yukawa binding becomes substantial. Although the relativistic expansion leads to qualitatively similar spectrum to the Bethe-Salpeter method, it nevertheless does not take into account the crucial contribution of the mixing from negative frequency solution.

In this paper, we mainly review the scenario pursued in [19], where the heavy quark mass m_Q was taken to be within 500 to 700 GeV interval. Based on the numerical analysis of Jain et al. the binding energy was chosen to be up to 200 GeV. The latest bound on 4G quark mass has now ruled out the lower half of this range and with the coming new set of data the bound will be further raised probing ever stronger Yukawa coupling domain. Keeping this in mind we will also touch on the discussion for higher values of m_Q up to $\sim$ TeV values.

Since π_1 and ω_1 are not produced by gg -fusion but by weak Drell-Yan processes, energywise they are not efficiently produced. π_8 can be produced via strong interaction in pairs from $\bar{q}q$ or gg scattering [26], which makes it low in luminosity at 7 and 8 TeV. Therefore, at least in the near future, ω_8 , which has the same quantum number as the gluon, is the most interesting object from LHC stand. It cannot be produced by gg -fusion since two massless vector particles cannot fuse to one massive vector particle according to the Landau-Yang theorem [27, 28]. This leaves $\bar{q}q$ annihilation as the principle production mechanism for ω_8 . The ordering of the spectra, according to [17] (which did not actually consider octet states), would be $M_{\pi_1} \leq M_{\omega_1} \leq M_{\pi_8} \leq M_{\omega_8}$. In the next section, we review the production and decay properties of the ω_8 meson.

3. Production and Decay of ω_8

In this section, we start the collider phenomenology discussion with the review of the ω_8 production and decay. We have concluded that it is the lowest lying colored object that can be singly produced. The other octet states are either pair produced due to the new *isospin* or heavier even potentially unbound. The following input parameters are chosen:

- (i) the ω_8 decay constant $\xi \equiv f_{\omega_8}/m_{\omega_8}$ defined by

$$\langle 0 | V^{\mu,a} | \omega_8^b(p) \rangle \equiv \frac{1}{\sqrt{2}} \delta^{ab} f_{\omega_8} m_{\omega_8} \varepsilon^\mu(p), \quad (2)$$

- (ii) the fourth-generation quark mass $m_Q = m_{t'} = m_{b'}$ (this degeneracy led to heavy isospin),
- (iii) the ω_8 mass m_{ω_8} (this is essentially the choice of binding energy $2m_Q - m_{\omega_8}$),
- (iv) the quark mixing elements $|V_{t'b}| = |V_{tb'}|$, where we assume that mixings to lighter generations are negligible.

In principle if the binding mechanism is quantitatively established the decay constant and binding energy are no longer free parameters. As mentioned before, the most reliable approach is the lattice field theory and it is far from reaching the required level at this point and therefore is beyond the scope of this report. What is done thus far is the qualitative estimates of the possible spectrum in Bethe-Salpeter and relativistic expansion approaches which fail to address these quantities. Therefore these quantities are chosen as free parameters for the phenomenological study.

3.1. Production. As it has been stressed earlier, the color-octet isospin singlet vector state ω_8 appears to be most prominent among the interesting objects for the LHC. ω_8 is dominantly produced by $\bar{q}q$ annihilation as shown in Figure 4 instead of gluon fusion. The latter is forbidden due to the Landau-Yang theorem [27, 28], which states that a massive vector particle cannot decay to two massless vector particles. One can explicitly check this fact.

At the parton level, the total cross-section is

$$\hat{\sigma}_{\bar{q}q \rightarrow \omega_8}(\hat{s}) = \frac{32\pi^3 \alpha_s^2}{9m_{\omega_8}^2} \xi^2 \delta\left(1 - \frac{m_{\omega_8}^2}{\hat{s}}\right). \quad (3)$$

Convoluting with the parton luminosity

$$\mathcal{L}(\tau; \mu_F^2) = \iint dx_1 dx_2 f_1(x_1; \mu_F^2) f_2(x_2; \mu_F^2) \delta(\tau - x_1 x_2), \quad (4)$$

we get the hadronic cross-section:

$$\sigma(s) = \int d\hat{\tau} \hat{\sigma}(\hat{\tau}s) L(\hat{\tau}; \mu_F^2). \quad (5)$$

The inclusive production cross-section σ_{ω_8} for ω_8 is given in Figures 5 and 6. The left- and right-hand sides of Figure 5 are the $\sqrt{s} = 7$ TeV and 14 TeV, respectively, while Figure 6 is for 8 TeV [19]. Here CTEQ6L [29] parton distribution functions (PDFs) are used with the renormalization and factorization scales to $\mu_R = \mu_F = m_{\omega_8}$. The decay constant parameter defined in (2) is chosen at three different values $\xi = f_{\omega_8}/m_{\omega_8} = 0.1, 0.03$, and 0.01 for illustration. One could intuitively argue that one expects stronger binding for large decay constant ξ . Because one lacks a reliable numerical tool to probe the highly relativistic system, in [19] the authors

FIGURE 5: Production cross-section of ω_8 at the LHC running at 7 TeV (a) and 14 TeV (b) for various $\xi = f_{\omega_8}/m_{\omega_8}$ values. The open $\text{Q}\bar{\text{Q}}$ cross-sections at LO and NLO are also plotted for comparison.

FIGURE 6: Production cross-section of ω_8 at the LHC running at 8 TeV.

chose these values, which one encounters in QCD mesons [30–32]. One should keep in mind that the system we are discussing here is not a confining QCD-like model. The binding energy was chosen constant to be 100 or 200 GeV

treating it as a free parameter. Therefore the production cross-section ranges from pb to fb in magnitude depending only on the choice of ξ . Due to higher gluon density at the LHC one might think that the production via higher order scattering process $gg \rightarrow \omega_8 g$ could be important. It turns out, upon explicit calculation, that for the range $m_{\omega_8} > 0.8$ TeV the rate is negligibly small.

Another question is whether the production of the heavy meson state is completely dominated by the open production of $\text{Q}\bar{\text{Q}}$. For this, the pair production cross-section is calculated at LO and NLO [33, 34] and the results are shown in the same figure. For NLO calculation CTEQ6M PDFs [29] are used. The dotted (hatched) band shows the range of uncertainty when the renormalization scale is varied from $\mu_R = \mu_F = m_Q$ to $4m_Q$ at LO (NLO). From this we see that the heavy meson ω_8 production can be at the same order or even bigger than the open pair production if the decay constant, therefore binding energy, is larger even though $q\bar{q}$ annihilation is the mechanism for the former while the latter occurs dominantly through gg -fusion at the LHC. At 14 TeV the production cross-section is an order of magnitude larger than the lower energy runs and open production is increased at higher m_Q values compared to ω_8 production due to higher gluon luminosity.

3.2. Decay. Due to the lower lying states, ω_8 has possible meson transition-like decays. Such channels are experimentally more distinct which will be a clear signal of the bound

states. In the following, possible leading channels are given with corresponding formulae. Using these it has been shown in [19] that such a meson transition-like decay indeed is sizable and often times dominates. The decay channels of ω_8 considered are as follows.

- (i) Annihilation decay: $\omega_8 \rightarrow q\bar{q}, t\bar{t}; t\bar{t}', b\bar{b}'$; the decay rates for these processes are

$$\begin{aligned}\Gamma(\omega_8 \rightarrow q\bar{q}) &= \xi^2 \frac{\pi\alpha_s^2}{3} m_{\omega_8} n_f, \\ \Gamma(\omega_8 \rightarrow t\bar{t}) &= \xi^2 \frac{\pi\alpha_s^2}{3} m_{\omega_8} \sqrt{1 - \frac{4m_t^2}{m_{\omega_8}^2}}, \\ \Gamma(\omega_8 \rightarrow t\bar{t}') &= \xi^2 |V_{tb}^* V_{t'b'}|^2 \frac{G_F^2 m_{\omega_8}^5}{192\pi} E(\widehat{m}_t, \widehat{m}_{t'}), \\ \Gamma(\omega_8 \rightarrow b\bar{b}') &= \xi^2 |V_{tb}^* V_{t'b'}|^2 \frac{G_F^2 m_{\omega_8}^5}{192\pi} E(\widehat{m}_b, \widehat{m}_{b'}),\end{aligned}\quad (6)$$

where $n_f = 5$ is the number of light quark flavors that enhances this rate compared to $t\bar{t}$ and $\widehat{m} \equiv m/m_{\omega_8}$. The first two are just inverses of the production process as shown in Figure 7, which only depend on the decay constant parameter. On the other hand, the latter two proceed with an exchange of W boson as depicted in Figures 8(a) and 8(b), where $V_{tb'}$ necessarily enters. The final configuration for this case is always $bWbW$. While the as same as for $t\bar{t}$, it is nevertheless kinematically distinct. Analogous to $gg \rightarrow \omega_8 g$ production, the three-body $\omega_8 \rightarrow ggg$ decay rate, following the tree-level calculation of [35], is calculated and found to be negligible.

- (ii) Free-quark decay: $\omega_8 \rightarrow bWt', tWb'$;

$$\Gamma_{\text{free}} \simeq \frac{1}{2} [2\Gamma_{t'} + 2\Gamma_{b'}] = \Gamma_{t'} + \Gamma_{b'}, \quad (7)$$

where $\Gamma_{t'}, \Gamma_{b'}$ are given at Born level as

$$\Gamma_{t'}(m_{t'}) = |V_{tb}|^2 \frac{G_F^2 m_{t'}^3}{8\sqrt{2}\pi} F(\widehat{m}_W, \widehat{m}_b), \quad (8)$$

$$\Gamma_{b'}(m_{b'}) = |V_{tb}|^2 \frac{G_F^2 m_{b'}^3}{8\sqrt{2}\pi} F(\widehat{m}_W, \widehat{m}_t), \quad (9)$$

with $\widehat{m} = m/m_{t'}$ or $m/m_{b'}$. These occur when one of the heavy constituent quarks decays as a free quark, which requires the flavor change via $V_{tb'}$. There are two distinct final state configurations depending on whether it decays either from the $t'\bar{t}'$ or $b'\bar{b}'$ mode of the bound state. If it is via t' decay, the final decay product would be $bWbW$, and $bWWbWW$ if it is from b' decay. Since these signals are the same as the final state of the open production, the search can be pursued along the standard 4G quark search strategy [36]. Due to the binding energy, the reconstructed

heavy quark mass would be lower than the open production case [37]. For example, one of the bW pairs in the b' decay case would have a lower invariant mass roughly by the binding energy.

- (iii) Meson transition: $\omega_8 \rightarrow \omega_1 g; \pi_8 W; \pi_8 \gamma$,

$$\begin{aligned}\Gamma(\omega_8 \rightarrow \pi_8 W) &= \frac{G_F m_{\omega_8}^3}{32\sqrt{2}\pi} \frac{m_{\omega_8}}{m_{\pi_8}} W(\widehat{m}_{\pi_8}, \widehat{m}_W), \\ \Gamma(\omega_8 \rightarrow \omega_1 g) &= \frac{\alpha_s}{18} \frac{m_{\omega_8}^2}{m_{\omega_1}} G(\widehat{m}_{\omega_1}), \\ \Gamma(\omega_8 \rightarrow \pi_8 \gamma) &\simeq \frac{\alpha}{3} \frac{(\Delta m)^3}{m_Q^2}.\end{aligned}\quad (10)$$

The meson transition decays shown in Figure 9 give the clearest experimental signal for the heavy bound states. One more possibility $\omega_8 \rightarrow \pi_1 g$ is forbidden by the heavy isospin symmetry. The partial width for these decays depends on the mass difference as well as the transition amplitude of these resonances. For example, the meson transition to $\pi_8 W$ is allowed only if $m_{\omega_8} - m_{\pi_8} > m_W$. For $m_{\omega_8} - m_{\pi_8} < m_W$, ω_8 can decay into $\pi_8 \ell \nu_\ell$ or $\pi_8 q \bar{q}'$ through the off shell W boson. However, the partial width would be negligibly small. Previous numerical studies on Yukawa bound states suggest that the binding energy is indeed large if the m_Q is beyond the unitarity bound. For this reason, it is expected that the on-shell transition is to be the most likely and was pursued in the phenomenological study of [19].

The functions appearing in the rate formulae are given as follows:

$$\begin{aligned}E(x, x') &= \frac{\lambda(1, x^2, x'^2)}{2(1 - 2x^2 - 2x'^2)} \\ &\quad \times [2 - 9x^2 + 15x^4 - 8x^6 - 9x'^2 + 18x^2 x'^2 \\ &\quad - 8x^4 x'^2 - 16x^6 x'^2 + 15x'^4 - 8x^2 x'^4 \\ &\quad + 32x^4 x'^4 - 8x'^6 - 16x^2 x'^6], \\ F(x, y) &= (1 + x^2 - 2x^4 - 2y^2 + x^2 y^2 + y^4) \lambda(1, x^2, y^2), \\ W(x, y) &= (1 - 2x^2 + x^4 + 3y^2 + 2xy^2 + 3x^2 y^2 - 4y^4) \\ &\quad \times (1 + 2x + x^2 - y^2) \lambda(1, x^2, y^2), \\ G(x) &\simeq \frac{(1 - x^2)^3}{x}, \\ \lambda(a, b, c) &= \sqrt{a^2 + b^2 + c^2 - 2(ab + bc + ca)}.\end{aligned}\quad (11)$$

For more details of the derivation of the rates and approximations and simplifications taken we refer readers to [19] and references therein.

FIGURE 7: QCD-induced annihilation decay of ω_8 into light quark pair (dijets) or $t\bar{t}$.

3.2.1. Numerical Estimates. For the numerical study we need to choose following parameters: the 4th and lighter generation mixings from which we take 34 mixings only $V_{t'b}$ ($= -V_{tb'}$), the decay constant $\xi = f_{\omega_8}/m_{\omega_8}$ of ω_8 , the mass difference between mesons $\Delta m = m_{\omega_8} - m_{\pi_8}$ for $\omega_8 \rightarrow \pi_8 W$, $\Delta m = m_{\omega_8} - m_{\omega_1}$ for $\omega_8 \rightarrow \omega_1 g$ which is essentially the binding energy. Numerical studies based on Bethe-Salpeter equation and relativistic expansion of variational method suggest much higher values for the binding energy than the QCD case at few 100s of GeV. So we choose Δm to be 100 or 200 GeV. The following four cases were examined which represent distinct features in the large parameter space:

- (i) Case 1: $\xi = 0.1$, $\Delta m = 100$ GeV, $V_{t'b} = 0.1$;
- (ii) Case 2: $\xi = 0.03$, $\Delta m = 100$ GeV, $V_{t'b} = 0.1$;
- (iii) Case 3: $\xi = 0.1$, $\Delta m = 200$ GeV, $V_{t'b} = 0.1$;
- (iv) Case 4: $\xi = 0.1$, $\Delta m = 100$ GeV, $V_{t'b} = 0.01$.

Case 1 can be called as a “nominal” choice for the parameters. For the upper bound of the mixing the results of [6, 7] were taken. The choices for smaller decay constant, stronger binding, and small mixing correspond to Cases 2, 3, and 4, respectively. The resulting branching fractions for the above choice of parameters are plotted in Figure 10 using expressions in (6)–(10). In Case 1, we see that the dominant decay mode in the low mass range is $\pi_8 W$ and the constituent quark decay (free quark) in the higher values. In the latter case the 4G quark mass enhances the rate. The other decay modes are in the range of 1% ($\omega_1 g$) to 10% level ($q\bar{q}$). We see that the $q\bar{q}$ is 5 times larger than $t\bar{t}$ due to the number of light generations.

In Case 2, the dominant modes are the same as the above. The smaller decay constant makes all the annihilation modes suppressed.

In both Cases 3 and 4, the dominant modes are $\pi_8 W$, due to the large mass differences among the mesons and small mixing (suppressing the free and $t'\bar{t}(b'\bar{b})$), respectively.

In all cases, $m_{\omega_8} - m_{\pi_8}$ is taken to be larger than M_W . If the binding is not as strong as the Bethe-Salpeter approach suggests, the decay becomes off shell therefore subdominant. In this case the dijet channel will be the dominant. Also if the spectrum turns out to have $m_{\omega_8} \simeq m_{\pi_8} > m_{\omega_1}$ pattern upon a closer numerical examination $\omega_1 g$ could become important especially in the low mass range. In summary, the dominant channel is $\pi_8 W$ in all the choices considered, while free quark

decay can become competitive for large mass range if the mixing is close to its upper bound.

The estimated total decay widths for ω_8 and π_8 are shown in Figure 11. For ω_8 it is at most 20 GeV which, considering its mass being above TeV, makes it a narrow resonance.

3.2.2. The Decay of π_8 . From the decay rates in Figure 10, we see that $\pi_8 W$ is one of the leading channels if not the most leading one. Therefore, for complete phenomenological discussion, one has to address the decay of π_8 further. In the rest of the paper we assume $m_{\omega_8} - m_{\pi_8} > M_W$. Such a high mass splitting is suggested by Bethe-Salpeter approach [17]. Hence one should keep in mind that its validity has to be further checked by lattice calculation [39].

The first thing to note is that $\pi_8 \rightarrow \pi_1 g$ is forbidden by angular momentum sum rule. Further, isospin forbids the annihilation via W -exchange, while s -channel via g is forbidden by its color and isospin. This leaves only two possibilities: the free quark decay and annihilation to gW final state, where the W is transverse. The gW decay rate is given by

$$\Gamma(\pi_8 \rightarrow Wg) = 2\xi_\pi^2 \alpha_s \alpha_W m_{\pi_8} \frac{(1 - M_W^2/m_{\pi_8}^2)^3}{(1 + 4m_Q^2/m_{\pi_8}^2 - 2M_W^2/m_{\pi_8}^2)^2}, \quad (12)$$

where $\xi_\pi = f_{\pi_8}/m_{\pi_8}$ is the π_8 decay constant normalized by π_8 mass. The decay rate is plotted in Figure 12 for the four different cases we considered earlier. In Cases 1 to 3, where $V_{t'b}$ mixing is large, the free quark decay channel is the dominant. For Case 2, $\pi_8 W$ is suppressed even more due to small decay constant. On the other hand, Case 4 gives $\pi_8 W$ as the dominant decay, since the free quark decay is suppressed by the choice of small mixing.

4. Phenomenology Discussion

The searches for a dijet resonance at the Tevatron and LHC put constraint on theories with TeV range particles. The ω_8 production cross-section times dijet branching fraction at 8 TeV is plotted in Figure 13 for the four cases along with the dijet search at CMS [38] with $\mathcal{L} = 4 \text{ fb}^{-1}$ data at $\sqrt{s} = 8$ TeV. This is a conservative estimate since, unlike the data, no cut was used in the dijet rate from ω_8 . The open $Q\bar{Q}$ pair productions via gg -fusion are also calculated at LO and NLO and are shown in orange and red bands, respectively. Although ω_8 is through $q\bar{q}$, it is of the same order or even larger than the open production. We see that all the rates are at least an order of magnitude lower than the CMS constraint.

Since the dominant decay is $\omega_8 \rightarrow \pi_8 W$ it is necessary to examine π_8 decay modes to get eventual signals at the LHC. Observe that $\pi_8 \rightarrow \pi_1 g$ is forbidden by angular momentum sum rule. The estimate for $\omega_8 \rightarrow \pi_8 W$ rate in Figure shows that in Cases 1 to 3 it dominantly decays through free quark channel. From the isospin one expects

FIGURE 8: Weak decay modes: (a) exchange diagram with off-diagonal quark mixing element; (b) free quark decay.

FIGURE 9: Meson transition currents for $\omega_8 \rightarrow \pi_8$ and $\omega_8 \rightarrow \omega_1$.

$\omega_8^0 \rightarrow \pi_8^\pm W_s^\mp, \pi_8^0 Z_s$ to have a 2 : 1 ratio, where the subscript s signifies that W is soft. The final state configuration is then $bW^+ \bar{b}' W_s^-$, $tW^- \bar{t}' W_s^+$, and $\{bW^+ \bar{t}', tW^- \bar{b}'\} Z_s$ with 7 and 6/8-jet systems, respectively. The associated soft W (Z) can be used as a tag. With additional soft W , the signal appears even more complicated than the open production. While they appear difficult, the boosted products can be exploited to analyze the system using their jet substructures. If the total mass can be constructed with good enough resolution the system adds up to a narrow resonance potentially giving signals for both ω_8 and π_8 .

This brings us to the discussion of Case 4, for which one has $\pi_8 W$ as dominant mode with final signal as $W_s W g$. ($W_s^\pm W^\mp g$ and $Z_s Z g$). This signal is very unique and highly detectable at the LHC with $p_T > 500$ GeV gluon and boosted W with soft W . The whole system also gives a narrow resonance.

The $\pi_8 \gamma$ channel is at most a percent level for all the Cases. Although the rate is low, due to high photon detection efficiency it can be searched and further studied in detail. Finally, the W -exchange channel to $t\bar{t}'$ and $b\bar{b}'$, while appears similar to the single $t'(b')$ production, and the associated boosted t (high p_T b) can give distinct resonant signal.

We conclude the phenomenological discussion emphasizing the difference compared to technicolor (TC) models. The first distinction is the presence of colored objects. For generic “walking” TC models [40], the technipion π_T tends to be closer to the technirho ρ_T in mass such that $\rho_T \rightarrow \pi_T \pi_T$ is absent, while (near) degeneracy of ω_T and a_T with ρ_T is also often invoked. The signature for these technimesons are typically WZ , $W\gamma$, and $Z\gamma$. Thus, not only the spectrum is rather different—absence of ρ and a mesons—the decay

signature is also in strong contrast. The bound states due to strong Yukawa coupling, which follow simply from the existence of new heavy chiral quarks without assuming new dynamics, should be easily distinguished from technicolor. The Yukawa-bound ultraheavy mesons are also quite distinct from QCD bound states. Not only there is the absence of η (where gg fusion would be possible) and ρ type mesons, but also they have a much larger binding energy and are much smaller in size. This is brought about not only by a strong coupling constant, but also by isospin-dependent Goldstone coupling.

5. Extension to Two-Higgs Doublet

With the discovery of ~ 125 GeV boson, a simple 4G extension of the SM appears unattainable if the signal is indeed confirmed as the SM Higgs. If an extra generation does exist at least the SM Higgs sector has to be extended further. Indeed a recent analysis has shown that the precision test for 4G with two Higgses is allowed by the current constraint if the lighter Higgs is indeed 125 GeV [22]. In the following, we discuss how an extra new doublet Higgs affects the phenomenology we considered so far in this paper.

We take, for simplicity, the case of Type-II of two-Higgs doublet (2HD) model, where up- (down-) type fermions couple only to H_1 (H_2). We also take the simplified limit, for illustration, wherein $\tan \beta = v_1/v_2 \simeq 1$ and $\sin(\beta - \alpha) \sim 1$. Here α is $h - H$ mixing angle. In the presence of a light Higgs, new ω_8 decay channels open up (i) from free quark decay with Higgs in the final state, (ii) to gh/H , and (iii) from the annihilation via Higgs exchange analogous to the W -exchange discussed earlier. These are as follows.

FIGURE 10: Branching ratio of ω_8 as a function of m_{ω_8} for given parameter sets in Cases 1 to 4.

(i) Free quark decay: $\omega_8 \rightarrow t'\bar{t}'(b'\bar{t}'H^+)$:

$$\begin{aligned} \Gamma_{\text{free}}^H &\simeq \Gamma(t' \rightarrow H^+ b) + \Gamma(b' \rightarrow H^- t), \\ \Gamma(t' \rightarrow H^+ b) &= \frac{G_F |V_{t'b}|^2 m_{t'}^3}{8\sqrt{2}\pi} \lambda(1, \hat{m}_b^2, \hat{m}_{H^\pm}^2) \\ &\quad \times \left\{ \kappa_A^{u^2} (1 + \hat{m}_b^2 - \hat{m}_{H^\pm}^2) + \hat{m}_b^2 \kappa_A^{d^2} + 4\hat{m}_b^2 \kappa_A^u \kappa_A^d \right\}, \\ \Gamma(b' \rightarrow H^- t) &= \Gamma(t' \rightarrow H^+ b) \Big|_{(t', b, d) \leftrightarrow (b', t, u)}, \end{aligned} \quad (13)$$

where $\kappa_A^{u(d)}$ is $\tan \beta = v_1/v_2(1/\tan \beta)$ and $\hat{m} = m/m_{t'(b')}$ for $t'(b') \rightarrow H^\pm b(t)$. For numerics we take $m_{H^\pm} = 200$ GeV as suggested by [22]. Note that $t'\bar{t}'A(b'\bar{b}A)$ is absent since the CP-odd A does not violate flavor at tree level in Type-II 2HD model. The loop induced amplitude are small at low $\tan \beta$.

(ii) Annihilation to $t'\bar{t}'(b'\bar{b})$ through Higgs exchange:

$$\begin{aligned} \Gamma^{H^\pm}(\omega_8 \rightarrow t'\bar{t}') &= \xi^2 \frac{G_F^2 |V_{t'b}|^2 m_Q^4 m_{\omega_8} \lambda(1, \hat{m}_t^2, \hat{m}_Q^2)}{24\pi} \\ &\quad \times \frac{(1 - \hat{m}_t^2 - \hat{m}_Q^2)}{(1 + 4\hat{m}_{H^\pm}^2 - 2\hat{m}_t^2 - 2\hat{m}_Q^2)^2}, \quad (14) \\ \Gamma^{H^\pm}(\omega_8 \rightarrow b'\bar{b}') &= \Gamma^{H^\pm}(\omega_8 \rightarrow t'\bar{t}') \Big|_{(t't) \rightarrow (b'b')}. \end{aligned}$$

As in the previous case, there is no A -exchange contribution since A does not break flavor in Type-II case at tree level and for low $\tan \beta$ the quantum effects are ignorable.

(iii) Annihilation with Higgs in the final state $\omega_8 \rightarrow gh$ and gH . In Type-II 2HD case with $\tan \beta \simeq 1$ and

FIGURE 11: Total decay width of ω_8 and π_8 for the four parameter sets. The upper curves are for ω_8 , where Case 3 is enhanced by a larger $m_{\omega_8} - m_{\pi_8}$ mass difference. The lower curves, plotted 100 GeV less in mass, are for π_8 , where Cases 1 and 2 are barely distinguishable.

FIGURE 12: Branching ratio $\mathcal{B}(\pi_8 \rightarrow Wg)$ for the four parameter sets, ranging from predominance (Case 4), to below 10^{-2} (Case 2).

$\sin(\beta - \alpha) \sim 1$, the rates for the light Higgs channel are given as

$$\Gamma(\omega_8 \rightarrow gh) \simeq \xi_h^2 \frac{2\sqrt{2}\alpha_s}{3} \frac{G_F m_Q^2 m_{\omega_8} (1 - \hat{m}_h^2)^3}{(1 + 4\hat{m}_Q^2 - 2\hat{m}_h^2)^2}, \quad (15)$$

where $\hat{m} = m/m_{\omega_8}$ and ξ_h is the decay constant for this channel taken to be equal to ξ . We take $m_h = 125$ GeV. In the simplified limit we are considering that $\omega_8 \rightarrow gH$ is suppressed. ω_8 and g have the same quantum numbers except mass which makes the $\omega_8 g$ system CP even. Since A does not mix with h/H in Type-II case, it has definite odd CP-nature.

FIGURE 13: Production cross-section times branching ratio for ω_8 into dijets, for the four parameter sets at the LHC 8 TeV. CMS upper limit [38] for the cross-section of dijet resonance production is also plotted.

FIGURE 14: Branching ratios for 2HD model for the parameter choices of Case 4. Here we take $m_h = 125$ GeV and $m_{H^\pm} = 200$ GeV.

Therefore $\omega_8 \rightarrow gA$ is also absent. One might think that any of $\omega_8 \rightarrow \pi_8 h/H/A$ could become important if kinematics allows. These are absent for the same reason that W from $\omega_8 \rightarrow \pi_8 W$ is purely transverse.

The rates for channels $t'\bar{b}H^- (b'\bar{t}H^+)$ and gh are plotted in Figure 14. In Type-II case, there is no free quark to a final state with neutral Higgs. Since we find $t'\bar{t}$ and $b'\bar{b}$ via charged Higgs exchange to be very small we do not plot them here. The free quark decays with charged Higgs are present if $m_{H^\pm} < m_{\omega_8} - m_Q$. The rates are at the same order as the free quark decay given in (9), which is expected since the heavy quark decay rates to tW are dominated by the Goldstone components.

For Type-II 2HD extension, the annihilation through charged Higgs exchange is also at the same order as the W -exchange counterpart. Therefore they have no qualitative and a very minimal quantitative effect on the discussion of the leading $\pi_8 W$ signal.

The decay to gh and gH could have large branching ratios at high ω_8 mass range. For our simplified discussion, gh reaches at most $\sim 20\%$. Therefore our $\pi_8 W$ discussions remain robust. The search for this channel appears challenging. If enough ω_8 is produced and its mass is high enough; boosted Higgs analysis [41] may be used.

In this section we have taken Type II 2HD model as a simple illustration and have shown that the main phenomenology expected at the LHC from the bound states may not have to be drastically different. A more thorough numerical study is needed for clarifying whether this or any other variants of 2HD model are phenomenologically viable.

6. Conclusion and Outlook

In this short paper we revisited current phenomenological study of possible bound states of heavy 4G quarks. The direct search limits now have reached the perturbative unitarity bound making their Yukawa couplings well in the strong coupling regime, where one expects a meson-like $Q\bar{Q}$ bound states in $\sim \text{TeV}$ range. Existing studies in Bethe-Salpeter and relativistic expansion approaches enable one to study their phenomenology at the LHC.

Electroweak Precision tests constrain the top and bottom-like 4G quark masses to be nearly degenerate leading to a “heavy isospin” symmetry, which in turn can be used to classify spectrum. Among these, one finds that color-octet isosinglet vector resonance ω_8 is to be the easiest produced one at early runs of the LHC via $q\bar{q}$ -annihilation. The parameters needed for such study are the meson mass splittings, 4G quark mixing to lighter generations, the decay constant, and 4G quark mass.

Possible final decay products have rich phenomenology with challenging signals. The leading decay channel of ω_8 in the cases considered is $\pi_8 W$ followed by free quark decay inside the bound state. With four different representative choices for the parameters their rates are studied. Among these, Case 4 with small mixing provides a clearest signal of possible bound state at the LHC, wherein the whole process chain is $\omega_8 \rightarrow \pi_8 W_s \rightarrow gW_s W$ with highly boosted W and g with one soft W_s as the final state. Other cases, with enough total energy resolution, also may provide a resonant signal.

With $\sim 125 \text{ GeV}$ object found at the LHC being most likely a Higgs, if proven, simple 4G extension is unlikely to be viable. In light of this, we also briefly explore how this phenomenology is affected in presence of additional Higgs. The most significant change comes in the form of opening up new large annihilation decay channels gh and gH . Numerical estimates of these rates indicate that in Case 4 $\pi_8 W \rightarrow gWW$ mode remains the leading channel, leaving the conclusions with single Higgs robust. Another easily detectable signature is $q\bar{q}$ dijet which becomes dominant if $\Delta m = m_{\omega_8} - m_{\pi_8} < M_W$. While this is unlikely considering

what Bethe-Salpeter study suggests, this can be eventually resolved by Lattice effort.

We have reviewed signatures for Yukawa-bound heavy $Q\bar{Q}$ mesons in the 1 to 2 TeV range. We also briefly touched upon two Higgs doublet case. What is covered in this review is qualitative in nature. As the LHC energy increases, and with higher luminosity, it could uncover new chiral quarks above the unitarity bound, with new unusual bound states. One could probe into the truly nonperturbative regime, in which these results only offer a glimpse of what might happen.

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