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Collider signals of a superlight gravitino

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Supergravity theories with a superlight gravitino (of mass $\sim m_W^2 M_{\rm Pl}^{-1}$) are shown to give rise to monojet and dijet events with large missing p_T which should be observable at the Fermilab Tevatron collider. For a gluino mass less than 200 GeV, the signal is much bigger than expected from the standard model or the usual supergravity theory. The observation of such events will thus be a clear signal of supersymmetry.

In the last decade, a great deal of effort has been devoted to the search for supersymmetry in the leptonic and hadronic colliders. For example, the Collider Detector at Fermilab (CDF) group at the Tevatron collider, from their analysis of 25.3 nb^{-1} of data, has set interesting bounds on the masses of the gluino ($\tilde{g}$) and the scalar quarks ($\tilde{q}$);[1] for $m_{\tilde{q}}=\infty$ ($m_{\tilde{g}}=\infty$), they obtain $m_{\tilde{g}}>73$ GeV ($m_{\tilde{q}}>75$ GeV), assuming the gluino decays dominantly to $q\bar{q}\tilde{\gamma}$ and the squark dominantly to $q\tilde{\gamma}$. Most of these searches have been restricted to the signatures arising from the usual supergravity theory[2] (USG) in which the gravitino ($\tilde{G}$) has a mass of the order of m_W, i.e., the same order as the masses of the other superpartners ($\tilde{g},\tilde{q},\tilde{W},\tilde{Z}$, etc.), and is not the lightest supersymmetric particle (LSP). There exists another class of local supersymmetry (SUSY) theory in which the gravitino is superlight, with its mass given by $m_{\tilde{G}}\sim m_W^2 M_{\rm Pl}^{-1}\sim 10^{-15}$ GeV.[3] In such theories, $\tilde{G}$ is the LSP, and the signal for SUSY changes drastically; for example, a gluino decays to $g+\tilde{G}$, instead of the traditional decay mode in the USG theory, $\tilde{g}\rightarrow q_i\bar{q}_j\tilde{x}$, where $\tilde{x}$ is a chargino or a neutralino. The object of this work is to point out that if nature chooses a locally supersymmetric theory with a superlight gravitino, then the prospect of discovering SUSY in the Tevatron collider is far greater than for the USG theory. We find that, at the Tevatron energy, the production of a gluino pair (by the usual QCD processes plus supergravitational interactions) or $\tilde{g}+\tilde{G}$ (by the supergravitational interaction), and the subsequent decay, $\tilde{g}\rightarrow g+\tilde{G}$, gives rise to dijet and monojet signals which are an order of magnitude or more larger than that expected from the standard model or USG theory.

To motivate our work, let us first discuss why the supergravity interactions can be interesting in laboratory experiments. The gravitino is the spin-$\frac{3}{2}$ superpartner of the graviton in the $N=1$ local SUSY theory. Its interactions are all gravitational, so one might think that those are completely negligible in the laboratory when compared to the weak or strong interactions. This is indeed true in the USG theory, but not true in supergravity theories with a superlight gravitino. If a superlight gravitino appears as an internal particle, its propagator gives a factor $2p_\mu p_\nu/2m_{\tilde{G}}^2$. If it appears as an external particle, then after the spin sum, it gives the same factor. The net effect is that its effective coupling is enhanced from the gravitational coupling κ to $\kappa_{\rm eff}=\kappa(m_{\tilde{g}}/m_{\tilde{G}})$. For $m_{\tilde{g}}=100$ GeV, $m_{\tilde{G}}\sim 10^{-15}$ GeV, this enhancement factor is $m_{\tilde{g}}/m_{\tilde{G}}\sim 10^{17}$, making its gravitational interaction comparable to or stronger than the strong interaction. [As first pointed out by Fayet,[4] the longitudinal part of such a superlight gravitino Ψ_μ effectively behaves like a spin-$\frac{1}{2}$ Goldstino (χ) with the replacement $\Psi_\mu=i(\frac{2}{3})^{1/2}m_{\tilde{G}}^{-1}\partial_\mu\chi$.] In fact, it has been pointed out that the superlight gravitino can be pair produced with an observable cross section in γ-γ collisions from e^+e^- annihilation.[5]

What is the plausible mass range for a gravitino? The gravitino acquires its mass from the super Higgs effect by absorbing the fermionic partner of the chiral superfield z, in the hidden sector. In a local SUSY theory, the interactions of gravity with matter and the Yang-Mills fields are completely specified in terms of the usual gauge and gravitational couplings, plus two unknown functions of the chiral superfields (z), $\mathcal{G}(z,z^*)$ and $f_{ab}(z)$.[6] $\mathcal{G}(z,z^*)$ multiplies the scalar kinetic term, and is the so-called Kähler potential, whereas $f_{ab}(z)$ multiplies the gaugino kinetic term. If we assume that the Kähler potential is a polynomial, and the gaugino kinetic term is minimal, i.e., $f_{ab}(z)=\delta_{ab}$, then, after the super Higgs mechanism, the masses of the gravitino and the usual superpartners become comparable (with $m_{\tilde{G}}\sim m_{\tilde{q}}\sim m_{\tilde{W}}\sim m_W$). However, in a general theory (i.e., a general Kähler potential and/or nonminimal kinetic terms), the two mass scales are unrelated. Their mass ratio has the form

$$\frac{m_{\tilde{g}}}{m_{\tilde{G}}}=\frac{1}{2}\,\mathcal{G}'(v)\,\mathcal{G}''(v)^{-1}\frac{{\rm Re}f'(v)}{{\rm Re}f(v)}\,, \qquad (1)$$

where v is the vacuum expectation value of the field z, and $m_{\tilde g} \sim m_W$. The right-hand side (RHS) of Eq. (1) is arbitrary. Thus, in a general theory, the gravitino mass is essentially arbitrary with many possibilities: (i) $m_{\tilde G} \sim m_W^2/M_{\rm Pl}$ (superlight gravitino); (ii) $m_{\tilde G} \sim m_W$ (usual supergravity), (iii) $m_{\tilde G} \sim M_{\rm Pl}$ (ultraheavy gravitino); (iv) $m_{\tilde G} \sim (m_W/M_{\rm Pl})^n M_{\rm Pl}$, $1 < n < 2$ (light gravitino). In our phenomenological consideration, we shall consider $m_{\tilde G}$ in the range 10^{-16}–10^{-10} GeV. For $m_{\tilde G} > 10^{-10}$ GeV, the theory behaves like the USG theory.

Now, we are ready to consider the implications of the superlight or light gravitino scenario in collider experiments. Our considerations are general; but in this work, we shall restrict ourselves to hadronic colliders, in particular to the Tevatron. The processes we consider are

$$\bar p p \rightarrow \tilde g \tilde g \,, \tag{2a}$$

$$\bar p p \rightarrow \tilde g \tilde G \,, \tag{2b}$$

$$\bar p p \rightarrow g S + g P \,, \tag{2c}$$

where S and P are the scalar and the pseudoscalar fields left from the hidden sector after the super Higgs phenomena. (The masses of S and P are essentially zero.) In the USG theory, the first process (2a) is the important one, the others are negligible. (This process has been experimentally studied by the CDF group to set the bound, $m_{\tilde g} > 73$ GeV for $m_{\tilde q} = \infty$.) But, even for this process, there are two important differences between the USG and the superlight-gravitino scenario. (1) Significant contributions arise from virtual superlight-gravitino exchange. This is in addition to the usual super QCD contributions. (2) There is an additional decay mode $\tilde g \rightarrow g + \tilde G$ which dominates over the usual three-body modes for a large region of the $m_{\tilde g}$-$m_{\tilde G}$ mass plane. Thus, the signature of SUSY is drastically altered, and at the Tevatron energy, we obtain a much larger dijet and monojet signal than the USG theory.

The relevant interactions for the processes under consideration are

$$e^{-1}\mathcal{L} = \tfrac{1}{4}\kappa\bar\lambda^a\gamma^\rho\sigma^{\mu\nu}\Psi_\rho F^a_{\mu\nu} + \frac{i}{2}\bar\lambda^a \not{D}\lambda^a + \tfrac{1}{4}\kappa\alpha S(F^a_{\mu\nu}F^{\mu\nu a} + \bar\lambda^a \not{D}\lambda^a) + \tfrac{1}{8}\kappa\alpha P[F^a_{\mu\nu}\tilde F^{\mu\nu a} - \tfrac{1}{2}e^{-1}D_\mu(e\bar\lambda^a\gamma_5\gamma^\mu\lambda^a)]$$
$$+\beta S\bar\lambda^a\lambda^a + \text{usual super QCD terms}\,. \tag{3}$$

Even in the superlight gravitino scenario, the couplings α and β are model dependent, in the sense that their values depend on the specific choices of the functions $\mathcal{G}(z,z^*)$ and $f_{ab}(z)$. For some models, α is enhanced by the $m_{\tilde G}^{-1}$ factor; but β is always small. For example, if $f_{ab}(z) = \delta_{ab}f(z)$, and $\mathcal{G}(z,z') = -3\ln[\kappa(z+z^*)]$, we get

$$\alpha = -\sqrt{2/3}\, m_{\tilde g}/m_{\tilde G}, \quad \beta \sim O(\kappa m_{\tilde g})\,. \tag{4}$$

We shall discuss the contributions of S and P interactions separately in order to separate out this model-dependent part of our results. In the above interactions (3), we have assumed strong CP conservation in the observable sector.

Now consider the new gluino decay mode

$$\tilde g \rightarrow g + \tilde G \tag{5}$$

and compare it with the usual decay mode

$$\tilde g \rightarrow q_i \bar q_j x, \quad x = \tilde W_i, \tilde Z_j\,. \tag{6}$$

In Fig. 1 we plot the branching fraction $B(\tilde g \rightarrow g\tilde G)$ as a function of $m_{\tilde G}$ for given gluino and squark masses, assuming the gluino decays either to $\tilde g \rightarrow g\tilde G$ or $\tilde g \rightarrow q\bar q\tilde Z_1$, where $\tilde Z_1$ is the LSP. [This result is also true when the other three-body decay modes in Eq. (6) are included.] The figure shows clearly that, for $m_{\tilde G} < 10^{-12}$ GeV, the gluino decays almost exclusively to a gluon gravitino; thus in this mass regime, we expect major modifications of the usual event topology of gluino pairs, irrespective of any role the gravitino may play in SUSY production.

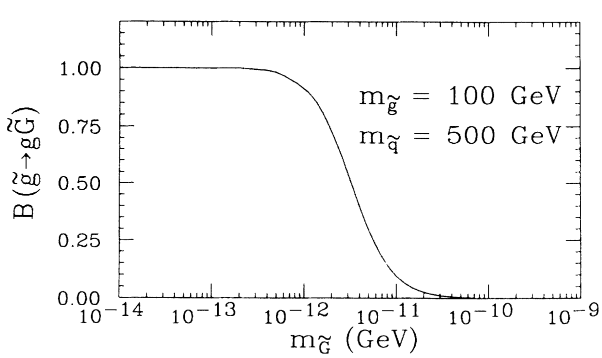

FIG. 1. Branching fraction for the $\tilde g \rightarrow g\tilde G$ decay.

Next, we consider the production of gluinos via the processes (2a) and (2b). At the Tevatron energy, we find production by gluon-gluon fusion to be dominant. For the process (2a), the Feynman diagrams are the t- and u-channel gravitino exchange, plus the usual supersymmetric QCD interaction diagrams. For the process (2b), the diagrams are s-channel gluon exchange, t- and u-channel gluino exchanges, plus the contact interaction. We have included all these contributions, and express our results in terms of the cross section for a specific number of jets.

These n-jet topological cross sections are displayed in Figs. 2 and 3. In these curves, we sum the contributions from both the $\tilde g\tilde g$ and $\tilde g\tilde G$ production processes and impose the following experimental cuts on the missing p_T, pseudorapidity of the leading jet, and jet topology of the events.

$$\not{p}_T > \max(40\text{ GeV}, 2.8\sqrt{\Sigma p_T})\,,$$
$$|\eta_{J_1}| < 1, \quad \Delta\phi(J_1,J_i) < 150^\circ\,.$$

We employ a jet-counting algorithm[7] and coalesce into a single jet partons for which $\Delta R = (\Delta\phi^2 + \Delta\eta^2)^{1/2} < 0.7$. We also assume that $m_{\tilde q} \geq 500$ GeV, so that signals due to squark production may be neglected.

Two factors affect event topology. The branching fraction $B(\tilde g \rightarrow g\tilde G)$ as mentioned earlier, and the relative

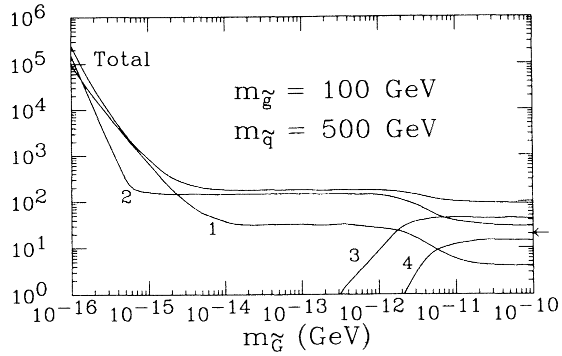

FIG. 2. n-jet topological and the total cross sections satisfying the cuts given in the text; the results are for the processes (2a) and (2b) for $\sqrt{s}=1.8$ TeV.

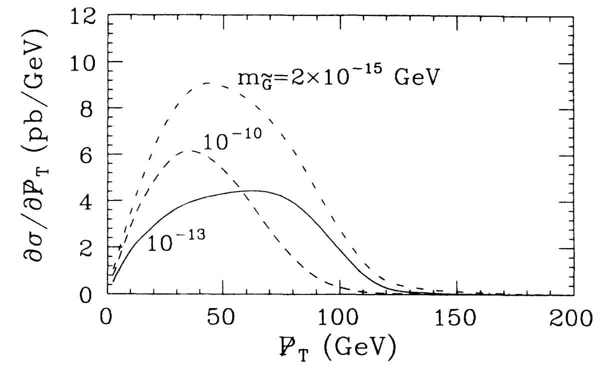

FIG. 4. The $\not{p}_T$ distribution of the total cross section for the processes (2a) and (2b) for $\sqrt{s}=1.8$ TeV.

rates of production of $\tilde{g}\tilde{g}$ vs $\tilde{g}\tilde{G}$. If, for instance, $m_{\tilde{G}} < 10^{-12}$ GeV, we know that $\tilde{g}\to g\tilde{G}$ dominates gluino decay. Thus in this range $\tilde{g}\tilde{g}$ production should lead to a predominance of dijet events, while $\tilde{g}\tilde{G}$ final states would preferentially produce monojets. These effects are clearly reflected in Figs. 2 and 3 which exhibit three distinct ranges in $m_{\tilde{G}}$. For $m_{\tilde{G}} > 10^{-11}$ GeV, one has $B(\tilde{g}\to g\tilde{G})\sim 0$, while $\tilde{g}\tilde{g}$ production totally dominates the cross section. Thus assuming $\tilde{g}\to q\bar{q}\tilde{Z}_1$, we find the standard $\not{p}_T$+multijet SUSY signal in which we have three- or four-jet dominance for $m_{\tilde{g}}=100$ to $m_{\tilde{g}}=200$ GeV. In the approximate range $10^{-14} < m_{\tilde{G}} < 10^{-12}$ GeV, however, $B(\tilde{g}\to g\tilde{G})\simeq 1$ and although $\tilde{g}\tilde{g}$ still dominates production, the event topology changes dramatically to exhibit dijet dominance. As we go still lower in gravitino mass, the $m_{\tilde{G}}$ dependence of the $\tilde{g}\tilde{G}$ cross section begins to have a visible effect. Below $m_{\tilde{G}}=10^{-14}$ for $m_{\tilde{g}}=100$ GeV and $m_{\tilde{G}}\simeq 5\times 10^{-14}$ for $m_{\tilde{g}}=200$ GeV, the monojet cross section rises sharply and quickly comes to dominate the dijet rate. At still lower gravitino masses ($m_{\tilde{G}}\lesssim 10^{-16}$ GeV) and $m_{\tilde{G}}^{-4}$ dependence of the gravitino-exchange contribution to $\sigma(\tilde{g}\tilde{g})$ begins to dominate and again, correspondingly, the dijet signal is largest. The general rise in the total cross section as $m_{\tilde{G}}$ decreases is understood in terms of the same factors. The $\not{p}_T$ distributions of the total cross sections are also given in Fig. 4 for three values of $m_{\tilde{G}}$. The harder $\not{p}_T$ distributions

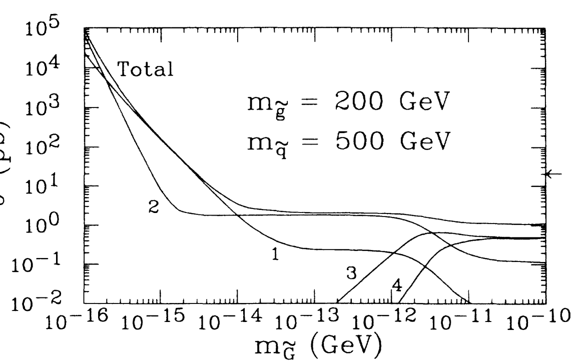

FIG. 3. Same as Fig. 2 for $m_{\tilde{g}}=200$ GeV.

for the lower gravitino masses are again due to the dominance of the new two-body decay over the usual three-body decay mode.

A shower Monte Carlo calculation of potential standard-model backgrounds to the $\not{p}_T+n$ jet signal has been performed[8] and yields the result of 22 and 21 pb, respectively, for the 1 jet and 2 jet backgrounds. The specific processes considered were W+jets, Z+jets, and $t\bar{t}$ production with $m_t=75$ GeV. The background falls to 17 and 15 pb, respectively, if $m_t=125$ GeV. We have indicated this background by the arrows on Figs. 2 and 3. For $m_{\tilde{g}}=100$ GeV, and $10^{-14} < m_{\tilde{G}} < 10^{-12}$ GeV, the dominant dijet signal is nearly an order of magnitude above background and the monojet signal below $m_{\tilde{G}}\simeq 3\times 10^{-15}$ GeV is even greater. For $m_{\tilde{g}}=200$ GeV, the cross section is of course suppressed; nevertheless, if $m_{\tilde{G}} < 3\times 10^{-15}$ the monojet signal remains above background.

The CDF result of $m_{\tilde{G}} > 73$ GeV, mentioned in the introduction, also implies an upper limit to the cross section for SUSY-particle production subjected to the cuts specified before. This result can be applied to any model in an effort to exclude regions of parameter space. To find the cross-section limit appropriate to our Monte Carlo simulation, we ran our program for $m_{\tilde{g}}=73$ GeV and $B(\tilde{g}\to q\bar{q}\tilde{Z}_1)=1$, making use of a simple relation between parton energies and measured jet cluster energies suggested by the CDF.[9] The result is $\sigma_0=121$ pb, in good agreement with the CDF bound of $\sigma_0 < 110$ pb. We then calculate the cross section $\sigma(m_{\tilde{g}},m_{\tilde{G}})$ passing the signal cuts given above and plot the contour $\sigma(m_{\tilde{g}},m_{\tilde{G}})=121$ pb in the $m_{\tilde{g}}$-$m_{\tilde{G}}$ plane in Fig. 5 (the solid curve). The figure exhibits the same general features as the topological cross sections. For $m_{\tilde{G}} > 10^{-11}$ GeV, the gravitational interaction is essentially negligible and $m_{\tilde{g}} > \sim 73$ GeV as in the CDF analysis. However, as $m_{\tilde{G}}$ decreases, maintaining a constant $m_{\tilde{g}}$ would lead to an increase in σ for the reasons cited earlier. Thus to maintain the cross section at the limit of 121 pb, the gluino mass must rise. The effect is most dramatic below $m_{\tilde{G}}=10^{-14}$ GeV, where the monojets from $\tilde{g}\tilde{G}$ production begin to dominate. One notes that if $m_{\tilde{G}}=10^{-15}$ GeV the Tevatron data already rules out gluino masses up to 220 GeV.

Now, we briefly discuss the effect of process (3a) on our

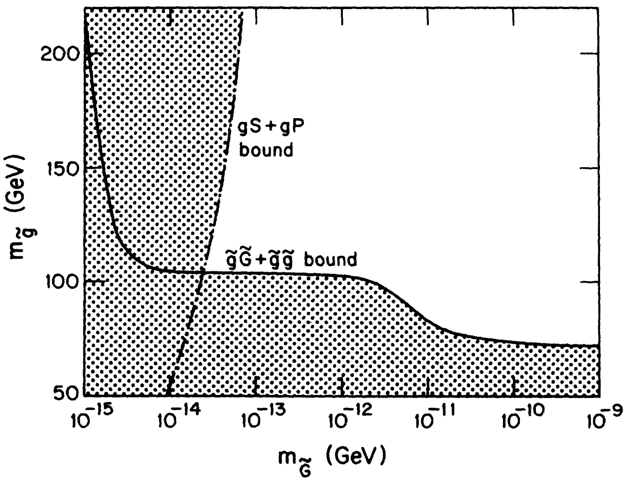

FIG. 5. $m_{\tilde{g}}$-$m_{\tilde{G}}$ mass bound using the CDF data (Ref. 1) from the processes (2a) and (2b) (solid curve) and from the process (2c) (dash-dotted curve). The shaded region is excluded.

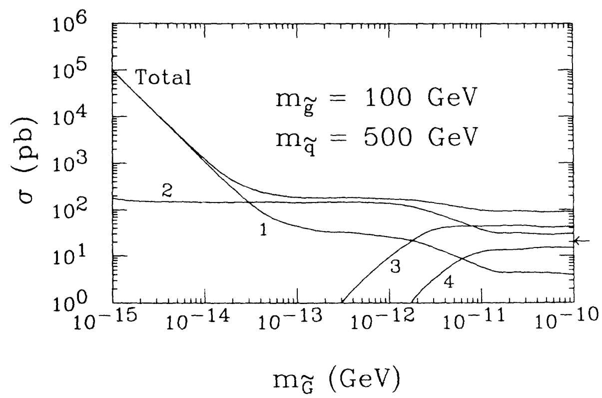

FIG. 6. Same as in Fig. 2, except the contributions from the processes (2c) are also included.

results, for models in which the coupling α is given by Eq. (4). The results for the topological cross sections are shown in Fig. 6. At the smaller gravitino mass ranges, the cross section is dominated by the monojets arising from this process, since there is no phase-space suppression in this case. The $m_{\tilde{g}}$-$m_{\tilde{G}}$ mass bound due to this process alone is shown by the dash-dotted line in Fig. 5. Note that the process $p\bar{p}\rightarrow gS+gP$ yields an upper bound on $m_{\tilde{g}}$ because the cross section is proportional to $m_{\tilde{g}}^2$. Thus, in Fig. 5, the region above the dash-dotted curve is excluded by $(gS+gP)$. It is also interesting to note that combining the two excluded regions yields an absolute lower bound on $m_{\tilde{G}}$ of about 2.2×10^{-14} GeV. This is a significant improvement over the previous lower bound, 2.3×10^{-15} GeV (not absolute) derived by Fayet[10] from the e^+e^- annihilation results.

Finally, we point out that for a given energy we cannot choose the gravitino mass as small as we please. For each gluino mass, there is a lower bound on the gravitino mass below which the tree-level amplitude violates unitarity. Such unitarity constraints and further details of our work (including the cross-section formulas and the light-squark case) will be presented elsewhere.[11]

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