Phenomenology of iQuarkonium

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Abstract

Phenomenology of uncolored iquarks – hypothetical fermions charged under a new confining unbroken non-abelian gauge group as well as the standard electroweak gauge group – is investigated for the iquark mass in the range near and above 100 GeV. If the new confining scale turns out to be higher than MeV but much less than the iquark mass, the iquark-antiquark pair produced in the collider will promptly relaxed into the ground state of the iquarkonium. Subsequent pair annihilation into standard model particles gives useful information of the iquark dynamics. We formulate in details production and decays of the iquark-antiquark bound states. Decay patterns of the iquarkonium can be distinguished from the superheavy quarkonium of a sequential fourth generation of quarks with degenerate mass.
A new confining unbroken non-abelian gauge interaction, in mimic of quantum chromodynamics (QCD) of strong interaction, may exist in models beyond the standard model (SM). Okun \cite{1} pioneered in this theoretical curiosity long time ago and named these new particles “thetons”, including $\theta$-strings, $\theta$-leptons, $\theta$-quarks, $\theta$-hadrons, etc associated with the gauge group $SU_{\theta}(N)$. This theory can also be regarded as a certain limit of QCD with light quarks removed such that the scale $\Lambda$ where QCD gets strong is much smaller than the heavy quark masses $\Lambda_{QCD} \ll M_Q$. Phenomenology of this hypothetical QCD can be drastically different from the real world where light quarks do exist and one must have to worry about the spontaneous chiral symmetry breaking. For instance, heavy quarks are connected by unbreakable long and stable string flux tube since there is no light quark pairs popped up from the vacuum as the heavy quarks are being pulled apart. However, the $\theta$-fermions $Q$ and $\bar{Q}$ of the new gauge interaction, if also carry standard model quantum numbers, have to be massive ($M_Q \gtrsim 100$ GeV). Otherwise, they would have been observed in current collider experiments.

These heavy $\theta$-fermions were recently revived and renamed as “quirks” by Kang and Luty \cite{4}, who emphasized their fantastic event structures due to the long and stable string flux tube connecting them. We will briefly mention some of their observations here. For definiteness, let us assume that this new gauge group is $SU_{\theta'}(N_{IC})$ with its characteristic scale $\Lambda' \ll M_Q$. The size of the string flux tube is of order $\Lambda'^{-1}$ and it can be macroscopic for $100 \text{ eV} \ll \Lambda' \ll 10 \text{ keV}$, or mesoscopic for $10 \text{ keV} \ll \Lambda' \ll \text{MeV}$, or microscopic for $\text{MeV} \ll \Lambda'$. In general, one also assume $\Lambda'$ is smaller than the $\Lambda_{QCD}$ such that the new color degree of freedom bears the name infracolor (IC). We call this infracolor (a.k.a. iQCD) gluon fields igluons, and the fermions iquarks. \footnote{Kang and Luty \cite{4} called this quirk. To some extent it is easy to mix up “quark” and “quirk” when listening to talks of quirks. So we introduce “iquark” (pronounced as i-quark), and similarly “igluon” (pronounced as i-gluon), etc.} It has been noticed that the infracolor is an example of the “hidden valley” model in Ref.\cite{5}. The igluons do not carry SM quantum numbers. Thus, iglueballs can only couple to SM gauge bosons through a heavy iquark loop assuming the heavy iquark is not a SM singlet. At two-loop level, the iglueballs can also couple to the SM fermions. Thus, these iglueballs may have escaped detection. In the macroscopic
The infracolor dynamics can allow the $Q\overline{Q}$ bound state to survive for distances of order centimeter and preventing them from annihilation. In the mesoscopic $Λ'$ scenario, the $Q\overline{Q}$ pair will appear as a single entity inside the detector. On the other hand, if $Λ'$ is microscopic, the $Q\overline{Q}$ pair will lost most of their kinetic energy and angular momenta by emitting iglueballs and/or light QCD hadrons like pions before annihilation. In the latter case in which iquarks are QCD-colored, it leads to a hadronic fireball along with the other SM decay products of the iquarkonium. In particular, in the context of folded supersymmetry [6] it was pointed out in Ref.[7] that production of the “squirk-antisquirk” pair $\tilde{Q}\tilde{Q}^*$ at the large hadron collider (LHC) would quickly lose their excitation energy by bremsstrahlung and relax to the ground state of the scalar iquarkonium. However, the energy loss due to iglueball emissions is harder to estimate.

In this work, we consider vector-like iquarks with respect to the electroweak gauge group but without carrying the QCD color. However, iquark carries a new color degree of freedom of $SU_{C'}(N_{IC})$. Thus, iquarks do not mix with SM quarks or leptons since the latter do not carry the new color degrees of freedom. We also assume $\text{MeV} \leq Λ' \ll M_{Q}$ so that the strings are microscopic but yet unbreakable. In analogous to the case of folded supersymmetry [7], the bound states formed by the iquark-antiquark pairs will annihilate promptly into SM particles. Let $Q$ denotes a heavy iquark doublet. The quantum number assignment for the iquark doublet $Q$ under $SU_{C'}(N_{IC}) \times SU_C(3) \times SU_L(2) \times U_Y(1)$ is given by

$$Q_{L,R} = \left( \begin{array}{c} U \\ D \end{array} \right)_{L,R} \sim (N_{IC}, 1, 2, \frac{1}{3}) .$$

These iquarks are the fractionally charged $θ$-leptons in Okun’s terminology [1]. The gauge interactions are given by

$$\mathcal{L}_{\text{gauge}} = - g_s' G^\mu_{\mu} \overline{Q} \gamma^\mu T^a_{\mu} Q - e A^\mu \left( e_U \overline{U} \gamma^\mu U + e_D \overline{D} \gamma^\mu D \right) - \frac{g}{\cos \theta_W} Z_{\mu} \left( v_U \overline{U} \gamma^\mu U + v_D \overline{D} \gamma^\mu D \right) - \frac{g}{\sqrt{2}} \left( W^+_{\mu} \overline{U} \gamma^\mu U + W^-_{\mu} \overline{D} \gamma^\mu D \right)$$

where we have suppressed generation indices and ignored possible mixings among iquarks. $T^a (a = 1, \cdots, N_{IC}^2 - 1)$ are the generators of the $SU_{C'}(N_{IC})$ in the defining representation where each iquark lives and $g_s'$ is its coupling. For vector iquark $Q = U$ or $D$, we have

$$v_Q = \frac{1}{2} (T_3(Q_L) + T_3(Q_R)) - e_Q \sin^2 \theta_W .$$
Here $T_3(Q_{L,R})$ is the third-component of the weak isospin for the left- (right-) handed iquark $Q$. Since we assume vector iquarks, $T_3(Q_L) = T_3(Q_R)$, they have the same value for each component of $Q$. For each vectorial iquark doublet, we can also have a Dirac bare mass term. Due to the assumption of vectorial iquarks, precision electroweak data from LEP do not provide any constraints even for a TeV iquark doublet. There is no Yukawa coupling between the SM Higgs doublet and the vector iquarks.

Iquarks in the above model can be copiously produced at the LHC, not via normal QCD interactions, but via electroweak interactions. The iquarks are stable with respect to the collider time scale, and linked by a string. Depending on the value of $\Lambda$ the iquarks will lose energy by emitting iglueballs and bremsstrahlung until the iquark pair linked by the string come together and form a iquarkonium \[2\]. The life time of the iquarkonium then depends on the annihilation rates into SM particles. In accord with the normal quarkonium, the life time of iquarkonium is inversely proportional to the square of the wave-function at the origin. The life time and decay patterns of the iquarkonium can therefore provide useful information about the quantum numbers of iquarks and the dynamics of the new strong interacting gauge group.

In this article, we set up detailed formulas for production and decay properties of the neutral $Q\overline{Q}$ and charged $Q\overline{Q}'$ iquarkonia. Our result extends a previous study on superheavy-quarkonia to our current interest in iquarks and iQCD. We will consider only one generation of iquark doublet. Extension to multiple generations of iquark doublets is straightforward.

The organization of the paper is as follows. In the next section, we present the cross sections for open production of iquarks. In Sec. III, we describe briefly the motion of the iquark once produced and the radiation. In Sec. IV, we present the results on the annihilation of the $S$-wave iquarkonium and discuss the decay patterns. Conclusions are made in Sec. V.

II. OPEN UNCOLORED IQUARK PRODUCTION

Even though we assume iquarks do not carry the usual color of the strong QCD interaction, they can still be pair produced at CERN LHC through electroweak interactions. In the following, we list the formulas for the production amplitude squared of these hard partonic
A. Open production of $\U\U$ and $\D\D$

1. $q(p_1)\pi(p_2) \to \U(k_1)\U(k_2)$ and $\D(k_1)\D(k_2)$

There are two s-channel diagrams from the $\gamma$ and $Z$ exchanges with the following result

$$\sum |M|^2 = 2 \epsilon^4 \frac{N_{IC}}{3} \left[ \left( \hat{t} - M_Q^2 \right)^2 + \left( \hat{u} - M_Q^2 \right)^2 + 2\hat{s}M_Q^2 \right]$$

$$\times \left[ \epsilon_q \epsilon_Q \frac{1}{\hat{s}^2} + \frac{v_Q^2 \left( g_V^2 + g_A^2 \right)}{\sin^2 \theta_W \cos^2 \theta_W} \frac{1}{\left( \hat{s} - M_Z^2 \right)^2 + \Gamma_Z^2 M_Z^2} \right.$$  
$$\left. + \frac{2\epsilon_q \epsilon_Q g_V^2 v_Q}{\sin^2 \theta_W \cos^2 \theta_W} \frac{1}{\hat{s} \left( \hat{s} - M_Z^2 \right)^2 + \Gamma_Z^2 M_Z^2} \right]$$

(4)

where $\hat{s} = (p_1 + p_2)^2 = (k_1 + k_2)^2$, $\hat{t} = (p_1 - k_1)^2 = (p_2 - k_2)^2$ and $\hat{u} = (p_1 - k_2)^2 = (p_2 - k_1)^2$; $g_V^2 = \frac{1}{2} \left( T_3^Q \right)_L - \epsilon_q \sin^2 \theta_W$, $g_A^2 = \frac{1}{2} \left( T_3^A \right)_L$ and $v_Q$ is given by Eq. (3) with $Q = \U$ or $\D$. All the initial-state quark masses are set to be zero.

B. Open production of $\U\D$ and $\D\U$

It turns out that the $\U\D$ and $\D\U$ pairs can be frequently produced via the virtual $W$ boson in the annihilation of $u\bar{d}$ and $d\bar{u}$, respectively.

1. $u_i(p_1)\overline{d}_j(p_2) \to \U(k_1)\D(k_2)$ and $d_j\pi_i \to \D\U$

There is only one s-channel diagram from the $W$ exchange.

$$\sum |M|^2 = \frac{1}{4} \epsilon^4 \frac{N_{IC}}{3} |V_{CKM}|^2 \frac{1}{\left( \hat{s} - M_W^2 \right)^2 + \Gamma_W^2 M_W^2} \times \left\{ \left( \hat{t} - M_D^2 \right) \left( \hat{t} - M_D^2 \right) + \left( \hat{u} - M_U^2 \right) \left( \hat{u} - M_U^2 \right) + 2\hat{s}M_D M_D \right\}$$

(5)

where $\hat{s} = (p_1 + p_2)^2 = (k_1 + k_2)^2$, $\hat{t} = (p_1 - k_1)^2 = (p_2 - k_2)^2$ and $\hat{u} = (p_1 - k_2)^2 = (p_2 - k_1)^2$. All the initial-state quark masses are set to be zero.
C. Cross sections at the LHC

With the above formulas Eqs. (4)-(5) we present in Fig. 1 the production cross sections of iquarks at the LHC. The production rates are significant even though the iquarks are produced via electroweak interactions rather than QCD. For $M_{U,D}$ around $100 - 200$ GeV the cross sections are of order $O(1) - O(10)$ pb and are about an order of magnitude larger than the scalar iquarkonium case in folded supersymmetry [6]. In contrast with QCD, due to the unbroken string flux tube, the two iquarks do not hadronize individually to form isolated jets. Since we assume the iquark doublet is vector-like, the $\beta$-decay between doublet members is suppressed by small mass splittings due to radiative corrections. Instead these initially flying-apart iquark-antiquark pair will come back close to each other to form iquarkonium after losing their kinetic energies by radiating off iglueballs and photons [4]. Essentially, in the case of microscopic $\Lambda'$ all open iquark pairs will, at the end, come together to form a iquarkonium. This is in sharp contrast to the normal quarkonium, in which we have to force them to go together in the same direction (e.g. by radiating one or more gluons) and in roughly the same velocity in order to form a quarkonium. Therefore, the iquarkonium production rates are not inferior to the quarkonium, although iquarks are only produced via electroweak interactions. We will discuss more about these interesting phenomena in the next section.

III. BEHAVIOR OF THE OPEN UNCOLORED IQUARKS

A. Macroscopic $\Lambda'$

The iquarks are produced via electroweak interactions. Once they are pair produced, they are still connected by a string or an iQCD flux tube. Recall in usual QCD the flux tube gets broken by creation of light $q\bar{q}$ pairs. However, in iQCD the energy density stored in the string cannot exceed $\Lambda'^2$, which is way smaller than the mass of iquarks. Therefore, the iquark-antiquark pair production is exponentially suppressed. The iglueball production from the vacuum is also expected to be suppressed due to its finite mass gap. Thus, the string is relatively stable.

Nevertheless, the iquarks carry electric and weak charges so that it will undergo bremsstrahlung and ionization with the detector materials. In Ref. [7], bremsstrahlung
is treated semi-classically as two massive charged particles connected a string of tension $\Lambda'^2$. The rate energy loss is proportional to $\alpha \Lambda'^4 / M^3$. On the other hand, the ionization energy loss when transversing the detector is given by the Bethe-Blöch equation \[9\]. Essentially, the penetrating particles lose energy by exciting the electrons of the material. Ionization energy loss $dE/dx$ is a function of $\beta \gamma \equiv p/M$ and the charge of the penetrating particle \[9\]. The signature of iquarks due to the above two energy loss mechanisms is very spectacular. The ionization energy loss in the detector could give rise to observable tracks in the tracking sector of the detector. In addition, the bremsstrahlung photons may be, though soft, identified along with the iquarks. Since the iquarks are roaming around in the detector connected by a string, there may be random tracks plus bremsstrahlung photons going along with the tracks. At the end of kinetic-energy loss, the iquarks and antiquarks could be still far apart so that they cannot annihilate. They would then stop inside the detector.
B. Microscopic $\Lambda'$

In this case, the size of the string is so small that when the iquarks lose their kinetic energy, they will come together to form a iquarkonium. The iquarkonium will then annihilate into SM particles. We will consider the annihilation channels in the next section. Whether the iquarkonium can survive a period of time depends on (i) the annihilation rate of the iquarkonium, and (ii) the rate of energy loss of the fast moving iquarks. It is easy to show that the annihilation rate into SM particles is proportional to \( \alpha_W^2 |R_S(0)|^2 / M^2 \) where \( \alpha_W = \alpha_{em} / \sin^2 \theta_W \), \( R_S(0) \) is the radial wave-function at the origin, and \( M \) is the mass of the iquarkonium. For heavy iquarkonium \( |R_S(0)|^2 \) scales as \( \alpha_s^2 M^3 \), so that the overall annihilation rate scales as \( \alpha_W^2 \alpha_s^2 M \). On the other hand, the rate of energy loss of the fast moving iquarks depends on the bremsstrahlung rate and the rate of radiating iglueballs. While the rate of radiating the iglueballs is largely unknown, the bremsstrahlung rate is proportional to \( \alpha_{em} \Lambda^4 / M^3 \). For \( M \sim 100 \text{ GeV} \) the bremsstrahlung rate is faster than the annihilation rate. The iquark pair loses the kinetic energy to almost stationary and form the iquarkonium, which then survives a short period of time before annihilates. In this case, \( P \)-wave annihilation is negligible. For the case of very heavy iquarks \( M \sim 1 \text{ TeV} \), the bremsstrahlung rate is comparable to the annihilation rate so that the iquarkonium annihilates right away while the iquarks are losing the kinetic energy. In this case, we expect the \( P \)-wave annihilation is also relevant. We will delay the \( P \)-wave annihilation to a future publication while we focus on the \( S \)-wave case in this paper. If the rate of radiating iglueball dominates over the bremsstrahlung, the discussion here is very straight-forward. The iquark and antiiquark quickly loses their kinetic energy, and the size of the string is so small that the iquark and antiiquark are readily coming together to form a iquarkonium, which then annihilates promptly. In this case, the annihilation rate is dominated by the \( S \)-wave states.

IV. IQUARKONIUM DECAY

In this section, we present the formulas for \( S \)-wave iquarkonium decays. We have checked that some of the results are consistent with a previous similar calculations for a superheavy quarkonium decays by Barger et al \cite{Barger90} with appropriate modifications. However, parts of
the results are genuinely new.

A. Neutral Iquarkonium

1. $ff$

$$\Gamma(\eta_{Q\bar{Q}} \to ff) = 0$$  \hspace{1cm} (6)

$$\Gamma(\psi_{Q\bar{Q}} \to ff) = \frac{4 N_{IC} N_{Cf} \alpha_{em}^2 \beta_f}{3} \left( 1 + 2 R_f \right) \left( e_{Qf} + \frac{v_Q g_V}{x_W (1 - x_W) (1 - R_Z)} \right)^2$$

$$+ \beta_f^2 \frac{v_Q^2 g_A^2}{x_W (1 - x_W)^2 (1 - R_Z)^2} \left| R_S(0) \right|^2 \frac{1}{M^2}$$  \hspace{1cm} (7)

Here, $N_{Cf}$ is the color factor for the fermion $f$ (1 for leptons and 3 for quarks), $R_i = M_i^2/M^2 \ (i = f, Z)$ with $M$ being the mass of the iquarkonium, $x_W = \sin^2 \theta_W$ and $\beta_f = (1 - 4 R_f)^{1/2}$.

2. $\gamma\gamma$

$$\Gamma(\eta_{Q\bar{Q}} \to \gamma\gamma) = 4 N_{IC} \alpha_{em}^2 e_Q^4 \left| R_S(0) \right|^2 \frac{1}{M^2}$$  \hspace{1cm} (8)

$$\Gamma(\psi_{Q\bar{Q}} \to \gamma\gamma) = 0$$

3. $Z\gamma$

$$\Gamma(\eta_{Q\bar{Q}} \to Z\gamma) = 8 N_{IC} \alpha_{em} \alpha_Z e_Q^2 v_Q^2 (1 - R_Z) \left| R_S(0) \right|^2 \frac{1}{M^2}$$  \hspace{1cm} (9)

$$\Gamma(\psi_{Q\bar{Q}} \to Z\gamma) = 0$$

Here, we have defined $\alpha_Z = \alpha_{em}/(\sin^2 \theta_W \cos^2 \theta_W)$ and $R_Z = M_Z^2/M^2$.
where $\beta_Z = (1 - 4R_Z)^{1/2}$. In addition to setting the axial vector couplings of the iquarks to be zero, we also exclude the Higgs exchange contribution in the formulas of Barger et al. because a vectorial iquark does not couple to the standard model Higgs boson.

5. $W^+W^-$

Using the Feynman rules defined by Eq. (2), one obtains

$$\Gamma(\eta_{Q\bar{Q}} \to W^+W^-) = \frac{N_{IC} \alpha_W^2 \beta_W^3}{2} \frac{1}{(1 - 2R_Z)^2} \frac{|R_S(0)|^2}{M^2}.$$  \hspace{1cm} (11)

Here $\alpha_W = \alpha_{em}/\sin^2 \theta_W$, $\beta_W = (1 - 4R_W)^{1/2}$ and only $t$-channel exchange of the $Q'$ iquark contributes for the $^1S_0$ state. However, for the $^3S_1$ state, both the $t$- and $s$-channels appear with the following amplitudes

$$M_t = \pm \frac{R_S(0) \sqrt{N_{IC}}}{M \sqrt{4\pi M}} \frac{2g_Q^2}{1 - R_{Q'W}} \left[(1 + r_{Q'}) (k_1^\mu g^{\mu\alpha} - k_2^\nu g^{\nu\alpha}) - g^{\mu\nu} k_1^\alpha \right] \epsilon_\alpha(P) \epsilon_\nu(k_1) \epsilon_\mu(k_2) \hspace{1cm} (12)$$

where the $+ \text{ or } -$ sign is for the case of $UU$ or $DD$ respectively and

$$M_s = -\frac{R_S(0) \sqrt{N_{IC}}}{M \sqrt{4\pi M}} \frac{2g_Q^2 g_Q'}{2(1 - R_Z)} \left[(k_2 - k_1)^\alpha g^{\mu\nu} - 2k_2^\mu g^{\nu\alpha} + 2k_1^\mu g^{\nu\alpha} \right] \epsilon_\alpha(P) \epsilon_\nu(k_1) \epsilon_\mu(k_2). \hspace{1cm} (13)$$

Here $g_Q = (T_3(Q_L) + T_3(Q_R)) - 2 e_Q \sin^2 \theta_W R_Z$, $r_{Q'} = M_{Q'}/M$ and $R_{Q'W} = (1 - 4R_{Q'} + 4R_W)/2$ with $(Q, Q') = (U, D)$ or $(D, U)$. Note that these two amplitudes partially cancel each other if $r_{Q'} = \frac{1}{2}$ when the iquarks $U$ and $D$ are degenerate in mass for the vectorial iquarks. The cancellation is expected to reduce the large decay rate into longitudinal polarizations of $W^\pm$ bosons. However, we allow the iquark masses to be different in our formulas. Our result of the partial width can be expressed compactly as

$$\Gamma(\psi_{Q\bar{Q}} \to W^+W^-) = \frac{N_{IC} \alpha_W^2 \beta_W^3 |R_S(0)|^2}{48R_W^2} \left[H^2 + 4 \left(H^2 + 3HG + G^2\right) R_W + 12G^2 R_W^2\right] \hspace{1cm} (14)$$

where

$$G = \frac{|g_Q|}{1 - R_Z} - \frac{1}{1 - R_{Q'W}}, \quad H = \frac{|g_Q|}{1 - R_Z} - \frac{2r_{Q'}}{1 - R_{Q'W}}. \hspace{1cm} (15)$$

\[10\]
6. $g'g', g'g'g', Zg'g', \gamma g'g'$

The iquarks couple to the igluon field of the $SU_C(N_{IC})$ with a coupling strength $g_s' = \sqrt{4\pi\alpha_s'}$. These igluon fields will escape the detection, giving rise to missing energies in the final state. We compute the leading order of these processes.

For the $^1S_0$ neutral quarkonium the leading decay mode involving $g'$ is $g'g'$:

$$\Gamma(n_{Q\bar{Q}} \to g'g') = \frac{N_{IC}^2 - 1}{N_{IC}^2} |R_S(0)|^2 \frac{\alpha_s'^2}{M^2},$$

where $\alpha_s'(Q) = 12\pi/[(11N_{IC} - 2n_Q)\ln(Q^2/\Lambda^2)]$. In our numerical works presented later, we will choose the number of infracolor $N_{IC} = 3$ and the number of iquark generation $n_Q = 1$ at the running scale $Q = M$.

For the $^3S_1$ neutral quarkonium the leading decay modes involving $g'$ are $g'g'g'$, $\gamma g'g'$, and $Zg'g'$. The formulas for $g'g'g'$ and $\gamma g'g'$ have simple closed forms:

$$\Gamma(\psi_{Q\bar{Q}} \to g'g'g') = \frac{\alpha_s'^3(N_{IC}^2 - 1)(N_{IC}^2 - 4)}{9\pi N_{IC}^2} \frac{|R_S(0)|^2}{M^2} (\pi^2 - 9),$$

$$\Gamma(\psi_{Q\bar{Q}} \to \gamma g'g') = \frac{4\alpha_s'^2 e_Q^2 \alpha}{3\pi} \frac{N_{IC}^2 - 1}{N_{IC}} \frac{|R_S(0)|^2}{M^2} (\pi^2 - 9).$$

The formula for $Zg'g'$ is shown in the appendix.

B. Charged iQuarkonium

The charged iquark-antiiquark pair once produced will settle down to a charged iquarkonium state and finally annihilates into SM particles. The decay formulae for each channel are listed in the following.

1. $f\bar{f}$

$$\Gamma(\psi_{u\bar{d}} \to u_i\bar{d}_j) = \frac{N_{IC} \alpha_W^2 |V_{ij}^{CKM}|^2}{12} \frac{1}{(1 - R_W)^2} \left[ 2 - R_i - R_j - (R_i - R_j)^2 \right] \frac{|R_S(0)|^2}{M^2},$$

where

$$\beta_{ij} = \left( (1 - R_i - R_j)^2 - 4R_iR_j \right)^{\frac{1}{2}},$$

where $R_{i,j} = r_{i,j}^2$ with $r_{i,j} = m_{i,j}/M$ and $M = M_d + M_D$. A similar decay formula can be written down for $\psi_{D\bar{d}} \to d_j\bar{u}_i$. 
2. \( W^\pm V \) with \((V = \gamma, Z)\)

For the spin singlet case, we have two contributions from the \( t \) and \( u \) channels while the \( s \) channel diagram vanishes for vector iquarks.

\[
\Gamma(\eta_{t\overline{t}} \rightarrow W^+V) = \frac{N_{IC}\alpha_W\alpha_V\beta_{WV}^3}{4} \left( \frac{c_D^V}{r - \tau R_W - r R_W} + \frac{c_U^V}{r - r R_W - \tau R_V} \right) \frac{|R_S(0)|^2}{M^2} \tag{21}
\]

Here we define \( r = M_t/M, \tau = M_D/M \) with \( M = M_t + M_D; \beta_{WV} = ((1 - R_W - R_V)^2 - 4R_W R_V)^{1/2}; \) \( c_D^V = e_u, D \) or \( v_{t, D} \) for \( V = \gamma \) or \( Z \) with \( v_{t, D} \) defined by Eq.(3); and \( \alpha_V = \alpha_{em} \) or \( \alpha_Z \) for \( V = \gamma \) or \( Z \) respectively.

For the spin triplet case, we have contributions from all \( s-, t- \) and \( u- \)channels. For the \( W\gamma \) case, we obtain

\[
\Gamma(\psi_{t\overline{t}} \rightarrow W^+\gamma) = \frac{N_{IC}\alpha_W \alpha_Z \beta_{WZ}^3}{12} \left( \frac{e_u}{r_u} - \frac{e_D}{r_D} - 2 \right) \frac{(1 - R_W^2)}{R_W} \frac{|R_S(0)|^2}{M^2}. \tag{22}
\]

It is interesting to note that all the \( s-, t- \) and \( u- \)channel amplitudes of the \( W^+\gamma \) mode completely cancel when \( M_t = M_D \). Such a cancellation is expected to avoid large contributions from the longitudinal mode of the \( W \) polarization.

Similarly, we work out the decay rate for the \( WZ \) case,

\[
\Gamma(\psi_{s\overline{s}} \rightarrow W^+Z) = \frac{N_{IC}\alpha_W \alpha_Z \beta_{WZ}^3}{24} \left( X_u - X_D - \frac{\cos^2 \theta_W}{1 - R_W} \right)^2 \frac{1}{R_W R_Z} \frac{|R_S(0)|^2}{M^2} \times \left\{ 8R_W + 2R_Z (1 + R_W) (2 + \Delta)^2 + (1 + R_W + R_Z + R_Z \Delta)^2 \right\} \tag{23}
\]

where we have defined

\[
\Delta = \left( X_u - X_D - \frac{\cos^2 \theta_W}{1 - R_W} \right)^{-1} \left( \frac{X_u}{r_u} - \frac{X_D}{r_D} - 2 (X_u - X_D) \right) \tag{24}
\]

with

\[
X_u = \frac{v_u}{1 - R_W - \frac{r_D}{r_u} R_Z} \quad \text{and} \quad X_D = \frac{v_D}{1 - R_W - \frac{r_u}{r_D} R_Z}. \tag{25}
\]

We have used a relation:

\[
e_d \left( \frac{r_D}{r_u} - 1 \right) - e_D \left( \frac{r_u}{r_D} - 1 \right) = \frac{e_u}{r_u} - \frac{e_D}{r_D} - 2 \tag{26}
\]

due to the identity \( e_u - e_D - 1 = 0 \) as a consequence of charge conservation. We take \( e_u = 2/3 \) and \( e_D = -1/3 \) as implied by our hypercharge assignment.

The CP conjugate processes give the same widths by symmetry, viz.

\[
\Gamma(\eta_{t\overline{t}} \rightarrow W^-V) = \Gamma(\eta_{t\overline{t}} \rightarrow W^+V), \quad \Gamma(\psi_{s\overline{s}} \rightarrow W^-V) = \Gamma(\psi_{s\overline{s}} \rightarrow W^+V). \tag{12}
\]
3. \( Wg'g' \)

For charged \( ^1S_0 \) quarkonium the decay into \( Wg'g' \) is zero on the amplitude level in the limit of degenerate \( M_U = M_D \). So we ignore this mode in the decay branching ratio of the charged \( ^1S_0 \) quarkonium. On the other hand, the leading mode for the charged \( ^3S_1 \) quarkonium is nonzero, and we list the formulas in the appendix. We include the \( Wg'g' \) in the decay branching ratio.

C. Decay patterns

We present the decay branching ratios of the \( S \)-wave \( ^1S_0 \) and \( ^3S_1 \) quarkonium of \( U \bar{U} \), \( D \bar{D} \), and the charged \( U \bar{D} \) in Figs. 2–4. In these plots, we have set \( N_{\text{IC}} = 3 \) and a small mass difference of \( M_U - M_D = 10 \) GeV for neutral and charged quarkonium.

The pseudoscalar state, \( \eta_{Q\bar{Q}} \), only decays into a pair of gauge bosons for both \( U \bar{U} \) and \( D \bar{D} \), as shown in Fig. 2. The dominant decay mode for \( \eta_{Q\bar{Q}} \) is \( g'g' \), which is still valid for \( \Lambda' \) down to 1 MeV. It gives rise to an invisible decay. The second largest decay mode is \( W^+W^- \) when the mass \( M \) of the quarkonium is above \( 2m_W \) threshold; otherwise \( \gamma\gamma \) and \( Z\gamma \) are large when \( M \) of the quarkonium is below \( 2m_W \) threshold. On the other hand, the fermion-antifermion modes are dominant in the decay of \( \psi_{Q\bar{Q}} \) states; especially the down-type quarks, followed by the up-type quarks and then lepton modes: see Fig. 3. The branching ratio into \( W^+W^- \) is small and so are the \( Zg'g' \), \( \gamma g'g' \), and \( g'g'g' \) modes. We show the decay branching
FIG. 3: Branching fractions of the iquarkonium of (a) $^3S_1(UU)$ and (b) $^3S_1(D\bar{D})$ versus the iquarkonium mass $M$. We have chosen $\Lambda' = 10$ MeV and $n_Q = 1$ in the running of $\alpha'_s$.

FIG. 4: Branching fractions of the charged iquarkonium of (a) $^1S_0(U\bar{D})$ and (b) $^3S_1(U\bar{D})$ versus the iquarkonium mass $M$. We have chosen $\Lambda' = 10$ MeV and $n_Q = 1$ in the running of $\alpha'_s$.

ratios for the charged $\eta_{U\bar{D}}$ and $\psi_{U\bar{D}}$ in Fig. 4. There are only two modes for $\eta_{U\bar{D}}$ state, namely $W^+\gamma$ and $W^+Z$. Whereas for the $\psi_{U\bar{D}}$ state the fermion pair modes dominate.

D. Comparison with 4-th generation quarkonium

A superheavy quarkonium made up of a pair of sequential 4-th generation quark and antiquark has QCD as well as Higgs interactions. We have to assume that each of the 4-th generation quark is stable against weak decay. In general it is not true, but just for comparison with iquarkonium we temporarily assume this is actually the case. The major
decay mode for the $^1S_0$ state is $gg$ while that of $^3S_1$ state is $ggg$. These QCD decays are far more efficient than those of electroweak decays. Therefore, we can make the following observations

1. The major decay mode of a $\eta_{Q\bar{Q}}$ iquarkonium is $g'g'$, followed by $\gamma\gamma, Z\gamma, WW$, and $ZZ$ so that it decays mainly into an invisible mode, which is very different from the 2-jet mode of a superheavy quarkonium.

2. The majority of the decay modes of a $\psi_{Q\bar{Q}}$ iquarkonium for all its mass range is $d\bar{d} + s\bar{s} + b\bar{b}$, which give rises to two jets in the final state. On the other hand, once the $WW$ channel is open, it will dominate over all other modes in the superheavy quarkonium decay while this mode occupies only about 1% for iquarkonium decay. Presumably the superheavy quarkonium decays dominantly into the longitudinal components of the $W$ bosons, while for the vectorial iquarks the longitudinal piece is largely cancelled in the limit of degenerate $M_U = M_D$ and $M \gg m_W$. Furthermore, the $3g$ mode of the superheavy quarkonium will give rise to a 3-jet final state while the $3g'$ mode of the iquarkonium has small branching ratio and is invisible.

3. For the $^1S_0$ charged quarkonium it decays into $W\gamma$ or $WZ$ in the leading order. This is similar to $\eta_{U\bar{U}}$, so it is hard to distinguish the charged $^1S_0$ state.

4. The $^3S_1$ charged quarkonium will decay into fermion pairs via a virtual $W$ boson, which is similar to the charged iquarkonium. However, the $^3S_1$ charged quarkonium also has the $Wgg$ mode, which is comparable to its $f\bar{f}'$ mode, while the $Wg'g'$ mode of the iquarkonium is very small. So, there is a chance to distinguish between the $^3S_1$ charged iquarkonium from the quarkonium.

V. CONCLUSIONS

Production of iquark-antiquark pair connected by unbreakable long and stable string associated with a new confining non-abelian gauge group may lead to spectacular events at the LHC. We presented in some details the phenomenology of uncolored iquarks in this work. Since the momentum transfer for the electroweak processes producing the iquark-antiquark pairs is typical of order $M_Q$ and one has to bring the pairs at a distance of order $1/M_Q$ in
order for them to annihilate, the bound states are thus formed with high excitation energies and large orbital momenta. These energies are quickly dissipated by emitting iglueballs and photons before they settle down in the ground state and annihilate into standard model particles.

We have studied open production of the uncolored iquarks at the LHC and the leading-order 2-body and 3-body decays of $S$-wave iquarkonium. The decay patterns of $S$-wave iquarkonium formed by a vector-like iquark doublet are found to be distinguishable from a superheavy quarkonium composed of a sequential 4-th generation quark doublet. With our choice of $\Lambda' \sim 10$ MeV, the emitted iglueballs due to energy loss as the iquark-antiquark crossing each other before annihilation may decay outside the detectors and become invisible \[4\]. Developing new search strategies at the LHC for detecting the decay products of iquarkonium together with many soft photons emitted due to the energy loss are very important to unravel this kind of new physics beyond the standard model. Some of these issues have been addressed in a recent article of Ref.\[10\]. For cosmological implications of iQCD, we refer the readers to the literature \[11\].

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APPENDIX A: THREE-BODY DECAY FORMULAS

For the process $\psi_{UD} \rightarrow W^+ g' g'$ we define the rescaled energy variables in the rest frame of the iquarkonium:

$$x_1 = \frac{2E_1}{M}, \quad x_2 = \frac{2E_2}{M}, \quad x_v = \frac{2E_W}{M}$$

so that $x_1 + x_2 + x_v = 2$. We also define

$$\xi = \frac{m_W^2}{M^2}.$$  

The differential width, in the limit of degenerate $M_U = M_D$, is given by

$$\frac{d\Gamma}{dx_v dx_1} (\psi_{UD} \rightarrow W^+ g' g') = \frac{2\alpha'^2\alpha}{3\pi \sin^2 \theta_w} \frac{N_{IC}^2 - 1}{N_{IC}} \frac{|R_s(0)|^2}{M^2} \frac{1}{x_1 x_2 (x_v - 2\xi)^2} \times \left[ 2\xi^4 + 2\xi^3 (6 - 4x_v + 2x_1 - x_v x_1 - x_1^2) 
\quad + 2\xi^2 \left( 11 - 16x_v + 6x_v^2 - (8 - 2x_v - x_v^2) x_1 + (4 + x_v) x_1^2 \right) 
\quad + \xi \left( 4(1 - x_v)(4 - 5x_v + 2x_v^2) - (32 - 44x_v + 14x_v^2)x_1 + (20 - 18x_v + x_v^2)x_1^2 
\qquad - 2(2 - x_v)x_1^3 + x_1^4 \right) 
\quad + 2 \left( 2 - 6x_v + 7x_v^2 - 4x_v^3 + x_v^4 - (6 - 13x_v + 9x_v^2 + 2x_v^3)x_1 
\qquad + (7 - 9x_v + 3x_v^2)x_1^2 - 2(2 - x_v)x_1^3 + x_1^4 \right) \right] .$$  

(A1)

The ranges of integration for $x_v$ and $x_1$ are

$$\frac{1}{2} \left( 2 - x_v - \sqrt{x_v^2 - 4\xi} \right) \leq x_1 \leq \frac{1}{2} \left( 2 - x_v + \sqrt{x_v^2 - 4\xi} \right) .$$  

(A2)

Note that $\eta_{UD} \rightarrow W^+ g' g'$ is zero on the amplitude level in the limit of degenerate $M_U = M_D$.

The substitutions needed to obtain $\psi_{UD} \rightarrow Z g' g'$ or $\psi_{D\overline{D}} \rightarrow Z g' g'$ from the above formula are

$$\frac{\alpha}{2 \sin^2 \theta_w} \rightarrow \frac{\alpha v_Q^2}{\cos^2 \theta_w \sin^2 \theta_w} , \quad m_W \rightarrow m_Z .$$

The decay of $\psi_{UD} \rightarrow \gamma g' g'$ or $\psi_{D\overline{D}} \rightarrow \gamma g' g'$ can be easily obtained from the above formula with the replacement

$$\frac{\alpha}{2 \sin^2 \theta_w} \rightarrow e_Q^2 \alpha , \quad \xi \rightarrow 0 .$$
So the differential partial width is given by
\[
\frac{d\Gamma}{dx_vdx_1}(\psi_{Q\bar{Q}} \rightarrow \gamma g'g') = \frac{4\alpha_s^2\alpha e_Q^2}{3\pi} \frac{N_{IC}^2 - 1}{N_{IC}} \frac{|R_S(0)|^2}{M^2} \frac{1}{x_1^2x_2^2(x_v - 2\xi)^2} 
\times \left[ 2 \left( 2 - 6x_v + 7x_v^2 - 4x_v^3 + x_v^4 - (6 - 13x_v + 9x_v^2 - 2x_v^3) \right) 
+ (7 - 9x_v + 3x_v^2)x_1^2 - 2(2 - x_v)x_1^3 + x_1^4 \right],
\]
(A4)

with the integration range

\[ 0 \leq x_v \leq 1, \quad 1 - x_v \leq x_1 \leq 1. \]

After integrating over \( x_1 \) and \( x_v \) we obtain
\[
\Gamma(\psi_{Q\bar{Q}} \rightarrow \gamma g'g') = \frac{4\alpha_s^2\alpha e_Q^2}{3\pi} \frac{N_{IC}^2 - 1}{N_{IC}} \frac{|R_S(0)|^2}{M^2} (\pi^2 - 9). \quad (A5)
\]

Finally, the decay width of \( \psi_{Q\bar{Q}} \rightarrow g'g'g' \) is
\[
\Gamma(\psi_{Q\bar{Q}} \rightarrow g'g'g') = \frac{\alpha_s^3}{9\pi} \frac{(N_{IC}^2 - 1)(N_{IC}^2 - 4)}{N_{IC}^2} \frac{|R_S(0)|^2}{M^2} (\pi^2 - 9). \quad (A6)
\]

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