Tests of substructure of heavy quarks

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We discuss the observability of flavor-changing transitions among heavy quarks in e^+e^- collisions: anomalous thresholds $e^+e^- \to Q+\bar{Q}'$, anomalous flavored-quarkonia formation $e^+e^- \to V(Q\bar{Q}')$, and anomalous decay modes of heavy quarkonia $(Q\bar{Q}) \to Q+\bar{Q}'$, $Q+\bar{Q}'+\gamma, \ldots$, etc. We express cross sections and partial widths in terms of a substructure scale and show what kind of limits can be obtained.

I. INTRODUCTION

The possibility of a composite structure for leptons, quarks, and weak bosons is a field of investigations which was recently intensively developed. For review and references see Ref. 1. In this paper we propose further tests of quark substructure.

Experimental limits on lepton (ν,e,μ,τ) substructure mainly come[2,3] from anomalous-magnetic-moment measurements $(g-2)$, $e^+e^- \to l^+l^-$, absence of excited lepton states, deep-inelastic lepton-nucleon scattering, and flavor-changing processes $\mu \to e\gamma$ and $\mu \to 3e$. They do not leave much room for observable effects at present energies as the typical mass scale (Λ) of the substructure lies high in the TeV range. In the case of quarks the situation is different. Anomalous magnetic moments and form-factor effects in $e^+e^- \to$hadrons, deep-inelastic scattering, and search for excited quark states have also been discussed[4] but the limits on Λ only lie in the 100-GeV range. The rare decay modes of K mesons and the K_L-K_S mass difference can be used[5] to also give Λ values in the TeV range but these features may be peculiar of the pseudoscalar-K-meson system. Here we are faced with a typical aspect of the subject. The absence of a signal in a given process can always be attributed to a particular dynamical reason. Cancellations may hide the existing substructures. For example, effects may appear much earlier in K^* vector states than in K pseudoscalar states. They may be functions of m/Λ where m is a quark mass such that the effects which disappear for light quarks will become sizable for heavy quarks. In addition high generations may present features of excited states with larger extensions (i.e., smaller values of Λ). The appreciable mixing angles between generations can be taken as an indication for quark substructure. In any case, because of the absence of a unique and firm theory of fermion substructure it is necessary to look for anomalous effects in various processes even if the result is negative in some of them.

With this motivation we looked for possible anomalous effects of heavy quarks. Magnetic-moment and form-factor effects in $e^+e^- \to$heavy quarks$\to$hadrons and in deep-inelastic scattering may be difficult to isolate because flavor identification is not easy. So it is not obvious that the limits on Λ can be substantially improved this way. Hence we concentrate on flavor-changing transitions among heavy quarks. We assume that compositeness (or any underlying new interaction) leads to effective $\gamma QQ'$,gQQ',ZQQ' flavor-changing couplings (see Sec. II). In Sec. III we start by computing the cross section of $e^+e^- \to \gamma,Z \to Q+Q'$ and discuss its particular aspects (threshold effects and vector-meson enhancements). We give the rates of the remarkable direct $D^{*0},B^{*0},T^{*0},\ldots$ formations. In Sec. IV we then consider peculiar decay modes of heavy quarkonia $(\psi,\Upsilon,\theta,\ldots)$ induced by these flavor-changing couplings: $(Q\bar{Q}) \to Q+\bar{Q}'$ (for example $\psi \to u\bar{c}$, $\Upsilon \to d\bar{b}$, $\theta \to u\bar{t},\ldots$) and $(Q\bar{Q}) \to Q+\bar{Q}'+\gamma$ or $(Q\bar{Q}) \to Q+\bar{Q}'+$gluon. In each case we express cross sections and partial widths in terms of the flavor-changing parameters. We show how measurements can give new limits on the substructure scale Λ. The radiative decays seem to be particularly promising.

II. FLAVOR-CHANGING COUPLINGS

We use the effective $\gamma QQ'$ couplings:

$$L=\frac{e}{2\Lambda}\bar{\psi}_{Q'}\sigma^{\mu\nu}F_{\mu\nu}(a^\gamma - ib^\gamma\gamma^5)\psi_Q + \text{H.c.} \tag{1}$$

Λ is a scale parameter introduced for convenience. The dimensionless parameters are a and b. C invariance would require both to be real and CP invariance would require a to be real and b to be purely imaginary. Here we do not want to rely on a given composite model. a and b could result from constituent-bound-state wave-functions overlap of Q and of Q' or from a new-heavy-boson exchange between different-flavor fermions (horizontal gauge bosons, Higgs bosons, etc.). We just notice that they may be functions of m_Q/Λ and $m_{Q'}/\Lambda$ and this would cut off their value in the case of light quarks. For timelike photons a and b can contain s-dependent form factors which can locally enhance or depress their value according to the substructure dynamics. In the case of gluon-QQ' couplings we just replace e by the QCD coupling constant g_s and a^γ,b^γ by a^g,b^g times $\lambda/2$ color matrices. For simplicity we shall also use the same form for ZQQ' couplings with a^γ,b^γ replaced by a^Z,b^Z although nonconserved currents could also be induced.

Nonzero values for a/Λ and b/Λ are at least expected from high-order electroweak effects[6,7] (diagrams with internal $W^\pm$ or Higgs-boson lines and quark lines). The order of magnitude is $(a/\Lambda,b/\Lambda) \simeq \alpha\Delta/M_W{}^2$. When no

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Glashow-Iliopoulos-Maiani (GIM) mechanism occurs (for example, when $m_Q \geq M_W$) Δ is proportional to quark-mass differences and mixing angles. This already gives a very low value $(a/\Lambda, b/\Lambda) \leq 1/10^3$ TeV for quark masses in the GeV range. If a GIM mechanism (cancellation between two generations of quarks) operates Δ is reduced by a factor $M_{\text{quark}}{}^2/M_W{}^2$ and this makes effects unobservable in the present range of energies.

III. APPLICATION TO $e^+e^- \to \bar{Q}+Q'$ PROCESSES

We now consider the reactions $e^+e^- \to \gamma, Z \to \bar{Q}+Q'$ $(e^+e^- \to \bar{d}+s, \bar{u}+c, \bar{d}+b, \bar{s}+b, \bar{u}+t, \bar{c}+t, \ldots)$ and their conjugate states $Q+\bar{Q}'$. Couplings of Eq. (1) lead to the differential cross section

$$\frac{d\sigma}{d\Omega} = \frac{6\alpha^2 p\sqrt{s}}{\Lambda^2}\{G_1[2EE' - \tfrac{1}{2}(s - m^2 - m'^2) - 2p^2\sin^2\theta] + G_3 mm' - G_{12}(E'-E)p\cos\theta\} . \tag{2}$$

m and m' are quark masses, $\sqrt{s}$ is the total e^+e^- c.m. energy,

$$E = \frac{s+m^2-m'^2}{2\sqrt{s}} , \quad E' = \frac{s+m'^2-m^2}{2\sqrt{s}} , \quad p = (E^2 - m^2)^{1/2} ,$$

$$G_{1,3} = \frac{1}{s^2}(|a^\gamma|^2 \pm |b^\gamma|^2) - \frac{2}{s}\text{Re}\frac{a_e^{Z*}(a^\gamma a^{Z*} \pm b^\gamma b^{Z*})}{D_Z^*} + \frac{(|a_e^Z|^2 + |b_e^Z|^2)(|a^Z|^2 \pm |b^Z|^2)}{|D_Z|^2} ,$$

$$G_{12} = \frac{2}{s}\text{Im}\frac{b_e^Z(a^{\gamma*}b^Z - a^Z b^{\gamma*})}{D_Z} + \frac{4\,\text{Re}(a_e^Z b_e^{Z*})\text{Im}(a^Z b^{Z*})}{|D_Z|^2} , \quad D_Z = s - M_Z{}^2 + iM_Z\Gamma_Z .$$

a_e^Z and b_e^Z are the standard vector and axial-vector Ze^+e^- couplings; see Ref. 8 for notations.

The cross section for $e^+e^- \to Q+\bar{Q}'$ is obtained by the replacement $a^{\gamma,Z} \to a^{*\gamma,Z}$ and $b^{\gamma,Z} \to b^{*\gamma,Z}$. The integrated cross section is

$$\sigma = \frac{24\pi\alpha^2 p\sqrt{s}}{\Lambda^2}\{G_1[2EE' - \tfrac{1}{2}(s-m^2-m'^2) - \tfrac{4}{3}p^2] + G_3 mm'\} . \tag{3}$$

The forward-backward asymmetry is

$$A_{\text{FB}} = -\frac{\tfrac{1}{2}G_{12}(E'-E)p}{G_1[2EE' - \tfrac{1}{2}(s-m^2-m'^2 - \tfrac{4}{3}p^2)] + G_3 mm'} . \tag{4}$$

The $Z \to \bar{Q}+Q'$ partial width is given by

$$\Gamma_{Z\to\bar{Q}Q'} = \frac{6\alpha p}{\Lambda^2}\{mm'(|a^Z|^2 - |b^Z|^2) + (|a^Z|^2 + |b^Z|^2)\tfrac{1}{3}[4EE' - \tfrac{1}{2}(s-m^2-m'^2)]\} . \tag{5}$$

In Fig. 1 we show the energy behavior of $\sigma(e^+e^- \to \bar{Q}+Q'$ and $Q+\bar{Q}')$ as given by (3) assuming constant values for a and b (no form factor). In this case the cross section would asymptotically tend to a constant. The threshold behavior for free quarks goes like p for a couplings and like p^3 for b couplings. However, strong interactions between low energy $\bar{Q}$ and Q' [for example, $(\bar{Q}Q')$ bound states discussed below] can locally modify this behavior.

It would be interesting to look for these new types of events in $e^+e^- \to$hadrons just above $\sqrt{s} = m_Q + m_{Q'}$ thresholds, for example, $m_u + m_c \simeq 1.8$ GeV, $m_d + m_b \simeq 5.4$ GeV, $m_s + m_b \simeq 5.7$ GeV, and $m_u + m_t, m_c + m_t$ above 20 GeV. These new thresholds will lead to usual effects (new event shapes, increase of sphericity, increase of leptonic rates, etc.) but should differ by asymmetrical flavor properties of the events. The $\bar{d}b$ and $\bar{s}b$ thresholds lie just above the rich $c\bar{c}$ region. The $\bar{u}t$ and $\bar{c}t$ thresholds would be a rather genuine way of finding the top quark.

Absence of a signal at a given level of cross section (for example, $\Delta R = \frac{1}{10}$) can be used to set limits on $(a/\Lambda, b/\Lambda)$ parameters. Assuming $m_u + m_t < 30$ GeV a run at $\sqrt{s} = 30$ GeV would give $(a^\gamma/\Lambda, b^\gamma/\Lambda) < 1/(100$ GeV). Once the t-quark mass will be known a longer run above the threshold could largely improve these limits. A future run at the Z peak excluding a branching ratio larger than 10^{-4} would give the limits $(a^Z/\Lambda, b^Z/\Lambda) < 1/(4$ TeV).

Around the thresholds one can look for enhancements

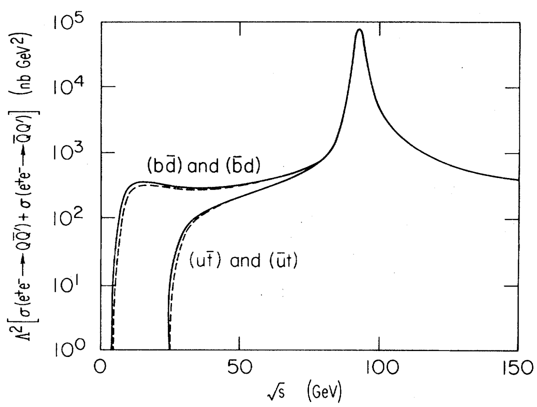

FIG. 1. Flavor-changing cross section $\sigma(e^+e^- \to Q\bar{Q}') + \sigma(e^+e^- \to \bar{Q}Q')$ scaled by factor Λ^2 in the cases $Q=b$, $Q'=d$ and $Q=t$, $Q'=u$. Solid curve $a^\gamma = a^Z = 1$, $b^\gamma = b^Z = 0$. Dashed curve, $a^\gamma = a^Z = 0$, $b^\gamma = b^Z = 1$.

due to $(\bar{Q}Q')$ and $(Q\bar{Q}')$ vector (or axial-vector) quarkonia formation. This means $e^+e^- \to \gamma, Z \to K^{*0}, D^{*0}, B^{*0}, T^{*0}, \ldots$ and their conjugate states. Except in the K^{*0} case the low-lying vector mesons are expected to be rather narrow (in the MeV range) so we discuss the magnitude of the integrated spectra:

$$\bar{\sigma} = \int d\sqrt{s} = \frac{6\pi}{M_V{}^2}\Gamma_{V\to e^+e^-} \qquad (6)$$

with[8]

$$\Gamma_{V\to e^+e^-} = \frac{4\pi\alpha^2}{3M_V{}^3}\left[|g_{V\gamma}|^2 + \frac{M_V{}^4}{|D_Z|^2}(|a_e{}^Z|^2 + |b_e{}^Z|^2)|g_{VZ}|^2 - 2M_V{}^2 \mathrm{Re}\left[\frac{a_e^Z g_{V\gamma}^* g_{VZ}}{D_Z}\right]\right] \qquad (7)$$

and

$$\Gamma_{A\to e^+e^-} = \frac{4\pi\alpha^2}{3}M_A\left[\frac{|a_e^Z|^2 + |b_e^Z|^2}{|D_Z|^2}\right]|g_{AZ}|^2 \qquad (8)$$

in the case of axial-vector mesons. In the nonrelativistic approximation (this may seem very crude for light-quark—heavy-quark bound states but it works reasonably well in several cases[7,9]) one gets

$$g_{V\gamma} = -\frac{2a^\gamma M_V\sqrt{M_V}}{\Lambda}\left[\frac{\phi(0)\sqrt{3}}{\sqrt{4\pi}}\right] \qquad (9)$$

(and similarly for g_{VZ}), and

$$g_{AZ} = \frac{ib^Z 3\sqrt{3}M_A\sqrt{M_A}}{\sqrt{2}mm'\Lambda}\left[\frac{\phi'(0)\sqrt{3}}{\sqrt{4\pi}}\right]. \qquad (10)$$

Values of wave functions and their derivatives at the origin can be taken from analysis of well-known pseudoscalar, vector, and axial-vector mesons.[7,9] This gives

$$\left[\frac{\phi(0)\sqrt{3}}{\sqrt{4\pi}}\right] \simeq 0.11 M_V \text{ GeV}^{3/2} \qquad (11)$$

and

$$\left|\frac{g_{AZ}}{M_A b^Z}\right| \simeq \left|\frac{g_{VZ}}{M_V a^Z}\right|. \qquad (12)$$

Applied to K^{*0}, D^{*0}, B^{*0}, and T^{*0} (assuming $M_V = 25$ GeV in the last case) one gets

$$\Gamma_{V\to e^+e^-} \simeq \left[\frac{a^\gamma}{\Lambda}\right]^2 \times (8 \text{ keV}, 30 \text{ KeV}, 250 \text{ KeV}, 6 \text{ MeV}),$$

where Λ is in GeV. Quick runs in the 20—30-GeV energy range excluding $\Gamma_{T^{*0}\to e^+e^-}$ larger than 0.1 keV already give the interesting limit

$$\left[\frac{a^\gamma}{\Lambda}\right] < \frac{1}{250 \text{ GeV}}.$$

Longer runs at the right energy position excluding $\Gamma_{V\to e^+e^-}$ larger than 1 eV (this may be possible in the low-energy range) would give the limit

$$\left[\frac{a^\gamma}{\Lambda}\right] < \left[\frac{1}{100 \text{ GeV}}, \frac{1}{200 \text{ GeV}}, \frac{1}{500 \text{ GeV}}, \frac{1}{2.5 \text{ TeV}}\right].$$

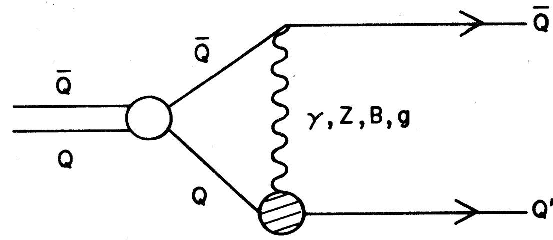

FIG. 2. γ, Z, g, and scalar-boson B exchange diagrams for flavor-changing decay of a heavy $(Q\bar{Q})$ quarkonium.

IV. APPLICATION TO HEAVY-QUARKONIA DECAYS

We now look for special decay modes of ordinary $(Q\bar{Q})$ quarkonia. The first one is $V(Q\bar{Q}) \to Q + \bar{Q}'$ or $\bar{Q} + Q'$. Examples of such modes are $\phi \to \bar{d} + s$ (i.e., $K+\pi$, $K+2\pi, \ldots$), $\psi \to \bar{u} + c$ (i.e., $D+\pi, \ldots$), $\Upsilon \to \bar{d} + b$, $\theta \to \bar{u} + t, \ldots$, etc. Final states (two-body, multibody, and jets) are similar to the ones quoted above for $e^+e^- \to \bar{Q} + Q'$. The process can go through an internal g, γ, or Z exchange where one effective flavor-changing coupling occurs (Fig. 2). An internal flavor changing due to the subquark dynamics may also directly occur (for example, through subconstituent rearrangement or exchange of a new kind of boson). The general form of the effective vertex $V(Q\bar{Q}) \to \bar{Q} + Q'$ is given by

$$R_{fi} = \bar{u}_{Q'}(k_1)[V(A - B\gamma^5) + V\cdot(k_1 - k_2)(C - iD\gamma^5)]v_Q(k_2) \qquad (13)$$

and the partial width by

$$\Gamma_{V\to\bar{Q}Q'} = \frac{k}{6\pi m_V 2}\{|A|^2[3(k_1\cdot k_2 + mm') - 2\vec{k}^2] + |B|^2[3(k_1\cdot k_2 - mm') - 2\vec{k}^2] + 4|C|^2\vec{k}^2(k_1\cdot k_2 - mm')$$
$$+ 4|D|^2|\vec{k}|^2(k_1\cdot k_2 + mm') - 4\vec{k}^2(m+m')(\mathrm{Re}AC^* + \mathrm{Im}BD^*)\}. \qquad (14)$$

The γ- or Z-exchange diagram would give

$$A = -\frac{e^2 a a_0(m' - m)\sqrt{M_V}}{2\Lambda D}\left[\frac{\phi(0)}{\sqrt{4\pi}}\right], \quad B = \frac{ie^2 b a_0(m + m')\sqrt{M_V}}{2\Lambda D}\left[\frac{\phi(0)}{\sqrt{4\pi}}\right],$$
$$C = -\frac{e^2(a a_0 - 2ibb_0)\sqrt{M_V}}{2\Lambda D}\left[\frac{\phi(0)}{\sqrt{4\pi}}\right], \quad D = -\frac{e^2(b a_0 + 2iab_0)\sqrt{M_V}}{2\Lambda D}\left[\frac{\phi(0)}{\sqrt{4\pi}}\right], \qquad (15)$$

where D is the value of the internal propagator [i.e., $(m'^2-m^2)/2$ or $(m'^2-m^2)/2-M_Z^2$ for γ or Z in the weak-binding limit] and a_0, b_0 are the standard vector and axial-vector γQQ or ZQQ couplings; k is the final c.m. momentum.

Adding $V \to \bar{Q}Q'$ and $V \to Q\bar{Q}'$ one gets the partial width

$$\Gamma=\frac{4\pi\alpha^2 k}{9M_V\,|D|^2\Lambda^2}\left[\frac{\phi(0)\sqrt{3}}{\sqrt{4\pi}}\right]^2 \times\{[3(k_1\cdot k_2+mm')-2\vec{k}^2](m-m')^2|a|^2|a_0|^2+[3(k_1\cdot k_2-mm')-2\vec{k}^2]|a_0|^2|b|^2(m+m')^2$$
$$+4\vec{k}^2(k_1\cdot k_2-mm')(|a_0|^2|a|^2+4|b_0|^2|b|^2)+4\vec{k}^2(k_1\cdot k_2+mm')(|a_0|^2|b|^2+4|b_0|^2|a|^2)\}\,. \tag{16}$$

In the case of gluon exchange one has to replace α^2 by $\frac{16}{9}\alpha_s^2$ and use $a_0=1$, $b_0=0$. A scalar-boson exchange (in place of γ or Z) of mass M_B and coupling constants ef for BBQ and ef' for BBQ' would give

$$A=\frac{e^2 ff'(3m+m')\sqrt{M_V}}{4mD}\left[\frac{\phi(0)}{\sqrt{4\pi}}\right],$$
$$C=-\frac{e^2 ff'\sqrt{M_V}}{D}\left[\frac{\phi(0)}{\sqrt{4\pi}}\right] \tag{17}$$

with $D=(m'^2-m^2)/2-M_B^2$ in the weak-binding limit.

Using again the empirical formula for $\phi(0)/\sqrt{4\pi}$ one gets in the case of single-γ exchange

$$\Gamma\simeq 4\times10^{-7}\left[\frac{a^\gamma}{\Lambda}\right]^2,\; 4\times10^{-5}\left[\frac{a^\gamma}{\Lambda}\right]^2,$$
$$7\times10^{-5}\left[\frac{a^\gamma}{\Lambda}\right]^2,\; 8\times10^{-3}\left[\frac{a^\gamma}{\Lambda}\right]^2$$

in GeV for ϕ, ψ,Υ, and θ(50 GeV), respectively. A limit of 10^{-3} on the ψ, Υ, and θ branching ratios for these decays would produce the bounds

$$\left[\frac{a^\gamma}{\Lambda}\right]<\frac{1}{30\text{ GeV}},\;\frac{1}{50\text{ GeV}},\;\frac{1}{700\text{ GeV}}\,.$$

In the case of internal gluon exchange the replacement of α^2 by $\frac{16}{9}\alpha_s^2$ will produce bounds on gluon flavor-changing couplings approximately 30 times large (using $\alpha_s\simeq0.15$) in the heavy-quark-mass range

$$\left[\frac{a^g}{\Lambda}\right]<\frac{1}{1.5\text{ TeV}},\;\frac{1}{2\text{ TeV}},\;\frac{1}{20\text{ TeV}}\,.$$

In the case of scalar-boson exchange the bounds are roughly obtained by the replacement

$$\left[\frac{a^\gamma a_0^\gamma}{2m\Lambda}\right]\to\frac{ff'}{m_B^2}\,.$$

The same kind of effects could induce an effective $(Q\bar{Q}')-Q-\bar{Q}$ vertex, for example, in the case of pseudoscalar $(Q\bar{Q}')$ states one gets contributions to the nonleptonic flavor-changing decays $K^0\to$pions, $D^0\to$pions, $B^0\to$pions and kaons, etc. These new contributions will compete with the ordinary weak nonleptonic decays due to internal $W^\pm$ exchange, for example, the so-called penguin diagrams for K^0 and D^0 decays.[7] Maybe they could be at the origin of the well-known problems for D^0 branching ratios and lifetime. As already noted in the Introduction, this will happen if Λ is of the order of 10^3 TeV. This is, however, very difficult to quantify more precisely because several types of other weak diagrams contribute to these decays at this level.[7]

Another very interesting process is the radiative mode $V(Q\bar{Q})\to\bar{Q}+Q'+\gamma$ (and its conjugate) described in Fig. 3 where the flavor changing occurs at the real photon vertex. With respect to the preceding case one gains one electromagnetic order, avoids possible form factors, and gets events easier to trigger with the hard photon. The final quarks $\bar{Q}$ and Q' (or Q and $\bar{Q}'$) can either (a) form jets or (b) bound into $\bar{Q}Q'$ (or $Q\bar{Q}'$) mesons like in the case of ordinary $V\to P+\gamma$ transitions.

In case (a) the shape of the photon spectrum is governed by the momentum dependence of the $V(Q\bar{Q})$ bound-state wave function

$$\frac{d\Gamma}{dq}=2\alpha(|a^\gamma|^2+|b^\gamma|^2)\frac{mq}{\Lambda^2}\times\int\frac{dk}{k^0}k\,|\tilde{\phi}(k)|^2\left[\frac{m_V^2+m^2-m'^2}{2}-m_V k^0\right]. \tag{18}$$

It peaks towards the value $q_0=(m^2-m'^2)/2m$ corresponding to Q and $\bar{Q}$ at rest in the $V(Q\bar{Q})$ rest system (weak-binding limit) and Q decaying into $Q'+\gamma$ (see Fig. 4).

The partial width $\Gamma_{V\to\bar{Q}Q'\gamma}$ is approximately equal to

$$\Gamma_{Q\to Q'\gamma}=\frac{\alpha(|a^\gamma|^2+|b^\gamma|^2)(m^2-m'^2)^3}{2m^3\Lambda^2}\,. \tag{19}$$

Assuming again that one can measure the modes $\phi\to\bar{d}+s+\gamma$ (i.e., $K+\pi+\gamma,\ldots$), $\psi\to\bar{u}+c+\gamma$ (i.e., $D+\pi+\gamma,\ldots$), $\Upsilon\to\bar{d}+b+\gamma$, and $\theta\to\bar{u}+t+\gamma$ down to a

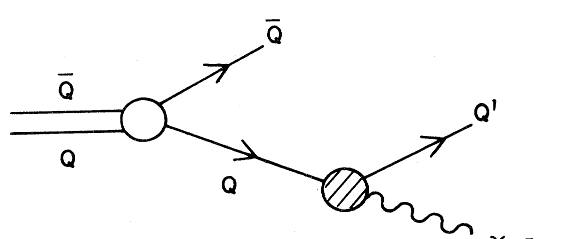

FIG. 3. Flavor-changing radiative decay of a heavy $(Q\bar{Q})$ quarkonium.

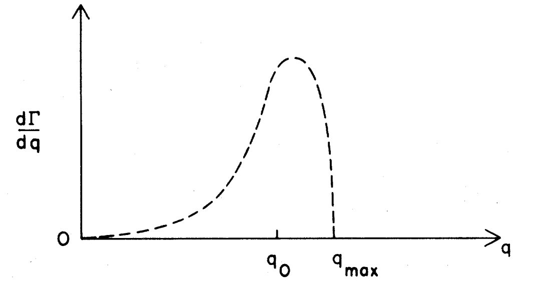

FIG. 4. Shape of photon spectrum in $(Q\bar{Q})\rightarrow\bar{Q}+Q'+\gamma$ decay. $q_0=(m^2-m'^2)/2m$, $q_{\max}=[M_V{}^2-(m+m')^2]/2M_V$.

10^{-3} branching ratio and using, respectively,

$$\Gamma=\left[\frac{|a^\gamma|^2+|b^\gamma|^2}{\Lambda^2}\right](4\times10^{-4},10^{-2},0.5,50)$$

(with Λ and Γ in GeV), one gets the very interesting limits

$$\frac{(|a^\gamma|^2+|b^\gamma|^2)^{1/2}}{\Lambda} < \frac{1}{10\text{ GeV}},\ \frac{1}{300\text{ GeV}},\ \frac{1}{3\text{ TeV}},\ \frac{1}{30\text{ TeV}}\ .$$

Here also one can replace the photon by a gluon and look for $V(Q\bar{Q})\rightarrow\bar{Q}+Q'+g$. Equations (18) and (19) should be modified with α replaced by $\frac{4}{3}\alpha_s$ and (a^γ,b^γ) replaced by (a^g,b^g). The increase in the rates will, however, surely be canceled by the difficulties in triggering on such events. In the low-energy range (ψ) these events will be mixed with the $(\bar{Q}+Q')$ ones studied above but appear to be more frequent by one α_s power.

In case (b) one can write the $V(Q\bar{Q})\rightarrow P(\bar{Q}Q')+\gamma$ partial width as

$$\Gamma_{V\rightarrow P\gamma}=\frac{q^3}{12\pi}(|g_{VP\gamma}|^2+|g'_{VP\gamma}|^2) \tag{20}$$

with

$$g_{VP\gamma}=\frac{iea^\gamma}{\Lambda}\left[\frac{2(m_V{}^2+m_P{}^2)}{m_V{}^2}\right]^{1/2}I_{VP} \tag{21}$$

(and $g'_{VP\gamma}$ in terms of b^γ). I_{VP} is the wave-function overlap, in the weak-binding limit,

$$I_{VP}=\frac{1}{4\pi}\int d_3k\,\tilde{\phi}_V(k)\tilde{\phi}^*_P(k)\ . \tag{22}$$

Its value is very model dependent; because of the large mass difference between $V(Q\bar{Q})$ and $P(Q\bar{Q}')$ I_{VP} is probably much less than one.

If the experimental accuracy for $\phi\rightarrow K^0\gamma$, $\psi\rightarrow D^0\gamma$, $\Upsilon\rightarrow B^0\gamma$, and $\theta\rightarrow T^0\gamma$ is also of the order of 10^{-3} in the branching ratio one will get approximately the same limits as above for the quantity $[(|a^\gamma|^2+|b^\gamma|^2)^{1/2}/\Lambda]I_{VP}$. Other contributions to $V(Q\bar{Q})\rightarrow P(\bar{Q}Q')+\gamma$ can occur in which the flavor changing does not appear at the photon vertex. For example, it can appear in an effective $V(Q\bar{Q})-\bar{Q}-Q'$ or $P(\bar{Q}Q')-Q-\bar{Q}$ vertex as discussed above (Fig. 5); if this last flavor changing is due to an internal gluon exchange one can just replace in Eq. (21) a^γ by $\alpha_s a^g$ in order to get the corresponding order of magnitude. All these uncertainties in the description of $\Gamma_{V\rightarrow P\gamma}$ will strongly affect the conclusions for Λ in absence of signal. In this case only a true signal would be a fruitful result.

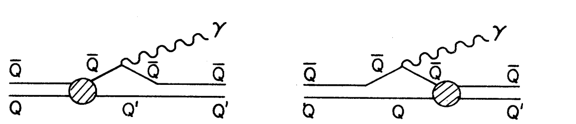

FIG. 5. Flavor-changing radiative decay of a heavy $(Q\bar{Q})$ quarkonium (other contributions).

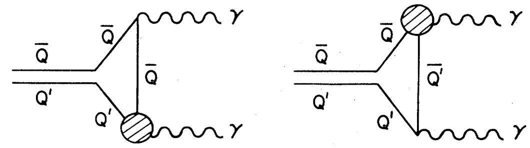

FIG. 6. Diagrams for flavor-changing $\gamma\gamma$ decays.

The real photon can also be replaced by a lepton pair $[V(Q\bar{Q})\rightarrow\bar{Q}+Q'+l^+l^-]$. In the case of very massive quarkonia (for example, a high-mass θ t-quarkonium) this process can be due to Z emission and decay. If the Z is itself composite the flavor-changing couplings ZQQ' may be larger than the $\gamma QQ'$ ones and this may compensate the additional α factor (or $B_{Z\rightarrow e^+e^-}$ branching ratio) controlling the rate. We have given elsewhere[10] other tests of Z compositeness.

Similar studies can be done for $\gamma\gamma$ collisions (for example, $\gamma\gamma\rightarrow Q+\bar{Q}'$ with jets or bound states) and deep-inelastic processes (anomalous heavy-quark production in neutral currents). However, for our purpose the experimental accuracy will probably not compete with that of direct e^+e^- collisions. Let us just estimate the rate of $P(Q\bar{Q}')\rightarrow\gamma\gamma$ using the diagrams of Fig. 6:

$$\Gamma_{P\rightarrow\gamma\gamma}\simeq12\alpha^2e_q{}^2\left[\frac{a^\gamma}{\Lambda}\right]^2|\phi(0)|^2\ . \tag{23}$$

With Eq. (11) one gets

$$\Gamma_{P\rightarrow\gamma\gamma}\simeq3\times10^{-5}e_q{}^2M_P{}^2\left[\frac{a^\gamma}{\Lambda}\right]^2\text{ GeV}\ .$$

Notice first that in the case of the K^0 meson a^γ/Λ could just compete with the effective value of the electroweak high-order effects quoted at the beginning (i.e. $\lesssim(1/10^5$ GeV in this particular case). In the case of heavy mesons taking $\Gamma_{P\rightarrow\gamma\gamma}\simeq1$ keV as a limit of detectability in $\gamma\gamma$ collisions one can hardly expect to get a bound better than 1/(100 GeV).

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