Can Flavor-Independent Supersymmetric Soft Phases Be the Source of All *CP* Violation?

M. Brhlik,[1] L. Everett,[1] G. L. Kane,[1] S. F. King,[2] and O. Lebedev[3]

[1]*Randall Laboratory, Department of Physics, University of Michigan, Ann Arbor, Michigan 48109*
[2]*Department of Physics and Astronomy, University of Southampton, Southampton, S017 1BJ, United Kingdom*
[3]*Department of Physics, Virginia Polytechnic Institute and State University, Blacksburg, Virginia 24061*

(Received 7 October 1999)

Recently it has been demonstrated that large phases in softly broken supersymmetric (SUSY) theories are consistent with electric dipole moment constraints, and are motivated in some (type I) string models. Here we consider whether large flavor-independent soft phases may be the dominant (or only) source of all CP violation. In this framework, ϵ and ϵ'/ϵ can be accommodated, and the SUSY contribution to the B system mixing can be large and dominant. An unconventional flavor structure of the squark mass matrices (with enhanced super–Cabibbo-Kobayashi-Maskawa mixing) is required for consistency with B and K system observables.

PACS numbers: 12.60.Jv, 11.30.Er, 13.25.Es, 13.25.Hw

Although the first experimental evidence of CP violation was discovered over thirty years ago in the K system [1], the origin of CP violation remains an open question. In the standard model (SM), all CP violation arises due to a single phase in the Cabibbo-Kobayashi-Maskawa (CKM) quark mixing matrix [2]. While the SM framework of CP violation provides a natural explanation for the small value of ϵ in the K system and is supported by the recent CDF measurement of $\sin 2\beta$ through the decay $B \to \psi K_S$ [3], it is not clear whether the SM prediction is in agreement with the observed value of ϵ'/ϵ recently measured by [4] (confirming the earlier results of [5]) due to theoretical uncertainties [6]. However, the SM cannot account for the baryon asymmetry [7], and hence new physics is *necessarily* required to describe all observed CP violations.

In this paper, we investigate the possibility of a unified picture of CP violation by adopting the hypothesis that all observed CP violations can be attributed to the phases which arise in the low energy minimal supersymmetric standard model (MSSM), as first suggested by Frère *et al.* [8]. The issue of CP violation in supersymmetric theories is not a new question [8–12]. However, much of our analysis is motivated from embedding the MSSM into a particular string-motivated D-brane model at high energies [13], which departs significantly from the standard results for CP violation in supersymmetric (SUSY) models (which we summarize for the sake of comparison). The CP-violating phases of the MSSM can be classified into two categories: (i) the flavor-independent phases (in the gaugino masses, μ, etc.) and (ii) the flavor-dependent phases (in the off-diagonal elements of the scalar mass squares and trilinear couplings). We focus here on the flavor-independent phases; these phases have traditionally been assumed small ($\lesssim 10^{-2}$) if the sparticle masses are $\mathcal{O}(\mathrm{TeV})$ as the phases are individually highly constrained by the experimental upper bounds on the electric dipole moments (EDM's) of the electron and neutron [14–16]. However, a reinvestigation of this issue [17,18] has demonstrated that cancellations between different contributions to the EDM's can allow for viable regions of parameter space with phases of $\mathcal{O}(1)$ and light sparticle masses.

Recently [13], we found a (type I) string-motivated model of the soft breaking terms based on embedding the SM on five-branes in which large flavor-independent phases can be accommodated. The large relative phases between the gaugino mass parameters in this model play a crucial role in providing the cancellations in the EDM's, yielding regions of parameter space in which the electron and neutron EDM bounds are satisfied simultaneously. In this model, the CP-violating phases in the soft breaking terms are due to the (assumed) presence of complex F-component vacuum expectation values (VEV's) of moduli fields. Complex scalar moduli VEV's can, in principle, also lead to phases in the superpotential Yukawa couplings; however, for simplicity we assume here that the phase of the CKM matrix is numerically close to zero [19]. The crucial feature of our scenario compared to previous work is that all *flavor-independent* phases in the soft SUSY breaking sector can be large, with the EDM constraints satisfied by cancellations motivated by the underlying theory. We will show that SUSY can account for all observed CP violations with large flavor-independent phases (including the relative phases of the gaugino masses, which are zero in many SUSY models) and a particular flavor structure of the squark mass matrices. We focus on the low $\tan\beta$ regime, distinguishing our results from other recent work [12]. The baryon asymmetry can be explainable in SUSY [20]; see [21] for a study of baryogenesis within this approach.

The CP-violating and flavor-changing neutral-current processes that we consider are presented in Table I (we do not list the electron EDM [13,18], but only consider parameter sets which satisfy the electron and neutron EDM constraints). Note that generically the matrices which diagonalize the quark mass matrices and those which diagonalize the squark mass matrices are not equivalent due to SUSY breaking effects. The sfermion mass matrices are

TABLE I. We list the CP-violating observables and our dominant one-loop contributions (we work within the decoupling limit and hence neglect the charged Higgs). The third column schematically shows the flavor physics. Basically, the δ's are elements of the squark mass matrices normalized to some common squark mass, and the $\tilde{K}$'s are related to the Γ^U matrices defined in the text (with the stop mixing factored out, so they represent the family mixing only). Subscripts label flavor or chirality. The table is designed to demonstrate symbolically which observables are related (or not) to others. (More technically, in the down-squark sector, we utilize the $(\delta_{ij})_{AB}$ parameters of the mass insertion approximation. $\tilde{K}_{ij}$ labels the flavor factors which enter in diagrams involving up-type squarks. The flavor factors which enter the $b \to s\gamma$ and the n EDM amplitudes are different from the $\tilde{K}$ matrices but the flavor structure is similar (analogous statements apply for $D - \bar{D}$ mixing).

Observable	Dominant Contribution	Flavor Content
n EDM	$\tilde{g}, \tilde{\chi}^+, \tilde{\chi}^0$	$(\delta_{dd})_{LR}, \sim\tilde{K}_{ud}\tilde{K}^*_{ud}$
ϵ	$\tilde{g}$	$(\delta_{ds})_{LR}$
ϵ'	$\tilde{g}$	$(\delta_{ds})_{LR}$
Δm_K	SM, $\tilde{g}$	SM, $(\delta_{ds})_{LR}$
$K_L \to \pi\nu\bar{\nu}$	$\tilde{g}$	$(\delta_{ds})_{LR}$
Δm_{B_d}	$\tilde{\chi}^+$	$\lvert\tilde{K}_{tb}\tilde{K}^*_{td}\rvert$
Δm_{B_s}	SM, $\tilde{\chi}^+$	$\lvert\tilde{K}_{tb}\tilde{K}^*_{ts}\rvert$
$\sin 2\beta$	$\tilde{\chi}^+$	$\tilde{K}_{tb}\tilde{K}^*_{td}$
$\sin 2\alpha$	$\tilde{\chi}^+$	$\tilde{K}_{tb}\tilde{K}^*_{td}$
$\sin 2\gamma$	$\tilde{\chi}^+$	$\tilde{K}_{tb}\tilde{K}^*_{ts}$
$\mathcal{A}_{CP}(b \to s\gamma)$	$\tilde{\chi}^+$	$\sim\tilde{K}_{tb}\tilde{K}^*_{ts}$
Δm_D	$\tilde{g}$	$\sim\lvert\tilde{K}_{tc}\tilde{K}^*_{tu}\rvert$
n_B/n_γ	$\tilde{\chi}^+, \tilde{\chi}^0, \tilde{t}_R$	–

expressed in the super-CKM (SCKM) basis, in which the squarks and quarks are rotated simultaneously. In this basis the sfermion mass matrices are nondiagonal, and the amplitudes depend on the matrices $\{\Gamma^{\mathrm{SKM}}_{U,D_{L,R}}\}$ which rotate the squarks from the SCKM basis into the mass eigenstates. As shown schematically in Table I, particular processes are sensitive to certain elements of the quark and squark diagonalization matrices. We find that an unconventional flavor structure at the electroweak scale of the Γ^{SKM} matrices in the up-squark sector is *required* to reproduce the observed CP violations in the K and B systems:

$$\Gamma^{\mathrm{SKM}}_{U_L} = \begin{pmatrix} 1 & \lambda' + \lambda & \lambda' c_\theta & 0 & 0 & -\lambda' s_\theta e^{i\varphi_{\tilde{t}}} \\ -(\lambda' + \lambda) & 1 & \lambda' c_\theta & 0 & 0 & -\lambda' s_\theta e^{i\varphi_{\tilde{t}}} \\ -\lambda' & -\lambda' & c_\theta & 0 & 0 & -s_\theta e^{i\varphi_{\tilde{t}}} \end{pmatrix}, \tag{1}$$

$$\Gamma^{\mathrm{SKM}}_{U_R} = \begin{pmatrix} 0 & 0 & 0 & 1 & 0 & 0 \\ 0 & 0 & 0 & 0 & 1 & 0 \\ 0 & 0 & s_\theta e^{-i\varphi_{\tilde{t}}} & 0 & 0 & c_\theta \end{pmatrix}, \tag{2}$$

where $\lambda' \sim lt\lambda \equiv \sin\theta_c$, θ and $\varphi_{\tilde{t}}$ denote the stop mixing parameter and its phase, respectively, and entries of $\mathcal{O}(\lambda^2)$ are neglected. Note that the mixing in the LL sector is enhanced as compared to that of the SM, while in the RR sector it is negligible (this is easily seen by setting $\theta = 0$).

We now estimate the SUSY contributions to the observables in Table I. We will be working in the framework similar to the one laid out in [18], except we will also assume significant flavor mixing in the trilinear soft terms already at the grand unified theory scale. In particular, we assume the A terms to be of the form $e^{i\phi_A} B Y^{u,d} B'$, where B, B' are *real* matrices with considerable off-diagonal elements. Further, we assume that the squarks (except for the lightest stop) are degenerate in mass, retain only the lightest stop except in the cases of ϵ and ϵ' (for which the first two generations give the leading contribution), and neglect all but top quark masses unless the other fermion masses give leading contributions. For the purpose of presentation, we separate the stop left-right mixing from the family mixing. The family mixing matrices $\tilde{K}^{L,R}$ are defined as $\tilde{K}^L_{ij} = (\Gamma^{\mathrm{SKM}}_{U_L})_{ij}|_{\theta=0}$, $\tilde{K}^R_{ij} = (\Gamma^{\mathrm{SKM}}_{U_R})_{i,j+3}|_{\theta=0}$ with $i, j = 1, \ldots, 3$. In accordance with the chosen form of the Γ's, we assume $\tilde{K}^L_{ij} \sim \lambda/3$ and $\tilde{K}^R_{ij} \sim 0$ for $i \neg = j$. These matrices are *real*, as the only source of CP-violating phases in the Γ's is the stop mixing. We assume maximal chargino and stop mixings, and the following parameter values: $m_{\tilde{t}} \sim 140$ GeV, $m_{\tilde{\chi}} \sim 100$ GeV, $m_{\tilde{q}} \sim m_{\tilde{g}} \sim 350$ GeV, and $A \sim 250$ GeV. Our estimates agree within better than an order of magnitude with the numerical results to be presented in Ref. [22].

Let us first turn to the discussion of ϵ and ϵ'. Here we utilize the mass insertion approximation and the associated $(\delta_{ij})_{AB}$ parameters (see, e.g., Ref. [9]). Since we study the impact of *flavor-independent* phases at high energies, the LL and RR insertions are essentially real (their phases are produced effectively at the two-loop level; see renormalization group equations in [10]). The LR insertion always occurs in combination with the gluino phase φ_3 due to reparametrization invariance; the physical combination of phases is $(\delta_{12})_{LR} e^{i\phi_3}$ (the gluino phase has generally been neglected in earlier work). Our numerical studies show that the observed values of ϵ and ϵ' can be reproduced for $|(\delta^d_{12})_{LR}| \approx 3 \times 10^{-3}$ and $\arg[(\delta_{12})^d_{LR} e^{i\phi_3}] \approx 10^{-2}$, in agreement with [9,11]. This value of $|(\delta^d_{12})_{LR}|$ can be obtained in models with a large flavor violation in the A terms. Note also that this value of $(\delta^d_{12})_{LR}$ leads to a significant gluino contribution to Δm_K.

The leading chargino contribution to $(M_K)_{12}$ is CP conserving, as can be seen from

$$(M_K)^{\tilde{t}\tilde{\chi}}_{12} \sim \frac{g^4}{384\pi^2} \frac{m_K f_K^2}{m_{\tilde{t}}^2} (\tilde{K}^L_{td}\tilde{K}^{L*}_{ts})^2 |V_{11}T_{11}|^4, \tag{3}$$

(recall that $\tilde{K}^{L,R}$ are real). V and T denote the chargino and stop mixing matrices; to simplify this expression we employed the approximation $m^2_{\tilde{t}} \gg m^2_{\tilde{\chi}}$. This contribution gives $\Delta m_K \sim 10^{-16}$ GeV, well below the experimental value. Therefore Δm_K is dominated by the standard model and gluino contributions (as in [9,11]).

In our approach, the SM tree diagrams for B decays are real, and there is negligible interference with the super-penguin diagrams. Therefore the B system is essentially superweak, with all CP violation due to mixing. In contrast to the case of $K - \bar{K}$ mixing, $B - \bar{B}$ mixing is dominated by the chargino contribution:

$$(M_B)_{12}^{\tilde{t}\tilde{\chi}} \sim \frac{g^4}{384\pi^2}\,\frac{m_B f_B^2}{m_{\tilde{t}}^2}\,(\tilde{K}_{td}^{L}\tilde{K}_{tb}^{L*})^2 |V_{11}T_{11}|^4 \times \left(1 - \frac{h_t}{g}\,\frac{V_{12}^{*}T_{12}^{*}\tilde{K}_{tb}^{R*}}{V_{11}^{*}T_{11}^{*}\tilde{K}_{tb}^{L*}}\right)^2. \qquad (4)$$

The corresponding Δm_B is of order 10^{-13} GeV, which is roughly the observed value. The SM contribution to Δm_B is significantly smaller since the CKM orthogonality condition forces V_{td} to take its smallest allowed value. The CP-violating gluino contribution requires two LR mass insertions and, as a result, is suppressed by $(m_b/\tilde{m})^2$. Similar considerations hold for $B_s - \bar{B}_s$ mixing although the mixing phase is generally smaller than that in $B_d - \bar{B}_d$ due to a significant CP-conserving SM contribution.

Although the CP asymmetries and CKM entries are not related, $\sin 2\beta$ and $\sin 2\alpha$ can be defined in terms of the above asymmetries ($\sin 2\gamma$ can be defined via the CP asymmetry in $B_s \to \rho K_s$). The angles of the "unitarity triangle" given in this way need not sum to 180° as in the SM. Our results demonstrate that the chargino contribution alone is sufficient to account for the observed value of $\sin 2\beta$ reported in the CDF preliminary results [3]. This can be seen from (4) since the mixing phase can be as large as $\pi/2$ if $O(1)$ phases are present in V and T. In Fig. 1 we show contour plots of both $\sin 2\beta$ and Δm_B in the $\varphi_{\tilde{t}}$-φ_μ plane.

The CP asymmetries in $B \to \psi K_s$ and $B \to \pi^+\pi^-$ are related: $\sin 2\beta = -\sin 2\alpha$. This relation is characteristic of superweak models with a real CKM matrix [23,24], and is not consistent with the SM, as seen using the "sin" relation: $\sin\beta/\sin\alpha = |V_{ub}|/|V_{cb}\sin\theta_c|$. The left-hand side implies $|V_{ub}|/|V_{cb}\sin\theta_c| = 1$, while the experimental upper bound on the right-hand side is 0.45, verifying the nonclosure of the unitarity triangle.

We now turn to the $b \to s\gamma$ CP asymmetry $\mathcal{A}_{CP}(b \to s\gamma)$. The dominant contribution is due to mixing between the magnetic penguin operator Wilson coefficients C_7 and C_8 [25],

$$\mathcal{A}_{CP}(b \to s\gamma) \sim \frac{-4\alpha_s(m_b)}{9|C_7|^2}\,\mathrm{Im}(C_7 C_8^{*}). \qquad (5)$$

Both C_7 and C_8 receive real SM contributions, and hence the SUSY contribution from the chargino-stop loop has to be competitive while at the same time respect the experimental limits on $\mathcal{B}(b \to s\gamma)$. As a result, larger values of $\mathcal{A}_{CP}(b \to s\gamma)$ usually imply branching ratios farther away from the experimental central value. Typical results still predict asymmetries larger than in the SM (of an order of several percent).

We checked that the enhanced super-CKM mixing does not lead to an overproduction of $D - \bar{D}$ mixing. Since the chargino contribution is subject to strong Glashow-Iliopoulos-Maiani cancellations, the leading contribution is given by the gluino-stop loop:

$$(M_D)_{12}^{\tilde{t}\tilde{g}} \sim -\frac{\alpha_s^2}{27}\,\frac{m_D f_D^2}{m_{\tilde{g}}^2}\,\ln\frac{m_{\tilde{g}}^2}{m_{\tilde{t}}^2}\,(\tilde{L}_{tu}^{L}\tilde{L}_{tc}^{L*})^2 |T_{11}|^4, \qquad (6)$$

where $\tilde{L}$ is a real matrix which has roughly the same form as $\tilde{K}$. Δm_D is of order 10^{-14} GeV which corresponds to $x = (\Delta m/\Gamma)_{D^0}$ between 10^{-3} and 10^{-2}, which is in the range of the SM prediction and is consistent with recent CLEO measurements [26].

Next, consider the CP-violating decay $K_L \to \pi^0\nu\bar{\nu}$, which in the SM provides an alternate way to determine $\sin\beta$ [27]. It proceeds through a CP-violating Zds effective vertex, for which the dominant SUSY contribution is the chargino-stop loop [28]:

$$Z_{ds}^{\tilde{t}} \sim \frac{1}{4}\,\frac{m_{\tilde{\chi}}^2}{m_{\tilde{t}}^2}\,\ln\frac{m_{\tilde{\chi}}^2}{m_{\tilde{t}}^2}\,|V_{11}T_{11}|^2 \tilde{K}_{td}^{L*}\tilde{K}_{ts}^{L}. \qquad (7)$$

This contribution conserves CP, and thus we expect the branching ratio $K_L \to \pi^0\nu\bar{\nu}/K^+ \to \pi^+\nu\bar{\nu}$ to be $\mathcal{O}(\epsilon)$. This clearly violates the SM relation between the CP asymmetry in $B \to \psi K_s$ and the branching ratio of $K_L \to \pi^0\nu\bar{\nu}$. However, the CP-conserving (charged) mode of this decay is dominated by the SM and chargino contributions. Typically, we expect $Z_{ds}^{\tilde{t}}$ to be of order 10^{-4} which translates into the branching ratio for $K^+ \to \pi^+\nu\bar{\nu}$ of the order of 10^{-10}. In certain regions of the parameter space,

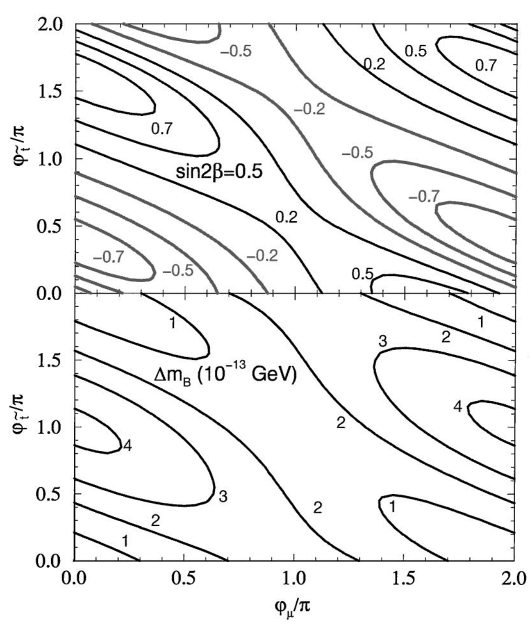

FIG. 1. Contours of $\sin 2\beta$ and Δm_B for $\lambda' = 0.07$, $\theta = \pi/5$, and the lightest stop mass $m_{\tilde{t}} \sim 140$ GeV. The absolute value of $\sin 2\beta$ can be as large as 0.78 for this choice of parameters. $\Delta m_B^{(\mathrm{exp})} \sim 3.1 \times 10^{-13}$ GeV and $\sin 2\beta^{(\mathrm{exp})} = 0.79 \pm 0.44$.

this branching ratio can be significantly enhanced (up to an order of magnitude) over the SM prediction.

In summary, our approach provides a unified view of all CP violation (including the baryon asymmetry [21]) which is testable at future colliders [29] and at the B factories, tying its origin to fundamental CP-violating parameters within a (type I) string-motivated context. CP violation in the K system is due mainly to the gluino-squark diagrams, with phases from the gluino mass M_3 and the trilinear coupling A. As the CKM matrix is by assumption (approximately) real, the B system is superweak: CP violation occurs mainly due to mixing. Therefore the unitarity triangle does not close, and we expect $\sin 2\beta/\sin 2\alpha \simeq -1$. Δm_K is dominated by the SM and gluino contributions, while Δm_B is dominated by the chargino-stop contribution. $K^+ \rightarrow \pi^+ \nu \bar{\nu}$ can be enhanced while $K_L \rightarrow \pi \nu \bar{\nu}$ is suppressed compared to the SM predictions. $D - \bar{D}$ mixing is expected to occur at a level somewhat below the current limit. The CP asymmetry in $b \rightarrow s\gamma$ can be considerably enhanced over its SM value. The electric dipole moments of the electron and neutron are suppressed by cancellations and should have values near the current limits. Our approach dictates an unconventional and interesting flavor structure for the squark mass matrices at low energies which is required for consistency with the preliminary experimental value of $\sin 2\beta$. An investigation of the connection of these matrices to the flavor structure of a basic theory at high energies is underway [22].

We would like to thank G. Good for helpful discussions and numerical work, J. Hewett for helpful suggestions, and S. Khalil for correspondence. This work is supported in part by the U.S. Department of Energy.

[1] J. Christenson, J. Cronin, V. Fitch, and R. Turlay, Phys. Rev. Lett. **13**, 138 (1964).
[2] M. Kobayashi and T. Maskawa, Prog. Theor. Phys. **49**, 652 (1973).
[3] CDF Collaboration, Report No. CDF/PUB/BOTTOM/CDF/4855, 1999.
[4] KTeV Collaboration, A. Alavi-Harati *et al.*, Phys. Rev. Lett. **83**, 22 (1999).
[5] NA31 Collaboration, G. D. Barr *et al.*, Phys. Lett. B **317**, 233 (1993).
[6] See, e.g., A. Buras, hep-ph/9905437; hep-ph/9908395, and references therein.
[7] M. Gavela *et al.*, Mod. Phys. Lett. A **9**, 795 (1994); Nucl. Phys. **B430**, 345 (1994); Nucl. Phys. **B430**, 382 (1994); P. Huet and E. Slather, Phys. Rev. D **51**, 379 (1995); G. R. Farrar and M. E. Shaposhnikov, Phys. Rev. D **50**, 382 (1994).
[8] J. M. Frère and M. Gavela, Phys. Lett. **132B**, 107 (1983); S. Abel and J. M. Frère, Phys. Rev. D **55**, 1623 (1997).
[9] F. Gabbiani *et al.*, Nucl. Phys. **B477**, 321 (1996).
[10] S. Bertolini *et al.*, Nucl. Phys. **B353**, 591 (1991).
[11] A. Masiero and H. Murayama, Phys. Rev. Lett. **83**, 9107 (1999).
[12] K. Babu and S. Barr, Phys. Rev. Lett. **72**, 2831 (1994); K. Babu *et al.*, hep-ph/9905464; S. Baek *et al.*, hep-ph/9907572, and references therein; S. Baek and P. Ko, hep-ph/9904283; S. Khalil, T. Kobayashi, and A. Masiero, Phys. Rev. D **60**, 075003 (1999); R. Barbieri *et al.*, hep-ph/9908255; A. Buras *et al.*, hep-ph/9908371; G. Eyal *et al.*, hep-ph/9908382; D. Demir, A. Masiero, and O. Vives, hep-ph/9909325; hep-ph/9911337; A. Kagan and M. Neubert, hep-ph/9908404; Y. Grossman *et al.*, hep-ph/9909297.
[13] M. Brhlik *et al.*, Phys. Rev. Lett. **83**, 2124 (1999); hep-ph/9908326.
[14] J. Ellis, S. Ferrara, and D. V. Nanopoulos, Phys. Lett. **114B**, 231 (1982); W. Buchmüller and D. Wyler, Phys. Lett. **121B**, 321 (1983); J. Polchinski and M. Wise, Phys. Lett. **125B**, 393 (1983); F. del Aguila *et al.*, Phys. Lett. **126B**, 71 (1983).
[15] M. Dugan, B. Grinstein, and L. Hall, Nucl. Phys. **B255**, 413 (1985).
[16] R. Garisto, Nucl. Phys. **B419**, 279 (1994).
[17] T. Ibrahim and P. Nath, Phys. Rev. D **58**, 111301 (1998); Phys. Rev. D **57**, 478 (1998); Phys. Rev. D **58**, 019901 (1998); Phys. Lett. B **418**, 98 (1998); T. Falk and K. Olive, Phys. Lett. B **439**, 71 (1998); S. Pokorski, J. Rosiek, and C. Savoy, hep-ph/9906206.
[18] M. Brhlik, G. Good, and G. L. Kane, Phys. Rev. D **59**, 115004 (1999).
[19] Within the type I string models, the matter fields arise from open strings which start and end on D-branes; in the heterotic language, they all have effective “modular weights” of -1, like untwisted sector fields. Therefore the trilinear superpotential couplings are $\mathcal{O}(1)$, and other mechanisms (e.g., higher-dimensional operators) will be required to obtain realistic fermion textures (in contrast to perturbative heterotic orbifold models, in which the moduli dependence of the twisted sector Yukawa couplings can allow for the generation of realistic fermion mass matrices at the trilinear order). The question of flavor physics within type I models is currently under investigation.
[20] See, e.g., M. Carena *et al.*, Nucl. Phys. **B503**, 387 (1997); J. Cline *et al.*, Phys. Lett. B **417**, 79 (1998); A. Riotto, Phys. Rev. D **58**, 095009 (1998), and references therein.
[21] M. Brhlik, G. Good, and G. L. Kane, hep-ph/9911243.
[22] M. Brhlik *et al.* (to be published).
[23] O. Lebedev, Phys. Lett. B **452**, 294 (1999).
[24] G. C. Branco, F. Cagarrinho, and F. Krüger, Phys. Lett. B **459**, 224 (1999).
[25] A. Kagan and M. Neubert, Phys. Rev. D **58**, 094012 (1998).
[26] CLEO collaboration, hep-ex/9908040.
[27] See, e.g., A. Buras, A. Romanino, and L. Silvestrini, Nucl. Phys. **B520**, 3 (1998).
[28] G. Colangelo and G. Isidori, JHEP **9809**, 009 (1998).
[29] G. L. Kane, S. Mrenna, and L. Wang, hep-ph/9910477.