F-flatness and universality in broken Supergravity

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Abstract

We show that in a broken supergravity theory any field that acquires a vacuum expectation value obtains at the same time an F-term of order $\langle\theta\rangle m_{3/2}$. These F-terms contribute unsuppressed to the trilinear couplings and have observable effects in the low energy phenomenology. We show that $\mu \to e\gamma$ decays can receive large contributions from this source at the level of current experimental bounds. Therefore this decay may provide the first clue of the structure of the soft breaking sector with the proposed experiments in the near future.

1 Motivations

Our knowledge of the Supersymmetry breaking sector in the Minimal Supersymmetric extension of the Standard Model (MSSM) is still very small due to the limited experimental information available. Only after the discovery of SUSY particles and the measurement of the Supersymmetric spectrum we will be able to explore in detail and improve our understanding on this fundamental piece of the MSSM. Nevertheless, we have already today a lot of useful information on this sector from experiments looking for indirect effects of SUSY particles in low-energy experiments [1]. In fact, it was readily realized at the beginning of the SUSY phenomenology era that large contributions to Flavour Changing Neutral Currents (FCNC) and CP violation phenomena were expected in Supersymmetric theories with a generic soft breaking sector. The absence of these effects much beyond the level most theorists considered reasonable became to be known as the SUSY flavour and CP problems. These problems are closely related to the flavour structure of the soft breaking sector and therefore to the source of flavour itself. In fact, the main solutions to these problems are either exact universality of the soft breaking terms or alignment to the Yukawa matrices. However, as we will show here, in any model where we intend to explain at the same time the strongly hierarchical structure of the Yukawa matrices, these two solutions are not perfect and the soft breaking terms are necessarily nonuniversal and not aligned with the Yukawa couplings. This is due to the fact that in a theory with broken local Supersymmetry, i.e. Supergravity, any field that acquires a vacuum expectation value (vev) obtains at the same time a non-vanishing F-term [2, 3].

To prove this we consider a generic supergravity scenario specified in terms of the superpotential W and the Kähler potential, $K = \hat{K}(\chi, \chi^*) + \sum_i K_i^i(\chi, \chi^*)|\varphi_i|^2$. Where K is a real function of the chiral superfields and χ are general hidden sector fields. In these conditions, the F-term contribution to the scalar potential is given by,

$$V = e^K \left[\sum_i (K^{-1})_i^j \tilde{F}^i \tilde{F}_j - 3|W|^2 \right] \tag{1}$$

with $\tilde{F}^i = \partial W/\partial\phi_i + K_i^i \phi^{i*} W$ and $\tilde{F}_i = \partial W/\partial\phi^{i*} + K_i^i \phi_i W^*$ related to the correctly normalized supergravity F-terms as $F^i = e^{K/2}\tilde{F}^i$ and $(K^{-1})_j^i$ is the inverse of the Kähler metric. We study a situation of broken supergravity where $m_{3/2} = e^K/2\langle W\rangle \neq 0$ and we have a non-vanishing vev for a certain field θ after minimization of the scalar potential. In particular we are interested in the case of fieds that obtain small vevs in units of the Plank mass and generate the hierarchy in the Yukawa couplings. So, we take this field as a singlet under the SM gauge symmetries but charged under a new flavour symmetry that controls the Yukawa structure. The natural scale of the F-terms with $m_{3/2} \neq 0$ is given by $m_{3/2}\langle\theta\rangle$ as this is the supergravity correction to the globally supersymmetric F-term. However this natural size could be reduced by some cancellation from other terms. In the case of a continuous flavour symmetry controlling the Yukawa hierarchy and when only the field θ acquires a vev, we have $\partial W/\partial\theta = 0$ at the minimum and the vev of θ will be driven by some scale dependent soft masses in analogy to the radiative symmetry breaking in the MSSM [4]. Hence it is strait-forward from the definition that $F^\theta = m_{3/2}\langle\theta\rangle$ as expected.

On the other hand in the presence of a discrete symmetry or a continuous symmetry with several fields ϕ_i which acquire vevs, the situation is more involved. The superpotential is of the kind, $W = a + b\ \theta^p \phi^q$ and the relevant part of the scalar potential is (after absorbing the Kähler metric in a redefinition of the fields),

$$\begin{aligned} V_{\theta\phi} \;\; = \;\; & \left|p\ b\ \theta^{p-1}\phi^q + \theta^* m_{3/2}\right|^2 + \left|q\ b\ \theta^p \phi^{q-1} + \phi^* m_{3/2}\right|^2 - 3\left|m_{3/2} + b\ \theta^p\phi^q\right|^2 + \\ & m_r^2(\theta) \mid \theta \mid^2 + A_r(\theta)\ b\ (p+q-3)\left(\theta^p\phi^q + \theta^{p*}\phi^{q*}\right) \end{aligned} \tag{2}$$

where we allowed the generation of D-terms at radiative order, $m_\theta^2(\theta)$, in order to generate a minimum with $\theta \neq 0$. D-flatness aligns the θ and ϕ vevs and we get $|\theta|^2/|\chi|^2 = p/q$ $\theta\ \partial V/\partial\theta \ - \ \theta^*\ \partial V/\partial\theta^* \ = 0$ implies that Im $\theta^p\chi^q = 0$ and we choose them as real. Using these relations and after some algebra, we obtain in terms of a new variable $X \equiv \frac{B}{m_{3/2}}\ p\ \left(\frac{q}{p}\right)^{q/2}\ \theta^{p+q-2}$,

$$\frac{\partial V}{\partial\theta} \;=\; m_{3/2}\ \theta\ \left[X^2(p+q-1) + X\ (1-\alpha)\ (p+q-3) + 1 - \beta\right] = 0 \tag{3}$$

where $\alpha \equiv \dfrac{A(\theta)}{m_{3/2}}$ and $\beta \equiv \dfrac{m_1(\theta)^2}{m_{3/2}^2}\dfrac{p}{p+q}$. The solutions with nonvanishing θ vev correspond to the roots in the quadratic equation in X. In terms of X the F terms are now,

$$F^\theta = \; m_{3/2}\ \langle\theta\rangle\ (1+X) \equiv \gamma\ \langle\theta\rangle\ m_{3/2} \tag{4}$$

and similarly for F^χ. Clearly there is no reason to expect a conspiracy between radiative terms and the Superpotential structure to obtain $X \simeq -1$ and so $\gamma = \mathcal{O}(1)$. Thus, terms involving θ also contribute to the soft SUSY breaking terms.

At this point, it is clear that any field θ that obtains a vev in the presence of broken SUSY, acquires simultaneously a non-vanishing F-term of order $\langle\theta\rangle m_{3/2}$ [1]. Thus, they also contribute to the SUSY soft breaking terms. However, these contributions are in principle small, suppressed by the small vev of θ, when compared with the dominant contributions of order $m_{3/2}$. Therefore, the main question now is whether these additional contributions can have any sizeable effect in the low energy physics. So we analyze the effective theory below the scale of flavour symmetry breaking with a W,

$$W = W^{hid}(\chi_k) + \left(\frac{\theta}{M}\right)^{\alpha_{ij}} H_a Q_{Li} q^c_{Rj} + \dots \tag{5}$$

where the hierarchical structure of the Yukawa couplings is generated through effective operators in terms of θ vevs suppressed by a heavy mediator scale which can be M_{Pl}.

Then, from Eq. (1) we get the soft breaking terms in the observable sector after SUSY breaking. At the minimum we obtain the trilinear terms as [5],

$$A_{ij}\hat{Y}^{ij} = F^{\chi_k}\hat{K}_{\chi_k}Y^{ij} + \alpha_{ij}\,\frac{e^{K/2}}{M}\left(\frac{\theta}{M}\right)^{\alpha_{ij}-1} m_{3/2}\theta \tag{6}$$

with $\hat{K}_{\chi_k} = \partial\hat{K}/\partial\chi_k$ and $Y^{ij} = e^{K/2}(\theta/M)^{\alpha_{ij}}$. The presence of the α_{ij} in the right hand side is due to the dependence of the effective Yukawa couplings on θ. Still the smallness of θ does not affect the trilinear couplings because is reabsorbed in the Yukawa coupling itself,

$$m_{3/2}\ \langle\theta\rangle\ \frac{\partial Y^{ij}}{\partial\theta} = \alpha(i,j)\ m_{3/2}\ Y^{ij}. \tag{7}$$

Therefore, in any framework explaining the hierarchy in the Yukawa textures through nonrenormalizable operators, the trilinear couplings are necessarily nonuniversal as in Eq. (6). In a similar way, we can also expect non-renormalizable contributions to the Kähler potential of the kind $(\theta\theta^*/M^2)^{\alpha(i,j)}$. Nevertheless these contributions appear only at order $2\alpha(i,j)$ in θ/M with respect to the dominant term $\mathcal{O}(1)$. So, in the following we concentrate in the nonuniversal trilinear couplings.

Next, we must check whether this breaking of universality does not contradict any of the very stringent bounds from low energy phenomenology. Unfortunately we do not have still a unique and complete theory of flavour that provides the full field dependence of the low energy effective Yukawa couplings [6]. Still, the analysis of the fermion masses and mixing angles seems to point to some definite textures for the Yukawa matrices [7],

$$\frac{M}{m_3} = \begin{pmatrix} 0 & b\epsilon^3 & c\epsilon^3 \\ b'\epsilon^3 & d\epsilon^2 & a\epsilon^2 \\ f\epsilon^m & g\epsilon^n & 1 \end{pmatrix}, \tag{8}$$

[1] As shown in [2] it is still possible to get a further reduction if the familon field has a vev in the SUSY limit. However, to obtain this vev requires a significantly more complicate theory and this is the correct order of magnitude for the F-terms in most simple models.

with $\epsilon_d = \sqrt{m_s/m_b} = 0.15$ and $\epsilon_u = \sqrt{m_c/m_t} = 0.05$. at the unification scale, and a, b, b', c, d, g, f coefficients $\mathcal{O}(1)$ and complex in principle. In this texture there are still two undetermined elements, $(3,1)$ and $(3,2)$, that would account for the unmeasured right-handed quark mixings although fermion masses require $m \geq 1$. Next, we assume that the hierarchy in the Yukawa matrices is generated by different powers in the vevs of scalar fields, $\epsilon_a = \langle\theta_a\rangle/M$, and we can immediately calculate the nonuniversality in the trilinear terms, $(Y^A)_{ij} \equiv Y_{ij} A_{ij}$,

$$(Y^A)_{ij} = A_0 Y_{ij} + m_{3/2}\; Y_{33} \begin{pmatrix} 0 & 3b\epsilon^3 & 3c\epsilon^3 \\ 3b\epsilon^3 & 2d\epsilon^2 & 2a\epsilon^2 \\ fm\epsilon^m & gn\epsilon^n & 0 \end{pmatrix} \tag{9}$$

with $A_0 = F^{\chi_k}\hat{K}_{\chi_k}$. The Yukawa texture in Eq. (8) is diagonalized by superfield rotations in the so-called SCKM basis, $\tilde{Y} = V_L^\dagger \cdot Y \cdot V_R$. However, in this basis large off-diagonal terms necessarily remain in the trilinear couplings, $\tilde{Y}^A = V_L^\dagger \cdot Y^A \cdot V_R$. The phenomenologically relevant flavour off-diagonal entries for FCNC contributions in the basis of diagonal Yukawa matrices are,

$$(\tilde{Y}^A)_{32} \simeq Y_{33}\; m_{3/2}\; g\; n\; \epsilon^n + \ldots \tag{10}$$
$$(\tilde{Y}^A)_{21} \simeq Y_{33}\; m_{3/2}\; \epsilon^3 (b' + a\; (\frac{b'}{d}\; g\; n\; \epsilon^n - f\; m\; \epsilon^{m-1}))$$

We take these matrices as the boundary conditions at a high scale, typically M_{GUT}. Then we must use the MSSM Renormalization Group Equations (RGE) [8] to obtain the corresponding matrices at the electroweak scale. The main effect in this RGE evolution is a large flavour universal gaugino contribution to the diagonal elements in the sfermion mass matrices (see for instance Tables I and IV in [9]). In this minimal supergravity scheme we take gaugino masses as $m_{1/2} = \sqrt{3} m_{3/2}$ and sfermion masses $m_0^2 = m_{3/2}^2$. So, the average squark and slepton masses,

$$m_{\tilde{q}}^2 \simeq 6 \cdot m_{1/2}^2 + m_0^2 \simeq 19\; m_{3/2}^2, \qquad m_{\tilde{l}}^2 \simeq 1.5 \cdot m_{1/2}^2 + m_0^2 \simeq 5.5\; m_{3/2}^2 \tag{11}$$

The RG evolution of the trilinear terms is also similarly dominated by gluino contributions and the third generation Yukawa couplings. However, the offdiagonal elements in the down and lepton trilinear matrices are basically unchanged for $\tan\beta \leq 30$ [9]. From here we can obtain the full trilinear couplings and compare with the experimental observables at low energies. The so-called Mass Insertions (MI) formalism is very useful in this framework. The left–right MI are defined in the SCKM basis as $(\delta_{LR})_{ij} = (m_{LR}^2)_{ij}/m_{\tilde{f}}^2$, with $m_{\tilde{f}}^2$ the average sfermion mass. We can estimate the value of $(\delta_{LR})_{21}$ as,

$$\left(\delta^d_{LR}\right)_{21} \simeq \frac{m_b\, \epsilon_d^3}{19\, m_{3/2}} (b' + a\, \frac{b'}{d}\, g\, n\, \epsilon_d^n - a\, f\, m\, \epsilon_d^{m-1}) \simeq$$
$$(b' + a\, \frac{b'}{d}\, g\, n\, \epsilon_d^n - a\, f\, m\, \epsilon_d^{m-1})\; 7.5 \times 10^{-6} \tag{12}$$

using $m_{\tilde{q}} \simeq 500$ GeV corresponding to $m_{3/2} \simeq 120$ GeV and $\epsilon_d \simeq 0.15$ [7]. We can compare our estimate for the mass insertion with the phenomenological bounds in Table

x	$\sqrt{\left\|\mathrm{Im}\left(\delta^d_{LR}\right)^2_{12}\right\|}$	$\sqrt{\left\|\mathrm{Re}\left(\delta^d_{13}\right)^2_{LR}\right\|}$	$\left\|\left(\delta^l_{LR}\right)_{12}\right\|$	$\left\|\left(\delta^l_{LR}\right)_{23}\right\|$
0.3	1.1×10^{-5}	1.3×10^{-2}	6.9×10^{-7}	8.7×10^{-3}
1.0	2.0×10^{-5}	1.6×10^{-2}	8.4×10^{-7}	1.0×10^{-2}
4.0	6.3×10^{-5}	3.0×10^{-2}	1.9×10^{-6}	2.3×10^{-2}

Table 1: MI bounds from ε'/ε, $b\to s\gamma$, $\mu\to e\gamma$ and $\tau\to\mu\gamma$ for an average squark mass of 500 GeV and different values of $x=m^2_{\tilde g}/m^2_{\tilde q}$ or an average slepton mass of 100 GeV and different values of $x=m^2_{\tilde\gamma}/m^2_{\tilde l}$. These bounds scale as $(m_{\tilde f}(\mathrm{GeV})/500(100))^2$ for different average sfermion masses

1 with $x=m^2_{\tilde g}/m^2_{\tilde q}\simeq 1$. Even allowing a phase $\mathcal{O}(1)$, necessary to contribute to ε'/ε, we can see the bound requires only $m\geq 1$ which is already required by fermion masses. Still, we can see that in the presence of a phase, ε'/ε naturally receives a sizeable contribution from the b' term [11].

The MI corresponding to the $b\to s\gamma$ decay are,

$$\left(\delta^d_{LR}\right)_{32}\simeq\frac{m_{3/2}\,m_b\,g\,n\,\epsilon_d^n}{19\,m^2_{3/2}}\simeq 2.2\times 10^{-3}\,g\,n\,\epsilon_d^n \tag{13}$$

again with $m_{3/2}\simeq 120$ GeV. This estimate is already of the same order of the phenomenological bound for any n and we do not get any new constraint on n.

The situation is more interesting in the leptonic sector. Here, the photino contribution is indeed the dominant one for LR mass insertions. In Table 1 , we show the rescaled bounds from Ref. [10] for the present limits on the branching ratio. In this case, it seems reasonable to expect some kind of lepton-quark Yukawa unification and we assume that the charged lepton Yukawa texture at M_{GUT} shares the same structure as the down Yukawa, except for the Georgi-Jarskog factors of 3 in (2,2), (2,3) and (3,2) entries. Therefore we obtain,

$$\begin{aligned}(\delta^e_{LR})_{12}\simeq\frac{m_\tau\,\epsilon_d^3}{5.5\,m_{3/2}}(b'+9\,a\,\frac{b'}{d}\,g\,n\,\epsilon_d^n-3\,a\,f\,m\,\epsilon_d^{m-1})\simeq\\ (b'+9\,a\,\frac{b'}{d}\,g\,n\,\epsilon_d^n-3\,a\,f\,m\,\epsilon_d^{m-1})\,8.7\times 10^{-6}\end{aligned} \tag{14}$$

where we take $m_{3/2}\simeq 120$ GeV corresponding to $m_{\tilde l}=280$ GeV. Notice that, at least, we have an unavoidable contribution from the b' entry, $(\delta^e_{LR})_{12}\simeq b'\,8.7\times 10^{-6}$. This must be compared with a bound, $(\delta^e_{LR})_{12}\leq 7\times 10^{-7}\;(280/100)^2=5.5\times 10^{-6}$. Clearly the estimate in Eq. (14) is still too large when compared with this experimental bound. Therefore this bound necessarily requires a larger value of the slepton mass. For an average slepton mass of 320 GeV ($m_{3/2}=136$ GeV, $m_{\tilde q}=600$ GeV) our estimate would be just below the MI bound. Again, there are no new constraints on the value of m and n. It is well worth to recall the very important results we get from this observable: Assuming a quark-lepton unification at M_{GUT}, nonuniversality in the trilinear terms predicts a large $\mu\to e\gamma$ branching ratio even beyond the values expected from other sources as SUSY seesaw [12].

This points out this decay as the most sensitive probe of SUSY and the soft breaking sector in the near future.

Another interesting constraint in this scenario is provided by Electric Dipole Moment (EDM) bounds. Even in the most conservative case, where all soft SUSY breaking parameters and μ are real, we know that the Yukawa matrices contain phases $\mathcal{O}(1)$. If the trilinear terms are nonuniversal, these phases are not completely removed from the diagonal elements of Y^A in the SCKM basis and hence can give rise to large EDMs [13]. However, it is possible to prove that the phase in the trilinear terms will be exactly zero at leading order in θ for any diagonal element. To see this we must take into account that the eigenvalues and mixing matrices of the Yukawa matrix depend on θ, $Y(\theta) = V_L^\dagger(\theta)\, D(\theta)\, V_R(\theta)$. The contribution to the trilinear terms proportional to $\theta\ \partial Y/\partial\theta$ is,

$$V_L \theta \frac{\partial Y}{\partial\theta} V_R^\dagger = V_L \frac{\theta \partial V_L^\dagger}{\partial\theta} D + \frac{\theta\partial D}{\partial\theta} + D \frac{\theta\partial V_R}{\partial\theta} V_R^\dagger \tag{15}$$

In this expression the dominant contribution in θ is controlled by the lowest power in D_{ii}. Clearly $\theta\ \partial V/\partial\theta$ always adds at least a power of θ and therefore the first and third terms in the above equation can only contribute to subdominant terms in the θ expansion for the diagonal elements. Hence the dominant term in a diagonal element is provided by the second term and is exactly proportional to the leading θ term in Y_{ii} with a coefficient equal to its power in θ. This implies that any observable phase in the diagonal elements will only appear at higher orders, for instance if $n \geq 1$ or $m \geq 2$ or with higher order contributions to entries of the Yukawa matrix. With real μ and soft breaking, EDMs are generated by,

$$\mathrm{Im}\left(\delta^{q,l}_{LR}\right)_{11} \simeq \frac{m_1}{R_{q,l}\, m_{3/2}} \left(\epsilon^n\, n + \epsilon^{m-1}\,(m-1)\right) \tag{16}$$

where we use $m_1 = m_3\, \epsilon^4 bb'/d$ and then take all unknown coefficients $\mathcal{O}(1)$. The coefficients $R_q = 19$ and $R_l = 5.5$ take care of the RGE effects in the eigenvalues as before. So, taking into account that $\epsilon_d = 0.15$, $\epsilon_u = 0.05$ and $m_d \simeq 10$ MeV, $m_u \simeq 5$ MeV and $m_e = 0.5$ MeV and assuming phases $\mathcal{O}(1)$, we get,

$$\begin{aligned}
\mathrm{Im}\left(\delta^d_{LR}\right)_{11} &\simeq \left(\epsilon_d^n\, n + \epsilon_d^{m-1}\,(m-1)\right)\ 3.9\times 10^{-6} \\
\mathrm{Im}\left(\delta^u_{LR}\right)_{11} &\simeq \left(\epsilon_u^n\, n + \epsilon_u^{m-1}\,(m-1)\right)\ 1.9\times 10^{-6} \\
\mathrm{Im}\left(\delta^e_{LR}\right)_{11} &\simeq \left(\epsilon_d^n\, n + \epsilon_d^{m-1}\,(m-1)\right)\ 6.7\times 10^{-7}
\end{aligned} \tag{17}$$

We compare these estimates with the phenomenological bounds in Table 2[14, 10]. The bounds from neutron and electron EDM do not provide any new information on the structure of the Yukawa textures. However, the mercury EDM bounds are much more restrictive and taking $x \simeq 1$ we find that, in the down sector, the case $n = 1$, $m = 2$ is not allowed by EDM experiments and we require $n \geq 2$, $m \geq 3$. As we said above, the same bound applies for subdominant corrections to Y_{22} and Y_{12}, Y_{21} where the first correction to the dominant terms can only be ϵ^4 or ϵ^5 respectively to satisfy EDM bounds.

In summary we have shown here that in a broken supergravity theory any field that acquires a vev obtains at the same time an F-term of order $\langle\theta\rangle m_{3/2}$. These F-terms contribute unsuppressed to the trilinear couplings and produce observable effects in the low

x	$\lvert\mathrm{Im}(\delta^d_{11})_{LR}\rvert$	$\lvert\mathrm{Im}(\delta^u_{11})_{LR}\rvert$	$\lvert\mathrm{Im}(\delta^d_{22})_{LR}\rvert$	$\lvert\mathrm{Im}(\delta^l_{11})_{LR}\rvert$
0.3	4.3×10^{-8}	4.3×10^{-8}	3.6×10^{-6}	4.2×10^{-7}
1	8.0×10^{-8}	8.0×10^{-8}	6.7×10^{-6}	5.1×10^{-7}
3	1.8×10^{-7}	1.8×10^{-7}	1.6×10^{-5}	8.3×10^{-7}

Table 2: MI bounds from the mercury EDM for an average squark mass of 600 GeV and for the electron EDM with an average slepton mass of 320 GeV and different values of $x = m^2_{\tilde{g}}/m^2_{\tilde{q}}$. The bounds scale as $(m_{\tilde{q}(\tilde{l})}(\mathrm{GeV})/600(320))$.

energy phenomenology. We have seen that $\mu \to e\gamma$ decays can receive large contributions from this source at the level of current experimental bounds. Therefore this decay may provide the first clue of the structure of the soft breaking sector with the proposed experiments in the near future.

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References

[1] A. Masiero and O. Vives, Ann. Rev. Nucl. Part. Sci. **51** (2001) 161; New J. Phys. **4**, 4 (2002).

[2] G. G. Ross and O. Vives, arXiv:hep-ph/0211279

[3] For an example of this in type I string models see S. A. Abel and G. Servant, Nucl. Phys. B **611**, 43 (2001).

[4] L. E. Ibanez and G. G. Ross, Phys. Lett. B **110**, 215 (1982).

[5] S. K. Soni and H. A. Weldon, Phys. Lett. B **126**, 215 (1983); V. S. Kaplunovsky and J. Louis, Phys. Lett. B **306**, 269 (1993); A. Brignole, L. E. Ibanez, and C. Munoz, Nucl. Phys. B **422**, 125 (1994) [Erratum-ibid. B **436**, 747 (1995)].

[6] S. F. King and G. G. Ross, Phys. Lett. B **520** (2001) 243; Z. Berezhiani and A. Rossi, Nucl. Phys. B **594**, 113 (2001); R. Barbieri and L. J. Hall, Nuovo Cim. A **110**, 1 (1997); G. G. Ross and L. Velasco-Sevilla, arXiv:hep-ph/0208218.

[7] R. G. Roberts, A. Romanino, G. G. Ross, and L. Velasco-Sevilla, Nucl. Phys. B **615** (2001) 358.

[8] N.K. Falck, Z Phys. **C30**, 247 (1986); S. Bertolini, F. Borzumati, A. Masiero, and G. Ridolfi, Nucl. Phys. **B353**, 591 (1991).

[9] A. Bartl *et. al*, Phys. Rev. D **64** (2001) 076009.

[10] F. Gabbiani, E. Gabrielli, A. Masiero, and L. Silvestrini, Nucl. Phys. **B477**, 321 (1996); J. Hagelin, S. Kelley, and T. Tanaka, Nucl. Phys. **B415**, 293 (1994).

[11] A. Masiero and H. Murayama, Phys. Rev. Lett. **83**, 907 (1999); S. Khalil, T. Kobayashi, and O. Vives, Nucl. Phys. **B580**, 275 (2000); A. Masiero and O. Vives, Phys. Rev. Lett. **86** (2001) 26.

[12] F. Borzumati and A. Masiero, Phys. Rev. Lett. **57**, 961 (1986); J. Hisano, T. Moroi, K. Tobe, and M. Yamaguchi, Phys. Rev. D **53** (1996) 2442; J. A. Casas and A. Ibarra, Nucl. Phys. B **618**, 171 (2001); S. Lavignac, I. Masina, and C. A. Savoy, Phys. Lett. B **520**, 269 (2001): A. Masiero, S. K. Vempati, and O. Vives, arXiv:hep-ph/0209303.

[13] S. Abel, S. Khalil, and O. Lebedev, Phys. Rev. Lett. **89**, 121601 (2002); Y. Nir and R. Rattazzi, Phys. Lett. B **382** (1996) 363.

[14] T. Falk, K. A. Olive, M. Pospelov, and R. Roiban, Nucl. Phys. B **560** (1999) 3; S. Abel, S. Khalil, and O. Lebedev, Nucl. Phys. B **606** (2001) 151.