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CP Violation and Moduli Stabilization in Heterotic Models[1]

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Abstract

The role of moduli stabilization in predictions for CP violation is examined in the context of four-dimensional effective supergravity models obtained from the weakly coupled heterotic string. We point out that while stabilization of compactification moduli has been studied extensively, the determination of background values for other scalars by dynamical means has not been subjected to the same degree of scrutiny. These other complex scalars are important potential sources of CP violation and we show in a simple model how their background values (including complex phases) may be determined from the minimization of the supergravity scalar potential, subject to the constraint of vanishing cosmological constant.

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Disclaimer

It has been argued by Dine et al. and Choi et al. [1] that CP is a gauge symmetry in string theory and that *explicit* breaking is therefore forbidden both perturbatively and nonperturbatively. They have shown that this is certainly true for weakly coupled heterotic orbifolds, the principal topic of this note. Thus, we envision *spontaneous* CP violation through complex scalar vacuum expectation values (*vevs*).

Compactification Moduli. One possible source is scalar fields corresponding to *compactification moduli* of the six-dimensional compact space. These are *Kähler moduli* (denoted with a "T") and *complex structure moduli* (denoted with a "U"). Ibáñez and Lüst [2] have enumerated the possibilities for Z_N and $Z_M \times Z_N$ orbifolds, based on the results of [3]:

$$\begin{aligned}
\text{case 1:} \quad \Gamma &= SL(2,Z)_T^3 \times SL(2,Z)_U^n \\
K &= -\sum_{I=1}^{3} \ln(T^I + \bar{T}^I) - \sum_{I=1}^{n} \ln(U^I + \bar{U}^I) \\
n &= 0: \quad Z_7, Z_8', \ldots \\
n &= 1: \quad Z_6, Z_8, \ldots \\
n &= 3: \quad Z_2 \times Z_2
\end{aligned} \tag{1}$$

$$\begin{aligned}
\text{case 2:} \quad \Gamma &= SL(2,Z)_T \times SL(2,2,Z)_T \times SL(2,Z)_U^n \\
K &= -\ln(T^1 + \bar{T}^1) - \ln\det(T + \bar{T}) - \sum_{I=1}^{n} \ln(U^I + \bar{U}^I) \\
& T \quad \text{is} \quad 2 \times 2, \qquad n = 0: \quad Z_6', \qquad n = 1: \quad Z_4
\end{aligned} \tag{2}$$

$$\begin{aligned}
\text{case 3:} \quad \Gamma &= SL(3,3,Z)_T: \quad Z_3 \\
K &= -\ln\det(T + \bar{T}), \quad T \quad \text{is} \quad 3 \times 3
\end{aligned} \tag{3}$$

Γ is the *target-space modular duality group.* It can be seen that each case has 3 *diagonal* Kähler moduli ($I = 1, 2, 3$):

$$T^I \equiv T^{II} \tag{4}$$

which transform under an $SL(2,Z)^3$ subgroup of the duality group Γ:

$$\begin{aligned}
T^I \quad \to \quad T'^I &= \frac{a^I T^I - ib^I}{ic^I T^I + d^I}, \\
a^I d^I - b^I c^I &= 1, \qquad a^I, b^I, c^I, d^I \in \mathbf{Z}
\end{aligned} \tag{5}$$

Enforcing this symmetry on the field theory limit leads to *modular invariant supergravity.* This symmetry constrains the stabilization of these moduli, and hence possible CP violating phases originating from $\arg(T^I)$ [4, 5, 6, 7].

FI induced vevs. Cancelation of the trace anomaly associated with an anomalous $U(1)_X$ by the GS mechanism [8] leads to an FI term [9] for the D-term of $U(1)_X$:

$$D_X = \sum_i K_i q_i^X \phi^i + \xi, \qquad \xi = \frac{g_H^2 \operatorname{tr} Q_X}{192\pi^2} m_P^2 \tag{6}$$

where g_H is the unified coupling at the string scale Λ_H and $m_P = 1/\sqrt{8\pi G} = 2.44 \times 10^{18}$ GeV is the reduced Planck mass. The fields ϕ^i which acquire nonvanishing vevs will be (following [5]) referred to as *Xiggs* fields in what follows.

In [10] it was shown that the presence of a $U(1)_X$ factor in the gauge group G is generic for semi-realistic orbifold models. For the class of standard-like orbifolds studied there, only 7 of 175 models did not have a $U(1)_X$. In the semi-realistic *free fermionic* models [11] a $U(1)_X$ is also generic.

In [12], the scalar potential V for SUGRA with a $U(1)_X$ was studied for vacuum configuations satyisfying $\langle V \rangle = \langle \partial V / \partial \phi^i \rangle = 0$. Supersymmetry breaking was characterized by

$$\langle |W|^2 \rangle = |\delta|^2, \quad \langle K_{i\bar{j}} F^i \bar{F}^{\bar{j}} \rangle = \alpha e^{\langle K \rangle} |\delta|^2 \tag{7}$$

According to expectations, it was found that $\langle V \rangle = \langle \partial V / \partial \phi^i \rangle = 0$ together with reasonable supersymmetry breaking scale requires

$$\langle D_X \rangle \sim |\delta|^2 \ll |\xi| \tag{8}$$

For canonical $K = \sum_i |\phi^i|^2$ and $\langle \phi^i \rangle = v^i$,

$$\langle D_X \rangle = \sum_i q_i^X |v_i|^2 + \xi \tag{9}$$

$$q_i^X \sim 1 \quad \Rightarrow \quad |v^i| \sim \sqrt{|\xi|} \tag{10}$$

Research in progress [13] has shown that (8,10) hold in cases more complicated than those studied in [12].

Based on (10), in [10] the $U(1)_X$ gauge symmetry breaking scale Λ_X was defined as

$$\Lambda_X = \sqrt{|\xi|} \tag{11}$$

For the class of models studied there it was found that for the 168 of 175 cases where $\xi \neq 0$,

$$\frac{g_H}{8.00} \leq \frac{\Lambda_X}{m_P} \leq \frac{g_H}{4.63} = \frac{\Lambda_H}{m_P} \tag{12}$$

where $\Lambda_H \approx 0.216 \times g_H m_P$ is the approximate string scale obtained in [14]. With $g_H \sim 1$ we have that $\Lambda_X \sim 0.1 \times m_P$ is a generic prediction.

The result of this is that nonrenormalizable operators should contribute *significantly* to the (effective) Yukawa couplings of the lighter quarks, since they are only down by $(\Lambda_X/m_P)^n \sim 10^{-n}, n > 0$. Given $\lambda_{u,d}/\lambda_t \sim 10^{-5}$ after running to the high scale, it is difficult to believe that nonrenomalizable operators would not play a role, generically speaking. Operators with $1 \leq n \leq 4$ would typically be present. Models with flat directions where this is not the case may be able to be found; for example, the 7 of 175 without a $U(1)_X$ found in [10]. However, if this is not the case generically, one can take the point of view that these models do not well represent "predictions" of string theory. That is, one can argue that in the absence of a dynamical vacuum selection mechanism in string theory, one should fall back on what is generic in extracting "predictions."

For these reasons, it was argued in [5] that FI induced vevs are more likely to be the dominant source of CP violation in string-derived models. These vevs are generically complex. Indeed, (9) is completely "phase-blind," leading to massless pseudoscalars termed *D-moduli* in [12]. Efforts are underway [13] to stabilize these moduli by including various terms in V (intentionally) neglected in [12]. Until this can be achieved, there is no *dynamical* reason to take the phases to be real. In [12] it was demonstrated that the generic case leads to nonzero KM phase.

Relevant orders. Renormalizing to $\Lambda_X \sim \Lambda_H \sim 10^{17}$ GeV,

$$\frac{m_u}{m_t} \sim 10^{-5} \tag{13}$$

for moderate values of $\tan\beta$. But for T^I stabilized at self-dual values, we do not get such a hierarchy from trilinear Yukawas. A natural source of such a hierarchy is the scale Λ_X. I.e., nonrenomalizable superpotential couplings can give effective Yukawa couplings from Xiggs vevs. The hierarchy is generically

$$\frac{\lambda_{eff}}{\lambda_{tri}} \sim \left(\frac{\Lambda_X}{m_P}\right)^n \tag{14}$$

where n counts dimensionality beyond trilinear. To get (13) we need $n \approx 5$. Thus, the structure of the Yukawas requires that we consider operators of rather large dimension in order to develop *predictions* in a realistic theory. The smallness of m_u, m_d, m_s requires that we work well beyond leading order to make firm conclusions about, say, V_{ub} vs. experiment. Not only should we work to higher orders in the superpotential couplings, but also in the Kähler potential since Xiggs vevs could give kinetic mixing which is not that small. E.g., $K \ni k_{ijk}\langle\phi_i\rangle u_j^\dagger u_k, j \neq k$. Such higher order terms are not well-understood and are not protected by nonrenormalization theorems. For the latter reason we will generally need to study supergravity loop corrections to any Kähler coupling extracted from string amplitudes at the string scale.

Note that all of these effects are highly dependent on the choice of Xiggs vevs. Which flat direction do we lead our expansion about? The results of [5] seem to suggest that the degeneracy is so great that we likely have the flexibility to tune the KM phase to any value we like. String theory has provided us with an empirical framework to match experimental data, while at the same time rendering our description quantum mechanically consistent in the ultraviolet limit with gravity included. But can we do better? Can we extract predictions?

Consider what Binétruy, Gaillard and Wu (BGW) did to extract predictions for the T-moduli phase [15]. They included nonperturbative effects from the hidden sector, coupled T^I to the gaugino condensate superfield U in the Veneziano-Yankielowicz lagrangian $\mathcal{L}_{VY}$ and imposed reparameterization symmetries. We must envision a similar analysis to stabilize the D-moduli if we are to stabilize them and predict Xiggs phases. The mixed anomaly

$$\text{tr } Q_X \neq 0 \quad \Rightarrow \quad \text{tr } T^a T^a Q_X \neq 0 \tag{15}$$

implies Xiggs couplings through gauge interactions with the hidden sector. A study of these effects is in progress.

Phase predictions from D-moduli stabilization. This is a modification of the linear multiplet toy model of [12]. The desire is to lift vacuum degeneracy by coupling D-moduli to matter condensates of the hidden sector condensing group G_C. As just mentioned, such couplings are expected from the mixed trace anomaly matching condition tr $T^a T^a Q_X \neq 0$, where T^a is a generator of G_C. The notation in what follows is defined in [12].

We continue to have

$$\begin{aligned} K &= k(L) + G(A, B, \Phi, \bar{A}, \bar{B}, \bar{\Phi}) \\ k(L) &= \ln L + g(L) \\ G &= \sum_i |A_i|^2 + \sum_i |B_i|^2 + \sum_i |\Phi_i|^2 \end{aligned} \tag{16}$$

On the other hand the superpotential now takes the form

$$\begin{aligned} W(A,B,\Phi,\Pi) &= \hat{W}(A,B,\Phi) + \breve{W}(\Phi,\Pi) \\ \hat{W}(A,B,\Phi) &= \lambda_{ijk} A_i B_j \Phi_k \\ \breve{W}(\Phi,\Pi) &= c_\alpha(\Phi)\Pi_\alpha \end{aligned} \tag{17}$$

The functional $c_\alpha(\Phi)$ is left unspecified at this point. The fields Π_α are chiral superfields corresponding to hidden sector matter condensate operators.

We implement dynamical supersymmetry breaking through the form of the Veneziano-Yankielowicz lagrangian assumed in [15]

$$\mathcal{L}_{VY} = \int \frac{E}{8R} U \left[b' \ln(e^{K/2}U) + \sum_\alpha b^\alpha \ln \Pi^\alpha \right] + \text{h.c.} \tag{18}$$

where U is the chiral superfield corresponding to the hidden sector gaugino bilinear condensate operator. We have no T-moduli appearing explicitly, no threshold corrections, and the only GS term is the one required to cancel the $U(1)_X$ anomaly. Following the BGW formulation with these simplifications one obtains for the scalar potential

$$\begin{aligned} V &= \frac{1}{2}\left(\frac{2\ell}{1+f}\right)\sum_a D_a D_a + (\ell g' - 2)\left|\frac{b' u}{4} - e^{K/2} W\right|^2 \\ &\quad + \left| e^{K/2}(\hat{W}_I + W G_I) - \frac{b' u}{4} G_I \right|^2 \\ &\quad + \left(\frac{1+\ell g'}{16\ell^2}\right)\left[(1+2\ell b')|u|^2 - \ell e^{K/2}(W\bar{u} + \bar{W} u)\right] \end{aligned} \tag{19}$$

Here, $\ell = L|, u = U|$.

We restrict our attention to the submanifold of the vacuum manifold where D_X is the only nonvanishing D-term, and $\langle a_i \rangle = \langle b_i \rangle = 0$, with $a_i = A_i|$, $b_i = B_i$. In this case the vacuum image of V is, after straightforward manipulations, given by

$$V = \frac{1}{2} g_H^2 D_X^2 + \hat{V} \tag{20}$$

$$\hat{V} = e^K \sigma^2 \left[b_c^2 (v^2 - 2 + \ell g') + \left(\frac{1+\ell g'}{2\ell^2}\right)(2 + 3\ell b' + \ell b_c)\right] \tag{21}$$

where we take ℓ at its vev and

$$\begin{aligned} v_i &= \langle \phi_i \rangle, \qquad v = \left[\sum |v_i|^2\right]^{1/2} \\ D_X &= \sum_i q_i |v_i|^2 + \xi, \qquad g_H^2 = \frac{2\ell}{1+f} \\ K &= k(\ell) + v^2, \qquad b_c = b' + \sum_\alpha b^\alpha \\ \sigma &= \frac{1}{4}\exp\left[-\frac{1}{b_c g_H^2} - \frac{b'}{b_c}\right] \prod_\alpha \left|\frac{4 c_\alpha(v)}{b^\alpha}\right|^{b^\alpha / b_c} \end{aligned} \tag{22}$$

From these expressions it is not hard to work out $V_i = \partial V/\partial v_i$:

$$V_i = \bar{v}_i \left[g_H^2 D_X q_i + \sigma^2 b_c^2 e^K + \hat{V} \right] + \mu_i \hat{V} \tag{23}$$

where

$$\mu_i = \sum_\alpha \frac{b^\alpha}{b_c} \frac{\partial}{\partial v_i} \ln c_\alpha(v) \tag{24}$$

Notice that all of the quantities in (23) are real except $\bar{v}_i$ and μ_i. Thus, the phase of v_i will be related to μ_i. We expect that this provides the necessary constraint to lift pseudoscalar D-moduli. More precisely,

$$\arg v_i = -\arg \mu_i \bmod \pi \tag{25}$$

For minimization with vanishing cosmological constant we require $V = V_i = 0$, which implies $\hat{V} \sim -D_X^2$. To have low scale supersymmetry breaking we demand $|D_X| \ll 1$, so from (23) $|D_X| \sim \sigma^2 \ll 1$. Then for $v_i \neq 0$ eq. (23) implies

$$g_H^2 D_X q_i + b_C^2 \sigma^2 e^K = \mathcal{O}(\sigma^4) \tag{26}$$

unless $v_i \lesssim \sigma^2$. Thus the delicate cancelation of $\mathcal{O}(\sigma^2)$ terms implied by (25) can, for reasonable splittings of q_i, only be achieved for *one choice* of charge q_i, just as was the case in [12]. Since the term $\mu_i \hat{V}$ in (23) is roughly $\mathcal{O}(\sigma^4)$, the result of [12] for *which* q_i gets vevs $v_i \gg \sigma^2$ is unchanged: the minimum q_i must be the set getting $v_i \gg \sigma^2$ vevs for positive mass-squared. To conclude,

$$\begin{aligned} |v_i| \quad &\sim \quad \begin{cases} \Lambda_X & q_i = -q \\ |\mu_i|\sigma^2 & q_i \neq -q \end{cases} \\ q \quad &= \quad -\min\{q_i\}, \quad \Lambda_X = \sqrt{\xi} \end{aligned} \tag{27}$$

since $D_X = \sum_i q_i |v_i|^2 + \xi \sim \sigma^2 \ll \xi$. From (23) we have

$$v^2 = \frac{\xi}{q} - \frac{b_C^2 \sigma^2 e^K}{q^2 g_H^2} + \mathcal{O}(\sigma^4) \tag{28}$$

We next suppose in (17)

$$c_\alpha(v) = \sum_A c_{\alpha A}(v), \qquad c_{\alpha A}(v) = \lambda_{\alpha A} \prod_i (v_i)^{p_{iA}^\alpha}. \tag{29}$$

Then it is easy to check that (24) yields

$$v_i \mu_i = \sum_\alpha \frac{b^\alpha}{b_c} \frac{\sum_A p_{iA}^\alpha c_{\alpha A}(v)}{\sum_A c_{\alpha A}(v)} \tag{30}$$

Consequently we can rewrite the minimization constraint which follows from (23) as (exactly)

$$0 = |v_i|^2 \left[g_H^2 D_X q_i + b_c^2 \sigma^2 e^K + \hat{V} \right] + \hat{V} \sum_\alpha \frac{b^\alpha}{b_c} \frac{\sum_A p_{iA}^\alpha c_{\alpha A}(v)}{\sum_A c_{\alpha A}(v)} \tag{31}$$

In the case where the sum in (29) has only a single term, the $c_{\alpha A}(v)$ cancel in (31) and no phase constraints exist. Thus, a non-monomial polynomial assumption for $c_\alpha(v)$ is required for phase stabilization.

As an example, consider the case of only two fields ϕ^1, ϕ^2 of charges $q_1 = q_2 \equiv -q$ and a single matter condensate field with superpotential coupling

$$\breve{W}(\phi, \pi) = c(\phi)\pi, \qquad c(\phi) = \lambda_1\phi_1 + \lambda_2\phi_2 \tag{32}$$

We can go to U-gauge by writing

$$\begin{pmatrix} \phi_1(x) \\ \phi_2(x) \end{pmatrix} = e^{i(\theta(x)+\varphi_+)} \begin{pmatrix} e^{i\varphi_-}\cos\eta & e^{i\varphi_-}\sin\eta \\ -e^{-i\varphi_-}\sin\eta & e^{-i\varphi_-}\cos\eta \end{pmatrix} \begin{pmatrix} d(x) \\ h(x)+v \end{pmatrix} \tag{33}$$

where $\langle\theta(x)\rangle = \langle d(x)\rangle = \langle h(x)\rangle = 0$ and

$$\begin{aligned} v_1 &= e^{i\varphi_1} v\cos\eta, \qquad v_2 = e^{i\varphi_2} v\sin\eta, \\ \varphi_\pm &= \frac{1}{2}(\varphi_1 \pm \varphi_2) \end{aligned} \tag{34}$$

Here, $\theta(x) + \varphi_+$ is eaten by the $U(1)_X$ vector boson, while the scalar $h(x) + v$ acquires a Λ_X mass and fills out the massive $U(1)_X$ vector multiplet. φ_- is the phase we would like to stabilize and $d(x)$ is the complex scalar D-modulus field. η is the mixing angle to the mass eigenstate basis which we would also like to stabilize. It is not hard to check that (23) gives

$$0 = b_c(\lambda_1|v_1|^2 + \lambda_2\bar{v}_1v_2)(-qg_H^2D_X + b_c^2\sigma^2e^K + \hat{V}) + b^\alpha\lambda_1\hat{V} \tag{35}$$

and a similar equation with $1 \leftrightarrow 2$, and then 2 conjugate equations. Manipulations on these four equations lead simply to

$$\frac{v_1\bar{v}_2}{\bar{v}_1v_2} = \frac{\bar{\lambda}_1\lambda_2}{\lambda_1\bar{\lambda}_2} \Rightarrow \varphi_1 - \varphi_2 = 2\varphi_- = \arg(\frac{\lambda_2}{\lambda_1}) \bmod \pi \tag{36}$$

It is also straightforward to check

$$\sin^2\eta = \frac{b^\alpha\hat{V}(|\lambda_1|^2 - |\lambda_2|^2) + b_cv^2|\lambda_1|^2(-qg_H^2D_X + b_c\sigma^2e^K + \hat{V})}{2b_cv^2|\lambda_2|^2(-qg_H^2D_X + b_c\sigma^2e^K + \hat{V})} \tag{37}$$

Thus $d(x)$ is stabilized and the phase and mixing are determined. In particular, the phase φ_- is not an independent source of CP violation, but is determined by whatever mechanism determines the phases of λ_i.

If we embed this toy into a string-inspired model, the source of $\arg(\lambda_2/\lambda_1)$ would be the phase of T-moduli. For example, suppose the modular weight of π is q_π^I while for the fields $\phi_{1,2}$ we have $q_{1,2}^I$. Then treating $\phi_{1,2}, \pi$ as untwisted fields we have

$$\lambda_i \sim \prod_I [\eta(t^I)]^{2(q_i^I + q_\pi^I - 1)} \qquad (i = 1, 2) \tag{38}$$

Then

$$\frac{\lambda_2}{\lambda_1} \sim \lambda_i \sim \prod_I [\eta(t^I)]^{2(q_2^I - q_1^I)} \tag{39}$$

From (36) we have

$$\varphi_1 - \varphi_2 = 2\sum_I (q_2^I - q_1^I)\arg(\eta(t^I)) \bmod \pi \tag{40}$$

If $\phi_{1,2}$ are untwisted with different compact space $SO(6)$ weights (*H-momenta*), say $(1,0,0)$ and $(0,1,0)$ resp., then

$$\varphi_1 - \varphi_2 = 2\left[\arg(\eta(t^2)) - \arg(\eta(t^1))\right] \bmod \pi \tag{41}$$

For $t^1 = e^{i\pi/6}$, $t^2 = 1$,

$$\arg(\eta(t^1)) = -\pi/24, \qquad \arg(\eta(t^2)) = 0 \tag{42}$$

and hence

$$\varphi_1 - \varphi_2 = \frac{\pi}{12} \bmod \pi. \tag{43}$$

In the case of twisted fields ϕ_1 or ϕ_2, the linear coupling in (32) implies that π is also twisted by the point group selection rule. In this case a major revision to (18) would be required because of mixing under $SL(2,Z)_T^3$. A mixing of π with other operators would not give the right sort of anomalous modular transformation for $\mathcal{L}_{YV}$ which could be canceled by $\mathcal{L}_{GS}$.

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References

[1] M. Dine, R. G. Leigh and D. A. MacIntire, Phys. Rev. Lett. 69 (1992) 2030; K. Choi, D. B. Kaplan, A. E. Nelson, Nucl. Phys. B (1993) 515 [hep-ph/9205202].

[2] L. E. Ibañez and D. Lüst, Nucl. Phys. B 382 (1992) 305.

[3] S. Ferrara, C. Kounnas and M. Porrati, Phys. Lett. B 181 (1986) 263; M. Cvetič, J. Louis and B. A. Ovrut, Phys. Lett. B 206 (1988) 227.

[4] D. Bailin, G. V. Kraniotis and A. Love, Phys. Lett. B 357 (1995) 387; B 414 (1997) 269; B 432 (1998) 343; Nucl. Phys. B 518 (1998) 92; Phys. Lett. B 483 (2000) 425 [hep-th/0004052].

[5] J. Giedt, Nucl. Phys. B 595 (2001) 3.

[6] T. Dent, hep-th/0011294; Phys. Rev. D 64 (2001) 056005 [hep-ph/0105285]; Nucl.Phys. B623 (2002) 73 [hep-th/0110110]; PhD Thesis, Univ. of Sussex, UK, 2000.

[7] O. Lebedev, Phys. Lett. B 521 (2001) 71 [hep-th/0108218]; S. Khalil, O. Lebedev and S. Morris, hep-th/0110063; O. Lebedev and S. Morris, hep-th/0203246.

[8] M. B. Green and J. H. Schwarz, Phys. Lett. B 149 (1984) 117.

[9] M. Dine, N. Seiberg and E. Witten, Nucl. Phys. B 289 (1987) 585; J. J. Atick, L. Dixon and A. Sen, Nucl. Phys. B 292 (1987) 109; M. Dine, I. Ichinose and N. Seiberg, Nucl. Phys. B 293 (1987) 253.

[10] J. Giedt, Ann. of Phys. (N.Y.) 297 (2002) 1 [hep-th/0108244].

[11] G. B. Cleaver, A. E. Faraggi and D. V. Nanopoulos, Phys. Lett. B 455 (1999) 135; Int. J. Mod. Phys. A 16 (2001) 425; G. B. Cleaver, A. E. Faraggi, D. V. Nanopoulos, and J. W. Walker, Mod. Phys. Lett. A 15 (2000) 1191; Nucl. Phys. B 593 (2001) 471; Nucl. Phys. B 620 (2002) 259 [hep-ph/0104091].

[12] M. K. Gaillard and J. Giedt, Phys. Lett. B 479 (2000) 308.

[13] M. K. Gaillard and J. Giedt, in preparation.

[14] V. S. Kaplunovsky, Nucl. Phys. B 307 (1988) 145; Erratum B 382 (1992) 436.

[15] P. Binétruy, M. K. Gaillard and Y.-Y. Wu, Nucl. Phys. B 481 (1996) 109; B 493 (1997) 27; Phys. Lett. B 412 (1997) 288.