

A model of Yukawa couplings with matter-messenger unification

J. Pawełczyk

*Institute of Theoretical Physics, Faculty of Physics, University of Warsaw,
Hoża 69, 00-681 Warsaw, Poland*

Abstract

We propose a GUT model in which visible matter and messengers are treated in unified way what unavoidably leads to messenger-matter mixed Yukawa interactions. Influence of this mixing on the fermion masses and the weak mixing angles is discussed.

GUT models are attractive scenarios for unification of all gauge forces [1]. The basic arguments supporting the idea are twofold: all known fermionic matter is organized in SU(5) multiplets and coupling constants seem to unify at some scale. It appears that supersymmetric versions of these models provide better coupling unification and shift the unification scale Λ_{GUT} to values acceptable for the proton stability under heavy gauge bosons exchange [2, 3]. Besides many nice features GUT models face serious issues among which there is the flavour problem i.e. the problem of generating appropriate hierarchy between masses of fermions. For example the simplest GUT models require $m_i^d = m_i^l$ at the GUT scale for each family ($i = 1, 2, 3$) of down quarks (d) and leptons (l). Calculations shows that one can hope to match masses of fermions of the 3rd family, $m_\tau = m_b$ [4, 5], but this can not be achieved for the other two families. The situation is readily apparent when one compares one-loop RG independent ratios of the light lepton and quark masses: $m_e/m_\mu \approx 0.004$, $m_d/m_s \approx 0.05$ what means that $m_e/m_d \sim 9 m_\mu/m_s$. It is clear that an extra source of flavour symmetry breaking is needed. The subject has very long history [6] and was studied very extensively in the past, see e.g. [7] and [8] and references contained therein.

The motivation behind our approach to the flavour problem lays in some F-theory constructions of GUT theories (the so-called F-GUTs) [9] where the both flavour visible chiral matter and messengers originate from the same D7-brane intersection [10] (see the Appendix). This means that matter and messengers¹ should be treated on equal footing e.g. should have common Yukawa matrix. Mixed couplings between fields is generated by fluxes on Calabi-Yau manifold [11, 12]. We shall claim that the mixing is responsible for the deviation from the naive relations between fermion masses.

We begin writing down all relevant superpotential terms

$$(y^u)_{ij} 10_i 10_j (5_H)_2 + (\tilde{y}^d)_{iJ} 10_i \bar{5}_J (\bar{5}_H)_2 + a_I \bar{5}_I 5 X \quad (1)$$

where the subscript 2 attached to the Higgs fields recalls that we take into account only the doublet part of the 5's. The above superpotential in principle could be extended with extra 10's but it is unlike in F-theory setup (see the Summary). Thus we let the flavour indices run as follows: $i, j = 1, 2, 3$; $I, J = 1, \dots, 4$ and we add an extra 5 which will form the messenger vector pair with one of the $\bar{5}$'s. The spurion v.e.v. $\langle X \rangle$ gives mass M_Y to the messengers (and could trigger SUSY breaking through its nontrivial

¹We shall abuse the name "messenger" here because the SUSY breaking will not discussed here.

F-term). It is clear that $\langle X \rangle$ provide mass M_Y only to one vector pair $(5, Y_{\bar{5}})$ where $Y_{\bar{5}} \sim a_J \bar{5}_J$ leaving the other $\bar{5}$ matter massless. Below the scale M_Y we are left with 3 massless families. Accordingly we need to redefine fields what results in modification of Yukawas \tilde{y}_d ².

F-theory GUTs [11]³ lead to Yukawa matrices of the Froggatt-Nielsen form [14] possessing naturally hierarchical structure. We write explicitly the general structure of these couplings in the case of three 10's and four $\bar{5}$ but the matrices can be easily made up for the other cases.

$$y_{NC} \sim \begin{pmatrix} h_\theta^5 & h_\theta^4 & h_\theta^3 & h_\theta^2 \\ h_\theta^4 & h_\theta^3 & h_\theta^2 & h_\theta \\ h_\theta^3 & h_\theta^2 & h_\theta & 1 \end{pmatrix} + \Phi_0 \begin{pmatrix} h_\theta^4 & h_\theta^3 & h_\theta^2 & h_\theta \\ h_\theta^3 & h_\theta^2 & h_\theta & 0 \\ h_\theta^2 & h_\theta & 0 & 0 \end{pmatrix} + \Phi_0^2 \begin{pmatrix} h_\theta^3 & h_\theta^2 & 0 & 0 \\ h_\theta^2 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \end{pmatrix} \quad (2)$$

where h_θ, Φ_0 are free parameters and as usual the order one constants have been dismissed. It is known that Φ_0 depends on the hypercharges of the corresponding matter. For relatively big hypercharges one expects $\Phi_0 = 0$. Thus this case should be applicable to leptons and to some extent to up-quarks⁴. We want to stress that \tilde{y}^d dependence on the hypercharge of matter breaks SU(5) symmetry so in fact Yukawa coupling are different for down quarks and leptons.

The model with just 3 families of matter and no matter-messenger mixing (thus disregarding advocated here messenger i.e. the last column of (2)) was discussed in [11]. For $h_\theta \sim \epsilon^2$, $\epsilon \sim 0.2$ this model is fine for the up quarks ($m_u : m_c : m_t \sim \epsilon^8 : \epsilon^4 : 1$) but fails for seemingly more favorable case of leptons due to $m_e/m_\mu \approx 0.005 \sim \epsilon^{3.5}$, $m_\mu/m_\tau \approx 0.05 \sim \epsilon^2$. Recall that in MSSM ratios of Yukawas for fermions differing only by the family index depend on scale mildly thus RG flow can not explain this deviation. For down quarks one puts instead $\Phi_0 \sim \epsilon \sim 0.2$ [11]. This yields some misfit for masses $m_1 : m_2 : m_3 = \epsilon^6 : \epsilon^3 : 1$ and correctly reproduces mixing angles $\theta_{23} \sim \epsilon^2$, $\theta_{13} \sim \epsilon^3$ although $\theta_{12} \sim \epsilon^2$ is too small.

Now we discuss how the situation changes if we go the model (1) with an extra $\bar{5}$. Its important to notice that the last term of (1) change the definition of messenger by mixing $\bar{5}$'s. One expects to get some kind of hierarchy between a_I s e.g. of the (2) type (take just the last row) i.e. one expects that $a_{I+1} \ll a_I$. Here we normalize

²For other GUT models utilizing mixing of matter with some extra fields see [15] and references therein.

³For different approaches to fermion hierarchies in F-theory GUTs see [12, 13].

⁴Recall values of the hypercharge for R component of fermions: $Y(u_R) = 2/3$, $Y(e_R) = -1$, $Y(d_R) = -1/3$.

$a_4 = 1$. The necessity to redefine messenger leads to redefinition of the other fields in multiplets of $\bar{5}$. This rotates Yukawa coupling matrix. On the other hand the leading terms of the couplings between the physical messenger $Y_{\bar{5}}$ and light matter has universal form.

$$(y_{14}\epsilon^4 10_1 + y_{24}\epsilon^2 10_2 + y_{34} 10_3) \cdot Y_{\bar{5}} \cdot (\bar{5}_H)_2 \quad (3)$$

After decoupling of the messenger $Y_{\bar{5}}$ we are left with 3×3 matrix of Yukawas for visible fermions. Below we analyze this matrix for the cases appropriate for up and down quarks and for leptons.

Leptons. Putting all the unknown constants in (2) in the leading order in ϵ we get for leptons ($\Phi_0 = 0$ case):

$$\tilde{y}^l = \begin{pmatrix} y_{11} \epsilon^{10} & y_{12} \epsilon^8 & y_{13} \epsilon^6 & y_{14} \epsilon^4 \\ y_{21} \epsilon^8 & y_{22} \epsilon^6 & y_{23} \epsilon^4 & y_{24} \epsilon^2 \\ y_{31} \epsilon^6 & y_{32} \epsilon^4 & y_{33} \epsilon^2 & y_{34} \end{pmatrix} \quad (4)$$

Properties of the resulting 3×3 matrix (called here y^l) for massless fermions depends on a_I coefficients. For the sake of illustration of possibilities at hand we display formulae for $a_1 \gg \epsilon^6$ and $a_2 \gg a_3 \epsilon^2$. The resulting Yukawa couplings for leptons is

$$y^l = \begin{pmatrix} y_{14} a_1^* \epsilon^4 & y_{14} a_2^* \epsilon^4 + y_{12} \epsilon^8 & y_{14} a_3^* \epsilon^4 + y_{13} \epsilon^6 \\ y_{24} a_1^* \epsilon^2 & y_{24} a_2^* \epsilon^2 + y_{22} \epsilon^6 & y_{24} a_3^* \epsilon^2 + y_{23} \epsilon^4 \\ a_1^* y_{34} & a_2^* y_{34} + y_{32} \epsilon^4 & a_3^* y_{34} + y_{33} \epsilon^2 \end{pmatrix} \quad (5)$$

Diagonalization of (5) yields the following eigenvalues (given here up to phases):

$$y_3 \approx a_3^* y_{34} + y_{33} \epsilon^2 \quad (6)$$

$$y_2 \approx \frac{(y_{24} y_{33} - y_{23} y_{34}) a_2 \epsilon^4}{a_3^* y_{34} + y_{33} \epsilon^2} \quad (7)$$

$$y_1 \approx \frac{\det(y_d)}{(y_{24} y_{33} - y_{23} y_{34}) a_2 \epsilon^4} \quad (8)$$

where $\det y^l \sim a_1^* \epsilon^{12}$. Assuming generic values of $y_{ij} \sim 1$ one gets

$$\frac{y_1}{y_2} \sim \frac{a_1 y_3 \epsilon^4}{a_2^2}, \quad \frac{y_2}{y_3} \sim \frac{a_2 \epsilon^4}{y_3^2} \quad (9)$$

Needed lepton hierarchy $y_1/y_2 \sim \epsilon^{3.5}$, $y_2/y_3 \sim \epsilon^2$ yields $a_2 \epsilon^2 \sim y_3^2$ and $a_1 y_3 \sim a_2^2 \epsilon^{0.5}$. For $y_3 \sim \epsilon^k$ this gives $a_2 \sim \epsilon^{2k-2}$ and $a_1 \sim \epsilon^{3k-4.5}$. Thus hierarchy in a_I 's is respected if $k \geq 2$, e.g. $y_3 \sim \epsilon^{2.5}$ produces $a_2 \sim \epsilon^3$, $a_1 \sim \epsilon^3$. This choice of y_3 requires some fine tuning i.e. $y_{33} \sim \epsilon^{0.5} \sim 0.5$ and $a_3 \sim \epsilon^{2.5}$.

Down quarks. Now we turn to the second example $h_\theta \sim \epsilon^2$ and $\Phi_0 \sim \epsilon$ which we use to fit to down quark Yukawas. Repeating the previous construction instead of (5) one obtains⁵:

$$y^d = \begin{pmatrix} y_{11}\epsilon^8 + y_{14}a_1^*\epsilon^3 & y_{12}\epsilon^6 + y_{14}a_2^*\epsilon^3 & y_{13}\epsilon^5 + y_{14}a_3^*\epsilon^3 \\ y_{21}\epsilon^6 + y_{24}a_1^*\epsilon^2 & y_{22}\epsilon^5 + y_{24}a_2^*\epsilon^2 & y_{23}\epsilon^3 + y_{24}a_3^*\epsilon^2 \\ y_{31}\epsilon^5 + y_{34}a_1^* & y_{32}\epsilon^3 + y_{34}a_2^* & y_{33}\epsilon^2 + y_{34}a_3^* \end{pmatrix} \quad (10)$$

The obtained diagonalized Yukawas are (for $a_1 \gg \epsilon^5$, $a_1 \geq \epsilon^2 a_2$, $a_1 \geq \epsilon^3 a_3$)

$$y_3 \approx a_3^* y_{34} + y_{33} \epsilon^2 \quad (11)$$

$$y_2 \approx \frac{a_2 y_{23} y_{34} \epsilon^3 + y_{32} y_{33} \epsilon^6}{y_3} \quad (12)$$

$$y_1 \approx \frac{\det(y_d)}{a_2 y_{23} y_{34} \epsilon^3 + y_{32} y_{33} \epsilon^6} \quad (13)$$

where $\det(y_d) \approx (y_{12} y_{23} y_{34} - y_{14} y_{23} y_{32}) a_1 \epsilon^9$. Phenomenological values of the ratios of the down quark masses $m_d : m_s : m_b = \epsilon^{2+s} : \epsilon^s : 1$ for $s = 2 - 3$. With $y_3 \sim \epsilon^k$, $a_2 \sim \epsilon^{2k+s-3}$, $a_1 \sim \epsilon^{3k+2s-7}$ we get correct relations for down quarks Yukawas. The coefficients a_I have hierarchical structure if $k+s > 4$. One can also calculate the mixing angles: $\theta_{12} \sim \epsilon^{6-(k+s)} + \epsilon^k$, $\theta_{23} \sim \epsilon^{3-k} + \epsilon^2$, $\theta_{13} \sim \epsilon^{5-k} + \epsilon^3$ (coefficients of the order one are suppressed). e.g. for $k = 1.5$, $s = 3$ one has : $\theta_{12} \sim \epsilon^{1.5}$, $\theta_{23} \sim \epsilon^{1.5}$, $\theta_{13} \sim \epsilon^3$. These values match phenomenology up to terms of the order $O(\epsilon^{0.5})$.

RG flow of masses. Finally we emphasize that the presented results hold at the GUT scale and are subject to renormalization what in case of matter-messenger coupling may have important influence on ordinary Yukawas. In order to analyze the problem we did numerical calculation based on one-loop RG equations for the wave functions renormalizations presented in [16] and corrected for the case at hand in [17]. It RG results depends mainly on the messenger mass M and the largest mixing coupling y_{34} . The RG flow modification of mixing angles is negligible. The results are presented in Table 1. It is clear that RG does not change the fermion masses ratios in a significant way.

⁵In general the coefficients y_{ij} in matrices y^l and y^d can be different.

	m_u	m_c	m_t	m_d	m_s	m_b	m_e	m_μ	m_τ
	[MeV]	[MeV]	[GeV]	[MeV]	[MeV]	[GeV]	[MeV]	[MeV]	[GeV]
$MSSM, \Lambda = 10^3$	1.15	560	161	2.20	42.0	2.23	0.41	88.3	1.50
$MSSM$	0.55	270	106	0.79	15.1	0.89	0.29	62.0	1.06
$y_{34} = 0.6, M = 10^{13}$	0.55	268	107	0.84	16.0	0.96	0.31	66.2	1.16
$y_{34} = 1.5, M = 10^7$	0.51	250	112	1.39	26.6	1.82	0.52	112	2.48

Table 1: One-loop fermion masses at Λ_{GUT} for different values of the messenger-matter mixing y_{34} and the messenger mass M . The initial values of the masses at 1 TeV are given in the first row.

Summary. Yukawa couplings for quarks and leptons can be modified by mixing with extra vector matter called messengers here. This mechanism opens up an extra path toward obtaining the desired structure of the Yukawa couplings of the visible matter. In the present paper we have worked out an example based on a model inspired by F-theory for which the possible Yukawa couplings have form of the Froggatt-Nielsen type. The considered vector pair of matter was $5 + \bar{5}$ of SU(5). Extending the model by an extra 10's is excluded if one wants to stay within perturbative physics. The reason is that due to hierarchical structure of the proposed Yukawa couplings and the couplings a_I 's to spurion (1) the extra 10's would have couplings hierarchically much bigger than that of the top quark i.e. much bigger than 1. Hence the messenger sector would be strongly interacting. This might be appealing in the F-theory context but usually is not welcomed feature.

From (11) one has $y_b \sim a_3^* y_{34} + y_{33} \ll y_{34} \sim 1$ what implies that $\tan \beta$ can not be large (typically smaller than 10) in this type of models. Adding more vector pairs of 5's would lower $\tan \beta$ much further what is excluded in MSSM-like models [18]. Thus we infer that our scenario prefers the single messenger in the representation 5 of SU(5) GUT. It is also interesting that [10] do not contain any model with Majorana neutrinos masses which would fit in our scheme but numerous with Dirac masses. This is shortly discussed in the appendix.

We want to emphasize here that the results obtained shows only the order of quantities of interest in terms of powers of ϵ . But because $\epsilon \sim 0.2$ and unknown coefficients $y_{ij} \sim 1$ this expansion may be misleading e.g. if one takes some $y_{ij} \sim 0.5$ than terms quadratic in these y 's are of the order ϵ what spoils the expansion we have made.

The presented model has various phenomenological consequences. The subject is very reach and requires separate studies so we shall be very brief here. Strong ($y_{34}^d \sim 1$) mixing of messengers to top quark produces large A_t term [19].

$$V = -\frac{F_X}{16\pi^2\langle X \rangle} (y_{34}^d)^2 y_t \tilde{Q} H_u \tilde{t} + c.c. \quad (14)$$

A_t of the order of several TeV gives positive contribution to the Higgs mass [20, 21] helping to achieve required value 126 GeV. Moreover it decreases the lighter stop mass [22]. Both contributions are very welcomed in view of LHC results: stop can still be relatively light (few hundred GeV) thus its discovery would favor large A_t term.

The other A-terms are much smaller than A_t but nevertheless are severely constrained. The reason is that flavour-messenger mixing yields non-diagonal soft masses [16, 17] which in turn generate FCNC. Phenomenological constrains on FCNC [23] put limits on mixing coefficients y_{i4} . The thorough analysis of this problem is left to our future work [24].

A About F-theory constructions

Here we discuss SU(5) F-theory local models of [10] which lead to the scenario proposed in the current work. In F-theory various matter is localized on different Riemann surfaces along which D7-banes intersect. Fluxes on these Riemann surfaces determines number of chiral species of given matter. For all models we need the flux to prevent existence of $Y_{\overline{10}}$ so all 10's are flavours. We also need appropriate Y_5 and the spurion X . In what follows we shall assume that the appropriate fluxes can be switched on.

The group theory properties of the matter follows from $SU(5)_\perp/\Gamma$ (which is the global symmetry group here) for various discrete groups Γ . In [10] scenarios for different Γ 's has been considered. Here we shall shot-cut the discussion to the relevant pieces of the models. We start with Dirac scenario models. For $\Gamma = \mathbb{Z}_2$ the global symmetry group is $U(1)_{PQ} \times U(1)_\chi$. All possible versions these models realizing scenarios proposed in the present work are enumerated in Table 2.

For $\Gamma = \mathbb{Z}_2 \times \mathbb{Z}_2$ only the last model of Table 2 survives. For $\Gamma = \mathbb{Z}_3$ there is one possibility for which all the relevant fields are the same as originally listed in [10]. On the other hand non of Majorana models of [10] can realize the scenario of the present paper.

		version 1		version 2		version 3	
	$\bar{5}$	$Y_5 = 5_{(1)}$	$X = N_R$	$Y_5 = 5_{(2)}$	$X = D_{(2)}$	$Y_5 = 5_{(3)}$	$X = D_{(4)}$
$U(1)_{PQ}$	+1	+2	-3	-5	+4	+6	-7
$U(1)_X$	+3	+2	-5	-3	0	+2	-5

Table 2: $U(1)_{PQ}$ and $U(1)_X$ charges for the candidates for messengers Y_5 and spurions X of Eq.(1) for different variants of $SU(5)_\perp/\mathbb{Z}_2$ model. The $\bar{5}$ has universal charges.

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References

- [1] P. Langacker, “Grand Unified Theories And Proton Decay,” Phys. Rept. **72** (1981) 185.
- [2] C. Kounnas, A. B. Lahanas, D. V. Nanopoulos and M. Quiros, Nucl. Phys. B **236** (1984) 438.
- [3] D. V. Nanopoulos, Phys. Rept. **105** (1984) 71.
- [4] M. S. Chanowitz, J. R. Ellis and M. K. Gaillard, Nucl. Phys. B 128, 506 (1977); A. J. Buras, J. R. Ellis, M. K. Gaillard and D. V. Nanopoulos, Nucl. Phys. B135, 66 (1978).
- [5] G. Ross and M. Serna, Phys. Lett. B 664 (2008) 97 [arXiv:0704.1248 [hep-ph]]; K. S. Babu, arXiv:0910.2948 [hep-ph]; G. Elor, L. J. Hall, D. Pinner and J. T. Ruderman, JHEP 1210 (2012) 111 [arXiv:1206.5301 [hep-ph]]; K. Bora, arXiv:1206.5909 [hep-ph].
- [6] H. Georgi and C. Jarlskog, Phys. Lett. B **86** (1979) 297.
H. Murayama, Y. Okada and T. Yanagida, Prog. Theor. Phys. **88** (1992) 791.

- [7] E. Witten, Phys. Lett. B 91 (1980) 81; Z. G. Berezhiani, Phys. Lett. B 150 (1985) 177; A. Davidson and K. C. Wali, Lett. 58 (1987) 2623; J. Hisano, H. Murayama and T. Yanagida, Phys. Rev. D 49 (1994) 4966; K. S. Babu and S. M. Barr, Phys. Rev. Lett. 75 (1995) 2088 [hep-ph/9503215]; Z. G. Berezhiani, Phys Lett. B 355 (1995) 178 [arXiv:hep-ph/9505384]; M. Malinsky, Phys. Rev. D 77 (2008) 055016 [arXiv:0710.0581 [hep-ph]]; Q. Shafi and Z. Tavartkiladze, Phys. Lett. B 459 (1999) 563 [hep-ph/9904249]; N. Oshimo, Phys. Rev. D 80 (2009) 075011 [arXiv:0907.3400 [hep-ph]].
- [8] L. Calibbi, Z. Lalak, S. Pokorski and R. Ziegler, JHEP **1206** (2012) 018 [arXiv:1203.1489 [hep-ph]].
- [9] R. Donagi, M. Wijnholt, [arXiv:0802.2969 [hep-th]]; C. Beasley, J.J. Heckman, C. Vafa, JHEP 0901:058 (2009), [arXiv:0802.3391 [hep-th]]; H. Hayashi et al., Nucl. Phys. B 806:224 (2009), [arXiv:0805.1057 [hep-th]]; C. Beasley, J.J. Heckman, C. Vafa, JHEP 0901:059 (2009), [arXiv:0806.0102 [hep-th]]; R. Donagi, M. Wijnholt, [arXiv:0808.2223 [hep-th]].
- [10] J.J. Heckman, A. Tavanfar, C. Vafa, [arXiv:0906.0581 [hep-th]].
- [11] J. J. Heckman and C. Vafa, Nucl. Phys. B 837 (2010) 137 [arXiv:0811.2417 [hep-th]]; S. Cecotti, M. C. N. Cheng, J. J. Heckman and C. Vafa, arXiv:0910.0477 [hep-th].
- [12] F. Marchesano and L. Martucci, Phys. Rev. Lett. 104, 231601 (2010) [arXiv:0910.5496 [hep-th]]; L. Aparicio, A. Font, L. E. Ibanez and F. Marchesano, JHEP **1108** (2011) 152 [arXiv:1104.2609 [hep-th]]; A. Font, L. E. Ibanez, F. Marchesano and D. Regalado, arXiv:1211.6529 [hep-th].
- [13] E. Dudas and E. Palti, JHEP 1001, 127 (2010) [arXiv:0912.0853 [hep-th]]; G. K. Leontaris and G. G. Ross, JHEP 1102, 108 (2011) [arXiv:1009.6000 [hep-th]].
- [14] C. D. Froggatt and H. B. Nielsen, Nucl. Phys. B 147, 277 (1979).
- [15] K. S. Babu, B. Bajc and Z. Tavartkiladze, Phys. Rev. D **86** (2012) 075005 [arXiv:1207.6388 [hep-ph]].
- [16] Z. Chacko and E. Ponton, Phys. Rev. D 66 (2002) 095004, [hep-ph/0112190].
- [17] J.A. Evans and D. Shih, JHEP **1308** (2013) 093, arXiv:1303.0228.

- [18] P. Draper, P. Meade, M. Reece and D. Shih, Phys. Rev. D **85** (2012) 095007 [arXiv:1112.3068 [hep-ph]].
- [19] G. Giudice and R. Rattazzi, Phys.Rept. 322 (1999) 419-499, [hep-ph/9801271].
- [20] S.P. Martin, in: G.L. Kane (Ed.), Perspectives on Supersymmetry, in: Advanced Series on Directions in High Energy Physics, vol. 18, World Scientific, 1998, pp. 1-98, hep-ph/9709356.
- [21] T. Jeliński, J. Pawełczyk and K. Turzyński, Phys. Lett. B **711** (2012) 307 [arXiv:1111.6492 [hep-ph]].
- [22] M.Drees, R.Goodbole and P.Roy, Theory and Phenomenology of Superparticles, World Scientific, 2004.
- [23] A. Masiero, S.K. Vempati, O. Vives, arXiv:0711.2903
- [24] T.Jeliński, J. Pawełczyk and K.Turzyński, work in progress.