ABSTRACT

Assuming that perturbative QCD is the dominant explanation for the narrowness of the vector quarkonia, we perform a $\chi^2$ minimization analysis of their hadronic decays as a function of two parameters, the mass of the gluino and the value of $\alpha_3(M_Z)$. A value below 1 GeV for the gluino mass is strongly preferred. Consequences for SUSY breaking scenarios are discussed.
Recently it has been pointed out that the quarkonia data can be made consistent with minimal supersymmetric grand unification if the gluino and photino masses are below half the $Z^0$ mass [1]. In addition, the quarkonia data becomes consistent with the world average measurements of the strong fine structure constant, $\alpha_3(M_Z)$, if these masses are not above the $\Upsilon$ region. The result of Ref. [1] was based on differences in the running of the strong coupling constant in the presence of gluinos. If, however the gluino mass is below half the $\Upsilon$ mass a second effect comes into play. Namely, the possibility for a quarkonium state to decay into gluino-containing final states affects the analysis of $\alpha_3$ at the relevant scale for that decay. In the current work we seek a best fit value to the gluino mass, $m_{\tilde{g}}$, taking both effects into account. This allows us to consider the full range of relevant gluino masses in a $\chi^2$ analysis and to determine the best fit value of the gluino mass $m_{\tilde{g}}$.

This work depends on the following basic assumption: \textit{The dominant explanation for the narrowness of the quarkonia states including the $J/\psi$ and $\phi$ is perturbative QCD.}

Any future models for the quarkonia data that rely on very large non-perturbative effects or relativistic corrections should be considered as alternative to the present analysis and judged on the basis of their relative physical plausibility. Small non-perturbative corrections, of course, would only modify our results by small amounts.

We analyze up to second order in QCD the hadronic decay rates of six quarkonia states, $\phi(1019)$, $J/\psi(3097)$, $J/\psi(3686)$, $\Upsilon(9460)$, $\Upsilon(10020)$, and $\Upsilon(10350)$. Ignoring possible higher order or non-perturbative effects and relativistic corrections each of these defines within errors the strong coupling constant at an appropriate scale $\mu_S(i)$

$$\alpha_3(\mu_S(i)) = \alpha_{3,i} \pm \delta_i, \quad (i = 1, \ldots, 6).$$

Assuming the gluino mass, $m_{\tilde{g}}$, lies above half the mass of the quarkonium state, these values are independent of the gluino mass and are given in Ref. [1] using the 1992 branching ratio averages of the Particle Data Group [2]. In the current work we consider also lower values of $m_{\tilde{g}}$ so that the $\alpha_{3,i}$ become functions of $m_{\tilde{g}}$.

If the gluino mass lies below half the quarkonium mass, there are decays into two gluinos ($\tilde{g}\tilde{g}$) and into two gluinos plus one gluon ($\tilde{g}\tilde{g}g$) which compete with the standard three gluon ($ggg$) decay. These gluino-containing final states [3] however are suppressed by four powers of $m_q/m_{\tilde{g}}$ and are negligible for currently allowed values of the squark mass $m_{\tilde{q}}$. The dominant gluino-related correction to the quarkonium width is therefore the two-gluon plus two-gluino ($gg\tilde{g}\tilde{g}$) final state with no intermediate squarks. The corresponding decay
rate as a function of the gluino to quark mass ratio $m_\tilde{g}/m_q$ was written down in Ref. [4] by applying a color+spin correction factor to the rate for the $ggq\bar{q}$ decay as calculated in Ref. [5]. The result is

$$\frac{\Gamma(3S_1(\bar{q}q) \to ggg\bar{g})}{\Gamma(3S_1(\bar{q}q) \to ggg)} = \frac{3\alpha_3(\mu_S, m_\tilde{g})}{\pi} R(r),$$  \hspace{1cm} (2)$$

where $r = 2m_\tilde{g}/m(3S_1) \simeq m_\tilde{g}/m_q$. For small $r$ (but $r > 0.1$),

$$R(r) = -\ln r + \frac{9}{32(\pi^2 - 9)} \left[ -3.56 - r^2(8\ln r + 0.95) + \frac{28}{27}\pi^2 r^3 
+ r^4\left(-\frac{16}{3}\ln^2 r + 1.1\ln r - 8.1\right) + \mathcal{O}(r^5) \right].$$  \hspace{1cm} (3)$$

Eq. (3) is infrared divergent as $r \to 0$. A correct treatment [3] shows that $R(0) \approx 1.57$.

We use the above expression for $R$ in region $r > 0.1$ and use a quadratic interpolating polynomial to join $R$ smoothly to its value at $r = 0$. The results presented here are not sensitive to the exact form of the cutoff. $R(r)$ falls rapidly for increasing $r$. For $r > 0.5$ it is necessary to use the exact integral expression for $R(r)$. It is possible [1] to choose the scale $\mu_S(i)$ so that the known (first order) standard model corrections to the 3 gluon decay rate vanish identically. The hadronic decay rate into states of dissimilar quarks is then

$$\Gamma(3S_1(\bar{q}q) \to \text{hadrons}) = \Gamma(3S_1(\bar{q}q) \to ggg) + \Gamma(3S_1(\bar{q}q) \to g\tilde{g}\tilde{g}) + \cdots$$  \hspace{1cm} (4)$$

We have then

$$\alpha^3_{3,i}(\mu_S, m_\tilde{g}) \left(1 + \frac{3\alpha_{3,i}(\mu_S, m_\tilde{g})}{\pi} R(r) \right) = \alpha^3_{3,i}.$$  \hspace{1cm} (5)$$

The $ggg$ and $ggq\bar{q}$ correction to Eq. (4) is absorbed into the three gluon term by the choice of scale $\mu_S$ [1]. The coupling $\alpha_{3,i}$ on the right hand side of Eq. (5) is the result obtained in Ref. [1] assuming no gluino contribution ($m_\tilde{g} > m(3S_1)/2$).

Our analysis proceeds as follows. For fixed value of the gluino mass we solve Eq. (5) to obtain six values of the strong coupling constant at six scales appropriate to each of the vector quarkonia. The extrapolation to these scales from the $Z^0$ mass is done using the three loop renormalization group equation,

$$Q \frac{d\alpha_3(Q)}{dQ} = -\frac{\alpha_3^2(Q)}{2\pi} \left[a + \frac{b}{4\pi} \alpha_3(Q) + \frac{c}{16\pi^2} \alpha_3^2(Q) \right],$$  \hspace{1cm} (6)$$
where

\begin{align}
a &= 11 - \frac{2}{3} n_f(Q) \left( 1 + \frac{11}{10} \frac{\alpha_1}{4\pi} + \frac{9}{2} \frac{\alpha_2}{4\pi} \right) - 2n_\bar{g}(Q), \quad (7a) \\
b &= 102 - \frac{38}{3} n_f(Q) - 48n_\bar{g}(Q), \quad (7b) \\
c &= \frac{2857}{2} - \frac{5033}{18} n_f(Q) + \frac{325}{54} n_f^2(Q). \quad (7c)
\end{align}

The effects of scalar quarks and Higgs bosons are neglected in accord with the decoupling theorem. The off-diagonal two loop effects are treated as electroweak corrections to the one-loop a coefficient. We neglect the running of the electroweak corrections, using instead as an average value

\[ \alpha_{em} = 1/133, \quad \sin^2\theta_W = 0.2333, \quad (8) \]

and

\[ \alpha_1 = \frac{5}{3} \frac{\alpha_{em}}{\cos^2\theta_W}, \quad \alpha_2 = \frac{\alpha_{em}}{\sin^2\theta_W}. \quad (9) \]

The gluino contribution to the 3-loop c coefficient is not known. We keep the zero gluino contribution (Eq. (7c)) but monitor the effect of the c term as an estimate of unknown perturbative effects. In the fits of reasonable \( \chi^2 \), the contribution of the c term is small.

We take into account the threshold dependence of quarks and gluinos according to the formula

\begin{align}
n_f(Q) &= \sum_i f\left( \frac{Q^2}{4m_{q_i}^2} \right), \quad (10a) \\
n_\bar{g}(Q) &= f\left( \frac{Q^2}{4m_{\bar{g}}^2} \right), \quad (10b)
\end{align}

with \[ f(x) = 1 + \frac{1}{2 \sqrt{x(1+x)}} \ln \left[ \frac{\sqrt{1+x} - \sqrt{x}}{\sqrt{1+x} + \sqrt{x}} \right]. \quad (11) \]

For each value of \( \alpha_3(M_Z) \) and \( m_\bar{g} \) we can extrapolate from the \( Z^0 \) mass down to the quarkonium region using Eq. (3) and calculate a \( \chi^2 \).

The basic assumption given above does not require the total absence of non-perturbative effects. In fact, in the case of the (most accurately measured) 1S quarkonia states the graphs shown in Ref. [1] do exhibit some scatter in the data around the QCD predictions at about the 10% level. To find the preferred values of the gluino mass we
follow two alternative customary procedures for the treatment of data in such cases. The results of these two procedures do not greatly differ.

**Method 1:** discard the data point of worst agreement and minimize the $\chi^2$ of the remaining data points. In the present case this is the $J/\psi(1S)$ decay. The remaining five vector quarkonia decays, which contain at least one entry from strange, charm, and bottom quarks agree well with the theory and provide a relatively sharp minimum $\chi^2$ as a function of two parameters $\alpha_3(M_Z)$ and the gluino mass $m_{\tilde{g}}$.

**Method 2:** retain all six vector quarkonia decays but add in quadrature with the experimental errors a “theoretical error” to take into account possible non-perturbative or binding effects, i.e.

$$\alpha_3(\mu_S(i)) = \alpha_{3,i}(\mu_S, m_{\tilde{g}}) \pm \sqrt{\delta_i^2 + \lambda_i^2 \alpha_{3,i}^2(\mu_S, m_{\tilde{g}})}.$$  

(12)

On general grounds one would expect such corrections to be more important for the lighter quarkonia than for the $\Upsilon$ states. The $\lambda_i$ parameterize our ignorance about higher order and non-perturbative corrections. They do not, of course, constitute a model for such effects (since they do not shift the central values of $\alpha_{3,i}$) and in fact no reliable model exists apart from lattice QCD which has not as yet attained sufficient numerical accuracy. Clearly for sufficiently large $\lambda_i$ all predictive power is lost. We take our basic assumption to imply $\lambda_i \ll 1$. An adequate $\chi^2$ is found with $\lambda_i = 0.05$ for the bottomonium states, and $\lambda_i = 0.10$ for the charmonium and strangeonium vector states. Our conclusions are qualitatively insensitive to increasing this value for the lighter quarkonia in the sense that the gluino mass of minimum $\chi^2$ remains low although the $\chi^2$ values increase more slowly away from the minimum. Similarly the favored light gluino also persists for smaller $\lambda_i$ although the minimum $\chi^2$ is then not a mathematically acceptable fit.

For comparison with the best fits allowing a light gluino, we show in Fig. 1 the best fit to the quarkonia data in method 2 assuming the gluino lies at high mass (400 GeV) where it essentially decouples. This fit seems surprisingly good to the naked eye specially considering that it relates pure perturbative QCD to the hadronic decay rates of six vector quarkonia of three species over mass scales varying by a factor of ten. However it is not a mathematically good fit since it corresponds to a $\chi^2$ per degree of freedom ($\chi^2$/DoF) of 3.7. In Fig. 2 we show the variation in the $\chi^2$/6 for this heavy gluino case as a function of $\alpha_3(M_Z)$. The minimum $\chi^2$ of 3.7 is several standard deviations worse than the best fits with a light gluino. In addition this best fit corresponds to a value of $\alpha_3(M_Z)$ that is many
standard deviations away from the world average value and is inconsistent with SUSY unification with a SUSY scale below 10 TeV. A similar attempt to fit the five quarkonia states as in method 1 but without light gluinos would yield a minimum $\chi^2$ per degree of freedom many times larger than the minimum $\chi^2$ of Fig. 2.

In Fig. 3 we show the $\chi^2$ contours for method 1 treating the gluino mass $m_{\tilde{g}}$ and $\alpha_3(M_Z)$ as variable. The $\chi^2/5 = 1$ contours define two acceptable regions:

$$\alpha_3(M_Z) = 0.1135 \pm 0.0005, \quad m_{\tilde{g}} = (0.32 \pm 0.05) \text{ GeV}; \quad (13a)$$

or

$$\alpha_3(M_Z) = 0.1145^{+0.0013}_{-0.0006}, \quad m_{\tilde{g}} = (0.01^{+0.04}_{-0.01}) \text{ GeV}. \quad (13b)$$

The best fit corresponding to the central values of $(13a)$ is shown in Fig. 4. Comparing with Fig. 1, it is clear that the fit has improved due to decay of $\Upsilon$ states into gluino containing hadrons and due to the slower falloff of the coupling constant as a function of energy. In the solution corresponding to $(13b)$, even the $\phi$ decay has significant contribution from gluino containing final states. In this case both gluinos must presumably hadronize into a single pion where they can readily mix with gluon pairs. With light gluinos one must expect that all hadrons have non-negligible gluino components just as there is a non-negligible probability to find strange quarks in the sea of non-strange hadrons.

Fig. 5 shows the $\chi^2$ contours for method 2. In this method the $\chi^2/6 = 1$ contour lies within the region:

$$\alpha_3(M_Z) = 0.1115^{+0.0018}_{-0.0013}, \quad m_{\tilde{g}} = (0.44 \pm 0.17) \text{ GeV}. \quad (14)$$

More conservative values (90% CL) can be read from the $\chi^2/\text{DoF} = 2$ contours in Fig. 3 or Fig. 5. In method 2 there is also a tendency for the $\chi^2$ to drop again toward zero gluino mass although in this case the $\chi^2/\text{DoF}$ does not fall below 1 outside of the region of Eq. (14). The $\chi^2/\text{DoF} < 2$ region is defined by gluino masses less than 1.2 GeV. The fit to the six vector quarkonia assuming the central values of Eq. (14) is shown in Fig. 6. All three values of $\alpha_3(M_Z)$ picked out by the quarkonia data with light gluinos are in excellent agreement with the world average value for this quantity:

$$\alpha_3(M_Z) = 0.113 \pm 0.003 \quad \text{(World Average) [7].} \quad (15)$$
There is at present no well established theory of supersymmetry breaking that would allow the unambiguous prediction of the gluino mass. Nevertheless, in the most realistic models that have been extensively studied, supersymmetry is softly broken, triggered by a super Higgs mechanism in the hidden sector of some minimal $N = 1$ supergravity theories [8]. Therefore, the possible SUSY breaking scenarios in such models can be parameterized in terms of only a few constants at the unification scale: the common gaugino mass $m_{1/2}$, scalar mass $m_0$, and the $A$ and $B$ parameters of dimension mass characterizing the cubic and quadratic soft-breaking terms that often exist as well. Although the analysis we presented above is independent of any specific SUSY breaking models, it is tempting to discuss the implication of our results to such models. For simplicity, we now consider such a SUSY GUTs model which assumes the low energy form of the minimal supersymmetric extension of the standard model (MSSM) [9]. For our purpose, it is enough to consider three soft-breaking parameters $m_{1/2}, m_0$ and $A$. It is interesting to note that, in this framework, the low gluino masses favored by our analysis are natural if the dominant SUSY breaking seed is the universal scalar mass $m_0$, i.e., $m_{1/2} \ll m_0$. Such a SUSY breaking pattern has been supported by recent considerations of proton stability in the context of (non-flipped) $SU(5)$ supergravity [10], and also favored by cosmological studies [11]. In fact, our results suggest

$$m_{1/2} = 0.$$  (16)

Such a model might be theoretically appealing since then SUSY breaking, like electroweak breaking, finds its origin in the scalar sector. On the other hand, SUSY breaking scenarios with $m_{1/2} \neq 0$ would lead to quite large gluino masses [12].

In the $m_{1/2} = 0$ scenario, generally, one expects a supersymmetric spectrum relatively light compared to what one would get in other scenarios. Besides the soft-breaking parameters and the currently unknown top quark mass, one needs two more parameters in order to specify the full spectrum: the ratio of the two Higgs vev’s $\tan\beta \equiv v_2/v_1$, and the Higgs mixing parameter $\mu$. In fact, all the gaugino masses then vanish at the tree level and the gauginos only receive masses through radiative corrections which are naturally small though dependent upon the masses of other particles [13]. In particular, the one-loop corrections to the gluino and photino (which is an exact mass eigenstate in this scenario)
masses are known, with the dominant contribution coming from graphs in which the top quark and its two superpartners circulate around the loop [13],

\[ \delta m_{\tilde{g}} = \frac{\alpha_3(m_t)}{8\pi} m_t F\left(\frac{m_{\tilde{t}_1}^2}{m_t^2}, \frac{m_{\tilde{t}_2}^2}{m_t^2}\right), \]  
\[ \delta m_{\tilde{\gamma}} = \frac{\alpha_{em}(m_t)}{3\pi} m_t F\left(\frac{m_{\tilde{t}_1}^2}{m_t^2}, \frac{m_{\tilde{t}_2}^2}{m_t^2}\right). \]  

where \( m_{\tilde{t}_1}, m_{\tilde{t}_2} \) are the masses of the two scalar top quarks (\( m_{\tilde{t}_1} < m_{\tilde{t}_2} \)), and

\[ F(x, y) = \sin 2\theta \left[ \frac{x}{1-x} \ln x - \frac{y}{1-y} \ln y \right]. \]  

The actual gluino and photino masses up to one-loop order are given by the absolute values of these mass corrections. The \( \theta \) in (19) is the mixing angle between the scalar partners of left- and right-handed top quarks, which rotates the \( \tilde{t}_{L,R} \) states into the mass eigenstates \( \tilde{t}_{1,2} \). The overall factor \( \sin 2\theta \) was omitted in Ref. [13], corresponding to the case where the difference between two diagonal terms of the \( 2 \times 2 \) mass-squared matrix of the scalar top quarks can be neglected (see Eq. 33 below). The more general result presented here has been recently calculated by one of us [14], and makes transparent the fact that these one-loop mass corrections vanish exactly if there is no left-right mixing, even if \( \tilde{t}_L \) and \( \tilde{t}_R \) (now mass eigenstates themselves) have non-degenerate masses. Since we are primarily interested here in the gluino mass, we neglect in Eq. (18) a contribution to the photino mass from the \( W^{\pm} \)-chargino loop diagrams [13]. A discussion of this contribution in its general form will be given elsewhere [14]. As shown below, with currently favored values for the top quark mass \( m_t \) and the two stop quark masses \( m_{\tilde{t}_1}, m_{\tilde{t}_2} \), the function \( F \) is such that gluino masses below 1 GeV are quite natural.

We now discuss the allowed regions for the relevant parameters in the \( m_{1/2} = 0 \) scenario. To simplify our approach we neglect the Yukawa couplings for all the fermions although, strictly speaking, this approximation is only very good for the first two generations. As a result, the diagonal elements of the sfermion mass-squared matrix can be written as [13]

\[ m_{\tilde{f}}^2 = m_0^2 + m_f^2 + M_Z^2 \cos 2\beta \left[ T_{3,f} - c_f \sin^2 \theta_W \right]. \]  

Here we have included the D-term contributions as well. The average mass-squared of the sfermions is seen to differ from the average mass-squared of the fermions by the parameter \( m_0^2 \),

\[ \langle m_{\tilde{f}}^2 \rangle = m_0^2 + \langle m_f^2 \rangle. \]  

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The average mass-squared of the SUSY partners is approximately equal to the effective SUSY scale $M_S^2$ that enters into grand unification considerations. In Ref. [1] (Eq. (3.11a)) it was shown that the assumption of minimal SUSY unification with a light gluino (below $M_Z/2$) requires the approximate relation

$$M_S = 150 \text{ GeV} \times e^{-518.5\sin^2\theta_W - 0.2336} e^{1.85\alpha_3^{-1}(M_Z) - 0.113^{-1}}.$$  \hspace{1cm} (22)

Equating $M_S^2$ with $\left\langle m_{\tilde{f}}^2 \right\rangle$ and substituting the world average values of $\sin^2\theta_W$ and $\alpha_3(M_Z)$ from Ref. [7] yields within errors the following range for $m_0$,

$$75 \text{ GeV} < m_0 < 270 \text{ GeV} \hspace{1cm} (23)$$

For the top quark mass, we will assume [22]

$$92 \text{ GeV} < m_t < 147 \text{ GeV} \hspace{1cm} (24)$$

In the $m_{1/2} = 0$ scenario, the tree-level masses of the two charginos, $\chi^\pm_{1,2}$, are given by

$$m_{\chi^\pm_{1,2}}^2 = \frac{1}{2} \left[ 2M_W^2 + \mu^2 \pm \sqrt{\mu^4 + 4M_W^2\mu^2 + 4M_W^4\cos^22\beta} \right]. \hspace{1cm} (25)$$

From Eq. (25) and the requirement that the lightest chargino ($\chi^\pm_1$) has to be heavier than about half the $Z^0$ mass [13], it is found that $\tan\beta$ is restricted from both sides, i.e.

$$0.441 < \tan\beta < 2.266. \hspace{1cm} (26)$$

And for each value of $\tan\beta$ in the above range there is an upper limit on the absolute value of the Higgs mixing parameter $\mu$. Furthermore, the lower limit of 41 GeV on the mass of the light CP-even Higgs boson yields the additional constraints [17]

$$0.55 < \tan\beta < 0.65; \hspace{1cm} \text{ or } \hspace{1cm} 1.5 < \tan\beta. \hspace{1cm} (27)$$

Combining (26) with (27) yields two allowed ranges for the $\cos2\beta$ factor of Eq. (20),

$$-0.674 < \cos2\beta < -0.385; \hspace{1cm} \text{ or } \hspace{1cm} 0.406 < \cos2\beta < 0.536. \hspace{1cm} (28)$$

The bounds of Eq. (27) change with the experimental lower limit on the Higgs mass, becoming inconsistent with Eq. (26) if this mass is required to be above 70 GeV. The
allowed parameter space of the $m_{1/2} = 0$ model also requires that the mass of the lightest chargino be below $M_W$.

The mass-squares of the two scalar top quarks entering into Eqs. (17) and (18) are given by

$$m^2_{t_1, t_2} = \frac{1}{2} \left[ m^2_{LL} + m^2_{RR} \mp \sqrt{(m^2_{LL} - m^2_{RR})^2 + 4m^4_{LR}} \right]$$

with the diagonal elements (see Eq. (20))

$$m^2_{LL} = m^2_0 + m^2_t + M_Z^2 \cos^2 \beta \left[ \frac{1}{2} - \frac{2}{3} \sin^2 \theta_W \right],$$

$$m^2_{RR} = m^2_0 + m^2_t + M_Z^2 \cos^2 \beta \left[ \frac{2}{3} \sin^2 \theta_W \right],$$

and off diagonal element

$$m^2_{LR} = m_t \left( A_t + \frac{\mu}{\tan \beta} \right) \equiv m_t m_0 \tilde{A}_t.$$  (32)

Here we have introduced the dimensionless mixing parameter $\tilde{A}_t$ as a useful combination of $\tan \beta$, $\mu$ and the low energy top soft-breaking parameter $A_t$. If there is a single source of SUSY breaking corresponding to a single scale $m_0$, we might expect values of $\tilde{A}_t$ to be either zero or of the order unity. In terms of $\tilde{A}_t$ the mixing angle factor in Eq. (19) is then

$$\sin 2\theta = \frac{-2m_0 m_t \tilde{A}_t}{\sqrt{(m^2_{LL} - m^2_{RR})^2 + 4m^4_0 m^2_t \tilde{A}_t^2}}.$$  (33)

Imposing the experimental constraint $m_{t_1} > M_Z/2$ together with (23), (24) and (28) requires that $\tilde{A}_t < 3.2$. We also use for $\alpha_3$ the value at the top mass $\alpha_3(m_t) \approx 0.1$ in evaluating the gluino mass according to Eq. (17). In Fig. 7 we show the range of gluino masses predicted by Eq. (17) as a function of $\tilde{A}_t$ for the allowed range of values of $\tilde{A}_t$, $m_0$, $m_t$, and $\cos 2\beta$ described above. The constraint from Eq. (29) that the stop quarks be above half the $Z^0$ mass is also required in this allowed range. For each value of $\tilde{A}_t$ there exist an upper and lower limit on the gluino mass $m_{\tilde{g}}$. Values of $\tilde{A}_t$ near zero are consistent with the near zero gluino masses of Eq. (13a). Values of $\tilde{A}_t$ near unity are consistent with the $\chi^2$ minima of Eqs. (13a) and (14). The entire range of values of $\tilde{A}_t$ assuming Eq. (23) yields gluino masses below 1.3 GeV in agreement with the quarkonia data at the two-standard-deviation level. Assuming the result of Eqs. (13) or (14), detailed predictions for the individual squark and slepton masses within narrow ranges can be made. Much of the allowed parameter space in Fig. 7 predicts one or more of the scalar quarks
and leptons to have a mass between $M_Z/2$ and $M_Z$. If the effect of a light gluino and a possible light squark is taken into account the anomalously large quoted values of $\alpha_3$ coming from the $Z^0$ hadronic decay can perhaps be reconciled with the world average value $^{[18]}$. In addition, using $\alpha(m_t) \simeq 1/127.9$ in Eq. (18) the photino mass $m_{\tilde{\gamma}}$ would then be about five times smaller than the gluino mass $m_{\tilde{g}}$. A stable photino of mass about 100 eV could provide enough dark matter to close the universe $^{[19]}$. More massive photinos would “overclose” the universe unless they could decay efficiently into photon plus gravitino or annihilate efficiently into photons. Such ultralight photinos have also been discussed as the explanation of other astrophysical observations $^{[20]}$. On the other hand photinos of such mass have been found to be disfavored $^{[21]}$ by data from Supernova 1987A unless the squark masses are outside the preferred range of 60 GeV to 2.5 TeV.

CONCLUSIONS

We have shown that the quarkonia data behave as if there were a gluino octet in the region below 1 GeV. Treating the data in either of two ways, such a light gluino is favored by at least several standard deviations over the best QCD fits without a light gluino. In addition the fits without a light gluino are in conflict with SUSY unification and in disagreement with the world average values of $\alpha_3(M_Z)$ as discussed in Ref. $^{[1]}$. However, one must remain aware of the usual possibility that any phenomenological fit to data could be coincidental. We can not rule out the possibility that the real explanation for the behavior of the quarkonia data might lie in relativistic binding corrections or other non-perturbative effects, although this would contradict the general assumption that the narrowness of the vector quarkonia is due to perturbative QCD and asymptotic freedom. Confirmation from other independent data will certainly be required before the results presented here could be considered compelling. We are presently investigating the possibility that supporting evidence may be present elsewhere in existing data or that definitive experimental tests can be proposed.

In addition we should address the question as to whether such light gluinos are ruled out by current bounds. The strongest bounds on the masses of supersymmetric particles come from the decay of the $Z^0$. Any particle with electroweak charge must lie above about one half the $Z^0$ mass.

Bounds on other particles such as gluinos or bounds from other processes are all to a greater extent model dependent. The status of these bounds is discussed in Ref. $^{[3]}$. Although many lower bounds on the gluino mass above the mass region indicated here
have been quoted, all of these are to some extent dependent on untested assumptions. For example, the stringent bounds from the hadron colliders on heavy gluinos allow windows for light gluinos below 50 GeV. The low energy windows are well illustrated in Ref. [23]. Ref. [2] confirms the lack of unanimous opinion concerning whether or not very light gluinos have been ruled out. Many of the purported bounds have obvious loopholes some of which are pointed out in Ref. [1] and elsewhere [24]. For example the CUSB [25] bound that disfavors gluino masses between 0.6 GeV and 2.2 GeV from the non-observation of γ+gluinoball final states in Υ decay is strongly dependent on the value of the wave function at the origin of the gluinoball for which only models can be made [26]. Although the gluino behaves like a quark with a different color charge, it does have quartic couplings to gluons and gluinos that do not affect quarkonia in the same way. The binding of gluinos into new hadrons is therefore more closely related to the binding of gluons into new hadrons which is a very poorly understood area of hadronic physics at present. Cosmological constraints on light photinos and gluinos are subject to similar uncertainties. Most of the range of photino masses between 100 eV and 2 GeV is disfavored by one or more cosmological arguments. However there is a window noted in Ref. [27] for a photino in the mass range from 4 to 15 MeV. The gluino would then be in the range from 20 to 75 MeV consistent with Eq. (13b). It is not clear whether there is sufficient uncertainty in the cosmological arguments to stretch this range by a factor of four to accommodate the results of our Eq. (13a) or (14).

Given the current situation we feel that prudence requires the adoption of a conservative, non-dogmatic attitude concerning the compatibility of light gluinos with existing data.

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Figure Captions

Fig. 1. Best fit to the 3-gluon decays of six quarkonia states assuming a gluino mass of 400 GeV. The experimental errors are increased according to “method 2”.

Fig. 2. $\chi^2/6$ as a function of $\alpha_3(M_Z)$ for the 3-gluon decays of six quarkonia states treated according to method 2 but assuming a heavy gluino (mass 400 GeV).

Fig. 3. $\chi^2$ contours for the fit to 5 vector quarkonia according to “method 1” as a function of $\alpha_3(M_Z)$ and the gluino mass $\tilde{m}_g$.

Fig. 4. Fit to the 5 vector quarkonia using the “best fit” values from Eq. (13a).

Fig. 5. $\chi^2$ contours for the fit to 6 vector quarkonia according to “method 2” as a function of $\alpha_3(M_Z)$ and the gluino mass $\tilde{m}_g$.

Fig. 6. Best fit to the 6 vector quarkonia with the expanded errors of method 2. The fit parameters are given by the central values of Eq. (14).

Fig. 7. The allowed values of the gluino mass in the soft SUSY breaking picture with $m_{1/2} = 0$ are bounded by the closed figure shown here. The range of allowed gluino masses at given $\tilde{A}_t$ corresponds to taking the full allowed range of the remaining parameters $m_0$, $m_t$, and $\tan\beta$. 