Phenomenology of gluino decays via loops and top-quark Yukawa coupling

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(Received 14 May 1990)

We have reanalyzed the decay patterns of the gluino ($\tilde{g}$) incorporating all its decays to charginos ($\tilde{W}_i$) and neutralinos ($\tilde{Z}_i$) as given by the minimal model. We have made an independent computation of the decay width for $\tilde{g}\to g\tilde{Z}_i$ which occurs at one loop and shown that the branching fraction for this can be up to 0.5. We have also computed the width for $\tilde{g}\to t\bar{t}\tilde{Z}_i$ including the effect of the top-quark Yukawa coupling, which significantly increases this width. We have studied the gluino decay modes as a function of parameter space, and clearly delineated the regions where the loop decays or $\tilde{g}\to t\bar{t}\tilde{Z}_i$ is large. The significance of these decays for gluino signatures in collider experiments is briefly commented on.

I. INTRODUCTION

The realization[1] that supersymmetry (SUSY) provides a potential solution to the gauge hierarchy problem[2] if the SUSY mass gap is smaller than ~ 1 TeV has provided experimentalists with ample motivation to search for supersymmetric partners of known particles (sparticles) at high-energy colliders. Negative results from these searches have been translated to lower bounds on sparticle masses, usually within the framework of the minimal supersymmetric model[3] (MSSM). The most straightforward bounds come from e^+e^- colliders since the production of heavy sparticles leads to striking signatures in a clean environment. The best limits from e^+e^- machines today come from the Z^0 factories in operation at SLAC and at CERN. Lower limits[4] of almost $M_Z/2$ have been accounced on the masses of charged scalar leptons ($\tilde{l}$), squarks ($\tilde{q}$), and the light chargino ($\tilde{W}_-$) under the assumption that all sparticles directly decay to the lightest supersymmetric particle (LSP) which escapes detection. It has also been shown that the measurement of the total Z^0 width and the peak hadronic cross section at the CERN e^+e^- collider LEP can be translated[5] to bounds on sparticle masses that are only slightly weaker than the bounds of Ref. 4, but are almost independent of how the sparticles decay. In the future, LEP II will probe[6] sparticle masses to near 100 GeV, but then there will be no further exploration until construction of a higher-energy e^+e^- machine becomes feasible.

However, the existing hadron colliders at CERN and Fermilab are already probing mass scales well in excess of $M_Z/2$. The UA2 experiment[7] at the CERN S$p\bar{p}$S and the Collider Detector at Fermilab (CDF) experiment[8] at the Tevatron have already announced lower limits of 70–80 GeV on the masses of squarks ($\tilde{q}$) and gluinos ($\tilde{g}$) under the assumption that they directly decay to the lightest neutralino ($\tilde{Z}_1$), which is taken to be the LSP. Since the publication of these bounds, the CDF experiment has accumulated almost 200 times as much data and is probing squark and gluino masses in the 120–150-GeV range.

It has been known for some time that gluinos and squarks have significant branching ratios for decays into all charginos and neutralinos that are kinematically accessible.[9,10] The daughter charginos and neutralinos further decay via a cascade until at the end of the chain the LSP ($\tilde{Z}_1$) is reached. These cascade decays lead to various promising experimental signatures, including (i) events containing jets plus missing p_T ($\not{p}_T$) due to nondetection of the $\tilde{Z}_1$'s,[11,12] (ii) isolated hard same-sign dileptons plus jets plus $\not{p}_T$,[13,12] and (iii) events containing one or more W and Z bosons plus jets plus $\not{p}_T$.[10,14] These signatures depend not only on the decay properties of the produced gluinos and squarks, but also on the decay properties of the daughter charginos and neutralinos.

To fix the various sparticle masses and couplings, which in turn specifies the allowed cascade decays and their rates, we adopt the masses and sparticle mixing angles as given by the MSSM. Within this framework, each chiral fermion f_i ($i=L,R$) of the standard model (SM) has a spin-zero partner $\tilde{f}_i$ with the same SU(3)$\times$SU(2)$\times$U(1) quantum numbers as f_i. Furthermore, there are two charginos ($\tilde{W}_-$ and $\tilde{W}_+$ with $m_{\tilde{W}_-} < m_{\tilde{W}_+}$), four Majorana neutralinos ($\tilde{Z}_{1,2,3,4}$ in order of increasing mass), and a color-octet gluino ($\tilde{g}$), all with spin $\frac{1}{2}$. If, as usual,[3] we assume that the SUSY-breaking Majorana masses of the SU(3), SU(2), and U(1) gauginos are equal at some unification scale, the masses and mixings of $\tilde{W}_i$ and $\tilde{Z}_j$ are fixed by three parameters which we take to be (i) the gluino mass, $m_{\tilde{g}}$ (ii) the supersymmetric Higgsino mass, $2m_1$ and (iii) the ratio v'/v of the vacuum expectation values (VEV's) of the Higgs field h' and h that give masses to the d- and u-type quarks, respectively. Just one more parameter,[15] the charged-Higgs-boson mass (m_{H^+}) suffices to fix the Higgs-boson sector of the MSSM, and thereby the decays of all the $\widetilde{W}_i$ and $\widetilde{Z}_j$. Specification of these four parameters together with the squark mass ($m_{\tilde{q}}$), assuming all squarks are approximately degenerate, suffices to fix all the tree-level decays of squarks and gluinos:

$$\begin{aligned} &\tilde{g} \to q\bar{q}\widetilde{W}_i \ \text{ or } q\bar{q}\widetilde{Z}_i \ \text{ or } q\tilde{q} \ , \\ &\tilde{q} \to q\widetilde{W}_i \ \text{ or } q\widetilde{Z}_i \ \text{ or } q\tilde{g} \ . \end{aligned} \tag{1.1}$$

Henceforth, we assume $m_{\tilde{q}} > m_{\tilde{g}}$, so $\tilde{g} \to q\tilde{q}$ is forbidden.

It has, however, been pointed out[16,17] that there are cases where the loop decays of the gluinos can dominate the tree-level decays. For instance, if $\widetilde{Z}_1$ is almost a Higgsino, the tree-level decay $\tilde{g} \to q\bar{q}\widetilde{Z}_1$ which is mediated by the virtual squark is strongly suppressed by the quark Yukawa coupling unless the decay $\tilde{g} \to t\bar{t}\widetilde{Z}_1$ is kinematically accessible. Then the decay $\tilde{g} \to g\widetilde{Z}_1$ mediated by the top-quark-top-squark loop may be competitive with the three-body decays of the gluinos into the heavier neutralinos and charginos. We will see that this is in fact the case in some regions of the parameter space of the MSSM.

The partial width for this decay has been computed by Ma and Wong[16] for $m_{\widetilde{Z}_1} = 0$ and by Barbieri *et al.*[17] in the limit $m_{\tilde{q}} \to \infty$. We have independently computed this decay width without any approximation for $m_{\tilde{q}}$ or $m_{\widetilde{Z}_1}$ and for general values of neutralino mixings as given by the MSSM. Our calculation generalizes the calculation of Haber and Kane[18] who computed it for the special case $\widetilde{Z}_1 = \tilde{\gamma}$ and showed that it was unimportant. Note also that the computation of the radiative decay of the gluino is analogous to earlier calculations[19] of the radiative decays of the photino (Higgsino) to a photon plus a Higgsino (photino). This calculation has recently been generalized for arbitrary neutralino mixings by Haber and Wyler.[20] This generalized calculation can be translated to the result for the gluino decay by abstracting out the fermion-sfermion loop contributions to the neutralino decay. We thought it useful, however, to present an explicit formula for $\Gamma(\tilde{g} \to g\widetilde{Z}_i)$ since one does not exist in the literature.

The two-body gluino loop decay can have important phenomenological implications. For instance, in a recent paper,[21] we pointed out that if $|2m_1|$ is not large, it may be possible to search for gluinos by looking for multijet $+\not{p}_T + Z^0$ events even at the Fermilab Tevatron, for a wide range of squark and gluino masses. For $m_{\tilde{g}} \gtrsim 250$ GeV, this signal was dominantly due to the decay $\tilde{g} \to g + \widetilde{Z}_4$, $\widetilde{Z}_4 \to \widetilde{Z}_1 + Z^0$, which occurred with a branching fraction of just 3–5%. Note also that if the two-body gluino events occur at a substantial rate, then gluino events can lead to much harder jet and $\not{p}_T$ activity than is usually expected.

Finally, since as remarked above, the Higgsino component of the neutralino may be significant for $\tilde{g} \to t\bar{t}\widetilde{Z}_i$, we have performed a new computation for this decay rate. We have also evaluated the branching fraction for this mode as the parameters of the MSSM are varied, since if it is significant, $\tilde{g}\tilde{g}$ production can lead to characteristic events with four hard t quarks in them. Since it is now known[22] that $m_t > M_W$, gluinos may then lead to four-W events at hadron supercolliders such as the Superconducting Super Collider (SSC) or the CERN Large Hadron Collider (LHC).

The remainder of this paper is organized as follows. In Sec. II, we present analytic formulas for the partial widths for the decays $\tilde{g} \to g\widetilde{Z}_i$ and $\tilde{g} \to t\bar{t}\widetilde{Z}_i$ (with the contribution of Higgsino included). In Sec. III, we present numerical results based on these formulas. We find that there are regions of parameter space of the MSSM that are not excluded by the bounds from LEP (and which are unlikely to be excluded in the near future) for which the loop branching fractions are almost 50%. Also, we find that for $m_{\tilde{g}} \gtrsim 0.6$ TeV, the branching fraction for the $\tilde{g} \to t\bar{t}\widetilde{Z}_i$ decay can be well in excess of 15%, reaching over 30% for certain ranges of parameters. As far as we are aware, neither of these decays have been taken into account in existing analysis[11,12,13,14] of gluino signals (except in Ref. 21 referred to above). We present some general conclusions and a summary of our results in Sec. IV.

II. FORMULAS FOR GLUINO WIDTHS VIA LOOPS AND YUKAWA COUPLINGS

The possibility of the two-body decay of the gluino, $\tilde{g} \to g\tilde{\gamma}$, was first pointed out by Haber and Kane[18] who noted that this decay rate vanishes if $m_{\tilde{q}_L} = m_{\tilde{q}_R}$. Since then, several authors[16,17] have observed that this is special to the case $\widetilde{Z}_i = \tilde{\gamma}$. The partial width for the radiative decay of the gluino has recently been recomputed by Ma and Wong[16] (for $m_{\widetilde{Z}_1} = 0$) and by Barbieri *et al.*[17] (for $m_{\tilde{q}} \gg m_{\tilde{g}}, m_q$). As discussed below, this width can also be obtained from the results of Haber and Wyler[20] using appropriate substitutions. Because these calculations[16,17,20] differ from one another, we have independently computed this width for arbitrary masses and mixing of the daughter neutralino and without any approximation on the squark or neutralino mass.

By gauge invariance, the amplitude for the radiative decay has the form

$$\begin{aligned} \mathcal{M}(\tilde{g}_A \to g_B \widetilde{Z}_i) = &(\tfrac{1}{2}\delta_{AB}) C_i \bar{u}(\widetilde{Z}_i) \\ &\times \not{k}\not{\epsilon}[P_R - (-1)^{\theta_{\tilde{g}} + \theta_i} P_L] u(\tilde{g}) \ , \end{aligned} \tag{2.1}$$

where $(\frac{1}{2}\delta_{AB})$ is the color factor and k and ϵ are the four-momentum and polarization vector of the gluon, $P_{L,R} = (1 \mp \gamma_5)/2$, and other particle labels denote their four-momenta. In terms of the constant C_i, the form of which is given below, the partial width for the radiative decay is given by

$$\Gamma(\tilde{g} \to g\widetilde{Z}_i) = \frac{|C_i|^2}{32\pi m_{\tilde{g}}^3} (m_{\tilde{g}}^2 - m_{\widetilde{Z}_i}^2)^3 \ . \tag{2.2}$$

From the calculation of the one-loop amplitudes[16,17,20] involving squark-quark loops we find that the constant C_i is given by

$$C_i = \frac{\sqrt{2}\alpha_s}{4\pi}(i)^{\theta_i+\theta_{\tilde g}-1}(-1)^{\theta_i+1} \times \sum_q \{ m_{\tilde Z_i}[(K+I_1)_L \tilde A^q_{\tilde Z_i} - (K+I_1)_R \tilde B^q_{\tilde Z_i}] + m_{\tilde g}(-1)^{\theta_i+\theta_{\tilde g}}(K_L \tilde A^q_{\tilde Z_i} - K_R \tilde B^q_{\tilde Z_i}) + m_q(-1)^{\theta_i}(I_L+I_R) f_q v_1^{(i)}\delta_{qt}\} , \quad (2.3)$$

where $\tilde A^q_{\tilde Z_i}$ ($\tilde B^q_{\tilde Z_i}$) are related to the couplings of the neutralino $\tilde Z_i$ to the $q\tilde q_L$ ($q\tilde q_R$) system via the gaugino content of $\tilde Z_i$. They are given in terms of the components $v_j^{(i)}$ of the neutralino in the $(\tilde h, \tilde h', \tilde\lambda_3, \tilde\lambda_0)$ basis (see Ref. 10 for details) by

$$\tilde A^u_{\tilde Z_i} = \frac{g}{\sqrt{2}}v_3^{(i)} + \frac{g'}{3\sqrt{2}}v_4^{(i)} ,$$
$$\tilde A^d_{\tilde Z_i} = \frac{-g}{\sqrt{2}}v_3^{(i)} + \frac{g'}{3\sqrt{2}}v_4^{(i)} , \quad (2.4)$$
$$\tilde B^u_{\tilde Z_i} = \frac{4}{3}\frac{g'}{\sqrt{2}}v_4^{(i)}, \quad \tilde B^d_{\tilde Z_i} = -\frac{2}{3}\frac{g'}{\sqrt{2}}v_4^{(i)} ,$$

and θ_i ($\theta_{\tilde g}$) equals 0 if the mass eigenvalue of $\tilde Z_i$ (the SUSY-breaking gaugino mass) is positive, and equal to 1 otherwise. The last term in Eq. (2.3) arises from the Higgsino content of the neutralino, and so is proportional to the quark Yukawa coupling f_q. This term, which is proportional to $f_q m_q$, has thus been retained only for the top family loops for which the coupling f_t is given by

$$f_t = \frac{g m_t}{\sqrt{2}M_W}\frac{(v^2+v'^2)^{1/2}}{v} . \quad (2.5)$$

Since the up-type quarks only couple to the H superfield, this also explains the mixing factor $v_1^{(i)}$ in this term.

Finally, the functions K, I, and I_1 which arises from the calculation of the loop are given by

$$K = \frac{-1}{m_{\tilde g}^2 - m_{\tilde Z_i}^2}\left[1 + \int_0^1 dx \left[\frac{m_{\tilde Z_i}^2 x(x-1) + m_{\tilde q}^2 x + m_q^2(1-x)}{(m_{\tilde g}^2 - m_{\tilde Z_i}^2)x(1-x)}\right]\ln X\right] , \quad (2.6a)$$

$$I = \frac{1}{m_{\tilde g}^2 - m_{\tilde Z_i}^2}\int_0^1 \frac{dx}{x}\ln X , \quad (2.6b)$$

$$I_1 = \frac{1}{m_{\tilde g}^2 - m_{\tilde Z_i}^2}\int_0^1 dx \ln X , \quad (2.6c)$$

where

$$X = \frac{-m_{\tilde g}^2 x(1-x) + m_{\tilde q}^2 x + m_q^2(1-x)}{-m_{\tilde Z_i}^2 x(1-x) + m_{\tilde q}^2 x + m_q^2(1-x)} .$$

The subscripts L and R in Eq. (2.3) denote the squark type whose mass enters into Eq. (2.6).

Equations (2.1)–(2.6) implicitly assume that $\tilde q_L$ and $\tilde q_R$ are mass eigenstates. Since the mixing angle (θ_t) between these states is[23] $O(m_q/m_{\tilde q})$, this assumption may be invalid for the top-squark eigenstates, $\tilde t_1$ and $\tilde t_2$. Equation (2.3) can easily be modified to take this mixing into account by multiplying the $\tilde q_L$ and $\tilde q_R$ contributions by $\cos\theta_t$ ($-\sin\theta_t$) and $\sin\theta_t$ ($\cos\theta_t$) for the $\tilde t_1$ ($\tilde t_2$) contribution and replacing the squark mass by $m_{\tilde t_1}$ ($m_{\tilde t_2}$).

Our result for the radiative decay width for gluino agrees with that of Ma and Wong[16] (who computed it for $m_{\tilde Z_1}=0$) aside from a color factor $\frac{1}{4}$ which was inadvertently omitted in their Eq. (18). Our result, for degenerate top-squark masses, is larger by a factor 4 as compared to Eq. (2.10) of Barbieri *et al.*[17] (who have computed it for $m_{\tilde q} \to \infty$). We have also checked that our amplitude (2.3) agrees with Eq. (53) of Haber and Wyler[20] with the substitutions $Z_{j1}=\cos\theta_W$, $Z_{j2}=\sin\theta_W$, $Z_{j3}=Z_{j4}=0$, and $ee_t^2 g\sin\theta_W \to g_s^2$ to convert their $\tilde Z_2 \to \tilde Z_1\gamma$ calculation to our calculation.

Note that if the squarks are degenerate only the Yukawa coupling term in (2.3) survives so that the radiative decay depends only on the h-Higgsino content of the neutralino and on the top-squark sector. The gaugino content of $\tilde Z_i$ is thus important only when $m_{\tilde t_L} \neq m_{\tilde t_R}$ since the other squarks are expected[23,3] to be degenerate.

Next we present a simplified formula for the decay $\tilde g \to t\bar t \tilde Z_i$, including the coupling of the Higgsino component of the neutralino. We have ignored the t-quark mass in the computation of the matrix element (except, of course, in the Yukawa coupling) but have retained it in the phase-space integration; this should be a good approximation since the $\tilde g \to t\bar t\tilde Z_i$ decay rate is significant only for very heavy gluinos. The partial width for this decay can be calculated from the couplings,

$$\mathcal{L}=-\sqrt{2}g_s(-i)^{\theta_{\tilde{g}}}\tilde{t}^{\dagger}_{L}\bar{\tilde{g}}_{A}\frac{\lambda_A}{2}\frac{1-\gamma_5}{2}t+\sqrt{2}g_s(i)^{\theta_{\tilde{g}}}\tilde{t}^{\dagger}_{R}\bar{\tilde{g}}_{A}\frac{\lambda_A}{2}\frac{1+\gamma_5}{2}t$$
$$+\bar{\tilde{Z}}_i\left[iA^t_{\tilde{Z}_i}\frac{1-\gamma_5}{2}-f_t v_1^{(i)}(i)^{\theta_i}\frac{1+\gamma_5}{2}\right]t\tilde{t}^{\dagger}_L+\bar{\tilde{Z}}_i\left[iB^t_{\tilde{Z}_i}\frac{1+\gamma_5}{2}-f_t v_1^{(i)}(-i)^{\theta_i}\frac{1-\gamma_5}{2}\right]t\tilde{t}^{\dagger}_R$$
$$+\bar{\tilde{Z}}_i\left[iA^b_{\tilde{Z}_i}\frac{1-\gamma_5}{2}-f_b v_2^{(i)}(i)^{\theta_i}\frac{1+\gamma_5}{2}\right]b\tilde{b}^{\dagger}_L+\bar{\tilde{Z}}_i\left[iB^b_{\tilde{Z}_i}\frac{1+\gamma_5}{2}-f_b v_2^{(i)}(-i)^{\theta_i}\frac{1-\gamma_5}{2}\right]b\tilde{b}^{\dagger}_R+\text{h.c.} \tag{2.7}$$

with

$$f_b=\frac{gm_b}{\sqrt{2}M_W}\frac{(v^2+v'^2)^{1/2}}{v'}$$

and the couplings $A^q_{\tilde{Z}_i}$ and $B^q_{\tilde{Z}_i}$ as given in Ref. 10, and particle labels denoting the fields. We obtain the $\tilde{Z}_i$ eigenvector components $v_j^{(i)}$ in the basis $(\tilde{h},\tilde{h}',\tilde{\lambda}_3,\tilde{\lambda}_0)$ by numerical diagonalization of the neutralino mass matrix.[10]

If the top-quark mass is neglected in the computation of the matrix element for the decay $\tilde{g}\to t\bar{t}\tilde{Z}_i$, but retained in the kinematics, the partial width is given by

$$\Gamma(\tilde{g}\to t\bar{t}\tilde{Z}_i)=\frac{1}{(2\pi)^5}\frac{1}{2m_{\tilde{g}}}(|\mathcal{M}_L|^2+|\mathcal{M}_R|^2)\ , \tag{2.8}$$

with

$$|\mathcal{M}_L|^2=2g_s^2[(|A^t_{\tilde{Z}_i}|^2+f_t^2v_1^{(i)2})\psi_t(m_{\tilde{g}},m_{\tilde{t}_L},m_{\tilde{Z}_i})+(-1)^{\theta_{\tilde{g}}+\theta_i}|A^t_{\tilde{Z}_i}|^2\phi_t(m_{\tilde{g}},m_{\tilde{t}_L},m_{\tilde{g}_i})]\ . \tag{2.9}$$

The functions ψ_t and ϕ_t are given by

$$\psi_t(m_{\tilde{g}},m_{\tilde{t}_L},m_{\tilde{Z}_i})=\pi^2m_{\tilde{g}}\int dE\frac{E(E^2-m_t^2)^{1/2}(m_{\tilde{g}}^2-m_{\tilde{Z}_i}^2-2m_{\tilde{g}}E)}{(m_{\tilde{g}}^2+m_t^2-2m_{\tilde{g}}E-m_{\tilde{t}_L}^2)^2}\frac{\lambda^{1/2}(m_{\tilde{g}}^2+m_t^2-2m_{\tilde{g}}E,m_{\tilde{Z}_i}^2,m_t^2)}{m_{\tilde{g}}^2+m_t^2-2m_{\tilde{g}}E} \tag{2.10a}$$

and

$$\phi_t(m_{\tilde{g}},m_{\tilde{t}_L},m_{\tilde{Z}_i})=\tfrac{1}{2}\pi^2m_{\tilde{g}}m_{\tilde{Z}_i}\int dE\frac{1}{m_{\tilde{g}}^2+m_t^2-m_{\tilde{t}_L}^2-2m_{\tilde{g}}E}$$
$$\times\left[-(E_{\max}-E_{\min})-\frac{2Em_{\tilde{g}}+m_{\tilde{Z}_i}^2-m_t^2-m_{\tilde{t}_L}^2}{2m_{\tilde{g}}}\ln\left[\frac{m_{\tilde{g}}^2+m_t^2-2m_{\tilde{g}}E_{\max}-m_{\tilde{t}_L}^2}{m_{\tilde{g}}^2+m_t^2-2m_{\tilde{g}}E_{\min}-m_{\tilde{t}_L}^2}\right]\right] \tag{2.10b}$$

with $E_{\min}$ and $E_{\max}$, the limits on the $\bar{t}$ energy in the gluino rest frame, given by

$$E_{\min,\max}=\frac{(m_{\tilde{g}}^2+2m_t^2-m_{\tilde{Z}_i}^2-2Em_{\tilde{g}})(m_{\tilde{g}}-E)\mp(E^2-m_t^2)^{1/2}m_{\tilde{g}}^2\lambda^{1/2}(1+m_t^2/m_{\tilde{g}}^2-2E/m_{\tilde{g}},m_t^2/m_{\tilde{g}}^2,m_{\tilde{Z}_i}^2/m_{\tilde{g}}^2)}{2(m_{\tilde{g}}^2+m_t^2-2Em_{\tilde{g}})}$$

and the limits of integration on the t-quark energy E in Eq. (2.10) ranging from m_t to $(m_{\tilde{g}}^2-2m_tm_{\tilde{Z}_i}-m_{\tilde{Z}_i}^2)/2m_{\tilde{g}}$. Here, $\lambda(a,b,c)=a^2+b^2+c^2-2ab-2ac-2bc$. Finally, $|\mathcal{M}_R|^2$ in Eq. (2.8) is obtained from $|\mathcal{M}_L|^2$ by the substitutions $A^t_{\tilde{Z}_i}\to B^t_{\tilde{Z}_i}$ and $m_{\tilde{t}_L}\to m_{\tilde{t}_R}$. We have checked that if the quark Yukawa coupling is neglected, this formula reduces to that given in Ref. 10. Finally, the width for the decay $\tilde{g}\to b\bar{b}\tilde{Z}_i$ can be read off from Eqs. (2.8)–(2.10) with obvious substitution in the couplings [see Eq. (2.7)] and kinematics.

For the decays $\tilde{g}\to q\bar{q}\tilde{W}_i$, we have used the partial widths as given in Ref. 10 except that we have multiplied the result by the appropriate phase-space suppression factor to obtain an estimate of $\Gamma(\tilde{g}\to tb\tilde{W}_i)$. This ignores the coupling of the charged-Higgsino component of the chargino. With γ_R as defined in Ref. 10 the magnitude of this coupling is $f_t\sin\gamma_R$ ($f_t\cos\gamma_R$) for the decay to $\tilde{W}_+$ ($\tilde{W}_-$) compared to $g\cos\gamma_R$ ($g\sin\gamma_R$) from the gaugino component of the chargino. This approximation thus affects only the $\tilde{g}\to tb\tilde{W}_i$ mode by at most a factor of about 2 so that its affect on the branching fraction for the other decays is just a few per cent. Note also that unlike the decays to neutralinos (for which the decay via the Higgsino components may be the dominant tree-level contribution to the gluino width), the Higgsino com-

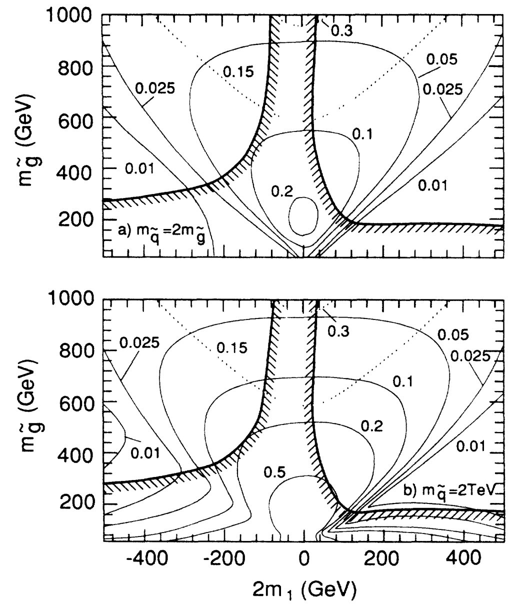

FIG. 1. Contours of fixed branching fraction for the one-loop decays $\tilde g \to g\tilde Z_i$ (solid lines) and the tree-level decays $\tilde g \to t\bar t\tilde Z_i$ (dotted lines) where the decays to all four neutralinos are included. In this plot, we have taken $v'/v=1/2$, $m_t=150$ GeV, and all squark masses to be equal with (a) $m_{\tilde q}=2m_{\tilde g}$ and (b) $m_{\tilde q}=2$ TeV. The cross-hatched area shows the region excluded by the constraints in Eq. (3.1) and denotes the region that can be probed in Z^0 decays at LEP I. Most of this region is already excluded.

ponents never dominate the decays to the chargino. For this reason, we have retained the Higgsino components in our computation of both $\tilde g \to t\bar t\tilde Z_i$ and (in case this channel is closed) also the $\tilde g \to b\bar b\tilde Z_i$ widths.

III. DECAY PATTERNS OF THE GLUINO

In this section, we present numerical results for the branching fractions of the gluino via its tree-level decays, $\tilde g \to q\bar q\tilde Z_i$ or $\tilde g \to q\bar q\tilde W_j$, or via the radiative decays $\tilde g \to g\tilde Z_i$ using the formulas of Sec. II. We plot in Fig. 1 the contours of constant branching fraction for the sum of all the radiative decay modes (solid lines) and the sum of the $\tilde g \to t\bar t\tilde Z_i$ modes (dotted lines) in the plane of $2m_1$ vs $m_{\tilde g}$, for $m_t=150$ GeV, $v'/v=\frac{1}{2}$, and with all squarks degenerate in mass. We will consider these as default parameters unless otherwise specified. In Fig. 1(a) we take $m_{\tilde q}=2m_{\tilde g}$, and in Fig. 1(b), we use $m_{\tilde q}=2$ TeV.

The cross-hatched region indicates the parameter space for which either

$$m_{\tilde W_-} < 45 \text{ GeV} , \tag{3.1a}$$

or

$$\Gamma(Z\to\tilde Z_1\tilde Z_1) > 20 \text{ MeV} , \tag{3.1b}$$

or

$$B(Z\to\tilde Z_i\tilde Z_j) > 10^{-6}, \quad i \text{ and } j\neq 1 . \tag{3.1c}$$

Chargino masses smaller than 45 GeV have nearly been excluded by the experiments[4] at LEP. Also, the measurement of the peak hadronic cross section at LEP already implies an upper limit[24] on the nonstandard model invisible width of the Z^0 boson of 38 MeV (under the assumption that the hadronic decay width of the Z^0 is as given by the SM). Further, the ALEPH experiment[25] has recently excluded the decays $Z\to\tilde Z_1\tilde Z_2,\tilde Z_2\tilde Z_2$ for branching fractions $\gtrsim 10^{-4}$ from a nonobservation of $\not p_T$ events. These bounds are expected to improve with the rapidly increasing size of the Z^0 decay sample, and Eq. (3.1) represents an estimate of the bounds that may be attained at LEP with a sample of a few $\times 10^6$ Z^0 events. The cross-hatched area, therefore, represents roughly the region[26] of SUSY parameter space that can be probed at LEP I.

Even in the region of parameter space that cannot be probed at LEP I, the branching fraction for the sum of the loop decays exceed $\tilde g \to g\tilde Z_i$ can be as large as 50%. For typical values of parameters, the loop decays exceed 5–10 %. To understand the behavior of Fig. 1, we must examine the structure of the coupling constant (2.3) for gluino radiative decays. If up-type ($T_3=+\frac{1}{2}$) and down-type ($T_3=-\frac{1}{2}$) left squarks are mass degenerate, then the terms in Eqs. (2.3) and (2.4) proportional to $v_3^{(i)}$ cancel. Further, if left and right squarks are mass degenerate, terms proportional to $v_4^{(i)}$ cancel as well, leaving no dependence on the gaugino components of the neutralino in (2.3). The cancellation of the gaugino components follows from the tracelessness of the diagonal generators of the electroweak gauge group. This yields the result of Haber and Kane,[18] that the decay $\tilde g \to g\tilde \gamma$ vanishes for degenerate squarks. Thus, for degenerate squarks, we see that the only significant contribution to the radiative decay comes from top loops.

The last term of Eq. (2.3) is proportional to the $\tilde h$-Higgsino component ($v_1^{(i)}$) of the neutralino $\tilde Z_i$, and is nonvanishing unless $v_1^{(i)}=0$. For $|2m_1|\simeq 0$ and large $m_{\tilde g}$, $\tilde Z_1$, and $\tilde Z_2$ are both mostly Higgsino-like and relatively light. For larger values of $|2m_1|$, the Higgsino content of the lighter neutralino is reduced, whilst that of the heavier ones increases. This explains why the radiative decay widths in Fig. 1 decrease with increasing $|2m_1|$ if the gluino mass is fixed.

Also shown by dotted lines in Fig. 1 are contours of 15% and 30% branching fractions for $\tilde g \to t\bar t\tilde Z_i$. This decay mode is only substantial for very heavy gluinos ($m_{\tilde g}\gtrsim 600$ GeV). For small values of $|2m_1|$, where $\tilde Z_1$ and $\tilde Z_2$ are dominantly Higgsino, the decays $\tilde g \to t\bar t\tilde Z_{1,2}$ mostly occur via the Yukawa couplings of the $t\bar t$ system to the Higgsino $\tilde h$. For $v'/v<1$ and small $|2m_1|$,$\tilde Z_2$ contains the larger fraction of $\tilde h$, so the majority of these decays are $\tilde g \to t\bar t\tilde Z_2$. For larger values of $|2m_1|$, these decays mainly occur via the gaugino components of the neutralino whereas those to the heavier neutralinos occur via the Yukawa coupling and no particular mode is dominant.

To see what is happening more clearly, we show in Fig. 2 a slice out of Fig. 1(b), for $2m_1=100$ GeV. We show the total radiative decay rate, along with its individual

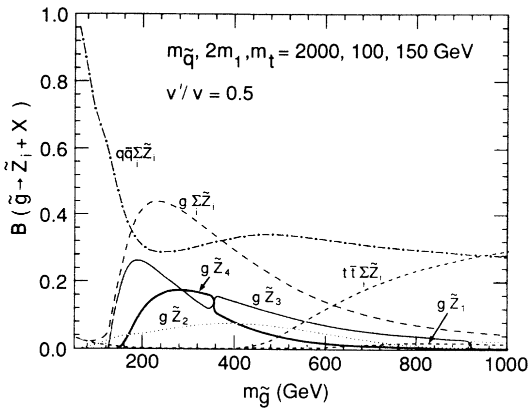

FIG. 2. A separation of the various components of the radiative decays shown in Fig. 1(b) for $2m_1=100$ GeV as a function of $m_{\tilde{g}}$. Also shown in the branching fraction for the tree-level decays into neutralinos summed over five quark flavors (dotted-dashed) and for the decay $\tilde{g}\rightarrow t\bar{t}\tilde{Z}_i$ (dashed), summed over all neutralinos. We have not shown the decays $\tilde{g}\rightarrow q\bar{q}\tilde{W}_i$ for clarity.

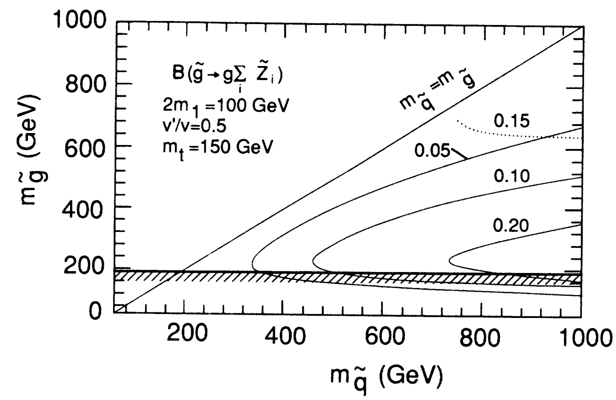

FIG. 3. Contours of the total branching fraction for gluino loop decays in the $m_{\tilde{q}}$ vs $m_{\tilde{g}}$ plane for degenerate squarks. The other parameters are as shown in the figure. Also shown is a contour for $\tilde{g}\rightarrow t\bar{t}\tilde{Z}_i$ summed over neutralinos. For $m_{\tilde{g}}>m_{\tilde{q}}$, the gluino decays via $\tilde{g}\rightarrow q\tilde{q}$. The cross-hatched region denotes the region to be explored at LEP I as given by Eq. (3.1).

components into the different neutralinos. The curve labeled $q\bar{q}\sum\tilde{Z}_i$ sums over all four neutralinos and five quark flavors. The contribution $\tilde{g}\rightarrow t\bar{t}\sum\tilde{Z}_i$ is shown separately, while decay rates to charginos are omitted for clarity. For small values of $m_{\tilde{g}}$, only decays to the lightest neutralino are allowed, and since it is a gaugino, it occurs via tree rather than loop decays. For the parameter choice in the figure, the radiative decay becomes dominant for $m_{\tilde{g}}\simeq 250$ GeV. This is partly due to the large squark mass chosen, but even $m_{\tilde{q}}=1$ TeV results in radiative branching fractions of 30%. For smaller values of $m_{\tilde{g}}$, the heavier neutralinos are dominantly Higgsino so that the decay $\tilde{g}\rightarrow g\tilde{Z}_{3,4}$ constitute the bulk of the loop decays. As the gluino mass increases the $\tilde{g}\rightarrow g\tilde{Z}_2$ gains at the expense of the $\tilde{Z}_4$ mode, but the radiative decays are less important. Finally, we see that for very heavy gluinos the decay $\tilde{g}\rightarrow t\bar{t}\tilde{Z}_i$ is very significant. Near the threshold for this decay the neutralinos $\tilde{Z}_3$ and $\tilde{Z}_4$ contain large $\tilde{h}$-Higgsino components so that this decay is suppressed either by phase space (for $i=3,4$) or small mixing angles. As $m_{\tilde{g}}$ becomes very large compared to $|2m_1|$ the lighter neutralinos contain the large Higgsino components so that the decays $\tilde{g}\rightarrow t\bar{t}\tilde{Z}_{1,2}$ are essentially unsuppressed. Of course, the decays to the neutralinos $\tilde{Z}_3$ and $\tilde{Z}_4$ which for large $m_{\tilde{g}}$ are dominantly the U(1) and SU(2) gaugino, respectively, are present via the gauge couplings. Note, however, that the top-quark Yukawa coupling that enters the Lagrangian (2.7), for $m_t=150$ GeV, is larger than even the SU(2) gauge coupling.

We plot in Fig. 3 the summed gluino loop decay branching fractions in the $m_{\tilde{q}}$ vs $m_{\tilde{g}}$ plane, for $m_{\tilde{q}}>m_{\tilde{g}}$ and $2m_1=100$ GeV, and the other parameters as in Fig. 1. For $m_{\tilde{q}}<m_{\tilde{g}}$, the decay $\tilde{g}\rightarrow\tilde{q}q$ dominates. The branching fraction for the radiative gluino decay is most substantial when $m_{\tilde{q}}\gg m_{\tilde{g}}$ as long as the gluino mass is not too light. Note that this branching fraction can easily exceed the $\sim 20\%$ maximum shown here if $|2m_1|$ is smaller.

To understand the increase of the radiative decay branching ratio for heavy squarks, we note that[16,17,20] the function I [given in Eq. (2.6b)] that enters the last term of Eq. (2.3) has the limit, as $m_{\tilde{t}}\rightarrow\infty$,

$$I\rightarrow\frac{1}{m_{\tilde{t}}^2}\left[1-\ln\frac{m_{\tilde{t}}^2}{m_t^2}\right] . \tag{3.2}$$

The prefactor $1/m_{\tilde{t}}^2$ provides the same suppression in the amplitude as the squark propagator in the tree-level three-body decays; the logarithm in Eq. (3.2), of course, enhances the radiative decay relative to the three-body tree decays when $m_{\tilde{q}}$ is large compared to $m_{\tilde{g}}$ and m_t, and accounts for the behavior of Fig. 3.

The loop decay rate depends sensitively on the top-quark mass via the Yukawa coupling as well as the explicit factor of quark mass in Eq. (2.3). This dependence is shown in Fig. 4 for various choices of SUSY parameters: in general, a heavier top mass enhances the loop decays. This enhancement can be very substantial depending on m_t so that the loop decays may be important even for gluinos in the mass range accessible at the Fermilab Tevatron.

In Fig. 5, we show the dependence of the loop branching fraction on variations in the ratio of Higgs-field VEV's, v'/v, for $m_{\tilde{g}}=150$ GeV, and 400 GeV. We see that there is a strong dependence on v'/v, especially when $m_{\tilde{g}}$ is small. The loop branching fraction is largest for $v'/v=1$. As v'/v is reduced from 1, the $\tilde{h}$ component (which couples to the top system) of the lighter neutralino is reduced while that of the heavier ones is increased. Furthermore, this reduces $m_{\tilde{W}_-}$ and $m_{\tilde{Z}_2}$, which allows significantly more phase space for the three-body decays.

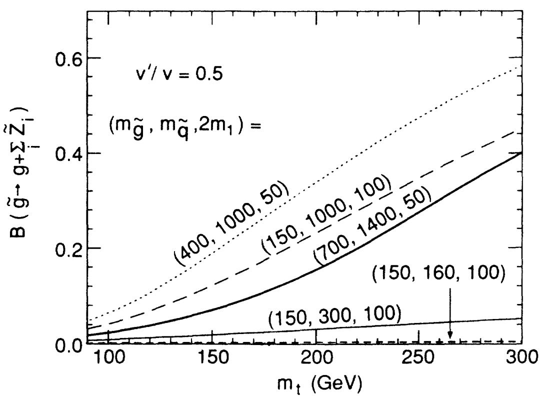

FIG. 4. The variation of the total branching fraction for the radiative decay of the gluino as a function of top-quark mass for degenerate squarks, and $v'/v = 1/2$ for representative values of SUSY parameters not excluded by the constraints (3.1).

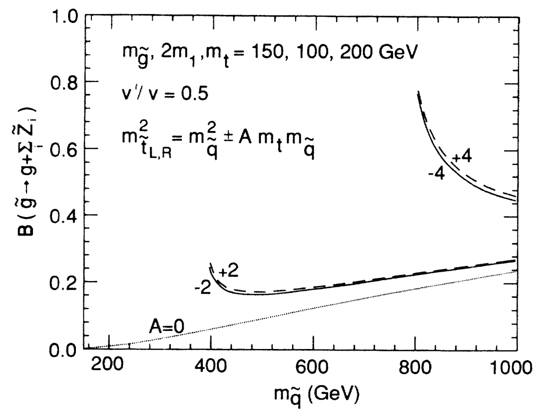

FIG. 6. The dependence of the total gluino radiative branching fraction on the splitting between the $\tilde t$-squark masses as parametrized by Eq. (3.3). For $A = 0$, the squarks are degenerate.

Thus, the loop decays of the gluino (which occur only via the $\tilde h$ component of the neutralino) tend to be largest when $v'/v = 1$ and get strongly suppressed by phase space if the gluino is light. For the MSSM, null searches for Higgs bosons at LEP exclude[27,28] $v'/v > 0.77$.

Up to now, we have confined our attention to the case of degenerate squarks and assumed that the scalar partners of chiral fermions are mass eigenstates. This is a good assumption in all models with a common scalar mass at the unification scale since, in the absence of strong Yukawa interactions, the dominant renormalization-group evolution of the squark mass is due to QCD interactions which are common to each flavor and handedness of squark. Since the t quark is known to be heavy,[22] soft SUSY-breaking terms can lead to a breaking of this degeneracy via $\tilde t_L$-$\tilde t_R$ mixing induced by the top-quark Yukawa coupling.[23] The resulting mass matrix depends on three unknown parameters.

Our interest in this stems from the earlier observation that if the squarks are not all degenerate the gaugino components of the neutralino do not exactly cancel so that their contribution can alter the analysis presented thus far. Note that only the third generation enters in

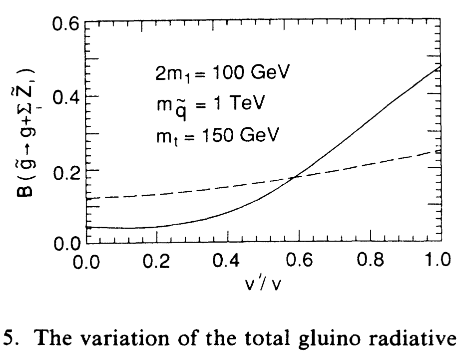

FIG. 5. The variation of the total gluino radiative branching with v'/v, the ratio of the Higgs-field vacuum expectation values, for degenerate squarks. The region $0.77 < v'/v < 1$ is excluded by the LEP constraints on the neutral-Higgs-boson mass, whereas Eq. (3.1) excludes $v'/v < 0.62$ for $m_{\tilde g} = 150$ GeV.

this analysis. In view of the theoretical uncertainties in the t-squark mass matrix, we parametrize the masses for the top-squark eigenstates by

$$m^2_{\tilde t_L} = m^2_{\tilde q} + A m_t m_{\tilde q}, \quad m^2_{\tilde t_R} = m^2_{\tilde q} - A m_t m_{\tilde q} . \tag{3.3}$$

In writing Eq. (3.3), we have assumed the extreme case that $\tilde t_L$ and $\tilde t_R$ are mass eigenstates, but that $m_{\tilde t_L} \neq m_{\tilde t_R}$. Further, $m_{\tilde q}$ is the common squark mass of the partners of the lighter quarks and A is a parameter ~ 1.

The resulting branching fraction for the radiative decay of the gluino is shown in Fig. 6 as a function of $m_{\tilde q}$ for the cases when (i) the squarks are degenerate, $A = 0$, (ii) $A = \pm 2$, which results in modest and perhaps realistic splitting in the stop system, and (iii) $A = \pm 4$. We have chosen $m_t = 200$ GeV in this figure. We see that the $A = \pm 2$ curves differ from the degenerate-squark case by only a few percent for most values of $m_{\tilde q}$, until $m_{\tilde q}$ decreases to 400 GeV, where $m_{\tilde t_R} \to 0$. In this case, the mass scale in this particular loop amplitude is m_t and not $m_{\tilde q}$, so this amplitude (unlike the tree-level amplitude), is unsuppressed by a large mass in the propagator. Similar behavior sets in even quicker in the $A = \pm 4$ case as shown since the top squark, in this case, is even lighter for the same value of $m_{\tilde q}$.

IV. SUMMARY AND CONCLUDING REMARKS

In this paper, we have made a detailed study of all the decay modes of the gluino within the framework of the minimal supersymmetric model.[3] In our computation, we have included all the tree-level decays $\tilde g \to q\bar q \tilde Z_i$ and $q\bar q \tilde W_i$ that are mediated by virtual squarks via the gaugino components of the charginos ($\tilde W_i$) or neutralinos ($\tilde Z_i$) using the minimal model as a guide to masses and mixing angles. For the decays $\tilde g \to t\bar t \tilde Z_i$ and $\tilde g \to b\bar b \tilde Z_i$, we have recomputed the partial widths [Eq. (2.8)–(2.10)] including the effects of the top- (bottom-) quark Yukawa coupling so that this decay can also occur via the $\tilde h$-Higgsino

($\tilde{h}'$-Higgsino) component of the neutralino. Finally, we have made an independent computation of the partial width for the radiative decay $\tilde{g}\to g\tilde{Z}_i$ for arbitrary values of neutralino mass and mixing angles and without making any approximation for the squark mass or neutralino masses [see Eqs. (2.2)–(2.6)].

Section III contains our results for the branching fractions for the various decay modes as a function of the unknown parameters ($m_{\tilde{g}}$, $m_{\tilde{q}}$, $2m_1$, v'/v, A, and m_t) defined in the text. We have seen that the radiative decay of the gluino can be a very significant fraction (up to 50%) of all gluino decays depending on the values of various parameters, and can be large even in regions of parameter space not accessible to LEP I. Our main results are summarized in Figs. 1–6 where we have shown the radiative gluino decay rate as functions of various combinations of the above six parameters. Typically, we find that the branching fraction for the radiative decay (i) is largest for small values of $|2m_1|$ (Fig. 1), (ii) increases with $m_{\tilde{q}}$ (Fig. 3) due to logarithms that enter the calculation of the loop graphs, (iii) increases with m_t since this decay is dominantly mediated by Yukawa interaction of the t quark (Fig. 4), (iv) depends on the ratio v'/v, especially for light gluinos (Fig. 5), and (v) can be significantly enhanced if the splitting between two top squarks is large so that one is considerably lighter than the other squarks (Fig. 6).

The large value of this two-body decay mode of the gluino, of course, has an impact on its expected signatures at hadron colliders. The obvious effect is to make gluino decays look more squark-like. Since the gluino has no two-body modes into charginos, these decays superficially resemble those of right-handed squarks and so will have lower jet multiplicity than is usually assumed. There is one important difference though, because unlike the squark case, the decay $\tilde{g}\to g\tilde{Z}_1$ is usually suppressed due to neutralino mixing. This is because for small values of $|2m_1|$ for which the lighter neutralinos are Higgsino-like, the $\tilde{h}$-Higgsino content of $\tilde{Z}_1$ is suppressed by v'/v relative to its $\tilde{h}'$ content. Thus unless v'/v is very close to its upper bound[27] of about 0.77 (in which case the discovery of the neutral Higgs boson at LEP should be imminent), the direct two-body decay rate of the $\tilde{g}$ to the LSP is suppressed by $(v'/v)^2$. This, of course, means that on an average, the $\not{p}_T$ in these events will be less than that due to direct two-body decays to the LSP. Of course, the same would apply to $\tilde{q}_R$'s if, in fact, the LSP is dominantly Higgsino-like.

The loop decays of the gluino can add new, interesting signatures for supersymmetry. An example that was recently pointed out[21] is to search for gluinos in events containing $Z^0+\not{p}_T+$jets, at the Fermilab Tevatron collider. Such events can come from $\tilde{g}\to g\tilde{Z}_4$ with $\tilde{Z}_4\to Z^0+\tilde{Z}_1$. Another interesting new signature for gluinos at hadron supercolliders results from the pair production of gluinos, with each gluino decaying via $\tilde{g}\to t\bar{t}\tilde{Z}_i$. In this case, one might look for four high p_T top quarks, or the four W bosons resulting from the top-quark decay. For instance, for $m_{\tilde{g}}=600$ GeV, and $m_{\tilde{q}}=2m_{\tilde{g}}$, one expects $\sim 2\times 10^6$ gluino pair events per year at the SSC. This would yield about 4×10^4 gluino pair events with four high $p_T W$'s, and if we require each W to decay to an e or a μ, we would get ~ 100 four-lepton events. Whether or not gluinos can be detected in this channel can only be answered after a detailed study.

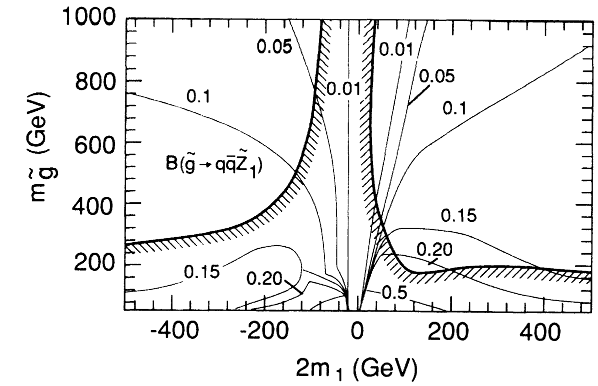

FIG. 7. Contours of fixed branching fraction for the decay mode $\tilde{g}\to q\bar{q}\tilde{Z}_1$ which has been assumed to be 100% in many analyses. The cross-hatched region is based upon Eq. (3.1). Squarks are assumed to be mass degenerate. We have taken $m_{\tilde{q}}=2m_{\tilde{g}}$, $v'/v=1/2$, and $m_t=150$ GeV.

In summary, we have seen that even in the minimal model, the decay patterns of gluinos are considerably more complex than those incorporated in many theoretical[9–14] and experimental[7,8] analysis of gluino signatures. If the minimal model is a guide for masses and mixings, then any analysis based on the assumption that the gluino always decays via $\tilde{g}\to q\bar{q}\tilde{Z}_1$ is highly suspect. This is illustrated in Fig. 7, where we have shown contours of constant branching fraction for the mode $\tilde{g}\to q\bar{q}\tilde{Z}_1$ in the $2m_1$ vs $m_{\tilde{g}}$ plane with v'/v fixed at 0.5 and $m_{\tilde{q}}=2m_{\tilde{g}}$. The cross-hatched area shows the region of this plane that can be excluded by a study of Z^0 decays at LEP I. As mentioned,[26] the bulk of this region is already excluded. In only a tiny region of parameter space of interest to hadron colliders does this branching ratio exceed 20%. It is clear that any serious search for supersymmetry at the Tevatron or hadron supercolliders should take into account all the various possible decay modes of the gluino.

ACKNOWLEDGMENTS

We thank D. Dicus, M. Drees, H. Haber, K. Hikasa, and E. Ma for discussions. This work was supported in part by the U.S. Department of Energy, Grants Nos. DE/GF05/85ER40215 and DE-AM03-76SF00235.

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