Searching for Supersymmetry in Hadrons

S. James Gates, Jr.

Department of Physics
University of Maryland
College Park, Maryland 20742-4111

and

Oleg Lebedev

Department of Physics
Virginia Polytechnic Institute and State University
Blacksburg, Virginia 24061-0435

March 27, 2022

Abstract

We discuss the possibility of the existence of a long-lived top squark ($\tau \ll 10^{10} \text{ years}$) and its motivation. If the stop is indeed metastable, it forms hadrons. We study properties of the low-energy stop-containing hadrons and their signatures in collider experiments.
PACS: 12.60.J, 14.80.Ly, 14.20
1 Introduction

A feature of most discussions on the possibility of observing supersymmetry in hadronic physics is that attention has been paid almost exclusively to “fragmentation physics.” That is, predictions have been made of what distinctive signature might occur in hadronic jets due to the production and subsequent decay of sparticles. This appears due to the expectation that the mass of the lightest squarks will be substantially greater than that of the lightest super-partners of standard model particles that do not carry color. This is the pattern within the conventional standard model and is so expected for its presumptive supersymmetric successor. Under this circumstance, bound states involving squarks and ordinary quarks would not undergo “hadronization,” superpartner hadrons (“shadrons”) would never form and hadron spectroscopy would be irrelevant.

Although unconventional, we can ask a simple question, “If the mass of squarks were comparable to that of the LSP (lightest supersymmetric particle), what type of hadrons might be expected?” Part of what prompts us to ask this unconventional question has been our formal investigations of possible structures that give naive supersymmetric extensions of the effective field theory describing low-energy pion physics. These results seem to indicate that supersymmetric extensions of these are formally possible.

The possibility of the existence of supersymmetric hadrons was first pointed out by Farrar and Fayet [2]. Historically, models containing light gluinos rather than squarks appeared to be better motivated. Subsequently, the gluino containing hadrons (“R-hadrons”) have been studied to a great extent [3]. On the other hand, another option - the squark containing hadrons - has not drawn similar attention. We believe that, in the ongoing search for supersymmetry, one cannot afford the luxury of ignoring this possibility however unlikely it may seem. In this letter, we attempt, at least partially, to fill this gap in the phenomenology of supersymmetry.

2 Metastable Stop

In this paper we study the possibility of a metastable top squark. It is well known that the top squark can be the lightest supersymmetric particle (LSP) if the left-right stop mixing is comparable to the squark masses (for
reviews, see [4]). Combined with the requirement that the interactions be strictly R-parity invariant, this results in the stability of the top squark as the stop decays into other superparticles are not allowed kinematically.

It is, however, also well known that such a scenario encounters serious cosmological problems. A stable stop would form stable hadrons. The abundance of such hadrons can be estimated in the standard cosmological model knowing the masses and cross sections of these particles. One finds that stable 100 $GeV$ hadrons should exist in concentrations of one in $10^{10} - 10^{12}$ nucleons [5]. Thus they could be detected as anomalously heavy protons or, if they get bound in nuclei, they could be found in anomalous isotopes [6]. However, years of experimental searches for anomalous protons and isotopes have resulted in extremely tight bounds on their concentration in matter [7], [8]:

\begin{align}
    n_{q=+1}/n_{\text{nucleon}} &< 10^{-29}, \\
    n_{q=-1}/n_{\text{nucleon}} &< 4 \times 10^{-20} \tag{1}
\end{align}

for $m=100$ $GeV$; the displayed bound on $n_{q=-1}$ results from the search for anomalous isotopes of carbon. Apparently, these numbers are in conflict with the expected concentration of the stop containing hadrons. Thus, the top squark cannot be absolutely stable.

The stability of the stop arises from the assumption of an exact R-invariance. However, recently, with the discovery of the neutrino mass, this assumption has lost some of its theoretical motivation. It seems more plausible that R-parity is not an exact symmetry as is not lepton (baryon) number (see [9] for a review). In this case the LSP decays into ordinary particles. If R-parity is broken only slightly, the decay time can be very long. This may lead to other cosmological problems. For example, if the decay products contain quarks, the balance between protons and neutrons would be broken during the nucleosynthesis leading to unacceptable abundances of helium and other light elements [10]. This places strict constraints on the lifetime of the dark matter candidates such as neutralinos and gravitinos. These considerations can also lead to significant constraints on the lifetimes of new strongly interacting particles if one assumes that the baryon asymmetry in the “new” sector is the same as it is in the regular matter sector. This assumption, however, is strongly model-dependent and needs not be the case. In particular, if the asymmetry in the stop sector is much smaller than that in the matter
sector, the concentration of the susy hadrons at the time of nucleosynthesis is too small ($\sim 10^{-10} n_{\text{bar}}$) to affect the formation of light elements. In this case the stop lifetime is left essentially unconstrained as long as

$$\tau_{\text{stop}} \ll 10^{10} \text{ years}.$$  

(2)

The top squark with such a lifetime is the subject of our paper.

The stop is allowed to decay into quarks and charged leptons via the R-breaking interactions:

$$W_R = \lambda'_{ijk} \hat{L}_i \hat{Q}_j \hat{D}_k + \lambda''_{ijk} \hat{U}_i \hat{D}_j \hat{D}_k,$$

(3)

where $\hat{L}_i$, $\hat{Q}_i$, $\hat{U}_i$, and $\hat{D}_i$ are the usual MSSM superfields. The corresponding decay width is given by

$$\Gamma = \frac{|\lambda', \lambda''|^2}{16\pi} m_{\tilde{t}}.$$  

(4)

The requirement that all the stops have decayed by now can then be translated into the lower bound on the couplings:

$$\lambda', \lambda'' \gg 10^{-21}$$

(5)

with $m_{\tilde{t}} = 100 \text{ GeV}$. The most recent upper bounds on these couplings from particle experiments can be found in [11].

A model containing such small numbers may seem ill-motivated. However, very small R-parity violation can naturally arise in some flavor models with a discrete symmetry [12],[13]. In particular, in the model of Ref. [13], the $(S_3)^3$ symmetry prohibits renormalizable R-breaking operators. R parity arises as an accidental symmetry and can be broken by nonrenormalizable operators induced at the Plank scale. Then the effective low-energy trilinear operators get suppressed by $(M_f/M_{Pl})^2$ and $M_f/M_{Pl}$, where $M_f$ is the flavor symmetry breaking scale. With a sufficiently low $M_f$, the R-breaking couplings in the range given by Eq. (3) can be naturally produced. For example, the lower bound on the $\lambda', \lambda''$ corresponds to $M_f \sim 10^9 \text{ GeV}$ if the R-breaking interactions are suppressed by $M_{Pl}^2$. This provides a theoretical motivation for studying a possibility of a long-lived top squark.
3 Lightest Stop Containing Hadrons

Since any observable object has to be a color singlet, a metastable top squark must be confined in hadrons. Within $\mathcal{O}(10^{-23})$ sec it will form a bound state with either an antiquark - mesonino, or a couple of quarks - sbaryon. Clearly, the lightest mesonino (sbaryon) is just as stable as the stop itself. Masses of the lightest shadrons can be approximated within 1 GeV by the stop mass which we will set to 100 GeV, however splittings between them will be very important for phenomenology. The natural candidates for the lightest hadron are listed in Table 1 (this is to be supplemented with the charge conjugate states).

In the presence of isospin symmetry breaking, $T^0$ is not a mass eigenstate. Indeed, due to the diagram in Fig.1, $T^0$ mixes with $\bar{T}^0$ forming Majorana fermions $T_1$ and $T_2$.

Let us estimate how large is the mixing. The leading contribution is given by the gluino diagram in Fig.1. Further, in the limit $m_{\tilde{g}}^2 \gg m_{\tilde{t}}^2$, the mixing is dominated by the diagram with a gluino mass insertion (Fig.2). It is also reasonable to assume that the lightest stop is predominantly right handed. Then, the generated effective Hamiltonian is given by

$$H_{eff} \approx \frac{g_2^2 V_{tu}^2}{3m_{\tilde{g}}} (\bar{u}_R \tilde{t}_R)^2,$$

where $\tilde{V}$ is the mixing matrix at the gluino vertex. To convert this interaction into a mass term, one needs to know the matrix element $\langle T^0 | H_{eff} | T^0 \rangle$. In the
case of the Standard Model $K - \bar{K}$ mixing, the needed matrix elements can be derived from the PCAC relation and the vacuum saturation assumption. In our case we will simply parametrize the matrix element assuming

$$\langle T^0 | (\bar{u}_R \tilde{t}_R)^2 | \bar{T}^0 \rangle = F^2,$$

(7)

where $F$ has dimension of mass\[^1\]. The size of $F$ can be estimated by the characteristic hadronic matrix element of a heavy quark meson $(f_B m^2_B)^{1/3} \sim \mathcal{O}(1) \text{ GeV}$ since this combination is constant in the heavy quark limit.

As a result, we get the following mass matrix in the $(T^0, \bar{T}^0)$ basis:

$$M = \begin{pmatrix} m & \epsilon \\ \epsilon^* & m \end{pmatrix},$$

(8)

where $\epsilon = \frac{g^2 V^2_{tu}}{3m_\tilde{g}} F^2$. In principle, $\tilde{V}_{tu}$ may contain a complex phase, leading to CP violation in mixing. However, in this paper, we restrict ourselves to the CP conserving case. The consequent mass eigenstates are maximal mixtures of $T^0$ and $\bar{T}^0$:

$$T_1 = \frac{1}{\sqrt{2}}(T^0 - \bar{T}^0), \quad m_1 = m - \epsilon,$$

$$T_2 = \frac{1}{\sqrt{2}}(T^0 + \bar{T}^0), \quad m_2 = m + \epsilon.$$

(9)

It is interesting to compare the mass splitting between the two eigenstates to a typical mass splitting in the neutral meson systems. For example, in the Standard Model, the $B^0 - \bar{B}^0$ mixing arises only through a box diagram resulting in $(\Delta m/m)_{B^0} \sim 10^{-13}$. In contrast, the $T^0 - \bar{T}^0$ mixing exists already at the tree level and

$$\frac{\Delta m}{m} = \frac{2}{3} g_s V^2_{tu} F^2 \frac{m_\tilde{g} m_{\tilde{t}}}{m_\tilde{g} m_{\tilde{t}}} \sim 10^{-9},$$

(10)

for $\tilde{V}_{tu} \sim V^{CKM}_{31}, \ m_\tilde{g} \sim 300 \text{ GeV}, \ m_{\tilde{t}} \sim 100 \text{GeV}$ and $F \sim 1 \text{ GeV}$. The result, of course, is sensitive to $F$, but it is clear that $\Delta m/m$ is enhanced by orders of magnitude over that in the Standard Model.

Let us now turn to the discussion of the isospin partner of $T^0 - T^+$. Since the top squark can be treated as a spectator, the mass splitting in this

\[^1\]In our convention $\langle T^0 | T^0 \rangle = \delta^{(3)}(p - p')$.  

6
isospin doublet can be attributed to the $m_d - m_u$ mass difference and the corresponding Coulomb interactions. As a result, $T^+$ is unstable and $\beta$-decays:

$$T^+ \rightarrow T^0 + e^+\nu \quad (11)$$

(It is understood that $T^+$ actually decays into a linear combination of the mass eigenstates $T_1$ and $T_2$.) Clearly, this process is governed by the usual Fermi interaction in analogy with a neutron $\beta$-decay:

$$H = \frac{G}{\sqrt{2}} T^+ \gamma^\mu (1 - \alpha \gamma_5) T^0 \gamma_\mu (1 - \gamma_5) \nu , \quad (12)$$

where $\alpha \sim 1$ accounts for hadronic effects. The decay width is given by that of a neutron $\beta$-decay:

$$\Gamma \sim \frac{G^2 \Delta^5}{60\pi^3} \quad (13)$$

with $\Delta = m_{T^+} - m_{T^0} \approx m_{D^+} - m_{D^0} \approx 5 MeV$. The corresponding decay time is about 3 sec. So within a few seconds all of the $T^+$'s will decay into the $T^0$'s.

One can go further in eliminating candidates for the lightest shadron. Let us consider the mass splitting between the vector isotriplet $V$ and the scalar isosinglet $S$. The mass difference arises due to the spin flip of one of the quarks. This effect is accounted for in the quark potential model via the spin-spin interaction (see, for example, [15]). Such an interaction is part of the one gluon exchange potential and is responsible for the $\rho - \pi$ and $\Delta - N$ mass differences. Again, treating a spin 0 stop as a spectator, we can write the hyperfine interaction as follows:

$$H_{h.f.} = k_{qq} \mathbf{s}_1 \cdot \mathbf{s}_2 \delta^{(3)}(\mathbf{r}). \quad (14)$$

The constant $k$ depends on the quark masses and on whether two quarks or a quark and antiquark are interacting, $k_{qq} = \frac{1}{2} k_{q\bar{q}}$. The scalar product $\mathbf{s}_1 \cdot \mathbf{s}_2$ is 1/4 for the vector state and -3/4 for the scalar state. This situation is very similar to what we encounter in the $\rho - \pi$ system up to the substitution $k_{q\bar{q}} \rightarrow k_{q\bar{q}}$. Therefore,

$$m_V - m_{S^+} \sim \frac{1}{2} (m_\rho - m_\sigma) \approx 300 \text{ MeV} \quad (15)$$

\footnote{This estimate is sensitive to the exact value of the mass difference. The actual decay time may be an order of magnitude larger. See [14] for similar considerations.}
Another way to estimate the $V - S^+$ mass difference is to exploit the heavy quark symmetry of the strong interactions [16]. In the heavy quark limit, the interactions of a heavy quark are independent of its spin and flavor. Thus the heavy quark spectroscopy depends on the light quark configuration only. Using this fact, one can approximate the $V - S^+$ mass difference by that of the $\Sigma$ and $\Lambda$ baryons containing a heavy quark. This gives

$$m_V - m_{S^+} \sim 200 \text{ MeV},$$

which is about 30% away from the previous estimate [3].

In both cases the mass splitting is well above the pion threshold, so $V$ will decay into $S^+$ within about $10^{-23} \text{ sec}$:

$$V \rightarrow S^+ + \pi$$

Thus we are left with only two potentially long-lived hadrons - $T^0$ and $S^+$. In principle, there is a possibility that one of them will decay into another. Since they carry different baryon numbers, a proton must be emitted in such a decay: $S^+ \rightarrow T^0 + p$. It is, however, quite unlikely that the $S^+ - T^0$ mass difference is above 1 GeV. To see this, we can again exploit the spin independence of the heavy quark interactions and use an analogy with the heavy quark hadrons. The $S^+ - T^0$ mass splitting can be compared to that for the $\Lambda_{b(c)}$ baryons and $B(D)$ mesons. This estimate yields

$$m_{S^+} - m_{T^0} \sim 400 \text{ MeV}.$$  

Thus, the decay $S^+ \rightarrow T^0 + p$ is not allowed kinematically.

To summarize, we have argued that there are four long-lived hadrons (including the charge conjugate states): $T_{1,2}$ and $S_{\pm}$. Their lifetimes are given by the stop lifetime. Other light hadrons decay into these states either strongly or via $\beta$-decay.

4 Possibility of Detection

Supersymmetric particles have been the subject of collider searches for many years (for recent reviews see [17] and references therein). In most of

\footnote{To obtain this number, we use the mass of $\Lambda_c$ and the average mass of $\Sigma_c$, spin 1/2 and $\Sigma_c$, spin 3/2 (which should be the same in the heavy quark limit) [21]. Analogous information for the $b$-baryons currently is not available.}
these searches it was assumed that the LSP is a neutral particle which eludes
detection. This assumption leads to the celebrated signature of supersym-
metry - missing transverse momentum and energy. The recent LEP and
Tevatron bounds on masses of the supersymmetric particles can be found in
[18].

In the scenario which we are discussing, the LSP is a charged strongly
interacting particle. If it were sufficiently long-lived, it would leave a track in
the tracking chambers and thus could be detected directly. In particular, if
a heavy charged particle is produced at the Tevatron, it can be identified in
the CDF detector due to its relatively low velocity and, consequently, large
ionization deposition, $dE/dx$. CDF measures $dE/dx$ in the silicon vertex
detector as well as in the central tracking chamber. In Run I, no events in
excess of the expected background have been found, setting a limit of 85 GeV
on the mass of a unit-charged scalar triplet [19]. A more recent analysis [20]
increases this bound up to about 100 GeV.

In the previous subsection we have argued that there are four metastable
shadrons. This is the case when these shadrons are in isolation. However,
when interactions with matter are turned on, only one of them survives. Let
us consider the states $\tilde{t}\bar{u}$ and $\tilde{u}t$ propagating in matter. Apparently, $\tilde{t}\bar{u}$ has
a larger interaction cross section since the annihilation channel is open:

\begin{align}
\tilde{t}\bar{u} + p &\rightarrow \tilde{t}ud + \pi^0, 2\pi, \ldots \quad (19) \\
\tilde{u}t + p &\rightarrow \tilde{u}dt + \gamma, 2\gamma, \ldots \quad (20)
\end{align}

As a result, the interaction eigenstates are different from the mass or flavor
eigenstates. This is a well known phenomenon in K-meson physics: the same
effect is used to regenerate the $K_S$ component via interactions with matter.

The effective mass matrix in the $(T^0, \tilde{T}^0)$ basis now takes the form

\begin{equation}
M' = \begin{pmatrix}
m & \epsilon \\
\epsilon & m
\end{pmatrix} \cdot \begin{pmatrix}
1 - i\gamma & 0 \\
0 & 1
\end{pmatrix} = \begin{pmatrix}
m - im\gamma & \epsilon \\
\epsilon & im\gamma
\end{pmatrix}, \quad (21)
\end{equation}

where we have taken into account only the largest absorptive effect. The
constant $\gamma > 0$ depends on the material and energy of the shadron. The size

\footnote{Many of these bounds are based on signatures specific to certain models and do not
hold in general. For example, the Tevatron bound on the stop mass is based on the
decay $\tilde{t} \rightarrow c\tilde{\chi}^0$, which cannot occur in our model.}

\footnote{Interactions with the material of the detector may weaken this bound by 30-40 GeV [19].}
of $\gamma$ can be estimated \textit{very roughly} using the nuclear interaction length $\lambda_I$. For protons propagating in iron, $\lambda_I \sim 17 \text{ cm}$ \cite{21}. This number triples for the $T^0$ as it contains only one light (anti)quark \cite{22}, yielding $m\gamma \sim 10^{-15}\text{ GeV}$. Therefore, $m\gamma \ll \epsilon$ and the interaction eigenstates can be well approximated perturbatively in $\gamma$. The corresponding effective masses are

$$m_1' = (m - \epsilon) \left[ 1 - \frac{i\gamma}{2} \right],$$

$$m_2' = (m + \epsilon) \left[ 1 - \frac{i\gamma}{2} \right].$$

(22)

Apparently, both mesonino states eventually vanish in a medium converting into sbaryons. For example, in iron, the attenuation length would be about 1 meter. At very low energies, the mesoninos may interact with matter primarily through the electromagnetic annihilation of $u$ and $\bar{u}$ if the strong channel (19) is kinematically prohibited.

Similar considerations can be applied to the sbaryon $S^-$ which contains light antiquarks. As a result, the only stop-containing particle which does not get destroyed in matter is $S^+$. All other shadrons propagating in a medium will go through a series of reactions and eventually end up as $S^+$.

Nevertheless, if the $\lambda'$ and $\lambda''$ are anywhere within 10 orders of magnitude from the lower bound (3), all of the charged shadrons but $V$ (which decays too quickly) would leave tracks in the tracking chambers. Since they are expected to be very massive, they would pass through the electromagnetic calorimeter undetected. With some energy loss (which may be considerable for slow particles \cite{22}) in the hadron calorimeter, a shadron will penetrate the detector and will be triggered on as a muon.

Let us consider the interaction of shadrons with the hadron calorimeter in more detail. In collisions with the nucleons, the momentum transfer at typical Tevatron energies is very small and the energy available in the center of mass system is likely to be below the pion threshold \cite{19}. Thus, the inelastic scattering would primarily go through the annihilation of the light quarks and antiquarks as in Eq. (13). Since the $S^+ - T^0$ mass difference is estimated to be around 400 $MeV$, a few soft pions (or an equivalent amount of radiation) can be produced in the collision of $T_{1,2}$ and $p$. Note that, at low energies, only the shadron states involving light antiquarks ($T^+, S^-$, etc.) would deposit any tangible amount of energy in the hadron calorimeter. However, at the Tevatron, the deposited energy is unlikely to exceed a few
GeV in any case and thus the signal would be hard to distinguish from that
generated by a charged colorless LSP such as a stau, especially taking into
account that the energy resolution is only $80 - 100\% / \sqrt{E}$ [17].

The situation will improve at the LHC. The energy deposition in the
calorimeter will be sufficient to distinguish between a strongly and a weakly
interacting LSP. The energy loss in a material as a function of the mass and
energy of a strongly interacting particle was studied in Ref. [23] based on
kinematical considerations. For instance, it was shown that a 100 GeV par-
ticle at $E = 1 \text{ TeV}$ would trigger a few 10 GeV hadron showers. Also, it was
noted that, due to a slower velocity, the opening angle of the shower would
be significantly wider than that of the shower induced by regular hadrons.
These calorimeter data combined with the tracking (and timing) information
would allow to eliminate the SM background and determine the nature of a
new (meta)stable particle.

Let us now briefly discuss other implications of the scenario. The stop may
be sufficiently long-lived as to be treated as stable. One may ask whether
there could be non-collider experiments which can detect shadrons produced
in the collisions. In principle, this can be done via mass spectroscopy. If a
produced shadron is sufficiently slow, it loses most of its energy through
ionization and thus can be trapped in the metal plates of the detector [22],
eventually ending up as an $S^+$. The isosinglet $S^+$ is very unlikely to be bound
in nuclei due to the Coulomb repulsion and absence of the one-pion-exchange
potential. Therefore, it should be found in isolation. Having picked up an
electron, it will form an anomalously heavy hydrogen atom. If such atoms
existed in an appreciable concentration, they would be detectable by mass
spectrometry [4] as they possess an anomalously small charge/mass ratio.

We would like to conclude with a remark:

“Curiosity is a delicate little plant which, aside from
stimulation, stands mainly in need of freedom.”

— Albert Einstein

We have benefited from very helpful correspondence with T. Cohen, H.
Dreiner, X. Ji, A. Kronfeld, D. Stuart as well as from discussions with mem-
bers of Virginia Tech and U. of Maryland high energy groups. We are par-
icularly thankful to Lay N. Chang, S. Eno, R.N.Mohapatra and G. Snow
for enlightening discussions.
References

[1] S.J. Gates, Jr., in the Proceedings of the 5th International Conference on Supersymmetries in Physics (SUSY’97), Philadelphia, USA, May 27-31, 1997, North-Holland (1998), pp. 171 - 181; in the Proceedings of the 2nd International Conference on Quantum Field Theory and Gravity, Tomsk, Russia, I. L. Buchbinder and K. E. Osetrin, pp. 79–92 Tomsk State Pedagogical Univ. Press, 1998.

[2] G. Farrar and P. Fayet, Phys. Lett. 76B (1978) 575; Phys. Lett. 79B (1978) 442.

[3] G.R. Farrar, Phys. Rev. Lett. 76 (1996) 4111; Phys. Rev. D 51 (1995) 3904; G.R. Farrar and G.T. Gabadadze, Phys. Lett. 397B (1997) 104; and references therein.

[4] H.E. Haber and G. L. Kane, Phys. Rep. 117 (1985) 75; S. Dawson, hep-ph/9712464; S.P. Martin, hep-ph/9709356.

[5] S. Wolfram, Phys. Lett. 82B (1979) 65; E. Nardi and E. Roulet, Phys. Lett. 245B (1990) 105.

[6] C.B. Dover, T.K. Gaisser, G. Steigman, Phys. Rev. Lett. 42 (1979) 1117.

[7] P.F. Smith et al., Nucl. Phys. B 206 (1982) 333.

[8] T.K. Hemmick et al., Phys. Rev. D 41 (1990) 2074.

[9] H. Dreiner, in “Perspectives on Supersymmetry”, ed. G.L. Kane, World Scientific, p. 462-479; G. Bhattacharyya, hep-ph/9709393.

[10] M.H. Reno, D. Seckel, Phys. Rev. D 37 (1988) 3441; J. Ellis, G.B. Gelmini, J.L. Lopez, D.V. Nanopoulos, S. Sarkar, Nucl. Phys. B 373 (1992) 399.

[11] B.C. Allanach, A. Dedes, H.K. Dreiner, Phys. Rev. D60 (1999) 075014; O. Lebedev, W. Loinaz, T. Takeuchi, hep-ph/9910435; O. Lebedev, W. Loinaz, T. Takeuchi, hep-ph/9911479.
[12] L. Hall, in *Proceedings of the Theoretical Workshop on Cosmology and Particle Physics*, Berkeley, CA, ed. I. Hinchliffe, World Scientific, 1987, p. 106-115.

[13] C.D. Carone, L.J. Hall, H. Murayama, Phys. Rev. D 54 (1996) 2328.

[14] R.N. Cahn, Phys. Rev. Lett. 40 (1978) 80; R.N. Cahn, S.D. Ellis, Phys. Rev. D 16 (1977) 1484.

[15] J.F. Donoghue, E. Golowich, B.R. Holstein, *Dynamics of the Standard Model*, Cambridge University Press, 1992.

[16] N. Isgur and M. B. Wise, Phys. Lett. 232B (1989) 113; Phys. Lett. 237B (1990) 527; Phys. Rev. Lett. 66 (1991) 1130.

[17] M. Carena, R. Culbertson, H. Frisch, S. Eno, S. Mrenna, Rev. Mod. Phys. 71 (1999) 937.

[18] J.F. Grivaz [ALEPH Collaboration], *talk given at SUSY’99*, Batavia, IL, USA, 1999, available at http://fnal.gov/pub/hep_descript.html; A. Favara [L3 Collaboration], *ibid.*; J. Qian [D0 Collaboration], *ibid.*

[19] F. Abe *et al.*, Phys. Rev. D 46 (1992) R1889; K. Maeshima [CDF Collaboration], in *Proceedings of ICHEP’96*, Warsaw, Poland, 1996; FERMILAB-Conf-96/412-E.

[20] A. Connolly [CDF Collaboration], hep-ex/9904010; D. Stuart, private communication; the bound on the mass of a scalar color triplet can be obtained using cross sections of [22].

[21] Particle Data Group, Eur. Phys. J. C 3 (1998) 1; updated information can be found at http://www-pdg.lbl.gov/.

[22] M. Drees and X. Tata, Phys. Lett. 252B (1990) 695.

[23] R.N. Mohapatra, F. Olness, R. Stroynowski, V.L. Teplitz, Phys. Rev. D 60 (1999) 115013; R.N. Mohapatra, V.L. Teplitz, Phys. Rev. Lett. 81 (1998) 3079.
Figure 1: $T^0 - \bar{T}^0$ mixing.

Figure 2: The dominant contribution to $T^0 - \bar{T}^0$ mixing.