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SEARCH FOR GLUINOS

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ABSTRACT

Simply supersymmetric gauge theories predict the existence of gluinos, massive and long-lived colour octet spin 1/2 partners of the vector gluons. We discuss their phenomenology in e^+e^- annihilation - focusing on quarkonium decays as well as continuum production - and in deep inelastic scattering experiments - focusing on the energy-momentum sum rule and bremsstrahlung production at high Q^2 . The best lower limit on the gluino mass comes from quarkonium decays and is about 3 GeV.

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Supersymmetry is such a beautiful idea¹⁾ that it would be particularly pleasing to have some experimental evidence for its validity. Unfortunately, the known elementary particles do not fit conveniently into supermultiplets, suggesting that the supersymmetric partners (squarks, sleptons, gluinos, etc.) of the known states (quarks, leptons, gluons, etc.) must have significantly higher masses. This means that supersymmetry must be broken, and various proposals have been made for the scale and manner of its breaking. These include explicit supersymmetry breaking²⁾, spontaneous breaking through elementary scalar fields³⁾, and most recently dynamical symmetry breaking^{4),5)}. In the first of these schemes²⁾ the magnitudes of the superpartner masses are put in by hand and largely arbitrary, whereas the other schemes offer some hope of predicting their masses. Spontaneous³⁾ and dynamical symmetry breaking schemes^{4),5)} differ on the masses of the squarks and sleptons, but both seem to like the gluinos (superpartners λ of the conventional massless gluons g) to have low masses of a few GeV^{5),6)}. This would make them accessible to experimental searches with present-day accelerators, and their phenomenology has been examined previously^{7),8)}. In particular, a lower limit on the gluino mass has been derived from the absence of a signal in beam dump experiments⁹⁾.

In this paper we extend the previous studies of gluino phenomenology, focusing mainly on the pair production of gluinos in e^+e^- collisions, on and off heavy quarkonium resonances, on pair production in deep inelastic lepton-nucleon scattering and on the indirect effects of gluinos on the energy-momentum sum rule for deep inelastic scattering. A striking feature of the recently proposed schemes^{4),5)} for dynamical breaking of supersymmetry is that they suggest⁵⁾ a relatively long lifetime $\tau_\lambda = O(10^{-10})$ seconds. This is considerably longer than the previous "spontaneous" estimate of $O(10^{-12} \text{ to } 10^{-15})$ seconds^{6),7),9)}, and provides the highly distinctive signature of an observable decay path, thus simplifying experimental searches. This relatively long lifetime alters the interpretation⁹⁾ of beam dump results, as we also discuss.

We start with a recap of some salient properties of gluinos⁶⁾⁻⁹⁾. They are Majorana spinors λ transforming as an octet of colour $SU(3)$. In schemes of spontaneous breaking of supersymmetry they may acquire masses from quantum effects⁶⁾, and an example suggests

$$m_\lambda = O(m_q) \approx \text{few GeV} ? \quad (1a)$$

Gluino masses arise naturally in models of dynamical symmetry breaking^{4),5)}, and a recent estimate⁵⁾ is

$$m_\lambda \approx \left(\frac{\alpha_3}{\pi}\right)\left(\frac{\alpha_2}{\pi}\right) N_f (g_2^2 \Lambda_{SC}) \ln g_2^2 \quad (1b)$$

where $\alpha_2 = g_2^2/4\pi$ and α_3 are the weak SU(2) and strong SU(3) couplings [the latter evaluated here at a scale $O(g_2^2 \Lambda_{SC})$], N_f is the number of quark flavours and Λ_{SC} is the scale of dynamical supersymmetry breaking. It has been estimated^{4),5)} that $\Lambda_{SC} = O(10 \text{ to } 100 \text{ TeV})$ and as a result the mass estimate (1b) becomes

$$m_\lambda = O(5) \text{ GeV} \times \left(\frac{\Lambda_{SC}}{10 \text{ TeV}}\right) \quad (2)$$

which should be accessible to contemporary experiments. Since gluinos are coloured, one expects them to be confined into "shadrons" containing gluons and/or quarks: $\Lambda \rightarrow (\lambda g)_0$ or $(\lambda \bar{q} q)_{0,\pm}$ where the subscripts denote electric charges. One might expect most such "shadrons" to be the electromagnetically neutral $(\lambda g)_0$, but who knows? Estimates have been made of the shadron lifetime in terms of the gluino lifetime:

$$\tau_{sh} \approx \tau_\lambda \approx O(10^{-12} \text{ to } 10^{-15}) \text{ sec.} \quad (3a)$$

in spontaneous schemes^{6),7),9)}, and

$$\tau_{sh} \approx \tau_\lambda \approx O(10^{-10}) \times \left(\frac{10 \text{ TeV}}{\Lambda_{SC}}\right) \text{ sec.} \quad (3b)$$

in dynamical schemes⁵⁾. The dominant decays are expected to be into conventional hadrons plus a "nuino", a light, colourless and neutral object resembling a neutrino, affording the distinctive pair production and decay signature illustrated in Fig. 1.

Armed with the expectations (1), (2) and (3), we turn first of all to gluino production in e^+e^- collisions. The rates for gluino pair production in heavy quarkonium decays via the diagrams of Fig. 2 can easily be computed using the known¹⁰⁾ rates for their decays into heavy quarks. For 3S_1 orthoquarkonium we have from Fig. 2a in leading order in α_s

$$R_{\lambda\lambda}(^3S_1) \equiv \frac{\Gamma(^3S_1(\bar{q}q) \rightarrow (\lambda\lambda)gg)}{\Gamma(^3S_1(\bar{q}q) \rightarrow gg\bar{g})} \approx \frac{9\alpha_3}{\pi} \int_0^1 \frac{dx}{x^2} \pi^{(2)}\left(\frac{x^2}{4x}\right) \Gamma(^3S_1; x) \quad (4)$$

where

$$Z \equiv \frac{2m_\lambda}{m(^3S_1)} \approx \frac{m_\lambda}{m_q} \quad (5)$$

The origin of formula (4) is clear: $\pi^{(2)}$ is the imaginary part of the vacuum polarization due to gluinos and Γ is a fictitious width for the decay of the 3S_1 into one massive and two massless gluons. The integration is performed over $x \equiv q^2/4m_q^2$ where q^2 is the (mass)² of the internal gluon. The expressions for $\pi^{(2)}$ and $\Gamma(^3S_1; x)$ can be found in Ref. 10). A plot of $R_{\lambda\lambda}(^3S_1)$ as a function of Z (5) is shown in Fig. 3. We see that if experiments can establish an upper limit which seems feasible

$$R_{\lambda\lambda}(^3S_1) < 10^{-3} \Rightarrow Z > 0.75 \quad (6)$$

one will be able to exclude gluinos with masses almost as large as that of the heavy quark q . We expect the 3P_1 decay into gluinos to be especially favourable since the competing decay modes are of the same order in α_s . Using formulae (3), (4) and (5) of Ref. 10), we calculate

$$R_{\lambda\lambda}(^3P_1) \equiv \frac{\Gamma(^3P_1(\bar{q}q) \rightarrow (\lambda\lambda)g)}{\Gamma(^3P_1(\bar{q}q) \rightarrow (\bar{q}q)g + (\lambda\lambda)g)} \quad (7)$$

which exceeds 0.1 for $Z < 0.9$ when the number of light flavours $N_L = 4$, the case for $\bar{b}b$ decays. A plot of $R_{\lambda\lambda}(^3P_1)$ as a function of $Z \equiv 2m_\lambda/m_{^3P_1}$ for $N_L = 4$ and 5 is shown in Fig. 4a. Since the radiative decays of 3S_1 radial recurrences into $^3P_1 + \gamma$ are expected to have branching ratios of a few per cent, it may be better to look for gluinos in inclusive N^3S_1 ($N \geq 1$) decays such as ψ' , T' and T'' (not even bothering to detect the photons from the radiative decays) rather than in the lowest-lying 1^3S_1 decays ($J/\psi, T$). We have also used Fig. 2b to estimate branching ratios $R_{\lambda\lambda}$ for other P wave quarkonium states¹⁰⁾:

$$R_{\lambda\lambda}(^3P_{0,2}) \equiv \frac{\Gamma(^3P_{0,2} \rightarrow (\lambda\lambda)g)}{\Gamma(^3P_{0,2} \rightarrow gg)} = \frac{6\alpha_s}{\pi} \int_{Z^2}^1 \frac{dx}{x} \pi^{(2)}\left(\frac{Z^2}{4x}\right) \Gamma(^3P_{0,2}; x) \quad (8)$$

where the notation was introduced after Eq. (5) above. Despite the external factor of α_s in Eq. (8), these branching ratios are expected to be large, as shown in Fig. 4b. We therefore conclude that the cascades $N^3S_1 \rightarrow N^3P_{0,1,2} + \gamma$, $N^3P_{0,1,2} \rightarrow (\lambda\lambda)g$ may be prime prospects for gluino hunting in quarkonium decays. An upper limit of 10^{-3} on gluino production in T or T' decays would imply $m_\lambda \gtrsim 4\frac{1}{2}$ GeV. Care should be exercised in using the results (4), (6), (7) and (8) in the charmonium system since lowest order perturbation theory may not be a good approximation at these low energies and masses. Also, for very heavy quarkonium systems there may be other important non-gluonic (weak?) decay modes which increase the total widths and hence dilute any gluino signal.

We now turn to gluino pair production in the e^+e^- continuum via the lowest order diagram of Fig. 5a which should dominate close to threshold. Using the calculations of Refs. 11) and 12) we obtain the curve shown in Fig. 6 for

$$R_{\lambda\lambda}(e^+e^-) \equiv \frac{\sigma(e^+e^- \rightarrow \bar{q}q\lambda\lambda)}{\sigma(e^+e^- \rightarrow \bar{q}q)} = \begin{cases} 10^{-3} & \text{for } E_{cm}^2/4m_\lambda^2 = 15 \\ 10^{-2} & \text{for } E_{cm}^2/4m_\lambda^2 = 130 \end{cases} \quad (9)$$

If experiments at PETRA could obtain the feasible(?) upper limit of 10^{-3} on $R_{\lambda\lambda}(e^+e^-)$ they would exclude gluinos weighing less than about 5 GeV.

Another set of plausible reactions to look for gluinos is provided by deep inelastic lepton-nucleon scattering. We have evaluated the basic diagram¹³⁾ of Fig. 5b for both charged current weak interactions and for electromagnetic interactions, with a particular view to ep colliding ring experiments. The ratio

$$R_{\lambda\lambda}(CC) \equiv \frac{\sigma(e_L^- + p \rightarrow \nu + (\lambda\lambda) + X)}{\sigma(e_L^- + p \rightarrow \nu + \text{all})} \quad (10)$$

is plotted as a function of $s[\equiv(p_e + p_p)^2]/4m_\lambda^2$ in Fig. 7a, while the analogous quantity $R_{\lambda\lambda}(\text{em}) \equiv d^2\sigma(\lambda\lambda x)/dx dy / (2\pi\alpha^2 s/Q^4)$, calculated with quark distributions $\propto (1-x)^3$ for representative values of the conventional deep inelastic scaling variables x and y is plotted in Fig. 7b. If experiments at HERA¹⁴⁾ [$\sqrt{s} \equiv (p_e + p_p)^2 = 314$ GeV] or CHEER ($\sqrt{s} = 200$ GeV) could exclude $(\lambda\lambda)$ production at the level of 10^{-4} (recall the distinctive signature of Fig. 1) they would imply $m_\lambda \gtrsim 4$ GeV.

It has been proposed previously⁸⁾ that one look indirectly for gluinos through their effect on the energy momentum sum rule for deep inelastic scattering. The Altarelli-Parisi equations¹⁵⁾ for the leading order evolution of singlet parton densities with increasing momentum transfers Q^2 are modified as follows by the lowest order diagrams of Fig. 8:

$$Q^2 \frac{\partial S(x, Q^2)}{\partial Q^2} = \frac{\alpha_s(Q^2)}{2\pi} \int_x^1 \frac{dy}{y} \left[S(y, Q^2) P_{q \rightarrow q} \left(\frac{x}{y} \right) + G(y, Q^2) P_{q \rightarrow g} \left(\frac{x}{y} \right) \right] \quad (11a)$$

$$Q^2 \frac{\partial G(x, Q^2)}{\partial Q^2} = \frac{\alpha_s(Q^2)}{2\pi} \int_x^1 \frac{dy}{y} \left[S(y, Q^2) P_{q \rightarrow g}\left(\frac{x}{y}\right) + G(y, Q^2) P_{g \rightarrow g}\left(\frac{x}{y}\right) + L(y, Q^2) P_{\lambda \rightarrow g}\left(\frac{x}{y}\right) \right] \quad (11b)$$

$$Q^2 \frac{\partial L(x, Q^2)}{\partial Q^2} = \frac{\alpha_s(Q^2)}{2\pi} \int_x^1 \frac{dy}{y} \left[G(y, Q^2) P_{g \rightarrow \lambda}\left(\frac{x}{y}\right) + L(y, Q^2) P_{\lambda \rightarrow \lambda}\left(\frac{x}{y}\right) \right] \quad (11c)$$

where $S(x, Q^2)$ is the singlet quark density, $G(x, Q^2)$ is the gluon density and $L(x, Q^2)$ is the gluino density. The new splitting functions $P_{\lambda \rightarrow g}$ (11b) and $P_{\lambda \rightarrow \lambda}$ (11c) are simply related to those for ordinary quarks:

$$P_{\lambda \rightarrow g}(z) = \frac{9}{4} P_{q \rightarrow g}(z) = 3 \left[1 + \frac{(1-z)^2}{z} \right] \quad (12a)$$

$$P_{g \rightarrow \lambda}(z) = 3 P_{g \rightarrow q}(z) = \frac{3}{2} [z^2 + (1-z)^2] \quad (12b)$$

$$P_{\lambda \rightarrow \lambda}(z) = \frac{9}{4} P_{q \rightarrow q}(z) = P_{\lambda \rightarrow g}(1-z) \quad (12c)$$

We will not discuss in detail the evolution of the gluino distribution $L(x, Q^2)$ and its implications for the observables $S(x, Q^2)$ and $G(x, Q^2)$. We will instead focus on the $n = 2$ moments of equations which can be written in the form:

$$Q^2 \frac{\partial}{\partial Q^2} \begin{pmatrix} S_2(Q^2) \\ G_2(Q^2) \\ L_2(Q^2) \end{pmatrix} = \frac{\alpha_s(Q^2)}{2\pi} \begin{pmatrix} -\frac{16}{9} & \frac{N_f}{3} & 0 \\ \frac{16}{9} & (-\frac{N_f}{3}-1) & 4 \\ 0 & 1 & -4 \end{pmatrix} \begin{pmatrix} S_2(Q^2) \\ G_2(Q^2) \\ L_2(Q^2) \end{pmatrix} \quad (13)$$

The evolution matrix (13) has three eigenvalues e :

$$e = 0, \quad e_{\pm} = \frac{-(3N_f+6) \pm \sqrt{9N_f^2 - 66N_f + 84}}{18} \quad (14)$$

corresponding to the eigenvectors

$$(1, 1, 1), \quad \left(\frac{1}{1 + \frac{9}{16}e_{\pm}}, 1, \frac{1}{1 + \frac{e_{\pm}}{4}} \right) \quad (15)$$

These describe combinations of $S_2(Q^2)$, $G_2(Q^2)$ and $L_2(Q^2)$ which evolve with Q^2 as

$$\left(\ln \frac{Q^2}{\Lambda^2} \right) \quad 0, \quad \frac{6e_{\pm}}{27-2N_f} \quad (16)$$

respectively. At infinite Q^2 one arrives at the steady state

$$\left. \begin{aligned} \int_0^1 dx \, x \, S(x, Q^2) &= \frac{3N_f}{20+3N_f} \quad , \quad \int_0^1 dx \, x \, G(x, Q^2) = \frac{16}{20+3N_f} \\ \int_0^1 dx \, x \, L(x, Q^2) &= \frac{4}{20+3N_f} \end{aligned} \right\} \quad (17)$$

or 47%, 42% and 11% respectively for $N_f = 6$. This compares with the conventional expectation

$$\int_0^1 dx \, x \, S(x, Q^2) = \frac{3N_f}{16+3N_f} \quad , \quad \int_0^1 dx \, x \, G(x, Q^2) = \frac{16}{16+3N_f} \quad (18)$$

or 53%, 47% respectively for $N_f = 6$ in the absence¹⁶⁾ of gluinos. The difference in $\int_0^1 dx \, x \, S(x, Q^2)$ is significantly smaller than previous estimates^{8), 17)}. Furthermore, the effect of gluinos on the singlet quark distribution is only indirect since gluinos only interact directly with gluons. This means that the asymptotic limit (17) is approached very slowly if one starts with $L(x, Q_0^2)$ at some initial value Q_0^2 of the momentum transfer. Calculations in the modified minimal subtraction scheme¹⁸⁾ suggest that the most appropriate Q_0^2 is equal to $2m_\lambda^2$. Figure 9 shows the effect of the modified evolution equations (13) assuming $Q_0^2 = 4.5 \text{ GeV}^2$ corresponding to $m_\lambda = m_c = 1.5 \text{ GeV}$. We see that the difference between $S_2(Q^2)$ in (17) and (18) develops rather slowly, and it is difficult to imagine this being detectable even at HERA¹⁴⁾. Fortunately, as we have seen earlier, there are plenty of more direct ways to look for gluinos.

Finally, we discuss gluino production in hadron-hadron collisions. Perturbative QCD estimates of the total cross-sections for heavy quark production in hadronic collisions are notoriously unreliable, generally falling short of the observed $(\bar{c}c)$ production rates. In the cases of pN collisions at 400 GeV and πN collisions at 300 GeV a reasonable rule-of-thumb guess seems to be

$$\sigma(h+h \rightarrow (ff)+X) = O(20) \mu b \times \left(\frac{m_c}{m_f} \right)^\alpha : \alpha = O(6) ? \quad (19)$$

for fermion masses m_f between $m_c \approx 1.5 \text{ GeV}$ and $m_b \approx 5 \text{ GeV}$. Thus if one could exclude gluino production at the level of the $\bar{b}b$ cross-section, which we guess to be of order 25nb at presently accessible energies, then one would exclude $m_\lambda \lesssim m_b \approx 5 \text{ GeV}$.

In their pioneering work, Farrar and Fayet⁹⁾ used the previous beam dump experiment limits on the anomalous prompt production of neutrino-like objects coming from shadron decays to infer that $m_b > m_c = 1.5$ GeV. While their arguments only apply directly to the short-lived shadrons expected (3a) in spontaneously broken supersymmetric theories^{3),6),9)}, one can now probably extend their limit to the longer lived shadrons expected (3b) in dynamically broken theories⁵⁾. These shadrons might have interacted in the target before decay, thus degrading the shadron energy. The eventual shadron decay would produce a low energy neutrino-like object whose interaction cross-section would be too small to get an observable rate in the beam-dump experiment detectors deployed so far. However, since the beam-dump experiments analyzed by Farrar and Fayet⁹⁾ there have been several experiments observing the hadronic production of charmed particles in visual detectors such as bubble chambers. They would surely have noticed the production of particles with cross-sections and masses comparable to those of charm, but which decayed after $O(10^{-10})$ seconds, and we are even aware of one intriguing observation¹⁹⁾. Presumably the previous lower limit of $O(m_c)$ on the gluino mass still stands.

Since no experiment has yet quoted any limits on new particles with the weird properties (2), (3b), it is difficult to be sure what is the most stringent limit on the mass of such a gluino. Almost certainly one has a limit of $O(1\frac{1}{2})$ GeV from the visual hadronic experiments discussed in the previous paragraph. We find it difficult to believe that long-lived gluinos could have been missed if they were produced in T or T' decays at the per cent level, and therefore suspect that a true lower limit on m_λ is about 3 GeV. However, this limit should be confirmed and improved by a dedicated search. Experiments at present energies could probably exclude $m_\lambda < 4$ or 5 GeV, which nudges the estimate (2) but does not rule it out. The best prospects for improving this limit in the near future and perhaps creating a conflict with the mass estimate (2) seem to come from toponium decay (gluinos with $m_\lambda < m_t \lesssim 20$ GeV could be excluded) and from Z^0 decays [$R_{\lambda\lambda}(e^+e^-) < 10^{-4}$ at $E_{cm} = 100$ GeV implies $m_\lambda \gtrsim 20$ GeV].

This paper has been written in a rather negative way, talking about establishing upper limits in gluino production rates and hence providing lower limits on the gluino mass. In fact, one should think more positively. Here is a relatively easily detectable, peculiar particle whose experimental discovery would finally demonstrate the relevance of supersymmetry¹⁾ and hence revolutionize physics.

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FIGURE CAPTIONS

- Figure 1 : Signature for gluino production in a high energy collision. Gluinos are produced in pairs, may live as long as 10^{-10} secs. and should decay into light neutrino-like particles plus hadrons.
- Figure 2 : Lowest order diagrams a) for $^3S_1 \rightarrow g\bar{g}\lambda\lambda$ and b) for $^3P_{0,1,2} \rightarrow g\lambda\lambda$.
- Figure 3 : A plot of the gluino production rate $R_{\lambda\lambda}(^3S_1)$ (4) as a function of $Z \equiv 2m_\lambda/m(^3S_1)$.
- Figure 4 : a) A plot of the gluino production rate $R_{\lambda\lambda}(^3P_1)$ (7) as a function of $Z \equiv 2m_\lambda/m(^3P_1)$. The solid line is calculated with four light flavours and a quark binding energy of 0.6 GeV, corresponding to the bottomonium system. The dashed line is calculated with five light flavours and assuming a Coulombic binding energy $m_t\alpha_s^2/9$ appropriate(?) for the toponium system.
b) The quantities $R_{\lambda\lambda}(^3P_{0,2}) \equiv \Gamma(^3P_{0,2} \rightarrow \lambda\bar{\lambda}g)/\Gamma(^3P_{0,2} \rightarrow g\bar{g})$, calculated using similar assumptions to those used for $R_{\lambda\lambda}(^3P_1)$, are plotted for both the bottomonium and toponium systems. The ratios $R_{\lambda\lambda}(^3P_{0,2})$ are rather sensitive to the assumption of a Coulombic binding energy for the toponium system.
- Figure 5 : The lowest order diagrams a) for $e^+e^- \rightarrow q\bar{q}\lambda\lambda$, and b) for $\ell + p \rightarrow \ell' + (\lambda\lambda) + X$.
- Figure 6 : A plot of the gluino production rate $R_{\lambda\lambda}(e^+e^-)$ (9) as a function of $Z \equiv 2m_\lambda/\sqrt{s}$. The solid line is an exact lowest order calculation¹¹⁾, while the dashed line is an approximation¹²⁾ neglecting constant terms. For rates at higher energies, use three times the value of the $(\bar{c}c)$ pair production rate plotted in Fig. 10 of Ref. 12) evaluated at $Q^2 = (m_c^2/m_\lambda^2)s$.
- Figure 7 : a) A plot of the gluino production rate $R_{\lambda\lambda}(CC)$ (10) as a function of $s/4m_\lambda^2$, neglecting effects due to the $W^\pm$ propagator and
b) the analogous quantity $R_{\lambda\lambda}(em)$ (defined in the text) for $x = 0.1, y = 0.2$ (solid line) and $x = 0.05, y = 0.1$ (dashed line).

Figure 8 : Additional lowest order diagrams contributing to the Altarelli-Parisi equations (11) when gluinos are present.

Figure 9 : The evolution of the fractions S_2 , G_2 and L_2 of the singlet nucleon momentum carried by quarks, gluons and gluinos respectively. The curves are evaluated assuming $\Lambda(\text{four flavours}) = 200 \text{ MeV}$. The solid lines apply in the absence of gluinos, while the dashed curves are calculated with gluinos of mass 1.5 GeV . The arrows indicate the asymptotic limits as $Q^2 \rightarrow \infty$.

Fig. 1

Fig. 2

Fig. 3

Fig.4a

Fig. 4b

Fig. 5

Fig. 6

Fig. 7a

Fig. 7 b

Fig . 8

Fig.9