WEAK EFFECTS IN e^+e^- PRODUCTION OF SUPERHEAVY FLAVOURS

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ABSTRACT

We investigate the experimental possibilities of detecting new superheavy flavours in e^+e^- annihilation when considering both electromagnetic and weak contributions and taking radiative corrections into account. For masses of the order of m_W or higher, the production of bound vector states is dominated by weak current effects, with cross-sections strongly peaked at the mass of the Z^0 . For continuum thresholds, also the sphericity behaviour is considerably distorted by the Z^0 peak and its radiative tail. A direct quarkonium search seems to be feasible only up to the gauge bosons mass region, whereas the detection of enhancements in the mean sphericity associated with either $q\bar{q}$ bound states or continuum thresholds allows the detection of superheavy flavours also above the gauge bosons mass region.

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1. INTRODUCTION

The excitation of new flavours, associated with superheavy quarks, may be one of the main objectives of the next generation of e^+e^- storage rings. This in fact would be of considerable relevance to our understanding of the group structure underlying the spectrum of hadronic constituents. In addition, bound vector states of the new flavours, for which a complex spectroscopy can be expected¹⁾, represent an almost ideal laboratory for studying some fundamental aspects of strong interactions, such as asymptotic freedom, confinement, and the perturbative applications of quantum chromodynamics²⁻¹⁰⁾.

At high enough energy the production of superheavy quarks, both in bound systems and in the continuum region, should reveal several new features because of the increasingly important role played by weak and radiative effects. In particular, weak effects, which become increasingly important with energy in quarkonium production and decay, may provide substantial new information on the dynamics of new flavours.

In Section 2 we first discuss the production of heavy, narrow vector mesons with the inclusion of weak contributions and electromagnetic radiative corrections in the framework of perturbative QCD and the Weinberg-Salam model. Taking into account the experimental energy resolution and the topology of the final states, this leads to specific predictions on the observability of new heavy quarkonium states. In Section 3 we discuss in the same framework the experimental implications for the detection of new flavour thresholds in the continuum region. Our conclusions are given in Section 4.

2. QUARKONIUM PRODUCTION

Interpreting heavy, narrow vector mesons as almost non-relativistic quark-antiquark bound systems, the dynamics of the lowest-lying 3S_1 states of superheavy quarkonium is greatly simplified, since a dominant part of it is described by an essentially Coulomb-like interaction potential¹¹⁻¹³⁾ arising from the exchange of massless vector gluons. Although logarithmic and linear confining terms are

necessary in the potential for describing^{12,14)} the known $q\bar{q}$ systems at large interquark distances, for quark masses of the order of 20 GeV or more, considerations about the ground-state Bohr radius show¹¹⁻¹³⁾ that these contributions can be regarded as just a small perturbation of the Coulomb part.

For superheavy quarkonium systems the Coulomb potential should actually govern the bound-state dynamics, thus allowing the use of the hydrogen-like wave functions for computing the leptonic widths of the lowest orthostate as a function of mass. The electromagnetic one-photon decay of the 3S_1 state (Fig. 1a) is assumed to be described by the QCD extension^{15,16)} of the Van Royen-Weisskopf formula¹⁷⁾, in terms of the bound-state wave function at the origin,

$$\Gamma_e(^3S_1) = \Gamma_{0e} f_R(\alpha_s) = 16\pi \alpha^2 e_q^2 \frac{|\psi(0)|^2}{m_V^2} f_R(\alpha_s), \quad (1)$$

with

$$\Gamma_{0e} = \frac{1}{2} e_q^2 \alpha^2 \left(\frac{4}{3} \alpha_s \right)^3 \frac{m_V}{2 - \left(\frac{2}{3} \alpha_s \right)^2}; \quad f_R(\alpha_s) = 1 - \frac{16\alpha_s}{3\pi},$$

where e_q is the charge of the constituent quark, m_V the mass of the quarkonium state and

$$\alpha_s(m_V^2) = \frac{12\pi}{(33 - 2N) \ln \left(\frac{m_V^2}{\Lambda^2} \right)}, \quad (2)$$

with⁷⁾ $\Lambda = 0.5$ GeV and N the number of quark flavours.

For a large quarkonium mass the leptonic width is also contributed by Z^0 exchange^{13,18,19)} (see Fig. 1b). In the standard model²⁰⁾

$$\Gamma_{Ze} = \left(\frac{G_F}{8\sqrt{2}\pi\alpha} \right)^2 \frac{[1 + (1 - 4 \sin^2 \theta_w)^2] [1 - 4 e_q \sin^2 \theta_w]^2}{e_q^2} \frac{m_V^4 m_Z^4}{(m_V^2 - m_Z^2)^2 + m_Z^2 \Gamma_Z^2} \Gamma_{0e}. \quad (3)$$

The mass dependence of the total electronic widths (weak plus electromagnetic) is shown in Fig. 2 for $\sin^2 \theta_w = 0.23$ and both values of the quark charge. Interference effects, which might slightly modify the results at the Z^0 peak, have been neglected.

It is worth noting that in this energy region the scaling quantity Γ_{0e}/e_q^2 for the Coulomb potential is almost equal to the measured values of the known vector mesons and very weakly energy dependent. This makes our considerations rather insensitive to the detailed assumptions on the form of the quark-antiquark potentials¹²⁾.

We proceed by evaluating, as a function of the vector meson mass m_V , the partial widths of the main decay channels, as shown in Fig. 1. The strong and electromagnetic decays of a heavy quarkonium yield^{6,15)}

$$\Gamma_q = 3 \sum_{j=1}^N e_{qj}^2 \left(1 + \frac{\alpha_s}{\pi}\right) \Gamma_{0e} \quad (4)$$

$$\Gamma_{3g} = \frac{5}{18} \frac{\pi^2 - 9}{\pi} \frac{\alpha_s^3}{\alpha^2} \frac{4}{9e_q^2} \Gamma_{0e} \quad (5)$$

$$\Gamma_{\gamma g g} = \frac{36}{5} e_q^2 \frac{\alpha}{\alpha_s} \Gamma_{3g} \quad (6)$$

The weak decays which proceed via Z^0 exchange^{16,18,21)} yield

$$\Gamma_{Zq} = N_G \frac{[1 + (1 - \frac{8}{3} \sin^2 \theta_w)^2] + [1 + (1 - \frac{4}{3} \sin^2 \theta_w)^2]}{1 + (1 - 4 \sin^2 \theta_w)^2} \Gamma_{Ze} \quad (7)$$

$$\Gamma_{\nu\bar{\nu}} = \left(\frac{G_F}{8\pi\alpha}\right)^2 N_V \frac{(1 - 4 e_q \sin^2 \theta_w)^2}{e_q^2} \frac{m_V^4 m_Z^4}{(m_V^2 - m_Z^2)^2 + m_Z^2 \Gamma_Z^2} \Gamma_{0e} \quad (8)$$

with N_G the number of quark generations and N_V the number of neutrinos.

The weak quarkonium decays can also proceed through the weak decay of one or both constituent quarks^{18,19,22)*} (Figs. 1g and 1h):

*) It has been observed²³⁾ that, for $m_q = O(m_W)$, Eq. (9) has to be modified to include the effect of the pole in the W propagator. Eventually for $m_q > m_W$ the dominant decay is the production of a real W , in which case

$$\Gamma_{W1} = \frac{\sum_q |\epsilon_q|^2 G_F^2 m_W^2}{2\pi\sqrt{2}} \frac{(m_q^2 - m_W^2)^2}{m_q^3} \left[2 + \left(\frac{m_q}{m_W}\right)^2\right],$$

where ϵ_q are angle suppression factors. This leads to a rapidly increasing rate of bound quark decay. This may only strengthen our conclusions on the relevance of weak quarkonium decays at very high energies.

$$\Gamma_{W1} = \frac{8N_G G_F^2 m_q^5}{192 \pi^3} \left(1 + \frac{m_q^2}{m_W^2} \right)^{-2}, \quad (9)$$

or, via t-channel W exchange,

$$\Gamma_{W2} = \frac{1}{2} \frac{m_V^4}{(m_W^2 + \frac{1}{4} m_V^2)^2} \frac{\Gamma_{0e}}{e_q^2}. \quad (10)$$

The close connection between the Higgs coupling strength and fermion masses suggests that heavy quarks may represent a practicable source of these scalars, if their mass is not too large²⁴⁾. The process in Fig. 1f has a partial width^{24,25)}:

$$\Gamma_{\gamma H} = \frac{G_F m_V^2}{4\sqrt{2}\pi\alpha} \left(1 - \frac{m_H^2}{m_V^2} \right) \left[1 + \frac{2m_V\Delta}{m_V^2 - m_H^2} \right]^{-2} \Gamma_{0e} \quad (11)$$

$$\Delta \approx 1 \text{ GeV}.$$

Other processes, involving Higgs bremsstrahlung from virtual Z^0 and vector mesons, yield very small contributions²⁶⁻³⁰⁾.

The branching fractions of the processes of Figs. 1d to 1h are shown in Fig. 3, normalized to the leptonic width Γ_{0e} in the one-photon channel (1).

The weak current couplings become clearly dominant over the electromagnetic and strong decays for $m_V \gtrsim m_W$.

The experimentally observed cross-section at the quarkonium peak can now be obtained, taking into account the radiative corrections and the effect of the machine energy resolution³¹⁾:

$$\begin{aligned} \tilde{\sigma}_V(m_V) = & \frac{6\pi^2(\Gamma_e + \Gamma_{Ze})\Gamma_h}{\sqrt{2\pi} \sigma_W m_V^2 \Gamma_V} \left(\frac{\Gamma_V}{m_V} \right)^{\beta_i} \exp \left(\frac{\Gamma_V}{2\sqrt{2}\sigma_W} \right)^2 \times \\ & \times \left[\operatorname{erfc} \left(\frac{\Gamma_V}{2\sqrt{2}\sigma_W} \right) + \frac{\beta_i}{2} E_1 \left(\frac{\Gamma_V^2}{8\sigma_W^2} \right) \right] (1 + C_h), \end{aligned} \quad (12)$$

where Γ_h and Γ_V are the hadronic and total widths, σ_W is the standard deviation of the total c.m. energy of the machine³²⁾, C_h the ultraviolet correction $O(10^{-1})$, and

$$\beta_i = \frac{4\alpha}{\pi} \left(\ln \frac{m_V}{m_e} - \frac{1}{2} \right). \quad (13)$$

With the LEP project parameters³³⁾ the peak quarkonium cross-sections are shown in Fig. 4. An enhancement of about two orders of magnitude is contributed by Z^0 exchange, corresponding to production rates $O(10^4)$ events per day. A similar enhancement is, though, present in the background cross-section, with a large contribution of the Z^0 pole and its radiative tail above the continuum cross-section. With radiative corrections included^{31,34)}, and without considering the heavy lepton contribution, the hadronic background is given by

$$\tilde{\sigma}_h = \tilde{\sigma}_{\text{cont}} + \tilde{\sigma}_{Zh} \quad (14)$$

$$\tilde{\sigma}_{\text{cont}} = \frac{4\pi\alpha^2}{m_V^2} \sum_{j=1}^N e_{qj}^2 \left(1 + \frac{\alpha_s}{\pi} \right) \frac{1}{\gamma^{\beta_i} \Gamma(1 + \beta_i)} \left(\frac{2\Delta\omega}{m_V} \right)^{\beta_i} (1 + C_h)$$

$$\tilde{\sigma}_{Zh} = C_i \sigma(Z^0 \rightarrow \text{hadrons}) (1 + C_Z),$$

where γ is Euler's constant, $(2\Delta\omega/m_V)$ is the fractional energy resolution of the experimental apparatus, C_h and C_Z the usual ultraviolet corrections³⁴⁾ $O(10^{-1})$, and

$$C_i = \left| \frac{\left(\frac{2\Delta\omega}{m_V} \right)}{1 + \left(\frac{2\Delta\omega}{m_V} \right) \frac{m_V^2}{m_Z \Gamma_Z} e^{i\delta} \sin \delta} \right|^{\beta_i} (1 - \beta_i \delta \cot \delta) \quad (15)$$

$$\tan \delta = \frac{m_Z \Gamma_Z}{m_Z^2 - m_V^2}.$$

Also the quantity $\tilde{\sigma}_h$ is averaged over the machine energy resolution. The signal-to-background ratio $\eta_V = \tilde{\sigma}_V(m_V)/\tilde{\sigma}_h(m_V)$ from Eqs. (12) and (14) is shown in Fig. 5a as a function of quarkonium mass for hadronic final states. The same ratio for

$\mu\bar{\mu}$ final states is shown in Fig. 5b. Some structure is generated by the Z^0 and its radiative tail. In comparing η_V for the two quark charges it is interesting to note how their relative ratio, due to electromagnetic effects, becomes modified and actually inverted in the Z^0 region by the onset of weak effects.

The values of the expected signal-to-background ratios given in Fig. 5 show that the excitation of heavy quarkonium states becomes rapidly unobservable above the Z^0 mass. More selective identification methods are therefore required in this energy region.

Among the topological variables describing jet production^{7,35-37)} we consider both sphericity and acoplanarity:

$$S = \left(\frac{4}{\pi}\right)^2 \min \left[\frac{\sum_i |p_{Ti}|}{\sum_i |p_i|} \right]^2 \quad (16)$$

$$A = 4 \min \left[\frac{\sum_i |p_{out\ i}|}{\sum_i |p_i|} \right]^2 \quad (17)$$

These quantities have different average values for different final states in quarkonium decay. In particular, for the diagrams of Figs. 1c, d, and h we have⁷⁾

$$\langle A \rangle \sim \frac{16}{\pi^2} \langle n_c(W) \rangle^2 \frac{\langle p_T^2 \rangle}{m_V^2} \quad (18)$$

as for the non-perturbative background^{*)}, and

$$\langle S \rangle \sim \frac{64}{\pi^2} \frac{\alpha_s}{3\pi} \left(64 \ln \frac{3}{2} - \frac{229}{9} \right). \quad (19)$$

In Eq. (17) the produced quarks are assumed to hadronize with an exponentially damped p_T distribution and an average multiplicity $\langle n_c \rangle$ rising as $\sim \ln^2 s$. The three-gluon decays of Fig. 1e have⁷⁾ $\langle A \rangle \sim O(\alpha_s/\pi)$ and $\langle S \rangle \approx 0.134$. The weak

*) This behaviour dominates over the perturbative QCD contribution $O(\alpha_s/\pi)^2$ up to c.m. energies $O(1000)$ GeV.

decay mechanism of Fig. 1g, on the other hand, is likely to give rise to nearly spherical events, because of the large number of quarks in the final state, with $\langle A \rangle \approx \langle S \rangle \approx 1$. The values of sphericity and acoplanarity, averaged over the different quarkonium decay channels, are shown in Fig. 6 as a function of mass. The strong rise of both quantities above the Z^0 region, owing to the dominant weak decay mechanism of Fig. 1g, allows an improvement in the signal-to-background ratio at high masses by simple cuts on sphericity or acoplanarity. In order to get an indication about the experimental possibilities of detecting heavy vector states with a LEP-type machine³³⁾, we show in Fig. 7 the running time needed for a 5 standard deviation significance as a function of quarkonium mass. It is clear that, above the Z^0 mass, heavy vector states can be reasonably detected only if suitable cuts are applied.

3. THRESHOLD BEHAVIOUR

The most effective way of identifying new flavour thresholds is given by the study of hadron jet cross-sections in the continuum region. It is known, in fact, that some topological variables³⁵⁻³⁷⁾ describing jet production are more sensitive to the excitation of a new quark than the corresponding cross-section step.

In this section we discuss the high-energy threshold behaviour of the sphericity in the presence of weak and radiative effects. Similarly to what was discussed in the previous section for the quark-antiquark bound states, we show that an important effect on the average value of the sphericity is contributed by the presence of the Z^0 and its radiative tail.

In order to study the behaviour of the average sphericity when crossing a new flavour threshold, we assume⁷⁾ for the background from the lower-lying flavours $\langle S \rangle_{\text{QCD}}$ the expression of Eq. (19). The step in R associated with the new flavour contributes above the threshold s_{th} with an average sphericity⁷⁾ decreasing as $\langle S \rangle_{\text{step}} \approx (s_{\text{th}}/s)$. The over-all value of the sphericity will be the average of these two contributions weighted on the corresponding background and step cross-sections. For the latter, one has²⁾ including Z^0 exchange,

$$\sigma_{\text{step}}(s) = \frac{1}{2} \beta(3 - \beta^2) \left[\frac{4\pi\alpha^2 e_q^2}{s} + \sigma(Z^0 \rightarrow q\bar{q}) \right] \left[1 + \frac{4}{3} \alpha_s f(\beta) \right], \quad (20)$$

where

$$f(\beta) = \frac{\pi}{2\beta} - \frac{3 + \beta}{4} \left(\frac{\pi}{2} - \frac{3}{4\pi} \right), \quad \beta = \left(1 - \frac{s_{\text{th}}}{s} \right)^{1/2}.$$

and $\sigma(Z^0 \rightarrow q\bar{q})$ refers to the specific new flavour being excited. This description, based on the Born approximation for free $q\bar{q}$ production, gives by duality arguments³⁸⁻⁴⁰⁾ the average over the resonant structures around threshold.

The radiative corrections³¹⁾ to σ_{step}

$$\tilde{\sigma}_{\text{step}}(W) = \beta_i \int_0^{\Delta\omega} \left(\frac{2\Delta\omega}{W} \right)^{\beta_i} \tilde{\sigma}_{\text{step}}(W - \omega) \left(1 - \frac{2\omega}{W} + \frac{2\omega^2}{W^2} \right) \frac{d\omega}{\omega} (1 + C_h) \quad (21)$$

turn out to have a small effect on the electromagnetic contribution to the sphericity, whereas there is a much larger effect on the weak contribution owing to the radiative tail^{34,41)} of the Z^0 . Following this discussion we evaluate the over-all average sphericity as a function of c.m. energy W as:

$$\langle S \rangle = \left[\left(\frac{W}{W_{\text{th}}} \right)^2 \tilde{\sigma}_{\text{step}}(W) + \langle S \rangle_{\text{QCD}} \tilde{\sigma}_h(W) \right] \left[\tilde{\sigma}_{\text{step}}(W) + \tilde{\sigma}_h(W) \right]^{-1}, \quad (W > W_{\text{th}}), \quad (22)$$

where the background contribution $\tilde{\sigma}_h(W)$ is given by (14).

The energy dependence of the average sphericity is shown in Fig. 8 for various threshold positions. In the Z^0 region the smooth decay of the average sphericity above threshold is modulated by the interplay of the excitation and background cross-sections. This effect is much more pronounced for the case $e_q = 1/3$ (Fig. 8a), which for some threshold positions displays a wide sphericity plateau.

The signal-to-background ratio in the mean sphericity at the maximum is shown in Fig. 9 for both quark charges. It indicates that measurable effects seem to be present throughout the whole energy range. This situation can be considerably improved with an appropriate cut in the differential sphericity distribution, which at the considered energies is almost totally confined⁷⁾ below $S = 0.5$ for

background events. With this cut and with a LEP-type machine³³⁾, the running time for a 5 standard deviation significance is given in Fig. 10, assuming the sphericity is measured just above threshold where the shoulder in R appears.

4. CONCLUSIONS

In this paper we discussed what the spectacular quarkonium peaking and the associated continuum threshold, already observed for light quarks, may look like in the case of superheavy quark production in e^+e^- annihilation. For masses of the quark-antiquark bound states not too different from those of the weak bosons, decays mediated by Z^0 and $W^\pm$ play a very important role and strong peakings occur at the Z^0 position. Above the Z^0 region the weak amplitudes of quasi-free quark decay contribute the dominant decay channels with a rising energy dependence. These decay mechanisms are characterized by high quark multiplicities and large sphericity, thus allowing the detection of quarkonium states also in the mass region above the Z^0 , where the peak cross-sections become unobservably small.

This high-energy behaviour of the sphericity at the quarkonium peak is of dynamical origin, and therefore physically very different from the sphericity step at the continuum thresholds of essentially kinematical nature. Also for the case of continuum thresholds we have shown that considerable distortion is introduced by the weak current terms. In the energy region dominated by electromagnetic interactions, the most efficient detection scheme⁴²⁾ of new flavours consists of a search for threshold effects in appropriate topological variables, followed by a fine scan in search of the associated bound states. We conclude that for e^+e^- production of superheavy quarks, in the energy region where weak current couplings become dominant, this strategy remains practically applicable.

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Figure captions

- Fig. 1 : Weak, electromagnetic, and strong contributions to heavy quarkonium decay.
- Fig. 2 : Mass dependence of the total leptonic width of 3S_1 quarkonium states for both values of quark charge e_q . The dotted lines show the one-photon electromagnetic contribution for the Coulomb-like $q\bar{q}$ potential.
- Fig. 3 : Mass dependence of the main strong and weak quarkonium widths normalized to the one-photon leptonic width.
- Fig. 4 : Peak cross-sections for weak and electromagnetic quarkonium production. Machine resolution effects and radiative corrections are included in the calculation. The fractional energy resolution of the apparatus is assumed to be 5×10^{-2} .
- Fig. 5 : a) Signal-to-background ratio at the quarkonium peak as a function of mass, for hadronic final states. Both quark charges are shown.
b) Same as in (a) for $\mu\bar{\mu}$ final states.
- Fig. 6 : a) Average sphericity of quarkonium decay as a function of mass. Average sphericity of the continuum is also shown.
b) Same as in (a) for the average acoplanarity.
- Fig. 7 : Running time at the nominal LEP luminosity to detect a quarkonium ground state with a 5 standard deviation significance as a function of mass. Both quark charges are shown. The dashed lines indicate the effect of a sphericity cut ($S \geq 0.5$) for high mass values.
- Fig. 8 : a) Average sphericity as a function of c.m. energy for different threshold positions, for $e_q = 1/3$. The background sphericity $\langle S \rangle_{QCD}$ is indicated by the dotted line.
b) Same as in (a), for $e_q = 2/3$.
- Fig. 9 : Signal-to-background ratio at the peak of the average sphericity as a function of threshold energy.

Fig. 10 : Running time at the nominal LEP luminosity to detect high-sphericity events ($S \geq 0.5$) with a 5 standard deviation significance as a function of threshold energy. The measurement is assumed just above threshold, at the shoulder in R.

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