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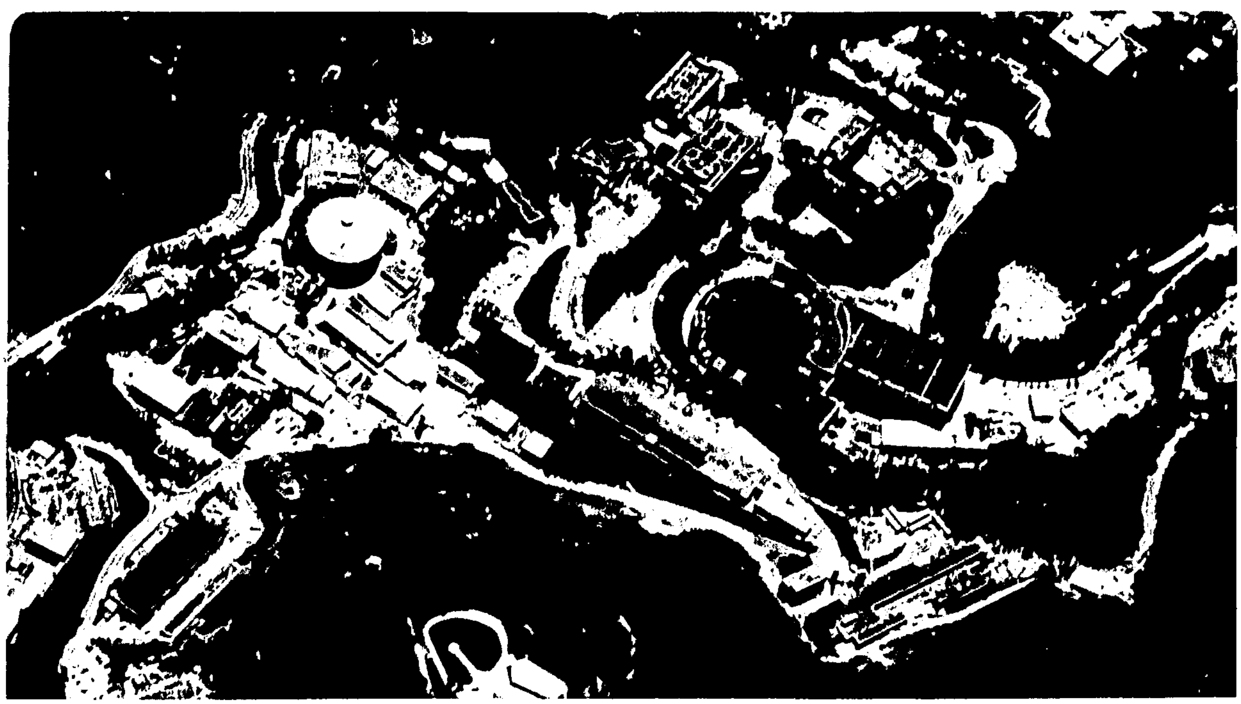

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Is $\zeta(8.3)$ a gluino-gluino bound state ?

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Abstract

We consider gluino-gluino ($\tilde{g}\tilde{g}$) interpretation of $\zeta(8.3)$. If ζ is a 1P, $\tilde{g}\tilde{g}$-state $B(T \to \gamma\zeta)$, $B(T' \to \gamma\zeta)$, $\Gamma(\zeta \to \text{all})$ are consistent with what is observed within theoretical ambiguity. We show $B(T \to \gamma\zeta) : B(T' \to \gamma\zeta) : B(T'' \to \gamma\zeta) = 1 : \sim 0 : 2$.

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The Cristal Ball group[1,2] has recently shown evidence for a state, which is called ζ, in the radiative decay of Υ.

$$
\begin{aligned}
&M = 8322 \pm 8 \pm 24 \text{ MeV}^{1},\ 8319 \pm 8 \pm 21 \text{ MeV}^{2}, \\
&\Gamma < 80 \text{ MeV} \quad (90\% \text{ C.L.}) \\
&B(\Upsilon(1S) \to \gamma\zeta) \sim 0.5\%^{1} \\
&\qquad\qquad\qquad = (0.47 \pm 0.11 \pm 0.26)\%^{2} \\
&\frac{B(\Upsilon(2S) \to \gamma\zeta)}{B(\Upsilon(1S) \to \gamma\zeta)} < 0.22 \quad (90\% \text{ C.L.})^{1}
\end{aligned} \tag{1}
$$

The CUSB group did not see this peak but obtained branching ratio $B(\Upsilon(1S) \to \gamma\zeta) < 0.22\%$ (90% C.L.)[2] is not inconsistent with above value.

If ζ is interpreted as the neutral Higgs expected in the standard model, the theoretical branching ratio $B(\Upsilon(1S) \to \gamma\zeta)$ becomes smaller than the experimental one by a factor 100. The predicted ratio $B(\Upsilon(2S) \to \gamma\zeta)/B(\Upsilon(1S) \to \gamma\zeta) \sim 1$ is not consistent with the experimental value either.

Supersymmetric particles are other candidates. Supersymmetry[3] is an only way of combining bosons and fermions in 4 dimension of space time and is expected to give a geometrical understanding of the massless fermion excitations. Supersymmetry also seems to play a rather fundamental role in solving several long-standing hierarchy problem, e.g., gauge hierarchy and strong CP hierarchy. If so, all the "known" particles should have superpartners. The gluino is the superpartner of the gluon.

Recently Barger et al[4] have proposed a supersymmetry scenario with a light gluino ($m(\tilde{g}) \lesssim 5$ GeV) in order to accomodate the observed UA1 missing-P_T events[5] (as for other scenarios see refs. 6 and 7). In this scenario the mass of the $\tilde{g}\tilde{g}$ state is likely to be below $\Upsilon(9.46)$.

One year ago it has been proposed[8] to look for $\tilde{g}\tilde{g}$-states, which is called glueballon, in the radiative decay of the heavy quarkonium states. Properties of $\tilde{g}\tilde{g}$ states were studied[8-10] by assuming $V_{\tilde{g}\tilde{g}}(r) = (9/4)V_{Q\bar{Q}}(r)$, where $V_{Q\bar{Q}}(r)$ is the quarkonium potential. So it is tempting to look for the possibility of $\zeta(8.3)$ being one of the $\tilde{g}\tilde{g}$ states.

In the following we consider two models, <u>Model 1</u>: ζ is a 1S, $\tilde{g}\tilde{g}$ state($\eta_{\tilde{g}}$), <u>Model 2</u>: ζ is a 1P, $\tilde{g}\tilde{g}$ state ($\chi_{\tilde{g}}$). We first use the Richardson potential[11] and compute $\tilde{g}\tilde{g}$ spectra in both models. Results are shown in Fig.1. The gluino mass is fixed to adjust $\tilde{g}\tilde{g}$,1S(1P) state mass to $\zeta(8.3)$ in model 1(model 2). We use $m(\tilde{g})$ = 4.91 GeV in model 1 and 4.282 GeV in model 2. Gluino masses are model dependent. For example, if the Kühn Ono potential(potential T in ref. 10) is used we need $m(\tilde{g})$ = 5.705GeV in model 1 and 5.19 GeV in model 2. The Richardson potential is more singular near origin than Kühn Ono potential. The main difference between two spectra in both models is 1P-1S level spacing i.e., E(1P) - E(1S) = 1165 MeV in Richardson model and 950 MeV in Kühn Ono model. Other higher levels above 1P-state is very similar in both models. Let us start with the first model.

(i) Model 1 ($\zeta = \eta_{\tilde{g}}(1S)$)

We first use the Richardson potential. ζ mainly decays into two gluons. The width is around 230 MeV (α_S = 0.157 is used, see ref. 10). This is couple of times larger than the experimental limit (< 80 MeV). We can use the following theoretical ambiguity to reduce the decay width. (i) Soften the short range potential singularity to reduce $\psi(0)$, (ii) reduce α_S , (iii) the higer order QCD corrections might change the width. However, if we assume Υ goes with single photon directly to $\eta_{\tilde{g}}$ through two gluons

(Fig.2) the obtained branching ratio $B(T \to \gamma\zeta) \sim 0.03\%$ is smaller than the experimental value by around factor ten. If $\Gamma(\eta_{\tilde{g}} \to gg)$ decreases, the already too small branching ratio $B(T(1S) \to \gamma\eta_{\tilde{g}})$ becomes even smaller.

In the present paper we compute the branching ratio $B(T \to \gamma\zeta)$ in another model. Since ζ happens to be rather near T the mixing between $\eta_{\tilde{g}}$ and η_b (Fig.3ab) might play an important role in the transition. Let us assume the following mixing

$$\zeta = \sqrt{1 - \sum_i \varepsilon_i^2}\, \eta_{\tilde{g}} + \varepsilon_{1S}\eta_b + \varepsilon_{2S}\eta_b' + \varepsilon_{3S}\eta_b'' \tag{2}$$

The mixing between $\eta_{\tilde{g}}$ and η_b (or η_b' , η_b'') can be explained by the two gluon intermediate state (Fig.3a) since the interaction is concentrated at short distance and the QCD perturbation is applicable. We estimate the imaginary part of the nondiagonal component of mass matrix

$$\varepsilon_{1S} = \frac{1}{2}\cdot\frac{1}{\Delta M}\left[\Gamma(\eta_b \to gg)\Gamma(\eta_{\tilde{g}} \to gg)\right]^{1/2} , \tag{3}$$

where ΔM is the mass difference $\Delta M = m(\eta_b) - m(\eta_{\tilde{g}})$ and other mixing parameters are given in the same way.

Decay widths of $\tilde{g}\tilde{g}$ states are given by[10]

$$\begin{aligned}
\Gamma(\eta_{\tilde{g}} \to gg) &= \frac{27}{4}\cdot\frac{8}{3}\cdot\alpha_s^2\,|\psi(0)|^2/m(\eta_{\tilde{g}})^2 \\
\Gamma(\chi_{\tilde{g}}(^3P_2) \to gg) &= \frac{27}{4}\cdot\frac{128}{5}\cdot\alpha_s^2\,|\psi'(0)|^2/m(\chi_{\tilde{g}})^4 \\
\Gamma(^3P_0) : \Gamma(^3P_1) &: \Gamma(^3P_2) = 15 : \sim 1.5 : 4
\end{aligned} \tag{4}$$

By using the theoretical estimate $\Gamma(\eta_b \to gg) = 5$ MeV ($\alpha_s(b\bar{b}) = 0.157$ is assumed) and the experimental upper limit $\Gamma(\eta_{\tilde{g}} \to gg) = 80$ MeV one obtains

$$B(\Upsilon(1S) \to \gamma\zeta) = 4.05\,eV/44.3\,keV = 9.1 \times 10^{-5}$$
$$B(\Upsilon(2S) \to \gamma\zeta) = 2.18\,eV/29.6\,keV = 7.4 \times 10^{-5} \quad (5)$$
$$B(\Upsilon(3S) \to \gamma\zeta) = 1.03\,eV/17.7\,keV = 5.8 \times 10^{-5}$$

where mixing parameters are $\varepsilon_{1S} = 0.00890$, $\varepsilon_{2S} = 0.00386$ and $\varepsilon_{3S} = 0.00273$.

The ratio $B(T(2S) \to \gamma\zeta)/B(T(1S) \to \gamma\zeta) \sim 0.8$, which does not depend on α_S, is larger than the experimental value <0.22. It was pointed out[12] that $|\psi(0)|$ for η_c determined from $\Gamma(\eta_c \to gg)$ is substantially larger than that for $^3S_1$ determined from $\Gamma(J/\psi \to e^+e^-)$. This discrepancy can be explained[12] if the spin spin force is taken into account. One might think $B(T(2S) \to \gamma\zeta)/B(T(1S) \to \gamma\zeta)$ will decrease if the spin-spin force is included in the potential because due to the force $\Gamma(\eta_b \to gg)$ increases more than $\Gamma(\eta_b' \to gg)$. However, the explicit calculation in this model[12] shows there is very little effect due to the force to this ratio. The reason is that the inclusion of this force breaks the outhogonality between 1S and 2S states and the transition matrix $\langle nS \mid j_0(kr/2) \mid mS \rangle$ changes much. Such effect improves[13] the agreement between theoretical and experimental decay rates in many radiative transitions. In the $b\bar{b} \to \gamma\zeta$ transition two effects (changes in $\psi_{2S}(0)/\psi_{1S}(0)$ and transition matrix) nearly cancell each other, so $B(T(2S) \to \gamma\zeta)/B(T(1S) \to \gamma\zeta)$ remains unchanged in the end. Thus we do not know the way how to improve the theoretical prediction for this ratio. This conclusion is true for any model in which ζ is assumed to be mixture of η_b, η_b', η_b'' and a certain 0^- particle (e. g., pseudoscalar Higgs, see ref.14).

One might already be able to rule out all such models including model 1 if one takes the discrepancy in

$B(T' \to \gamma\zeta)/B(T \to \gamma\zeta)$ seriously. A caveat[15] against this conclusion is that we cannot understand the small ratio $B(\psi' \to \gamma\eta')/B(\psi \to \gamma\eta')$ and $B(\psi' \to \gamma\iota)/B(\psi \to \gamma\iota)$ either.

The predicted branching ratio from imaginary part $B(T \to \gamma\zeta) \sim 9 \times 10^{-5}$ is substantially smaller than the experimental value $\sim 0.5\%$. To look at the real part of the mixing between two systems ($b\bar{b}$ and $\tilde{g}\tilde{g}$), we compute the box diagram (Fig. 3a) for the 0^{-+} channel. In the limit of the total energy $\sqrt{s} \approx 2\, m_b,\ 2\, m_{\tilde{g}}$, we get

$$\mathcal{E}_R = \frac{2}{\pi}(1 - \ln 2)\, \mathcal{E}_I \approx 0.195\, \mathcal{E}_I\ , \qquad (6)$$

where the imaginary part ε_I is given by (3). From eq.6 we conclude that the imaginary part dominates over the real part in the mixing via two gluons for the 0^{-+} channel. The similar conclusion is found by Körner, Kühn, Krammer and Schneider[16]. Thus due to the inclusion of the real part $B(T \to \gamma\zeta)$ increases only 4% which is far from enough.

In principle we can increase $\Gamma(\eta_b \to gg)$ if we assume that α_S for $\tilde{g}\tilde{g}$ is different from α_S for $b\bar{b}$ or that $V_{\tilde{g}\tilde{g}}(r)$ is not equal to $(9/4)V_{Q\bar{Q}}(r)$. If one uses $\Gamma(\eta_b \to gg) \simeq 50$ MeV one finds $B(T \to \gamma\zeta) \sim 10^{-3}$ which is around the experimental lower limit of this branching ratio. However, $\Gamma(\eta_b \to gg) \sim 50$ MeV might be unacceptably large.

Instead of studying this model further we consider the second possibility.

(ii) Model 2 ($\zeta = \chi_{\tilde{g}}(1P)$)

There are three $\tilde{g}\tilde{g}$, 1P-states ($^3P_{0,1,2}$). Since $^3P_1$ does not couple to two gluons the partial decay rate into $^3P_1$, $\tilde{g}\tilde{g}$ state from T is small and can be neglected. Thus ζ should be either $\chi_{\tilde{g}}(^3P_0)$ or $\chi_{\tilde{g}}(^3P_2)$. Through P wave $\tilde{g}\tilde{g}$ - $b\bar{b}$ mixing one finds

$$B(\Upsilon \to \chi_{\tilde{g}}(^3P_2)+\gamma)/B(\Upsilon \to \chi_{\tilde{g}}(^3P_0)+\gamma) = 5 \times \frac{4}{15} \times (\text{phase space}) \tag{7}$$

Thus both branching ratios are of comparable order. If the fine structure of $\chi_{\tilde{g}}(^3P_J)$ states is given simply by changing the color factor from $\chi_b(^3P_2) - \chi_b(^3P_0)$ one finds $\chi_{\tilde{g}}(^3P_2) - \chi_{\tilde{g}}(^3P_0)$ $\sim (9/4) \times 40$ MeV ~ 90 MeV. Since this is larger than the upper limit of the decay width 80 MeV one must have already seen two peaks while only one peak has been observed experimentally. However, the theoretical calculation of the spin-orbit splitting is full of ambiguities. It depends not only on the precise shape of the potential but also on the Lorentz structure (scalar or forth component of the vector). In some cases the contribution from the long range part and short range part cancell each other and the spin orbit force becomes very small. We here simply assume that the mass splitting $\chi_{\tilde{g}}(^3P_2) - \chi_{\tilde{g}}(^3P_0)$ is much smaller than 80 MeV and both peaks are observed unresolved.

In model 2 we assume $\zeta = \chi_{\tilde{g}}(1P)$ or more precisely

$$\zeta = \sqrt{1 - \sum_i \varepsilon_i^2}\, \chi_{\tilde{g}}(1P) + \varepsilon_{1P}\chi_b + \varepsilon_{2P}\chi_b' + \varepsilon_{3P}\chi_b'' \tag{8}$$

From imaginary part one finds

$$\varepsilon_{1P} = \frac{1}{2}\,\frac{1}{\Delta M}\left[\Gamma(\chi_b \to gg)\,\Gamma(\chi_{\tilde{g}} \to gg)\right]^{1/2} \tag{9}$$

After taking into account the size effect of $b\bar{b}$ the matrix element of the E1 transition from Υ to $b\bar{b}$ component in ζ is given by

$$R_{fi} = \sum_{f=1P,2P,3P} \varepsilon_f \langle f | j_1\left(\frac{kr}{2}\right) | i \rangle \quad (10)$$

We find the following branching ratios

$$B(\Upsilon \to \gamma\zeta) : B(\Upsilon' \to \gamma\zeta) : B(\Upsilon'' \to \gamma\zeta) \quad (11)$$

$= 1 : \sim 0 : 2.43$ for Richardson potential[10]

$= 1 : 0.031 : 2.08$ for Kuhn Ono potential[11]

here we use $(\varepsilon_{1P}, \varepsilon_{2P}, \varepsilon_{3P}) = (0.002651, 0.002125, 0.001858)$ and $\Gamma(\chi_b, \chi_b', \chi_b''(^3P_0) \to gg) = (852, 828, 807 \text{ keV})$ obtained by using Richardson potential with $\alpha_s = 0.25$. Above ratios does not depend on α_s. Similar numbers are found in Kühn Ono model.

Both models predict small $B(T' \to \gamma\zeta)$ because we find positive contribution to R_{fi} from $\varepsilon_{1P}\, \chi_b$ but negative ones from $\varepsilon_{2P}\chi_b'$ and $\varepsilon_{3P}\, \chi_b''$ and these terms cancell each other. There are no such cancellations for the processes $T \to \gamma\zeta$ and $T'' \to \gamma\zeta$. Thus in model 2 the absence if ζ peak in $T' \to \gamma\zeta$ can be understood naturally. Moreover, our model can clearly be checked by finding ζ in the process $T'' \to \gamma\zeta$ which has large branching ratio.

There are three further points which have to check in order to see if model 2 is consistent with experimental data. (i) $\Gamma(\zeta \to \text{all})$. (ii) $B(T \to \gamma\zeta)$ (iii) Why is only $\chi_{\tilde{g}}(1P)$ found ? The decay rates of various $\tilde{g}\tilde{g}$-states and partial decay widths $\Gamma(b\bar{b} \to \gamma + \tilde{g}\tilde{g})$ are computed by assuming $\alpha_s = 0.25$ and listed in table 1.

The real part of the mixing in the 0^{++} channel can be obtained in the same manner as the 0^{-+} channel and is given by

$$\mathcal{E}_R = \frac{1}{3\pi}\left[\frac{5}{9} + \frac{10}{3} \ln 2 + \frac{8}{3} \ln \frac{\sqrt{s}}{\lambda} \right] \mathcal{E}_I \qquad (12)$$

where ε_I is the imaginary part (eq. 9). The cut off parameter is introduced to avoid the singular integral at soft gluon and can be replaced by the binding energy of the systems. When we put $\sqrt{s} = m_\zeta = 8.3$ GeV and $\lambda = 300$ MeV into eq. 12, we get

$$\mathcal{E}_R = 1.244\ \mathcal{E}_I \qquad (13)$$

By taking the real part also into account, we find that the estimated branching ratio increases by the factor $(1 + \varepsilon_R^2/\varepsilon_I^2) \doteq 2.55$.

Thus the partial decay width becomes $\Gamma(T \to \gamma + \chi_{\tilde{g}}(^3P_{0,2}))$ = 2.55 × 4.20 = 10.7 eV. There are couple of factors which might increase this width. (i) The calculation of the real part of the mass matrix has ambiguity. (ii) The width $\Gamma(\chi_b(^3P_0) \to gg)$ = 852 keV which is used to compute the above width might be underestimated. Recent experimental data suggest $\Gamma(\chi_b(^3P_0) \to gg) \sim 1.5$ MeV. This can be adjusted e.g., by increasing $\alpha_s(b\bar{b})$ or $\psi'_{1P}(0)$. Most probably other decay rates e.g., $\Gamma(\chi_b(2P,3P) \to gg)$ also become large by the same amount. Thus $\Gamma(T \to \gamma + \chi_{\tilde{g}}(^3P_{0,2}))$ will increase by the factor 2 or 3. (iii) The $\tilde{g}\tilde{g}$ - $b\bar{b}$ mixing will considerably increase if the superpartner of the b quark is not too heavy (e.g., $m(\tilde{b}) \sim O(15$ GeV$)$).

Therefore, if we become optimistic, we can increase the branching ratio $B(T \to \gamma\zeta)$ up to around 40 eV i.e., ∿ 0.1 %, which is marginally consistent with the experimental datum.

The predicted branching ratio of $T \to \gamma + \eta_{\tilde{g}}(1S)$ is more than ten times larger than $T \to \gamma + \zeta$ in our model. However, the width of $\eta_{\tilde{g}}(1S)$ is as large as 550 MeV and $\eta_{\tilde{g}}(1S)$ is not easy to identify as a resonance. The branching ratio $\Upsilon \to \gamma + \eta_{\tilde{g}}(2S)$ is around the same order (note real part for the S-state is small) or even larger than that of $\Upsilon \to \gamma + \zeta$. $\eta_{\tilde{g}}(2S)$ is around 200 to 300 MeV heavier than ζ, and the width is twice as large as that of ζ. It will be a crucial test of this model to check the existence of such particle.

In summary we assume the ζ is a $\tilde{g}\tilde{g}$ state with small $b\bar{b}$ component. We obtains $B(\Upsilon \to \gamma\zeta) : B(\Upsilon' \to \gamma\zeta) : B(\Upsilon'' \to \gamma\zeta) = 1 : 0.8 : 0.6$ if ζ is a 1S, $\tilde{g}\tilde{g}$ state and $1 : \sim 0 : 2$ if ζ is a 1P, $\tilde{g}\tilde{g}$ state. Thus the latter model is preferred by the data, $B(\Upsilon' \to \gamma\zeta)/B(\Upsilon \to \gamma\zeta) < 0.22$. This conclusion is valid for any model which assume that the transition is mediated by small $b\bar{b}$ component inside ζ (Higgs is another example). If ζ is a 1P, $\tilde{g}\tilde{g}$ state it is possible to reproduce $\Gamma(\zeta \to \text{all})$, $B(\Upsilon \to \gamma\zeta)$, $B(\Upsilon' \to \gamma\zeta)$ within theoretical ambiguity. We have computed many quantitied which help to check the present model experimentally. For example, the processes $\Upsilon'' \to \gamma\zeta$ and $\Upsilon \to \gamma + \eta_{\tilde{g}}(2S)$ must be easy to observe, where $\eta_{\tilde{g}}(2S)$ will be heavier than ζ by 200 - 300 MeV. We expect that ζ consists of two resonances with $J^{PC} = 0^{++}$ and 2^{++} which are very close each other but have different decay widths.

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Table Caption

Table 1 (a) Two gluon decay width of $\tilde{g}\tilde{g}$-states. $\alpha_S = 0.25$ and Richardson potential is used to compute width.

(b) Partial decay width of radiative transition $b\bar{b} \to \gamma + \tilde{g}\tilde{g}$.

Table 1 (a)

$\tilde{g}\tilde{g}$-state	$\Gamma(\eta_{\tilde{g}} \to gg)$ (MeV)
1S	550
2S	153
3S	91
4S	66
5S	53
6S	45
7S	39
	$\Gamma(\chi_{\tilde{g}} \to gg)$
1^3P_2	21
1^3P_0	79
2^3P_2	14.5
2^3P_0	54.3
3^3P_2	11.3
3^3P_0	42.5

(b)

$b\bar{b}$	$\tilde{g}\tilde{g}$	$\Gamma(b\bar{b} \to \gamma + \tilde{g}\tilde{g})$ (eV)
$\Upsilon(1S)$	1S	105
	1P	4.2
	2S	16.7
	2P	0.47
	3S	2.52
$\Upsilon(2S)$	1S	41.2
	1P	5×10^{-7}
	2S	10.7
	2P	0.11
	3S	4.42
	3P	1.45
$\Upsilon(3S)$	1S	17.8
	1P	4.1
	2S	6.5
	2P	2.8
	3S	4.1
	3P	2.1

Figure captions

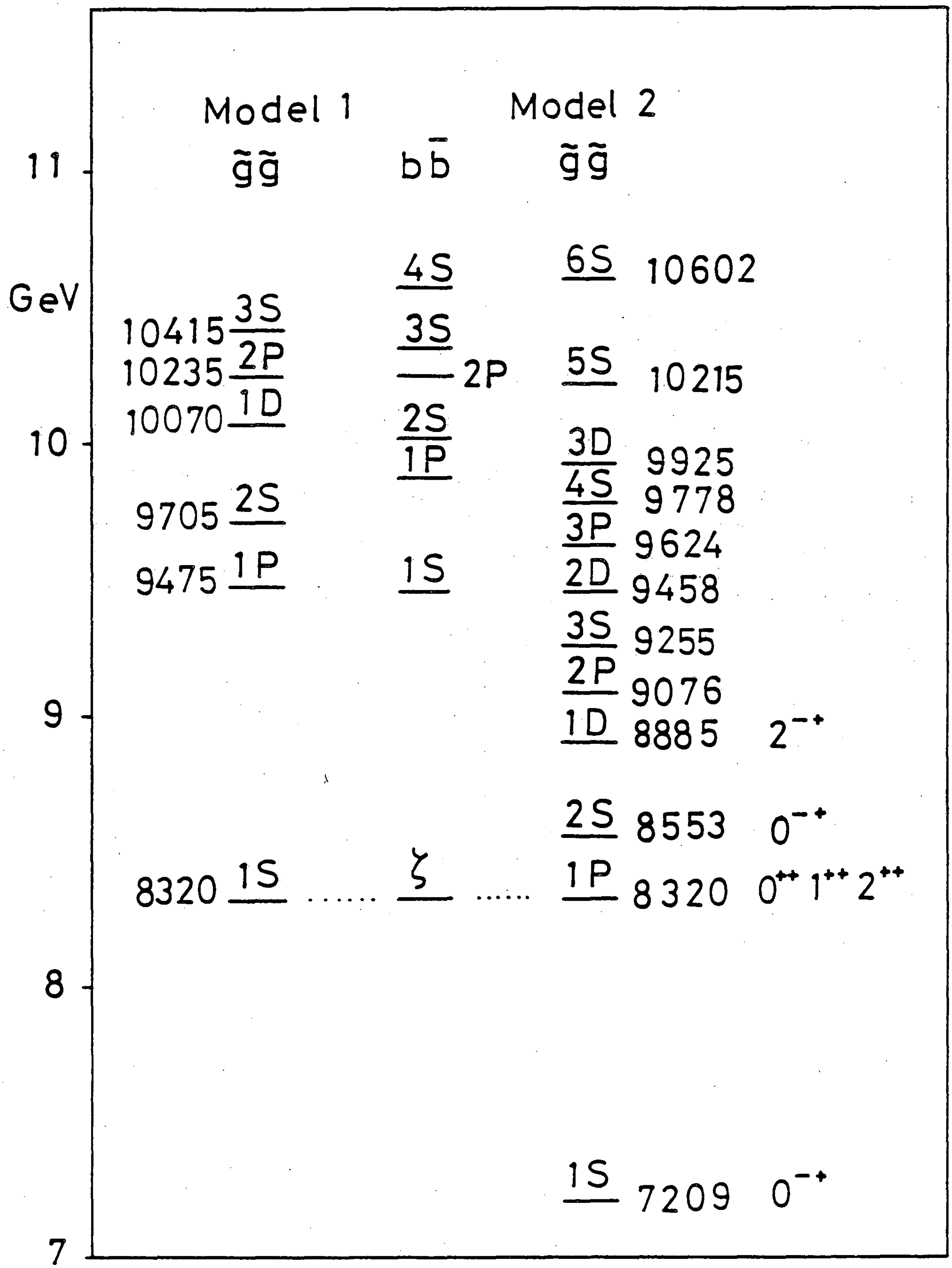
Fig. 1 The $\tilde{g}\tilde{g}$ spectra in model 1 ($\zeta = \eta_{\tilde{g}}(1S)$) and in model 2 ($\zeta = \chi_{\tilde{g}}(1P)$) are compared with $b\bar{b}$ spectra.

Fig. 2 The radiative decay of $b\bar{b}$ states into $\tilde{g}\tilde{g}$ states.

Fig. 3 (a) The $b\bar{b}$ - $\tilde{g}\tilde{g}$ mixing through two gluons.

(b) The $b\bar{b}$ - $\tilde{g}\tilde{g}$ mixing by a scalar quark exchange.

Fig.1



Model 1
Model 2
$\tilde{g}\tilde{g}$
$b\bar{b}$
$\tilde{g}\tilde{g}$
11
GeV
10
9
8
7
10415 3S
10235 2P
10070 1D
9705 2S
9475 1P
8320 1S
4S
3S
2P
2S
1P
1S
ζ
6S 10602
5S 10215
3D 9925
4S 9778
3P 9624
2D 9458
3S 9255
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1S 7209 0^{-+}

Fig.2

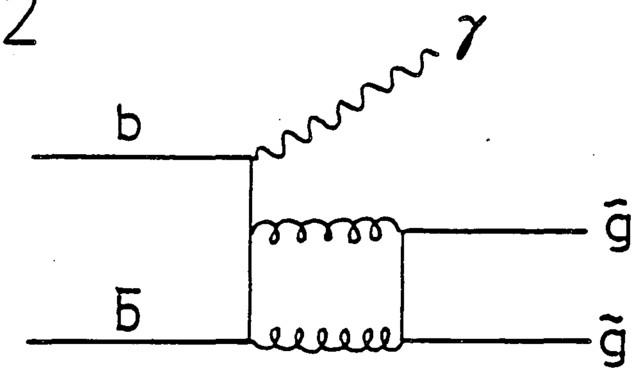

Fig.3a

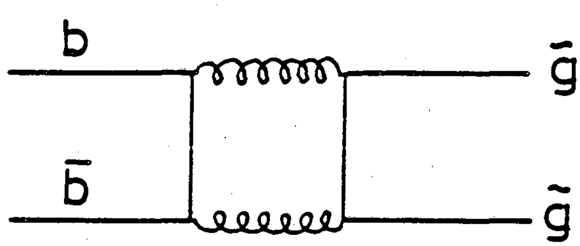

Fig.3b