Brief Reports

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Single-top-quark production in e^+e^- collisions via superstring Yukawa interactions

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(Received 30 November 1987)

The top quark may be singly produced in e^+e^- collisions via the exchange of leptoquarks. We present results for $e^+e^- \to t\bar{c}$, $c\bar{t}$, or $t\bar{u}$, $u\bar{t}$ open and bound-state production in terms of superstring Yukawa couplings.

In the E_6 theory obtained from compactification of the extra dimensions[1] in the ten-dimensional heterotic string,[2] totally new low-energy interactions are expected.[3] These interactions may have observable consequences at present or future colliders.[4,5] In this Brief Report we consider some new implications for e^+e^- collisions, pointing out that the top quark can be singly produced in the processes $e^+e^- \to t\bar{c}$ or $t\bar{u}$ via leptoquark exchange. In the standard model such flavor-changing neutral currents, occurring at the one-loop level in Z exchange, are highly suppressed and unobservable;[6] this conclusion remains essentially unchanged even in its supersymmetric extensions.[7]

In the E_6 theory, the matter particles of each generation are classified in a 27-dimensional representation and the low-energy superpotential consists of the Yukawa terms from an E_6-invariant $\mathbf{27}\times\mathbf{27}\times\mathbf{27}$ interaction. The couplings need not respect the E_6 relations,[3] but are invariant under the low-energy gauge group, taken here to be the rank-5 subgroup[3,5] $SU(3)_c\times SU(2)_L\times U(1)_Y \times U(1)_\eta$. The **27** representation of chiral superfields naturally incorporates both the 15 observed chiral fermions of each generation and a pair of Higgs doublets in one grand multiplet.

The **27** also contains an exotic color triplet, SU(2) singlet, electric charge $-\frac{1}{3}$ spinless particle S and its charge conjugate S^c, as well as their fermionic superpartners $\tilde{S}$ and $\tilde{S}^c$. These particles can be regarded as a color-triplet partner of the Higgs doublet in the SU(5) grand unified theory.[8] They can mediate rapid proton decays and the standard solution for proton longevity has been to give them very large masses ($\gtrsim 10^{14}$ GeV) while retaining light doublets.[9,10] The superstring theory opened up an alternative possibility that these particles are indeed as light as Higgs bosons but that their dangerous Yukawa couplings are somehow forbidden by the symmetry and/or topology of the compactified dimensions.[3,11,12]

The vanishing of certain Yukawa couplings that lead to proton decay allows us to assign baryon quantum numbers to all the light particles. Assuming also lepton-number conservation, we are left with three distinctive choices of their quark and lepton numbers: S can be a leptoquark, an antidiquark, or a squark.[13,14] In the simplest model of a rank-5 low-energy gauge group, the squark assignment is ruled out.[14] Here we are concerned with the case that S is a leptoquark.

From the superpotential the Lagrangian for the S couplings to electrons and quarks is[15]

$$-\mathcal{L} = (\lambda_1)_{ijk} S^*_{Ri} \overline{(e_{Lj})^c} u_{Lk} + (\lambda_2)_{ijk} S_{Li} \overline{e_{Rj}} (u_{Rk})^c + \text{H.c.} ,$$

where i,j,k are generation indices. Here we denote the scalars by S_{Li} and S_{Ri}, in the basis where their fermionic (Dirac) superpartners $\tilde{S}_i$ are diagonal. In this basis, the scalar mass eigenstates $S_{i\pm}$ are obtained by the rotation

$$\begin{bmatrix} S_{i-} \\ S_{i+} \end{bmatrix} = \begin{bmatrix} \cos\theta_i & \sin\theta_i \\ -\sin\theta_i & \cos\theta_i \end{bmatrix} \begin{bmatrix} S_{Li} \\ S_{Ri} \end{bmatrix} ,$$

where $m_{S_{i-}} < m_{S_{i+}}$. The left-right mixing is expected to be significant because of the mixing due to soft supersymmetry-breaking terms proportional to large Dirac mass $m_{\tilde{S}}$, and also because of the mixing due to the existence of supersymmetric four-point $S_L S_R$ couplings with Higgs bosons. Further intergenerational mixing from renormalization effects[16] do not affect our qualitative conclusions.

With the Yukawa couplings of the leptoquarks, generation-changing transitions such as $e \to c$ or $e \to t$ are possible. Hence these new interactions from the superstring may have striking physical consequences.[17,15] Of particular interest here is the possibility of single t-quark production in e^+e^- collisions via the diagram in Fig. 1. With this mechanism a t-quark mass up to 180 GeV can potentially be produced at CERN LEP II, which covers the upper limit on m_t allowed by neutral-current analyses.[18] Also if the t quark is light, single top production might be observed at KEK TRISTAN, the SLAC Linear Collider, and LEP I even if pair production is kinematically forbidden.

The amplitudes for the process

$$e^-(k,\eta)+e^+(\bar{k},\bar{\eta})\to t(P)+\bar{c}(p) ,$$

depend on the helicities η and $\bar{\eta}$. The $\eta=-,\bar{\eta}=+$ amplitude is given by

$$\mathcal{M}_{-+}=A(t)\bar{v}(k,\eta)P_L v(p)\bar{u}(P)P_R u(\bar{k},\bar{\eta}) ,$$

where $P_{L,R}$ are the chirality projections, $t=(k-p)^2=(\bar{k}-P)^2$, and

$$A(t)=\sum_{i=1}^{3}(\lambda_1)_{i12}(\lambda_1)^*_{i13}\left[\frac{\sin^2\theta_i}{t-m^2_{S_{i-}}}+\frac{\cos^2\theta_i}{t-m^2_{S_{i+}}}\right] .$$

The resulting differential cross section is

$$\frac{d\sigma_{-+}}{dt}=\frac{3}{16\pi s^2}\,|A(t)|^2(m_c^2-t)(m_t^2-t) .$$

In the heavy S limit the cross section is

$$\sigma_{-+}=\frac{s}{16\pi}\,|A(0)|^2\beta\left[\beta^2+\frac{3}{2}\frac{m_c^2+m_t^2}{s}\beta+\frac{3m_c^2m_t^2}{s^2}\right]$$
$$\simeq\frac{s}{16\pi}\,|A(0)|^2\left[1+\frac{m_t^2}{2s}\right]\left[1-\frac{m_t^2}{s}\right]^2 ,$$

where $\beta=\lambda^{1/2}(s,m_c^2,m_t^2)/s$ with $\lambda(a,b,c)=a^2+b^2+c^2-2ab-2bc-2ca$, and terms of order m_c^2/m_t^2 have been neglected in the last expression.

The cross section σ_{+-} for initial helicities $\eta=+$ and $\bar{\eta}=-$ is obtained by substitutions $\lambda_1\to\lambda_2$, $\cos\theta_i\to\sin\theta_i$, and $\sin\theta_i\to\cos\theta_i$ in $A(t)$. The σ_{++} and σ_{--} cross sections are proportional to the coupling product $\lambda_1\lambda_2\sin\theta_i\cos\theta_i$. However, such a nonvanishing product would give mass to the electron at one-loop-order proportional to $\lambda_1\lambda_2\sin\theta_i\cos\theta_i m_t$; see Fig. 2. Unless the product $\lambda_1\lambda_2$ is of order m_e/m_t, unnatural cancellations between tree-level and loop-level masses would be required. We therefore expect that at least λ_1 or λ_2 is small, but not necessarily both. This argument then implies that only one of σ_{-+} or σ_{+-} can be large, and the σ_{--} and σ_{++} cross sections should be negligible. With these constraints the resulting effective Lagrangian preserves approximate electron chiral symmetry, which also forbids large contributions to the electron anomalous magnetic moment.[19]

The cross-section expressions for the process $e^+e^-\to t\bar{u}$ ($\bar{t}u$) can be obtained from those for the $t\bar{c}$ ($\bar{t}c$) production by replacing m_c by m_u and changing generation indices of the couplings λ_1 and λ_2; e.g., the couplings $(\lambda_1)_{i12}$ in $A(t)$ for $\mathcal{M}_{-+}$ should be replaced by $(\lambda_1)_{i11}$.

If the Yukawa couplings are of the order of the electric charge, then an approximate estimate of the ΔR value of single-t cross section for unpolarized beams is

$$\Delta R=\frac{\sigma(e^+e^-\to t\bar{c}\text{ or }\bar{t}c)}{\sigma_{\rm pt}}$$
$$\simeq\frac{3}{8}\left[\frac{s}{m_S^2}\right]^2\left[1+\frac{m_t^2}{2s}\right]\left[1-\frac{m_t^2}{s}\right]^2 ,$$

where $\sigma_{\rm pt}=\sigma(e^+e^-\to\gamma^*\to\mu^+\mu^-)$, and we have set $A(0)=e^2/m_S^2$. Thus if $m_S\sim m_t$ the ΔR value near threshold is approximately $\frac{9}{16}(1-m_t^2/s)^2$. There is considerable ambiguity in this estimate; if the λ were instead of order $g=e/\sin\theta_W$ the rate would be a factor of 20 higher; if three almost degenerate S leptoquarks contribute constructively the rate would be increased by another factor of 9.

The experimental signature for such events would be spectacular. For $m_t<M_W$ production near threshold would give 3-hard jets of invariant mass m_t or a lepton, jet, and missing momentum. For $m_t>M_W+m_b$ the events would contain a single W boson. It would be hard to miss such signals.

Below threshold for $T(t\bar{u})+\bar{D}(\bar{c}u)$ meson pair production the $t\bar{c}$ system can form a quarkonium bound state.[20] Since only the $\mathcal{M}_{-+}$ or $\mathcal{M}_{+-}$ amplitudes can be nonnegligible, the S-wave $t\bar{c}$ bound state will be ${}^3S_1$, which we denote by T_c^*. The partial width for its decay to e^+e^- is given by

$$\Gamma(T_c^*\to e^+e^-)=\frac{m^2_{T_c^*}}{32\pi}\,|A(0)|^2|\psi(0)|^2 ,$$

where the $\psi(0)$ is the wave function at the origin. Since

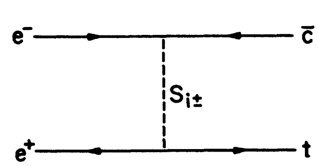

FIG. 1. Feynman diagram for $e^+e^-\to t\bar{c}$ via leptoquark exchange.

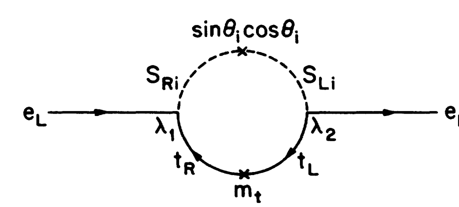

FIG. 2. One-loop contribution to the electron mass. There is also a corresponding diagram with $S\to\tilde{S}$ and $t\to\tilde{t}$ which is of the same order of magnitude.

the reduced mass μ of the $t\bar{c}$ system

$$\mu = \frac{m_t m_c}{m_t + m_c} \sim m_c$$

lies between that for the $J/\psi(m_c/2)$ and $\Upsilon(m_b/2)$, we may interpolate from their known leptonic widths to estimate that for $t\bar{c}$. By making use of the empirical relation, $|\psi(0)|^2/\mu^2 \simeq$ const for charmonium and bottomonium, we obtain

$$\frac{\Gamma(T_c^* \to e^+e^-)}{\Gamma(\Upsilon \to e^+e^-)} \simeq \frac{1}{2}\left[\frac{3m_t m_c\,|A(0)|}{4\pi\alpha}\right]^2 \simeq \frac{9}{2}\left[\frac{m_t m_c}{m_S^2}\right]^2,$$

where the last expression uses $A(0) \sim e^2/m_S^2$ as above. For $m_S \sim m_t$ the ratio becomes $\simeq 4.5(m_c/m_t)^2$ giving a width of $\Gamma(T_c^* \to e^+e^-) \simeq 5$ eV for $m_t = 50$ GeV; again this estimate could be 2 orders of magnitude low.

The produced T_c^* decays dominantly via the constituent t-quark decay, and the signal should be distinctive, as in the case of open $T\bar{D} + \bar{T}D$ production. For $m_{T_c^*} < m_W$, the T_c^* width is negligible compared to the beam spread, δW, and the total production rate is

$$\Delta R(T_c^*, \bar{T}_c^*) = \frac{2\sigma(e^+e^- \to T_c^*)}{\sigma_{\rm pt}} \approx \frac{\Gamma(T_c^* \to e^+e^-)}{5 \times 10^{-6}\delta W}.$$

Hence with a typical beam spread of $\delta W = 40$ MeV and the parameters above, $\Delta R \approx \frac{1}{40}$ and the resonances are detectable only when the couplings are significantly larger than for the estimate with $\lambda \sim e$. For $m_{T_c^*} > M_W + m_b$, $\Gamma_{\rm tot}(T_c^*) \approx \Gamma(t \to bW)$ becomes rapidly larger than the anticipated beam spread, making T_c^* detection still more difficult.

We also anticipate production of $T^*(t\bar{u})$ and $\bar{T}^*(\bar{t}u)$ about 1 GeV below the $T^*(t\bar{c})$ resonance position, although here we cannot use the potential model to reliably estimate the cross section.

We thank K. Igi for stimulating discussions. This research was supported in part by the University of Wisconsin Research Committee with funds granted by the Wisconsin Alumni Research Foundation, in part by the U.S. Department of Energy under Contract No. DE-AC02-76ER00881, in part by a 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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