Light scalar top quark at e^+e^- colliders

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(Received 13 April 1987)

We examine the production and decay of a supersymmetric scalar top $\tilde{t}_1$ at e^+e^- colliders such as TRISTAN at KEK, the Stanford Linear Collider, and LEP at CERN, based on supersymmetric models with soft breaking. Existing bounds for supersymmetric particles imply that $\tilde{t}_1$ in the energy range of these colliders would be lighter than other scalar quarks and than the top quark. We make a careful study of available $\tilde{t}_1$ decay modes in this scenario. The flavor-changing decay $\tilde{t}_1 \to c\tilde{\gamma}$ is found to proceed through the misalignment of the scalar-quark and quark mixing matrices. If the three-body decay $\tilde{t}_1 \to bl^+\tilde{\nu}$ is kinematically accessible it dominates the two-body decay. The decay $\tilde{t}_1 \to b\nu\tilde{l}^+$ is even larger due to different chirality structure, but is likely to suffer from a large phase-space suppression. Four-body decays are negligible. The $\tilde{t}_1$ pair production cross section including the γ and Z contribution and the left-right mixing is presented.

I. INTRODUCTION

The search for supersymmetry (SUSY) is one of the most active fields of new particle hunting. Various limits[1,2] have been obtained from e^+e^- and $p\bar{p}$ colliders. Strongly interacting SUSY particles can be copiously produced in $p\bar{p}$ interactions. The bounds[3] $m_{\tilde{q}} > 60$ GeV, $m_{\tilde{g}} > 70$ GeV were obtained assuming that five flavors and both chiralities of scalar quarks were degenerate. This (approximate) degeneracy is predicted by the minimal supergravity models[4-6] and is required by K^0-$\bar{K}^0$ phenomenology as well.[7] This requirement is, however, not very stringent for the third-generation scalar quarks because of the small intergeneration mixing. Current models allow for the possibility[8] that one of the scalar partners of the top quark ($\tilde{t}_1$) is lighter than other scalar quarks and also than the top quark. Implications of this scenario have been discussed in the literature.[8-12]

In this paper we update and refine these previous studies and clarify some existing confusion. We focus on the case in which the $\tilde{t}_1$ can be produced at the forthcoming e^+e^- colliders [TRISTAN at KEK, the Stanford Linear Collider (SLC), LEP at CERN], although most of our results should remain applicable if the $\tilde{t}_1$ is heavier. In this case, the existing limits for SUSY particles restrict the possible decay modes of the t_1 to a few. We compare the flavor-changing two-body decay $\tilde{t}_1 \to c\tilde{\gamma}$ with the possible three-body decays $\tilde{t}_1 \to bl^+\tilde{\nu}$, $\tilde{t}_1 \to b\nu\tilde{l}^+_L$, and also with the four-body decay $\tilde{t}_1 \to bl^+\nu\tilde{\gamma}$. Although these decay rates have been estimated, we find several errors in the calculation[11] and reevaluate the widths. The evaluation of the two-body mode $\tilde{t}_1 \to c\tilde{\gamma}$ reveals an interesting dynamical structure of the supersymmetric standard model with soft-breaking terms.

In Sec. II we review the mass matrix of scalar quarks and show that the lighter $\tilde{t}_1$ may have a mass smaller than the top and other scalar quarks if the left-right mixing is reasonably large. In Sec. III we review the existing experimental bounds for SUSY particles. These limits imply that if the $\tilde{t}_1$ is to be found at the newly built/forthcoming e^+e^- colliders (TRISTAN, SLC, LEP), the above scenario should be realized. Possibly lighter SUSY particles than the $\tilde{t}_1$ are then the photino, scalar neutrinos, and scalar leptons. (In addition, a very light gluino may still be allowed.) Available decay modes of the $\tilde{t}_1$ are thus quite restricted and different from those of other scalar quarks.

Section IV includes the study of these $\tilde{t}_1$ decay modes and comprises the main body of the paper. In Sec. IV A the flavor-changing two-body decay $\tilde{t}_1 \to c\tilde{\gamma}$ is studied. The decay can be regarded as occurring at the one-loop order. The examination of the $N=1$ supergravity models shows that the leading contribution to the decay actually comes from the large logarithm associated with the divergence due to the soft-supersymmetry-breaking terms. This can alternatively be interpreted as the mismatch of the quark and scalar-quark mass matrices in the flavor space, which introduces a small admixture of scalar charm into the scalar-top mass eigenstate $\tilde{t}_1$. The diagonalization of the scalar-quark states, including this flavor mixing, is worked out and the decay rate is estimated.

In Sec. IV B, the three-body decays $\tilde{t}_1 \to bl^+\tilde{\nu}$ and $\tilde{t}_1 \to b\nu\tilde{l}^+_L$ are calculated. These modes occur at the tree order and are dominant if kinematically allowed. The latter mode is seen to be much larger than the former if there is no phase-space suppression. This reflects the different chirality structure of the two modes and can be understood using effective-Lagrangian analysis. The phase-space reduction of the rates is much more severe for the three-body decays than the two-body decay. In Sec. IV C, we present an estimate of four-body decays which, however, turns out to be negligible.

In Sec. V, the production cross section for the process $e^+e^- \to \tilde{t}_1\bar{\tilde{t}}_1$ is presented. Both the photon and Z exchange are taken into account. The mixing of $\tilde{t}_L$ and $\tilde{t}_R$ induces an off-diagonal coupling of the Z boson to $\tilde{t}_1\bar{\tilde{t}}_2$.

Section VI includes comments and the conclusion. In the following we adopt mainly the notation of Gunion and Haber.[13]

II. THE SCALAR-TOP MASS MATRIX

The mass matrix of a flavor of the scalar quark takes the following form[8,6,13] (we ignore the flavor mixing for a while):

$$-\mathcal{L}_m=(\tilde{q}_L^* \ \ \tilde{q}_R^*)\begin{bmatrix} m_{\tilde{q}_L}{}^2 & am_q \\ a^*m_q & m_{\tilde{q}_R}{}^2 \end{bmatrix}\begin{bmatrix}\tilde{q}_L \\ \tilde{q}_R\end{bmatrix}, \tag{1}$$

where for a charge-$\frac{2}{3}$ scalar quark

$$m_{\tilde{q}_L}{}^2=\tilde{M}_Q{}^2+m_Z{}^2\cos2\beta(\tfrac{1}{2}-\tfrac{2}{3}\sin^2\theta_W)+m_q{}^2 , \tag{2a}$$

$$m_{\tilde{q}_R}{}^2=\tilde{M}_U{}^2+m_Z{}^2(\cos2\beta)\tfrac{2}{3}\sin^2\theta_W+m_q{}^2 , \tag{2b}$$

$$a=\mu\cot\beta+A_u^*\tilde{M} . \tag{2c}$$

Here, $\tilde{M}_Q{}^2$, $\tilde{M}_U{}^2$ are the soft-supersymmetry-breaking mass terms for the left-handed and right-handed up scalar quarks, $\tan\beta=v_2/v_1$ is the ratio of the vacuum expectation values of the two Higgs fields, H_1 (H_2) giving masses to down-type quarks and leptons (up-type quarks). μ is the supersymmetric Higgs-boson mass term, and $A_u\tilde{M}$ is the coefficient of the dimension-three soft-breaking term proportional to the superpotential. Hereafter, we neglect CP violation and assume that a is real. At the Planck mass scale we expect

$$\tilde{M}_Q{}^2=\tilde{M}_U{}^2=\tilde{M}^2 .$$

The second term of Eqs. (2a) and (2b) comes from the $SU(2)\times U(1)$ D term. For the first two generations, the left-right mixing is negligible[14] and the mass eigenstates are practically $\tilde{q}_L$ and $\tilde{q}_R$ with masses $\sim\tilde{M}=O(m_W)$. However, for the scalar-top sector the nondiagonal term may not be small if $m_t\sim\tilde{M}$, resulting in a large mixing of the two weak eigenstates and a large splitting of the two mass eigenstates. One of the states may become lighter than the top and also than all the other scalar quarks if the mixing is sufficiently large.

The mass matrix $\mathcal{M}_{\tilde{t}}{}^2$ in Eq. (1) is diagonalized by a rotation

$$R\mathcal{M}_{\tilde{t}}{}^2R^T=\begin{bmatrix} m_{\tilde{t}_1}{}^2 & 0 \\ 0 & m_{\tilde{t}_2}{}^2\end{bmatrix} \tag{3}$$

with

$$R=\begin{bmatrix}\cos\theta_t & -\sin\theta_t \\ \sin\theta_t & \cos\theta_t\end{bmatrix}, \tag{4}$$

$$\tan2\theta_t=\frac{2am_t}{m_{t_R}{}^2-m_{\tilde{t}_L}{}^2} , \tag{5}$$

and

$$(m_{\tilde{t}_1,\tilde{t}_2})^2=\tfrac{1}{2}\{m_{\tilde{t}_L}{}^2+m_{\tilde{t}_R}{}^2\mp[(m_{\tilde{t}_L}{}^2-m_{\tilde{t}_R}{}^2)^2+4a^2m_t{}^2]^{1/2}\} . \tag{6}$$

We take $m_{\tilde{t}_1}{}^2<m_{\tilde{t}_2}{}^2$. As an illustrative example we set $\tilde{M}_Q=\tilde{M}_U=\tilde{M}=m_W$, $\cos2\beta=\mu=0$, to obtain

$$m_{\tilde{t}_{1,2}}{}^2=m_W{}^2+m_t{}^2\mp|A_t|m_Wm_t .$$

We have $m_{\tilde{t}_1}<m_t$ if $|A_t|>m_W/m_t$. Of course, this is a very simplified case, but in general there is a substantial region in the parameter space where $m_{\tilde{t}_1}<m_t$ is satisfied. Also $m_{\tilde{t}_1}<m_{\tilde{q}}$ holds for $|A_t|>m_t/m_W$ which is a weaker condition if $m_t<m_W$.

III. SUSY-PARTICLE SPECTRUM LIMITS AND ITS IMPLICATION FOR $\tilde{t}_1$ DECAY MODES

Bounds on the scalar-quark mass for various decay possibilities have been obtained[1] at the DESY storage ring PETRA, assuming only one species of scalar quark. This bound applies to the $\tilde{t}_1$: $m_{\tilde{t}_1}>20$–21 GeV. We consider the mass range above this limit, for which a $\tilde{t}_1$ pair could be produced at e^+e^- colliding machines such as TRISTAN, SLC, or LEP. The UA1 Collaboration at the CERN $p\bar{p}$ collider has given bounds[3] on some SUSY-particle masses. In particular, from the analysis of monojet events they obtain $m_{\tilde{q}}>60$ GeV and $m_{\tilde{g}}>70$ GeV, barring the light-gluino scenario[15,16] (say $m_{\tilde{g}}\sim5$ GeV, $m_{\tilde{q}}\gtrsim100$ GeV) on which we will comment also. This bound was derived with the assumption that ten species of scalar quarks ($\tilde{u}_L,\tilde{u}_R,\tilde{d}_L,\ldots,\tilde{b}_R$) are degenerate. If only the $\tilde{t}_1$ is light, the production cross section of $\tilde{t}_1\bar{\tilde{t}}_1$ is less than one-tenth of that for the degenerate case, so any light $\tilde{t}_1$ we are considering could escape the bound. It has also been pointed out[17] that the monojet data should be able to exclude the weak gaugino ($\tilde{w},\tilde{z}$) masses less than ~35 GeV if the photino mass is not too close to those of $\tilde{w}$ and $\tilde{z}$. Thus, we assume that $\tilde{t}_1$ is lighter than all other scalar quarks and gauginos (except for the photino).

Limits on scalar leptons are derived from e^+e^- collider data. The single photon search[2,18] set a limit $m_{\tilde{e}}>65$ GeV (for degenerate scalar electrons) for vanishing photino mass, but the bound becomes much weaker for a massive photino. In fact for $m_{\tilde{\gamma}}\gtrsim12$ GeV the best limit comes from the direct production search,[1,2] $m_{\tilde{e}}\gtrsim22$ GeV. Other charged scalar leptons are expected to have similar masses as the scalar electron. (Direct limits[1] are $m_{\tilde{\mu}}>20$ GeV and $m_{\tilde{\tau}}>19$ GeV.) Limits for scalar neutrinos are not very significant. Neutrino-counting-type experiments are sensitive to $\tilde{\nu}$, but the $p\bar{p}$ collider limits[19] are not good enough to exclude light $\tilde{\nu}$'s. The e^+e^- single-γ limit[2,18] excludes the region $\lesssim10$ GeV only when the W gaugino is rather light ($\lesssim50$ GeV). In addition, the UA1 Collaboration has claimed[20] a limit for the scalar electron and electron scalar neutrino from the decay $W\to\tilde{e}\bar{\tilde{\nu}}_e$, $\tilde{e}\to e\tilde{\gamma}$. For equal masses the limit is $m_{\tilde{e}}\sim m_{\tilde{\nu}_e}\gtrsim25$ GeV, and for $m_{\tilde{\nu}_e}\sim0$ they give $m_{\tilde{e}}\gtrsim33$ GeV.

In the minimal supergravity models, the scalar leptons can be much lighter than the scalar quarks if the gaugino mass term is substantial compared to the "gravitino" mass term. In addition, the scalar neutrinos can be lighter than their charged counterparts when the two vacuum expectation values are not equal ($v_2 > v_1$). There is a model which predicts massless scalar neutrinos at the lowest order.[21]

Under the simplest grand unification condition (equal gaugino masses at the unification scale) the gluino mass bound[22] $m_{\tilde g} \gtrsim 70$ GeV implies $m_{\tilde\gamma} \gtrsim 10$ GeV, for which the scalar-electron mass bound from the single-photon search becomes quite weak.

Thus, the scalar top in the mass range 20–50 GeV which is reachable by the near-future e^+e^- colliding facilities is likely to be lighter than both the t quark and other scalar quarks. In this case, supersymmetric particles, possibly lighter than $\tilde t_1$, are $\tilde\gamma$, $\tilde\nu$, and $\tilde l$. One of these particles must be in the final decay products of $\tilde t_1$ provided that the R parity is not broken (which is assumed throughout the paper). The study of the available $\tilde t_1$ decay modes is the subject of the next section.

IV. DECAY OF LIGHT SCALAR TOP QUARK

In this section we study the various decay modes of the $\tilde t_1$. We shall treat the cases in which (1) all the scalar leptons are heavier than $\tilde t_1$ and (2) scalar neutrinos are lighter than $\tilde t_1$. Throughout the paper we assume that the photino is light ($\ll m_W$), in which case, it is a good approximation to regard it as a mass eigenstate of the neutralino sector. The possible decay modes of the $\tilde t_1$ we consider in this paper are the following.

(1) $\tilde t_1 \to c\tilde\gamma$. Although this decay is flavor changing, it can occur via radiative corrections. The main contribution comes from the mismatch of the quark and scalar-quark mass matrices due to the soft-supersymmetry-breaking term. This decay is actually suppressed by the Glashow-Iliopoulos-Maiani mechanism, but the phase space is large since the final state is two body. The decay $\tilde t_1 \to u\tilde\gamma$ should be smaller because $|K_{ub}| \ll |K_{cb}|$ (K is the flavor-mixing matrix of quarks).

(2) $\tilde t_1 \to bl^+\tilde\nu_l$. This occurs at the tree level if the scalar neutrino is light. If the scalar-neutrino mass is closer to the parent mass, the rate is much suppressed due to the three-body phase space. When $m_{\tilde\nu} < m_{\tilde t_1} < m_b + m_{\tilde\nu}$, the decay is kinematically forbidden; however $\tilde t_1 \to (u,s)l^+\tilde\nu_l$ is still allowed, but with an additional suppression coming from the small flavor mixing. We will also discuss the decay $\tilde t_1 \to b\nu_l\tilde l_L^+$.

(3) $\tilde t_1 \to bl^+\nu\tilde\gamma$. This four-body decay proceeds at the tree level. $\tilde t_1 \to bu\bar d\tilde\gamma$ has a similar rate. This will be compared to the two-body decay (1). When the three-body decay (2) is kinematically allowed, it dominates the four-body decay since both occur at the tree level. If the other three-body decay $\tilde t_1 \to b\nu\tilde l$ is allowed, most parts of the four-body decay (3) actually correspond to this three-body decay followed by $\tilde l \to l\tilde\gamma$.

A. The decay $\tilde t_1 \to c\tilde\gamma$

The soft-supersymmetry-breaking scalar mass terms split the scalars from their superpartners, quarks and leptons. It has been noted[7] that the supersymmetric contribution to the K^0-$\bar K^0$ mixing requires $\sum_q K_{qs}^* \Delta m_{\tilde q}^2 K_{qd} / \langle m_{\tilde q}^2 \rangle \ll 1$: the scalar quarks should be approximately degenerate. The supersymmetry-breaking mechanism via the hidden sector in $N=1$ supergravity models provides the flavor-independent scalar mass term proportional to the unit matrix. This diagonality holds, however, at the Planck scale and is violated by radiative corrections since it is not protected by symmetry. The scalar mass terms at the weak scale may be obtained by solving the renormalization-group equations[23] for these mass parameters. Off-diagonal masses are induced by the Yukawa couplings and the diagonal entries become no longer equal. Of course, a part of this off-diagonal mass terms for scalar quarks is common to those induced in the quark mass matrix, and the physical flavor mixing is that which remains in the scalar-quark sector after diagonalizing the quark mass matrix. In other words, if the scalar-quark mass matrix cannot be diagonalized simultaneously as with the quark matrix, the physical "scalar-top" states have some admixture of scalar charm which induces the $\tilde t$-c-$\tilde\gamma$ coupling.

Let us discuss this induced mixing from a different point of view. If there is no tree $\tilde t$-c-$\tilde\gamma$ coupling (as in the theory renormalized at the Planck scale), the decay $\tilde t_1 \to c\tilde\gamma$ goes through the one-loop diagrams shown in Figs. 1 and 2. There are three types of diagrams: scalar self-energy, quark self-energy, and proper vertex. Each set of diagrams has logarithmic divergences. If the soft-supersymmetry-breaking terms are set to zero, all divergences are related to the wave-function renormalization and cancel after adding all diagrams. When the soft breaking is turned on, no new divergences are introduced in the latter two types of diagrams, but there appear additional divergences in the scalar self-mass diagrams. This divergence must be subtracted using a soft counterterm at the Planck scale. A large logarithm $\ln(M_P{}^2/m_W{}^2)$ remains after renormalization. Because of this factor (~ 80), the remaining (nonlogarithmic) part of the one-loop diagrams is not important.[24] We can also see that the two-loop leading logarithm is smaller for the process under consideration (due to the small b-quark Yukawa coupling), so the procedure to calculate the one-loop logarithmic contribution should give a reliable estimate. In the renormalization-group approach this corresponds to solving the renormalization-group equation by one iteration. Both methods should give the same result. We have checked that this is actually true.

Before presenting the result we comment on the existing estimate of the process. The authors of Ref. 11 claim that one of the vertex-type diagrams shown in Fig. 2 (the first vertex graph with $\tilde\chi$, b, and $\tilde b$ comprising the loop) is suppressed by only one power of m_b. However, they overlooked the fact that the graph needs the mixing of $\tilde b_L$ and $\tilde b_R$, giving another factor of m_b. We have checked

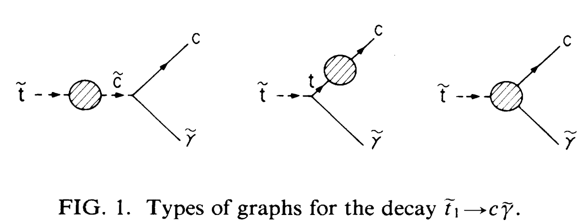

FIG. 1. Types of graphs for the decay $\tilde t_1 \to c\tilde\gamma$.

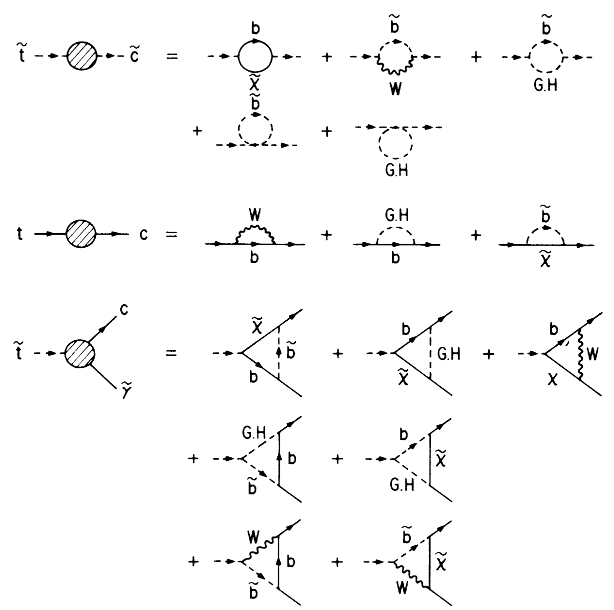

FIG. 2. Diagrams contributing to the decay $\tilde{t}_1 \to c\tilde{\gamma}$. Arrows represent the flow of the baryon number. G denotes the unphysical charged scalar needed in renormalizable gauges. We have adopted the 't Hooft–Feynman gauge in the calculation.

that there are no diagrams which give only one factor of m_b. This can be seen if we consider an appropriate supersymmetric limit where the answer must be zero. If there were a term with the b-quark loop which gives only one factor of m_b, it would be impossible to have cancellation since the scalar b loop cannot give an odd power of m_b. Thus, the finite part of the graphs is proportional to $m_b{}^2$ and is much smaller than the estimate of Ref. 11. (However the large logarithm in our estimate makes our resulting rate larger than that of Ref. 11.)

The logarithmic part gives the mixing of $\tilde{c}_L$ with $\tilde{t}_L$ and $\tilde{t}_R$. The $\tilde{c}_R$ does not mix with $\tilde{t}$'s in the approximation $m_c \to 0$. The mass matrix of these three scalars is in the left-right basis:

$$\mathcal{M}^2 = \begin{array}{c} \tilde{t}_L^* \\ \tilde{t}_R^* \\ \tilde{c}_L^* \end{array} \overset{\begin{array}{ccc} \tilde{t}_L & \tilde{t}_R & \tilde{c}_L \end{array}}{\begin{bmatrix} \multicolumn{2}{c}{\mathcal{M}_{\tilde{t}}{}^2} & \Delta_L \\ & & \Delta_R \\ \Delta_L^* & \Delta_R^* & m_{\tilde{c}_L}{}^2 \end{bmatrix}}, \tag{7}$$

where

$$\Delta_L = -\frac{g^2}{16\pi^2}\ln\left[\frac{M_P{}^2}{m_W{}^2}\right]\frac{K_{tb}^* K_{cb} m_b{}^2}{2m_W{}^2\cos^2\beta} \times(\tilde{M}_Q{}^2 + \tilde{M}_D{}^2 + \tilde{M}_{H_1}{}^2 + |A_d|^2 \tilde{M}^2), \tag{8a}$$

$$\Delta_R = -\frac{g^2}{16\pi^2}\ln\left[\frac{M_P{}^2}{m_W{}^2}\right]\frac{K_{tb}^* K_{cb} m_b{}^2}{2m_W{}^2\cos^2\beta} m_t A_d^* \tilde{M}. \tag{8b}$$

Here g is the SU(2) gauge coupling. Note that $\Delta_{L,R} \ll m_W{}^2$. We diagonalize (7) by first going to the $\tilde{t}_1$-$\tilde{t}_2$ mass basis

$$\begin{bmatrix} R & & 0 \\ & & 0 \\ 0 & 0 & 1 \end{bmatrix} \mathcal{M}^2 \begin{bmatrix} R^T & & 0 \\ & & 0 \\ 0 & 0 & 1 \end{bmatrix} = \begin{bmatrix} m_{\tilde{t}_1}{}^2 & 0 & \Delta_L\cos\theta_t - \Delta_R\sin\theta_t \\ 0 & m_{\tilde{t}_2}{}^2 & \Delta_L\sin\theta_t + \Delta_R\cos\theta_t \\ \Delta_L^*\cos\theta_t - \Delta_R^*\sin\theta_t & \Delta_L^*\sin\theta_t + \Delta_R^*\cos\theta_t & m_{\tilde{c}_L}{}^2 \end{bmatrix}, \tag{9}$$

then making perturbative diagonalization gives

$$\tilde{t}_1(\text{phys}) = \tilde{t}_1 + \epsilon\tilde{c}_L \tag{10}$$

with

$$\epsilon = -\frac{\Delta_L\cos\theta_t - \Delta_R\sin\theta_t}{m_{\tilde{c}_L}{}^2 - m_{\tilde{t}_1}{}^2}. \tag{11}$$

The decay $\tilde{t}_1 \to c\tilde{\gamma}$ proceeds through the $\tilde{c}_L$ component of the physical $\tilde{t}_1$. Using the photino interaction

$$\mathcal{L} = -\sqrt{2}ee_c(\bar{c}P_R\tilde{\gamma}\,\tilde{c}_L - \bar{c}P_L\tilde{\gamma}\,\tilde{c}_R) + \text{H.c.} \tag{12}$$

with $e_c = \frac{2}{3}$, $P_{L,R} = (1 \mp \gamma_5)/2$, we obtain the decay rate

$$\Gamma(\tilde{t}_1 \to c\tilde{\gamma}) = \tfrac{1}{2}\alpha e_c{}^2|\epsilon|^2 m_{\tilde{t}_1}\left[1 - \frac{m_{\tilde{\gamma}}{}^2}{m_{\tilde{t}_1}{}^2}\right]^2. \tag{13}$$

The decay is isotropic. Note that the decay time is far longer than the strong-interaction time scale, so a produced $\tilde{t}_1$ hadronizes into a $\tilde{t}_1$ hadron before decay. The lowest-lying states would be $(\tilde{t}_1\bar{u})$ and $(\tilde{t}_1\bar{d})$ with spin $\frac{1}{2}$, but they are not expected to be polarized.

As a numerical estimate, we may take

$$|\epsilon| = (3-10)\frac{g^2}{16\pi^2}\ln\left[\frac{M_P{}^2}{m_W{}^2}\right]\frac{K_{tb}^* K_{cb} m_b{}^2}{m_W{}^2} = (1-4)\times 10^{-4}.$$

We have used $|K_{tb}| = 1$, $|K_{cb}| = 0.05$, $m_b = 5$ GeV, and $m_W = 82$ GeV. This gives

$$\Gamma(\tilde{t}_1 \to c\tilde{\gamma}) = (0.3-3)\times 10^{-10} m_{\tilde{t}_1}\left[1 - \frac{m_{\tilde{\gamma}}{}^2}{m_{\tilde{t}_1}{}^2}\right]^2,$$

which corresponds to a lifetime of $10^{-15}-10^{-16}$ sec for $m_{\tilde{t}_1} \sim 20$ GeV.

If the gluino is lighter than the $\tilde{t}_1$, which is realized in the light-gluino scenario, the decay $\tilde{t}_1 \to c\tilde{g}$ dominates $\tilde{t}_1 \to c\tilde{\gamma}$ with the rate

$$\Gamma(\tilde{t}_1 \to c\tilde{g}) = \tfrac{2}{3}\alpha_s |\epsilon|^2 m_{\tilde{t}_1} \left[1 - \frac{m_{\tilde{g}}^2}{m_{\tilde{t}_1}^2}\right]^2 . \tag{14}$$

The decay rates for $\tilde{t}_1 \to u\tilde{\gamma}, u\tilde{\gamma}$ may be obtained by multiplying the rates for $c\tilde{\gamma}, c\tilde{g}$ by a factor $|K_{ub}/K_{cb}|^2$.

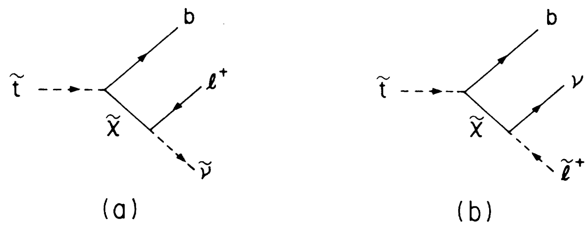

FIG. 3. The diagram for (a) $\tilde{t}_1 \to bl^+\tilde{\nu}$; (b) $\tilde{t}_1 \to b\nu\tilde{l}^+$.

B. The decay $\tilde{t}_1 \to bl^+\tilde{\nu}$

This three-body decay is allowed if $m_{\tilde{t}_1} > m_{\tilde{\nu}} + m_b$. The diagram for the decay is shown in Fig. 3(a). The matrix element for the decay $\tilde{t}_1 \to bl^+\tilde{\nu}_l$ is

$$\mathcal{M} = \sum_j \frac{g^2 V_{j1}^*}{(p_{\tilde{t}_1} - p_b)^2 - m_{\tilde{\chi}_j}^2} \left[\left[V_{ji}\cos\theta_t + \frac{m_t V_{j1}\sin\theta_t}{\sqrt{2} m_W \sin\beta} \right] \bar{u}(p_b) \not{p}_{\tilde{\nu}} P_L v(p_l) + \frac{m_b U_{j2}^* \cos\theta_t}{\sqrt{2} m_W \cos\beta} m_{\tilde{\chi}_j} \bar{u}(p_b) P_L v(p_l) \right] , \tag{15}$$

where $\tilde{\chi}_j$ $(j=1,2)$ is the chargino ($\tilde{W}$-gaugino–charged-Higgs-fermion) mass eigenstate, and U (V^*) is the mixing matrix for the right- (left-) handed charginos.[13] (Here the handedness refers to that of the positively charged fermions.) The first subscript of U, V corresponds to mass eigenstates $(j=1,2)$, and the second to weak eigenstates ($1 = \tilde{W}$ gaugino, $2 =$ Higgs fermion). We have neglected the charged-lepton mass and set $K_{tb}=1$.

The decay distribution can be calculated from Eq. (15). We show the distribution with the approximation $m_b = 0$, $m_{\tilde{t}_1} \ll m_{\tilde{\chi}_1} \ll m_{\tilde{\chi}_2}$, since there are many unknown parameters and the resulting expression for the general case is very complicated:

$$\frac{d\Gamma}{dz_b dz_l} = \frac{m_{\tilde{t}_1}}{256\pi^3} \sum_{\text{spins}} |\mathcal{M}|^2 = \frac{\alpha^2 |c|^2}{16\pi \sin^4\theta_W} \frac{m_{\tilde{t}_1}^5}{m_{\tilde{\chi}_1}^4} [(1-z_b)(1-z_l) - R_{\tilde{\nu}}] , \tag{16}$$

where z_b and z_l are scaled energy variables for b and l,

$$z_b = \frac{2p_{\tilde{t}_1} \cdot p_b}{m_{\tilde{t}_1}^2}, \quad z_l = \frac{2p_{\tilde{t}_1} \cdot p_l}{m_{\tilde{t}_1}^2} \tag{17}$$

with the range $z_b + z_l > 1 - R_{\tilde{\nu}}$ and $(1-z_b)(1-z_l) > R_{\tilde{\nu}}$,

$$R_{\tilde{\nu}} = m_{\tilde{\nu}}^2 / m_{\tilde{t}_1}^2 ,$$

and

$$c = |V_{11}|^2 \cos\theta_t + \frac{m_t}{\sqrt{2} m_W \sin\beta} V_{11}^* V_{12} \sin\theta_t . \tag{18}$$

The total rate is

$$\Gamma(\tilde{t}_1 \to bl^+\tilde{\nu}_l) = \frac{\alpha^2}{384\pi \sin^4\theta_W} \frac{m_{\tilde{t}_1}^5}{m_{\tilde{\chi}_1}^4} |c|^2 f(R_{\tilde{\nu}}) = \frac{G_F^2 m_{\tilde{t}_1}^5}{192\pi^3} \left[\frac{m_W}{m_{\tilde{\chi}_1}}\right]^4 |c|^2 f(R_{\tilde{\nu}}) , \tag{19}$$

where

$$f(R) = (1-R^2)(1-8R+R^2) - 12R^2 \ln R . \tag{20}$$

The function $f(R)$ has the property

$$f(0) = 1 ,$$

$$f(R) = \tfrac{2}{5}(1-R)^5 + O((1-R)^6) \quad (R \to 1) .$$

If the scalar-neutrino mass is close to the $\tilde{t}_1$ mass the rate is dramatically suppressed [for instance, $f(\frac{1}{2}) = 0.08$ already]. To obtain the total three-body rate, one should multiply Eq. (19) by 3 to take into account three generations of leptons. As a numerical estimate we take $m_{\tilde{t}_1} = 20$ GeV, $m_{\tilde{\chi}_1} = m_W$, $|c|^2 = \frac{1}{2}$, and $m_{\tilde{\nu}} = 0$. The result is

$$\Gamma(\tilde{t}_1 \to bl^+\tilde{\nu}_l) = 1.8 \times 10^{-9} m_{\tilde{t}_1} = (1.8 \times 10^{-17}\ \text{sec})^{-1} .$$

It is worth comparing this decay with $\tilde{t}_1 \to b\nu_l \tilde{l}_L^+$. ($\tilde{l}_R$ is forbidden for $m_l \to 0$.) A remarkable feature is that the $\tilde{l}$ mode is less suppressed than the $\tilde{\nu}$ mode as will be shown below. The contributing diagram is depicted in Fig. 3(b). The amplitude is[25]

$$\mathcal{M}=\sum_j \frac{-g^2U_{j1}}{(p_{\tilde{t}_1}-p_b)^2-m_{\tilde{\chi}_j}{}^2}\left[\left[V_{j1}\cos\theta_t+\frac{m_tV_{j2}\sin\theta_t}{\sqrt{2}m_W\sin\beta}\right]m_{\tilde{\chi}_j}\bar{u}(p_b)P_Rv(p_\nu)+\frac{m_bU_{j2}^*\cos\theta_t}{\sqrt{2}m_W\cos\beta}\bar{u}(p_b)\not{p}_{\tilde{l}}P_Rv(p_\nu)\right] . \tag{21}$$

The decay distribution with the same approximation as in Eq. (16) is

$$\frac{d\Gamma}{dz_b dz_\nu}=\frac{\alpha^2|c'|^2}{16\pi\sin^4\theta_W}\frac{m_{\tilde{t}_1}{}^3}{m_{\tilde{\chi}_1}{}^2}(z_b+z_\nu-1+R_{\tilde{l}}) , \tag{22}$$

where

$$c'=U_{11}\left[V_{11}\cos\theta_t+\frac{m_tV_{12}\sin\theta_t}{\sqrt{2}m_W\sin\beta}\right] . \tag{23}$$

The total rate is

$$\Gamma=\frac{\alpha^2|c'|^2}{96\pi\sin^4\theta_W}\frac{m_{\tilde{t}_1}{}^3}{m_{\tilde{\chi}_1}{}^2}g(R_{\tilde{l}}) \tag{24}$$

with

$$g(R)=(1-R)(1+10R+R^2)+6R(1+R)\ln R . \tag{25}$$

The $R\to 1$ behavior of this function is given by $g(R)\simeq\frac{1}{10}(1-R)^5$. Comparison of Eqs. (16) and (19) with Eqs. (22) and (24) reveals that the former is suppressed by $m_{\tilde{\chi}_1}{}^4$ whereas the latter only by $m_{\tilde{\chi}_1}{}^2$: The $\tilde{l}$ mode dominates the $\tilde{\nu}$ mode if both are not kinematically suppressed. (Note, however, that the existing bound for charged scalar leptons previously discussed implies that the $\tilde{l}$ mode is rather suppressed even if allowed.)

This fact can be easily understood in effective Lagrangian language. Ignoring the fermion masses and the Higgs fermion coupling (proportional to fermion masses), we notice that only the "left-handed" sector participates since the W gaugino does not couple to the "right-handed" matter. Now, we are to find the lowest-dimensional operator for the combinations $\tilde{t}_L\bar{b}_L l_L\tilde{\nu}_L^*$ and $\tilde{t}_L\bar{b}_L\bar{\nu}_L\tilde{l}_L$. For the former, chirality matching requires a gamma matrix between the fermion and antifermion fields, which must be contracted with a derivative for the scalar fields: $\bar{b}_L\gamma^\mu l_L(\tilde{\nu}_L^*\overleftrightarrow{\partial}_\mu\tilde{t}_L)$. This operator has dimension six and thus a coefficient $m_{\tilde{w}}{}^{-2}$. On the other hand, the operator for the latter can be constructed without derivatives since both fermion fields are antifermion: $\bar{b}_LC\bar{\nu}_L^T\tilde{l}_L\tilde{t}_L$ (C is the charge-conjugation matrix). This dimension-five operator only needs a factor $m_{\tilde{w}}{}^{-1}$ opposed to the other. In passing we note that if the charginos are much heavier than the W, the decay rate for the $\tilde{\nu}$ mode also develops a $m_{\tilde{\chi}_1}{}^{-2}$ behavior through the Higgs-fermion component, which is, however, suppressed by a factor $m_b{}^2/m_W{}^2$.

In any case, the three-body decays dominate the two-body decay $\tilde{t}_1\to c\tilde{\gamma}$ or $c\tilde{g}$ if energetically possible. The three-body phase-space suppression might be important, however, when the $\tilde{\nu}$ or $\tilde{l}$ mass is not negligible. The $\tilde{l}$ mode is unlikely to be fully open due to the bound for $m_{\tilde{l}}$.

C. Four-body decay

There are many diagrams[9,11] contributing to the decay $\tilde{t}_1\to bl^+\nu_l\tilde{\gamma}$ as shown in Fig. 4. We calculate the diagram (a) to make an order-of-magnitude estimate of the rate. The matrix element with the approximation $m_b, m_l=0$ is

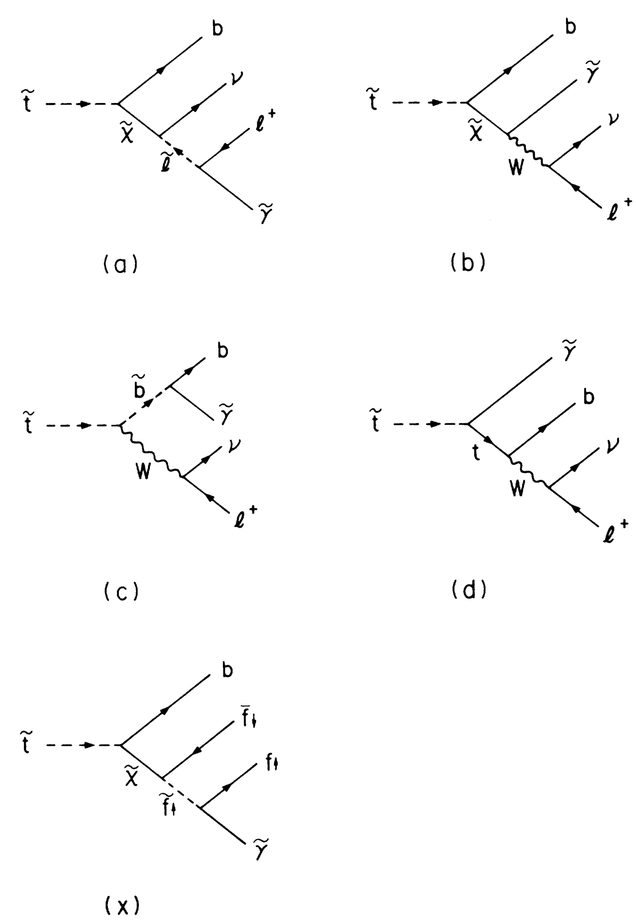

FIG. 4. Diagrams for the four-body decay $\tilde{t}_1\to bl^+\nu\tilde{\gamma}$ (a)–(d). Relabeling of these gives the diagrams for the other four-body decays $\tilde{t}_1\to b\bar{u}d\tilde{\gamma}$, $\tilde{t}_1\to b\bar{u}d\tilde{g}$, and $\tilde{t}_1\to bl^+\nu\tilde{g}$ [(c) and (d) only]. The last diagram (x) does not exist for the decay ($f_\uparrow=\nu$), but must be included for $f_\uparrow=u$, $f_\downarrow=d$.

$$\mathcal{M} = -\sum_j \sqrt{2} e g^2 U_{j1} \left[V_{j1} \cos\theta_t + \frac{m_t}{\sqrt{2} m_W \sin\beta} V_{j2} \sin\theta_t \right] \times \frac{m_{\tilde{\chi}_j}}{(p_{\tilde{t}_1} - p_b)^2 - m_{\tilde{\chi}_j}{}^2} \frac{1}{(p_l + p_{\tilde{\gamma}})^2 - m_{\tilde{l}_L}{}^2} \bar{u}(p_{\tilde{\gamma}}) P_L v(p_l) \bar{u}(p_b) P_R v(p_\nu) . \quad (26)$$

The total decay rate calculated in the limit $m_{\tilde{t}_1} \ll m_{\tilde{\chi}_1} \ll m_{\tilde{\chi}_2}$, $m_{\tilde{t}_1} \ll m_{\tilde{l}_L}$ is

$$\Gamma = \frac{\alpha^3 |c'|^2}{23\,040\pi^2 \sin^4\theta_W} \frac{m_{\tilde{t}_1}{}^7}{m_{\tilde{\chi}_1}{}^2 m_{\tilde{l}_L}{}^4}$$

with c' given by Eq. (23). If the W gaugino is essentially a mass eigenstate lighter than the Higgs fermion we have

$$\Gamma = \frac{\alpha^3 \cos^2\theta_t}{23\,040\pi^2 \sin^4\theta_W} \frac{m_{\tilde{t}_1}{}^7}{m_{\tilde{w}}{}^2 m_{\tilde{l}_L}{}^4} .$$

If one allows for a finite mass of the photino, the rate receives an extreme suppression factor of $(1-R_{\tilde{\gamma}})^8$.

In Ref. 11, a similar diagram [Fig. 4(x) with the gluino instead of the photino and with the W gaugino for the chargino] was evaluated and a rate proportional to $m_{\tilde{t}_1}{}^9 / m_{\tilde{\chi}_1}{}^4 m_{\tilde{f}}{}^4$ was found. The additional suppression of $m_{\tilde{t}_1}{}^2 / m_{\tilde{\chi}_1}{}^2$ is due to the different chiral structure of the diagrams (a) and (x), this situation is similar to that occurring between the two three-body diagrams in Fig. 3. [Inclusion of the Higgs-fermion component in the diagram (x) would remove the suppression but introduce a Yukawa-coupling factor $m_t{}^2/m_W{}^2$.] Thus, diagram (a) has a more important contribution to four-body decays than diagram (x). The other graphs (b)–(d) are not likely to give a much larger rate. Even with our larger estimate, the four-body decay cannot be competitive with the two-body decay

$$\frac{\Gamma(\tilde{t}_1 \to b l^+ \nu \tilde{\gamma})}{\Gamma(\tilde{t}_1 \to c \tilde{\gamma})} \approx (0.02\text{–}0.2) \frac{m_{\tilde{t}_1}{}^6}{m_{\tilde{w}}{}^2 m_{\tilde{l}_L}{}^4}$$

(and of course with the three-body decay) and may be ignored.

V. PRODUCTION RATE

In continuum e^+e^- annihilation, a $\tilde{t}_1$ pair can be produced via virtual photon or Z; see Fig. 5. The relevant gauge couplings of $\tilde{t}_1$ are given by

$$\mathcal{L} = -iee_t A^\mu (\tilde{t}_1^* \overleftrightarrow{\partial}_\mu \tilde{t}_1 + \tilde{t}_2^* \overleftrightarrow{\partial}_\mu \tilde{t}_2) - i(g/\cos\theta_W) Z^\mu [(\tfrac{1}{2}\cos^2\theta_t - e_t \sin^2\theta_W) \tilde{t}_1^* \overleftrightarrow{\partial}_\mu \tilde{t}_1 + (\tfrac{1}{2}\sin^2\theta_t - e_t \sin^2\theta_W) \tilde{t}_2^* \overleftrightarrow{\partial}_\mu \tilde{t}_2 + \tfrac{1}{2}\cos\theta_t \sin\theta_t (\tilde{t}_1^* \overleftrightarrow{\partial}_\mu \tilde{t}_2 + \tilde{t}_2^* \overleftrightarrow{\partial}_\mu \tilde{t}_1)] . \quad (27)$$

Note that there is an off-diagonal coupling connecting $\tilde{t}_1$ and $\tilde{t}_2$. The differential cross section for $e^+e^- \to \tilde{t}_1 \bar{\tilde{t}}_1$ is

$$\frac{d\sigma}{d\Omega} = \frac{3\alpha^2}{8s} \beta^3 \sin^2\theta \times \left[\left(-e_t + c_V \frac{s}{s - m_Z{}^2} \right)^2 + \left(c_A \frac{s}{s - m_Z{}^2} \right)^2 \right] , \quad (28)$$

where s is the c.m. energy squared, $\beta = (1 - 4m_{\tilde{t}_1}{}^2/s)^{1/2}$, θ the scattering angle, and

$$c_V = \frac{1}{\cos^2\theta_W \sin^2\theta_W} (\tfrac{1}{2}\cos^2\theta_t - e_t \sin^2\theta_W) \times (-\tfrac{1}{4} + \sin^2\theta_W) , \quad (29a)$$

$$c_A = \frac{1}{\cos^2\theta_W \sin^2\theta_W} (\tfrac{1}{2}\cos^2\theta_t - e_t \sin^2\theta_W)(-\tfrac{1}{4}) . \quad (29b)$$

The angular distribution has a form typical for scalar pair production. The total cross section is

$$\sigma = \frac{\pi\alpha^2}{s} \beta^3 \left[\left(-e_t + c_V \frac{s}{s - m_Z{}^2} \right)^2 + \left(c_A \frac{s}{s - m_Z{}^2} \right)^2 \right] \quad (30)$$

which corresponds to $R = \frac{3}{4} e_t{}^2 \beta^3 = \frac{1}{3}\beta^3$ for $s \ll m_Z{}^2$. The P-wave threshold rise is rather slow. The QCD correction should give an important contribution near threshold due

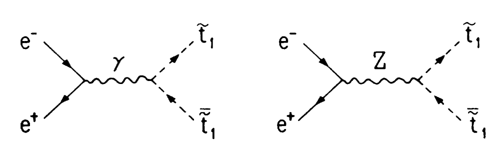

FIG. 5. Diagrams for the reaction $e^+e^- \to \tilde{t}_1 \bar{\tilde{t}}_1$.

to the Coulomb-type singularity; in fact, when $\beta \lesssim \alpha_s$ the $O(\alpha_s)$ "correction" is expected to be larger than the lowest-order result. However, in this region, the absolute cross section is tiny so the "large" correction does not have much significance unless one has a very large integrated luminosity.

On the Z resonance we have for the relative branching fraction for a $\tilde{t}_1$ pair

$$\frac{\Gamma(Z\to\tilde{t}_1\bar{\tilde{t}}_1)}{\Gamma(Z\to\nu\bar{\nu})}=6(\tfrac{1}{2}\cos^2\theta_t-e_t\sin^2\theta_W)^2\beta^3$$

$$=[0.13,0.064,0.75]\beta^3$$

for $\cos^2\theta_t=[0,\frac{1}{2},1]$. This corresponds to a branching ratio of $[0.8,0.4,4.5]\beta^3\%$ if $\tilde{t}_1$ is the only SUSY particle producible in Z^0 decay.

At the t-quarkonium resonance there could be a substantial branching into the $\tilde{t}_1$ pair if the gluino mass is not too large. Also, the single decay $(t\bar{t})\to\tilde{t}_1+\tilde{\gamma}+\bar{t}$ could be dominant if $m_{\tilde{t}_1}+m_{\tilde{\gamma}}<m_t$. In this case, the top quark produced in continuum $e^+e^-\to t\bar{t}$ also decays to $\tilde{t}_1+\tilde{\gamma}$. For further details see Ref. 8.

VI. COMMENTS AND CONCLUSION

At the highest PETRA energies ($\sqrt{s}=46.3$–46.8 GeV) the Mark J and JADE Collaborations have observed[26] anomalous hadronic events with low thrust and an isolated muon. Similar events with an electron were not found. These events could be interpreted as a $\tilde{t}_1$ pair production with the decay $\tilde{t}_1\to b\mu\tilde{\nu}_\mu$. If we assume that most of the decay gives muons, which would require that $\tilde{\nu}_\mu$ is lighter than other scalar neutrinos, the observed cross section would imply $m_{\tilde{t}_1}\sim 20$ GeV. (If all the scalar neutrinos are degenerate one would need a smaller $m_{\tilde{t}_1}$.) Possible difficulties with this interpretation include the fact that the slow rise of the cross section may be in conflict with the nonobservation of similar events at lower energies, and the fact that there must be a second muon in the event.

We have not discussed the $\tilde{t}_1$ production at $p\bar{p}$ colliders, which is studied in Refs. 9 and 12.

To summarize, the dominant decay mode of a light scalar top producible at the forthcoming e^+e^- colliders is $\tilde{t}_1\to c\tilde{\gamma}$ or $\tilde{t}_1\to bl^+\tilde{\nu}$, depending on whether the latter mode is kinematically allowed. If the available Q value for the latter is small the two modes may compete. More generally, we have obtained the following hierarchy for decays:

$$\Gamma(\tilde{t}_1\to b\nu\tilde{l}^+)\gg\Gamma(\tilde{t}_1\to bl^+\tilde{\nu})$$
$$\gtrsim\Gamma(\tilde{t}_1\to c\tilde{g})\gg\Gamma(\tilde{t}_1\to c\tilde{\gamma})$$
$$\gg\text{four-body decays}\ ,$$

assuming no phase-space suppression. Electron-positron colliding machines such as TRISTAN, SLC, and LEP have a good change of finding the scalar top if its mass lies within the reach of the machine.

ACKNOWLEDGMENTS

We would like to thank K. Hagiwara for discussions and comments on the manuscript. This work was supported in part by Grant-in-Aid for Scientific Research from the Ministry of Education, Science, and Culture, Japan and in part by the Japan-U.S. Cooperative Science Program.

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