Production and Decay Properties of Ultra-Heavy Quarks[*]

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ABSTRACT

The widths of ultra-heavy quarks that can decay into W, Z or Higgs bosons are discussed. If the lifetimes become much shorter than the typical strong interaction time scale $\Lambda_{QCD}^{-1} \sim 10^{-23}$ sec, then open-flavor hadrons and quarkonium bound states cannot be formed any more. Consequences for the jet evolution are investigated. On the other hand, if such quarks can decay only through tiny mixing angles - as it could happen for sequential down-type quarks and for SU(2) singlet quarks in E_6 models - then these bound states do form. Production rates for quarkonia in e^+e^- annihilation and in hadronic collisions are estimated and their decay signatures are discussed.

1. Introduction

It has often been conjectured that more leptons and quarks might exist than those presently known in the first three generations, with masses in the range of $\mathcal{O}$ (100 GeV):

(i) No explanation has been given so far for the replication of fermion generations in the Standard Model. The existence of a fourth family is not in conflict with any experimental facts. It has even been speculated recently that another generation might facilitate the description of CP violation in K_L decays.

(ii) Each generation itself might have a richer fermion content than presently observed. Grand unified models based on E_6, for example, predict the existence of isoscalar, charge - 1/3 quarks with small, quite possibly tiny mixing to the well-known doublet quarks.

The only constraint on heavy quark masses follows from the requirement that low energy observables must not be renormalized too strongly. In particular, a stringent bound on the mass splitting in isodoublets can be derived from fermion-

loop contributions to the ρ parameters[1]

$$\left| m_U^2 - m_D^2 \right|^{1/2} \lesssim 0.3 \text{ TeV} \tag{1}$$

The masses in nearly degenerate doublets can therefore be very large. Since high energy e^+e^- and hadron-hadron colliders will provide the opportunity to explore the $\mathcal{O}$ (100 GeV) domain experimentally [and hopefully the range beyond in the more distant future], a general discussion of the production of such particles and their decay properties is called for. In this note we shall discuss two topics:

(i) First we focus on the ultra-heavy quarks that decay semi-weakly into light quarks and W/Z bosons[2] $-Q \rightarrow q + W/Z-$ or Higgs bosons[3] $- Q \rightarrow q + H^\pm/H^0$. One finds that the Q lifetime will - for a sufficiently large quark mass m_Q - become much shorter than the time scale $\Lambda_{QCD}^{-1} \sim 10^{-23}$ sec typically needed to build up open-flavor hadrons $(Q\bar{q})$ and quarkonium bound states $(Q\overline{Q})$. In many aspects such quarks behave like (quasi)-free particles, their dynamical properties determined only by electroweak and perturbative strong interactions. The large width of these states requires a refined analysis of the threshold behavior in e^+e^- annihilation generally used to measure particle masses very accurately. A very similar discussion applies to corrections due to the finite W/Z width in Q decays throughout the transition region from virtual to real bosons.

(ii) In the second part we will elaborate on a scenario where the decays $Q \rightarrow q + W/Z$, $q + H$ are greatly suppressed due to tiny quark mixing angles as it might happen for sequential down-type quarks and for singlet quarks in E_6 models. In an extreme case such quarks Q might even be stable. Then $(Q\bar{q})$ and $(Q\overline{Q})$ bound states can form. We analyze their production and decay properties in more detail.

2. Q Decay

Adopting the Standard Model coupling with a mixing parameter $V(Qq)$ for the $Q \to q + W$ decay vertex, one finds for the width

$$
\Gamma(Q \to q + W) = \frac{G_F m_Q^3}{8 \, \pi \sqrt{2}} \, |V(Qq)|^2 \, \frac{2k}{m_Q}
$$

$$
\left\{ 1 - \left(\frac{m_q}{m_Q}\right)^2 + \left[1 + \left(\frac{m_q}{m_Q}\right)^2 \right] \left(\frac{m_W}{m_Q}\right)^2 - 2 \left(\frac{m_W}{m_Q}\right)^4 \right\}
\tag{2}
$$

where

$$
k = \left[m_Q^2 - (m_W + m_q)^2 \right]^{1/2} \left[m_Q^2 - (m_W - m_q)^2 \right]^{1/2} \bigg/ 2 \, m_Q
$$

denotes the W momentum in the Q rest frame. This width quickly reaches the asymptotic form

$$
\Gamma(Q \to q + W) \simeq 180 \text{ MeV} \, |V(Qq)|^2 \left(\frac{m_Q}{m_W}\right)^3
\tag{3}
$$

and the lifetime drops below 10^{-23} sec. This behaviour is illustrated in Fig. 1 (solid line) where the light quark mass is chosen to be 5 GeV. The final state consists of a q jet plus a lepton pair or two other jets from W decay. For Q an up-type quark we expect the mixing parameter $|V(Qq)|$ to be close to unity while for a down-type quark it should be much smaller, namely $|V(Qq)| \sim \theta_c$, θ_c^2 or even less.

Near W threshold Eq. (2) has to be modified due to the finite W width which provides a smooth transition to the conventional weak three-body decay below W threshold. The partial width into $qe\nu$, for instance, via virtual or real W

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emission can be written as

$$\Gamma(Q \to q + W(\to e\nu)) = \frac{G_F^2 m_Q^5}{192\pi^3} \; |V(Qq)|^2 \; f\left(\frac{m_Q^2}{m_W^2}, \; \frac{m_q^2}{m_Q^2}, \; \frac{\Gamma_W^2}{m_W^2}\right)$$

$$f(\rho, \; \mu, \; \gamma) = 2 \int\limits_0^{(1-\sqrt{\mu})^2} \frac{dx}{[(1 - x\rho)^2 + \gamma^2]^2} \left[(1 - \mu)^2 + (1 + \mu)x - 2x^2\right]$$

$$\times \sqrt{1 + \mu^2 + x^2 - 2(\mu + \mu x + x)}$$

$$\tag{4}$$

For quark masses m_Q up to ~ 3 GeV below the threshold for real W production one can ignore the effect of the non-vanishing W width; the W propagator, however, remains important for smaller masses as well, as shown in Fig. 2. The narrow width approximation becomes again adequate for m_Q roughly 10 GeV above the qW threshold. A similar calculation can be found in ref.17.

Isoscalar quarks might decay via Z° emission; the required flavor-changing neutral current coupling could be generated via mixing of the left-handed components of isoscalar and isodoublet quarks. The mass dependance of $\Gamma(Q \to Z + q)$ is the same as in Eq. (2) - m_W replaced by m_Z - while the overall normalization is reduced by a factor of two.

Charged Higgs particles - as for example predicted by SUSY extensions of the Standard Model - provide a similarly rapid decay mode of heavy quarks. Defining the coupling by $\sqrt{G_F/\sqrt{2}} \; M_Q\bar{q}\left[(1 - \gamma_5) \; v_L + (1 + \gamma_5) \; v_R\right]Q$ one obtains for the width of this decay mode[3]

$$\Gamma(Q \to q + H) = \frac{G_F m_Q^3}{8\pi\sqrt{2}} \; \frac{2k}{m_Q} \left\{ \left(|v_L|^2 + |v_R|^2\right) \times \right.$$

$$\tag{5}$$

$$\left. \left[1 + \left(\frac{m_q}{m_Q}\right)^2 - \left(\frac{m_H}{m_Q}\right)^2\right] + 4 \; Re v_L v_R^* \; \frac{m_q}{m_Q} \right\}$$

which rises asymptotically again like m_Q^3. The mass dependence is shown for a light and heavy charged Higgs by the dotted lines in Fig. 1. These widths get

reduced by mixing parameters if the $Q \to q$ transition connects two different generations; the relative strength of H versus W transitions, however, remains unchanged. At first thought it might be expected that a peculiar choice of the unknown couplings v_L, v_R could reduce this rate to an arbitrarily small level. However in a model with just two Higgs doublets[4] $v_L = \tan\beta$, $v_R = \dfrac{m_q}{m_Q}\cot\beta$ hold so that the minimal rate is obtained for $\tan^2\beta = m_q/m_Q$. Even for this extreme choice, the Higgs decay mode would still dominate the conventional weak decays for Higgs masses nearly up to the kinematic limit.

3. Short-lived Quarks: Hadron Formation and Jet Evolution

Rapid semiweak decays have very interesting consequences for strong interaction phenomena involving ultra-heavy quarks:

(i) The binding force in a heavy quarkonium state is essentially Coulombic. The revolution time of the $(Q\overline{Q})$ bound state is then estimated as $t_R \sim 9/(4m_Q\alpha_s^2)$. If the lifetime of the $(Q\overline{Q})$ system becomes shorter than the revolution time t_R, then the quarkonium bound state cannot be formed any more. Setting $\alpha_s = 0.15$ to illustrate the point one finds this to happen for

$$m_Q \gtrsim 125 \text{ GeV } \times |V(Qq)|^{-\frac{2}{3}} \qquad (6a)$$

Equation (6a) can be interpreted also in a different way: when the uncertainty in the quark masses becomes larger than the $(Q\overline{Q})$ level spacing of ~ 800 MeV estimated in potential models,[5] the non-perturbative binding forces become ineffective. Any resonance structure in $Q\overline{Q}$ production near threshold is washed out, and the properties of the process $e^+e^- \to Q\overline{Q}$ follow literally the predictions of the free quark model, modified only by perturbative QCD corrections. The apparent threshold will be lowered by about $2\Gamma_Q$, as shown in Fig. 3. This can be described by smearing out the quark mass values (in the phase space factor of $\sigma(e^+e^- \to Q\overline{Q})$) by Breit-Wigner expressions, an approximation valid as long

as the matrix element does not change rapidly over the threshold domain. [A similar prescription has been used for W^+W^- production[6]]. This pattern complicates an accurate measurement of m_Q and limits precision tests of potential models already at lower energies.

A somewhat stronger limit applies to $(Q\bar{q})$ bound states. We find a mass limit of

$$m_Q \gtrsim 100 \text{ GeV} \times |V(Qq)|^{-\frac{2}{3}} \tag{6b}$$

above which no more open-flavor hadrons can exist, i.e. Q decays before it can form a meson by picking up a light quark q: $\tau_Q < t_{\text{Had}} \sim \Lambda_{\text{QCD}}^{-1} \sim 10^{-23}$ sec.

(ii) Hadronization of a heavy quark jet with $E_Q \gg m_Q$, produced for example in e^+e^- collisions, is a complicated interplay of non-perturbative as well as perturbative QCD effects. In a parton picture, due to the inertia of the heavy quark, most of the energy resides in the $(Q\bar{q})$ open-flavor hadrons leaving only a small fraction for the light particles in the accompanying jet.[7-10] This leading particle effect and the truncation of the plateau has in fact been observed for charm and bottom jets, and the total energy of the light particles in a jet amounts to approximately $(0.6 \text{ GeV}/m_Q)E_Q$.[9] For ultra-heavy quarks the picture simplifies considerably: the energy <u>fraction</u> due to the non-perturbative QCD force approaches zero when the lifetime τ_Q becomes much less than the typical strong interaction time characterized by the confinement radius $R_{\text{conf}} \sim \Lambda_{\text{QCD}}^{-1} \sim 10^{-23}$ sec; on the other hand the degrading of the virtual quark Q due to perturbative gluon bremsstrahlung needs less than the decay time for any foreseeable energy.

At high energies $E \gg m_Q$, perturbative gluon bremsstrahlung softens the Q spectrum to[10]

$$D_Q(x_Q = E_Q/E) = A \frac{1 + x_Q^2}{2} (1 - x_Q)^{-1 + \frac{16}{3}\Delta\xi} \tag{7}$$

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where

$$\Delta\xi = \frac{1}{b}\log\left(\alpha_s(m_Q^2)\,/\,\alpha_s(E^2)\right)$$

$$b = 11 - \frac{2}{3}\,N_F$$

N_F is the number of active flavors, and A is chosen such that $\int dx_Q D_Q = 1$. For all practical cases $\Delta\xi$ remains small (for $\Lambda \sim 100$ MeV, $M_Q \gtrsim 100$ GeV: $\Delta\xi \lesssim 0.05$ even at $E = 5$ TeV). Note that the distribution turns into a δ-function for $\Delta\xi \to 0$. The distribution $D_Q(x_Q)$ is infrared stable for $(1 - x_Q) \gg (m_Q\Gamma_Q R_{conf}^2)^{-1}$. For $E \gg m_Q$ the average fraction of energy residing in the quark Q is given by

$$\langle x_Q \rangle = \left[\frac{\alpha_s(m_Q^2)}{\alpha_s(E^2)}\right]^{-\frac{32}{9b}}\left(1 + \frac{44}{27}\,\frac{\alpha_s}{\pi}\right) \tag{8}$$

The distributions of W (or H) are controlled by perturbative QCD, unaffected by hadronization of the heavy quark. Apart from a few low energy hadrons due to the remnant non-perturbative interaction energy, the final state jet distributions in $e^+e^- \to Q\overline{Q}$ after Q, $\overline{Q}$ fragmentation and decay are analogous to particle distributions in heavy lepton pair production. [In events like $e^+e^- \to Q(\to W + q) + \overline{Q}(\to W + \bar{q})$ with the W's decaying leptonically, the q, $\bar{q}$ jets will fragment independently of each other, their color fluxes connected to the remnants of Q, $\overline{Q}$ fragmentation]. Note that the emission of hard particles with $E_H > (E/m_Q)\Gamma_Q$ is suppressed at angles $\theta < m_Q/E$ due to the short lifetime of the heavy quark.

Near threshold when $(E - m_Q)/m_Q \ll 1$, the perturbative gluon bremsstrahlung is drastically suppressed due to the dipole character of the emission. The average energy loss through perturbative radiation is

$$1 - \langle x_Q \rangle_{\text{pert}} = \frac{16}{3}\,\frac{\alpha_s}{\pi}\left(\frac{E - m_Q}{m_Q}\right)^2 \tag{9}$$

If the emitted gluon energy is less than the Q width, the $Q\overline{Q}$ is produced practically without accompanying color quanta, and the production characteristics

will parallel those for heavy leptons near threshold. As a consequence, the quark is, for example, not depolarized through either fragmentation or perturbative multigluon emission in electroweak e^+e^- annihilation.[11] The primary quark polarization in $e^+e^- \to \gamma$, $Z^\circ \to Q\overline{Q}$ can therefore manifest itself via parity violations in the energy and angular distributions of the decay products [even for top production at LEP and SLC one expects at most a few light, low energy hadrons to accompany the top meson pair.][11,12]

4. Long-lived Quarks:Quarkonium Formation and Decay

$(Q\overline{Q})$ quarkonium resonances can be formed even for very large quark mass values if the mixing parameter is sufficiently small:$m_Q \leq 125 \; GeV \times |V(Qq)|^{-2/3}$. While for up-type quarks, $I_3 = 1/2$, $|V(Qq)| \sim 1$ should hold, we expect on the other hand $|V(Qq)| \sim \theta_c^n, n = 1, 2, 3, \ldots$ for down-type quarks, $I_3 = -1/2$, which reduces single quark decays significantly. For isoscalar quarks in E_6 type GUTs $|V(Qq)|$ might even be much smaller. Nevertheless such states are not easy to observe.

The production cross section in e^+e^- annihilation on the peak of the $(Q\overline{Q})$ resonance depends on the energy spread ΔW of the beams:[2]

$$R_{\mathrm{peak}} = \frac{9}{2\sqrt{2}\pi\alpha^2} \frac{\Gamma((Q\overline{Q}) \to e^+e^-)}{\Delta W} \tag{10}$$

Since $\Gamma((Q\overline{Q}) \to e^+e^-) \sim 10 keV$ is expected to hold away from Z,Z',... resonances, a decent cross section is obtained only if the beam spread could be kept sufficiently small in a dedicated effort.

The cross section for the production of the O^{-+} $(Q\overline{Q})$ bound state in hadronic collisions is estimated to be [13]

$$\sigma\left(p\overset{(-)}{p} \to O^{-+} + X\right) \sim \frac{\pi^2}{8m_{Q\overline{Q}}^2} \frac{\Gamma(O^{-+} \to gg)}{m_{Q\overline{Q}}} \left[\tau\frac{d\mathcal{L}}{d\tau}\right]_{gg} \tag{11}$$

where $\left[\tau\frac{d\mathcal{L}}{d\tau}\right]_{gg}$ denotes the gluon luminosity for $\tau = \frac{m_{Q\overline{Q}}^2}{s}$. Using the estimate

$\Gamma(0^{-+} \to gg) \simeq 2$ MeV and the luminosities computed in ref.13,14 we obtain the cross section estimates shown by the dashed lines in Figs. 4 and 5 for the three different energies $\sqrt{s}$ =0.63,2,40 TeV.Uncertainties in the cross sections are estimated to be of the order of three. Above $m_{Q\overline{Q}}$ ~400 GeV the width into two gluons increases slightly. Below $m_{Q\overline{Q}}$ ~400 GeV the only distinctive decay mode is the $\gamma\gamma$ decay with a branching ratio of 1 % . However the background due to π^0 production in jets would be very severe. At $m_{Q\overline{Q}}$ ~400 GeV the spectacular decay mode $(Q\overline{Q}) \to Z^0 H$ could become dominant for isodoublet down-type quarks[15].

To estimate the production of the 1^{--} quarkonium state we employ a different method: first one computes the continuum $Q\overline{Q}$ cross section near threshold (see, e.g., ref. 13)

$$\sigma(q\bar{q} \to Q\overline{Q}) \simeq \frac{4\pi\alpha_s^2}{9\, m_{Q\overline{Q}}^3} \sqrt{\hat{s} - m_{Q\overline{Q}}^2} \qquad (12a)$$

$$\sigma(gg \to Q\overline{Q}) \simeq \frac{7\pi\alpha_s^2}{48\, m_{Q\overline{Q}}^3} \sqrt{\hat{s} - m_{Q\overline{Q}}^2} \qquad (12b)$$

where $\sqrt{\hat{s}}$ denotes the CM energy of the $q\bar{q}$ or gg system respectively. Then we invoke the concept of duality to argue that integrating these parton cross sections from threshold over the level spacing $\Delta \sim 800$ MeV [5] should give a *rough estimate* of the $1^{--}(Q\overline{Q})$ production rate. The integration which contains as a weight factor the $q\bar{q}$ and gg luminosities [13,14] yields

$$\langle\sigma_q\rangle \sim \frac{8\pi\alpha_s^2}{27m_{Q\overline{Q}}^2} \left[\tau\frac{d\mathcal{L}}{d\tau}\right]_{gg} \left(\frac{2\Delta}{m_{Q\overline{Q}}}\right)^{3/2} \qquad (13a)$$

$$\langle\sigma_g\rangle \sim \frac{7\pi\alpha_s^2}{72m_{Q\overline{Q}}^2} \left[\tau\frac{d\mathcal{L}}{d\tau}\right]_{q\bar{q}} \left(\frac{2\Delta}{m_{Q\overline{Q}}}\right)^{3/2} \qquad (13b)$$

This estimate is shown by the solid curves in Figs. 4 and 5 for the three different CM energies $\sqrt{s} = 0.63, 2, 40$ TeV using $\alpha_S = 0.15$. It should be clear from the

preceding discussion that these curves represent ball-park estimates only. At the parton level the produced $Q\overline{Q}$ system is not yet in the proper colour and spin configuration for the production of a 1^{--} bound state; therefore the solid curves in Figs. 4 and 5 could represent overestimates.

A possible clean signature for the 1^{--} production is provided by the leptonic mode $1^{--} \rightarrow e^+e^-, \mu^+\mu^-$. In discussing the branching ratios for these modes we have to distinguish between three different cases: (i) Q belongs to an isodoublet with electric charge $e_Q = -\frac{1}{3}$ or (ii) $e_Q = +\frac{2}{3}$ or (iii) Q forms an isoscalar with $e_Q = -\frac{1}{3}$. Applying the formulae of Ref. 2 with $\sin^2\theta_w = 0.23$ one then finds:

$$\text{Case (i)} \quad R_i \equiv \frac{\Gamma(1^{--} \rightarrow e^+e^-, \mu^+\mu^-)}{\Gamma(1^{--} \rightarrow \sum f\bar{f}, 3g, \gamma H, \gamma Z)} \sim 0.064; \ 0.07; \ 0.06 \quad (14a)$$

$$\text{Case (ii)} \quad R_{ii} \simeq 0.11; \ 0.10; \ 0.077 \quad\quad\quad\quad\quad (14b)$$

where the three numbers each in (14a,b) refer to $m_{Q\overline{Q}} = 120, 200, 300 GeV$ and $\sum f\bar{f}$ denotes the sum over all fermion-antifermion final states. Note that these ratios change very little over this range in $m_{Q\overline{Q}}$.

Case (iii) has to be treated somewhat differently since isoscalar quarks couple to isoscalar Higgs fields and not to the conventional Higgs doublet, at least in the absence of quark mixing. Since such isoscalar Higgs are presumably quite heavy we ignore the transition $Q\overline{Q} \rightarrow \gamma H$ and thus find

$$\text{Case (iii)} \quad R_{iii} \equiv \frac{\Gamma(1^{--} \rightarrow e^+e^-, \mu^+\mu^-)}{\Gamma(1^{--} \rightarrow \sum f\bar{f}, 3g)} \simeq 0.10; \ 0.12; \ 0.13 \quad (14c)$$

For $m_{Q\overline{Q}} > 2 \ m_W$ one has to consider $(Q\overline{Q}) \rightarrow W^+W^-$ as well. Modifying the results of Ref. 16 appropriately one finds that this channel is very insignificant for Q being an isoscalar. This is not surprising since isospin arguments tell us that this transition has to be suppressed for $m_{Q\overline{Q}} \gg m_Z$.

If Q is an isodoublet quark then $(Q\overline{Q}) \to W^+W^-$ becomes significant once the threshold suppression fades away. For example the leptonic branching ratios (14a, b) get reduced for $m_{Q\overline{Q}} = 300$ GeV: $R_i \simeq 0.043$, $R_{ii} \simeq 0.043$.

So far we have ignored single quark decays (=SQD) $(Q\overline{Q}) \to Q\bar{q}$ "W" where "W" denotes an on or off shell W (or Z) boson, apart from noting that for $M_Q \gtrsim 125$ GeV $\times |V(Qq)|^{-\frac{2}{3}}$ no resonance can be formed. For quarks Q with $I_3 = +\frac{1}{2}$ one expects $|V(Qq)| \sim 1$. Then for $M_{Q\overline{Q}} \geq 110$ GeV SQD already dominates $(Q\overline{Q})$ decays and the branching ratio for $1^{--} \to e^+e^-$, $\mu^+\mu^-$ quickly decreases below the 1% level.

For quarks with $I_3 = -1/2$ on the other hand one expects $|V(Qq)| \sim \theta_c^n, n = 1, 2, 3, \ldots$ which reduces SQD significantly. If $|V(Qq)| \sim \theta_c^3$ - a conjecture made by several authors - then one reads off from Fig. 1: $\Gamma_{SQD}((Q\overline{Q})) \simeq 2\Gamma(Q) \sim 200 \ MeV \ |V(Qq)|^2 \sim 30 \ keV$ for $m_{Q\overline{Q}} = 200 \ GeV$; this reduces the leptonic branching ratio of 1^{--} to roughly 5% . The same argument applies to isoscalar quarks where one might actually expect even smaller mixing angles: $|V(Qq)| \leq \theta_c^4$; then SQD contributes not more than 24 % to $\Gamma(1^{--})$ for $m_{Q\overline{Q}} = 250 GeV$.

From Figs. 4 and 5 one reads off that the production cross section for a $(Q\overline{Q})$ vector resonance with $m_{Q\overline{Q}} \leq 250$ GeV exceeds 1 pb for $\sqrt{s} = 40$ TeV. An integrated luminosity $\int L dt = 10^{40}$ cm^{-2} then yields more than 10^4 produced 1^{--} states. For $BR(1^{--} \to e^+e^-$, $\mu^+\mu^-) > 1\%$ one obtains > 100 leptonic decays. As discussed above one expects $BR(1^{--} \to e^+e^-$, $\mu^+\mu^-) \sim 3 - 13\%$ for quarks with $e_Q = -\frac{1}{3}$; the argument becomes even stronger for isoscalar quarks. Thus the SSC would allow searches for 1^{--} resonances with $M_{Q\overline{Q}}$ at least up to 250 GeV if $e_Q = -\frac{1}{3}$, if a sufficiently good mass resolution and background rejection can be achieved.

The decays in particular of isoscalar quarks are expected to be suppressed by tiny mixing angles $V(Qq)$. If $|V(Qq)| \sim \theta_C^4$ for example then $(Q\overline{Q})$ resonances can form up to $M_{Q\overline{Q}} \sim 12$ TeV. Extrapolating our curves of Fig. 5 one finds a production cross section that could be as large as 10^{-1} pb for $M_{Q\overline{Q}} \sim 1$ TeV if

$\sqrt{s} = 40$ TeV. If the decay $(Q\overline{Q}) \rightarrow \gamma H$ is still forbidden by phasespace then $BR(1^{--} \rightarrow e^+e^-, \mu^+\mu^-) \sim 14\%$ is estimated leading to 14 $1^{--} \rightarrow e^+e^-, \mu^+\mu^-$ decays for $\int L dt = 10^{40}$ cm^{-2}. This presumably denotes the limit of sensitivity.

It is amusing to note that this leptonic branching ratio actually increases for stable quarks, i.e. $|V(Qq)| \rightarrow 0$ since SQD then becomes irrelevant. Thus the decays $(Q\overline{Q}) \rightarrow e^+e^-, \mu^+\mu^-$ would still provide a striking signature while it would be very difficult in this extreme case to identify the open flavor hadrons $(Q\bar{q})$.

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FIGURE CAPTIONS

Fig.1: The decay widths of ultra-heavy quarks Q

Fig.2: The effects of the finite W width on $\Gamma(Q)$

Fig.3: The impact of $\Gamma(Q)$ on the threshold behaviour for $Q\overline{Q}$ production.

Fig.4: 0^{-+} and 1^{--} production cross sections in $p\overset{(-)}{p}$ collisions for $\sqrt{s}$=0.63, 2 TeV. Parton luminosities are taken from ref. 13.

Fig.5: The same as Fig. 4 for $\sqrt{s}$=40 TeV. Parton luminosities from ref. 14.

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Fig. 1

$$\frac{\Gamma\ (Q \rightarrow q+e+\nu)}{G_F^2\ m_Q^5\ /192\ \pi^3}$$

$m_q = 5$ GeV
$m_w = 82$ GeV

——— $\Gamma_w = 2.3$ GeV
— — — $\Gamma_w = 0$

10^2

10^1

50   60   70   80   90   100

m_Q (GeV)

5520A3

Fig. 2

$\times 10^{-3}$

2 particle
phase space

——— Breit-Wigner

— — Narrrow width
approximation

$\Gamma_Q / m_Q = 10^{-3}$

$(E - 2\, m_Q)/\Gamma_Q$

5520A4

Fig. 3

Fig. 4

Fig. 5