EPJ Web of Conferences **125**, 02017 (2016) DOI: 10.1051/epjconf/201612502017
QUARKS-2016

Sgoldstino physics and flavor-violating Higgs boson decays

Sergei Demidov[1,⋆] and *Ivan* Sobolev[1,2,⋆⋆]

[1] *Institute for Nuclear Research of the Russian Academy of Sciences, 60th October Anniversary prospect 7a, Moscow 117312, Russia*
[2] *Department of Particle Physics and Cosmology, Physics Faculty, M. V. Lomonosov Moscow State University, Vorobjevy Gory, 119991, Moscow, Russia*

Abstract. In this talk we study effects of Higgs-sgoldstino mixing on flavour-violation in Higgs sector and in particular on $h \to \mu\tau$ decay.

1 Introduction

One of the most interesting anomalies in the LHC run-1 was a hint of Higgs flavour-violating (HFV) decay. CMS collaboration observed small excess of signal events which can be interpreted as $h \to \mu\tau$ decay with $\mathrm{Br(h} \to \mu\tau) = \left(8.4^{+3.9}_{-3.7}\right) \times 10^{-3}$ [1]. In the Standard Model (SM) in the absence of neutrino masses this process is forbidden since in this case muon and tau lepton numbers are conserved individually. When extending the SM to include neutrino masses, such a processes can occur via 1-loop diagrams with virtual neutrinos but their rate is suppressed by powers of $\frac{m_\nu^2}{M_W^2}$, so their effect is far too small to affect Higgs decay in practice. A lot of ideas was proposed in an attempt to explain the mentioned excess. [3–14]. In this contribution we briefly present results of our paper [15] and make some new comments and remarks on the results obtained.

2 Model description

2.1 Lepton nubmer violation in the Minimal Supersymmetric Standard Model (MSSM)

The SM is minimally flavour violating: the only sources of flavor violation are the Yukawa matrices. In this way, if assuming that the neutrino mass is zero, the SM posseses accidental $U(1)_e \times U(1)_\mu \times U(1)_\tau$ symmetry. The same holds for purely supersymmetric part of the MSSM Lagrangian (i.e. for Kähler potential and superpotential). However this is not the case for SUSY breaking soft terms, and in particular for slepton mass terms which in general can violate this accidental lepton number symmetry. The quadratic part of slepton lagrangian can be written as

$$\mathcal{L}_{slepton} \supset -\tilde{l}^{\dagger}_{Li} M_{\tilde{l}LLij} \tilde{l}_{Lj} - \tilde{l}^{\dagger}_{Ri} M_{\tilde{l}RRij} \tilde{l}_{Rj} - \tilde{l}^{\dagger}_{Li} M_{\tilde{l}LRij} \tilde{l}_{Rj} - \tilde{l}^{\dagger}_{Ri} M_{\tilde{l}RLij} \tilde{l}_{Lj}, \quad (1)$$

⋆e-mail: demidov@ms2.inr.ac.ru
⋆⋆e-mail: sobolev.ivan@physics.msu.ru

where

$$\tilde{l} = (\tilde{e}, \tilde{\mu}, \tilde{\tau}) \tag{2}$$

and $M^2_{\tilde{l}LL}$,$M^2_{\tilde{l}RR}$,$M^2_{\tilde{l}RL}$,$M^2_{\tilde{l}LR}$ are 3×3 matrices. Obviously any non-diagonal term in either of these matrices leads to flavour-violation in charged lepton sector. Several comments here are in order. Fist, HFV decays exclusively in the scope of the MSSM have been studied in [3, 16]. It was found that parameter space capable of explaining CMS excess is very narrow. In other words, one needs significant fine-tuning of the MSSM parameters to reach Br(h $\to \tau\mu$) of order of $\sim 1\%$. Second, we assume that non-diagonal terms in $M^2_{\tilde{l}LL}\left(M^2_{\tilde{l}RR}\right)$ are suppressed and thus can be neglected. We will elaborate on this assumption in the next section.

2.2 Models with low SUSY breaking scale

If supersymmetry is actually realized in Nature it should be spontaneously broken at low energies. In this fashion phenomenologically viable supersymmetric theory should possess some sector [17] (which is usually called "hidden" in contradistinction to "visible" or MSSM sector) where symmetry breaking occurs and then is transmitted to "visible" sector by means of some mediation mechanism. According to the analogue of Goldstone theorem, a massless fermion – goldstino – should appear in the spectrum of a model. We utilize the simplest case of a "hidden" sector described by the following lagrangian

$$\mathcal{L}_\Phi = \int d^2\theta\, d^2\bar{\theta}\left(\Phi^\dagger\Phi + \tilde{K}(\Phi^\dagger,\Phi)\right) - \left(\int d^2\theta F\Phi + \text{h.c.}\right), \tag{3}$$

where $\Phi = \phi + \sqrt{2}\theta\psi + \theta^2 F_\phi$ is a chiral superfield and $\tilde{K}$ can be chosen in the form

$$\tilde{K}(\Phi^\dagger,\Phi) = -\frac{m_s^2 + m_p^2}{8F^2}(\Phi^\dagger\Phi)^2 - \frac{m_s^2 - m_p^2}{12F}((\Phi^\dagger\Phi^\dagger\Phi) + (\Phi\Phi^\dagger\Phi)) \tag{4}$$

Chiral superfield is a singlet under the action of the SM gauge group and it contains massless godstino ψ and its superpartners – scalar (s) and pseudoscalar (p) sgoldstinos which form complex field ϕ.The last linear term in the superpotential of Eq. (3) forces the auxiliary field F_ϕ to acquire non-zero vacuum expectation value $\langle F_\phi\rangle = F + O\left(\frac{1}{F}\right)$ and hence triggers spontaneous supersymmetry breaking. The value $\sqrt{F}$ can be regarded as supersymmetry breaking scale. When $\sqrt{F}$ is not far away from the electroweak scale (say it is of order of $\sim O$(TeV)) "hidden" sector becomes not so hidden and sgoldstino interactions with MSSM fields should be included into the low-energy theory. Here we briefly recall a minimal set of interactions of goldstino multiplet with MSSM fields. It can be obtained by making use of spurion method [18]. We will restrict ourselves only to the simplest set of operators which reproduce soft-terms when auxiliary field F_ϕ acquires its vacuum expectation value. Furthermore we assume R-parity conservation and presume that all of the free parameters of the model are real and hence neglect any possible CP-violation.

$$\mathcal{L}_{\text{Kähler}} = -\int d^2\theta\, d^2\bar{\theta} \sum_k \frac{m_k^2}{F^2}\Phi^\dagger\Phi\Phi_k^\dagger e^{g_1V_1+g_2V_2+g_3V_3}\Phi_k, \tag{5}$$

where k runs over all chiral MSSM superfields.

$$\mathcal{L}_{\text{superpotential}} = \int d^2\theta \left\{\epsilon_{ij}\left(-\frac{B}{F}\Phi H_d^i H_u^j + \frac{A^L_{ab}}{F}\Phi L_a^j E_b^c H_d^i\right) + \sum_\alpha \frac{M_\alpha}{2F}\Phi \text{Tr}\, W^\alpha W^\alpha\right\} + \text{h.c.} \tag{6}$$

Several important remarks here are in order. First, lagrangians (4)–(6) contain operators with dimension higher than four. Thereby this theory should be regarded as *effective theory* valid under some cutoff scale. Possible values of the cutoff Λ in the low scale supersymmetry breaking models have been discussed some time ago in Refs. [18, 20, 21]. It has been found that the cutoff for this model can lie somewhere between the level of soft masses of matter scalars $\tilde{m}$ (the largest of which can not exceed $\sqrt{F}$ (see table 1)) and the value $\Lambda^2 = 16\pi F^2/\tilde{m}^2$ (which is the scale of perturbative violation of unitarity of the effective theory). We choose the latter for our numerical estimates. Second, non-renormalizable interactions cause divergences in loop diagrams which involve particles from the hidden sector. These divergences turn out to be of two types: logarithmic and quadratic. The latter are caused by box diagrams in $\tau \to 3\mu$ process and are proportional to non-diagonal entries in $M^2_{\tilde{l}LL}$ $(M^2_{\tilde{l}RR})$ blocks [20]. This is the reason why we assume that this entries are suppressed (in our analysis we set them to zero). Logarithmic divergences take place in $\tau \to \mu\gamma$ [1] due to sgoldstino-Higgs mixing. That is why they are proportional to $\sim \sin^2 \frac{\theta}{2} \log \frac{\Lambda}{m_S}$, where θ is the mixing angle, and in general are suppressed (for more detailed discussion see [15]).

2.3 Higgs-sgoldstino mixing

From (5)-(6) one can derive the following Higgs-sgoldstino mixing term

$$\mathcal{L}_{s-h} = -\frac{X}{F} s \cdot h, \tag{7}$$

with [2]

$$X = 2\mu^3 v \sin 2\beta + \frac{v^3}{2}\left(g_1^2 M_1 + g_2^2 M_2\right)\cos^2 2\beta \tag{8}$$

which was obtained in so-called "decoupling limit" (i.e. when all the MSSM Higgs bosons are considerably heavier than the lightest one) and after the EWSB. The mixing angle can be written as

$$\tan 2\theta = \frac{2X}{F(m_s^2 - m_h^2)}. \tag{9}$$

We stress that m_s and m_h in this expression are *non-physical* masses of scalar sgoldstino and Higgs respectively. This remark is important since corresponding *physical* masses can differ drastically from the mentioned ones in case of large of mixing angles and this will be the case for most of the interesting points in the parameter space of our model. We will denote mass states as $\tilde{s}$ and $\tilde{h}$ keeping notations s and h for non-physical states. Next, the second term in (7) is suppressed compared to the first one. In this way the sign of the mixing angle is dictated by the sign of μ parameter. Unlike Higgs boson sgoldstino has flavour-violating couplings to the SM fermions

$$\mathcal{L}_{s-fermions} \supset \frac{A^L_{ab}}{\sqrt{2}F}\bar{e}_b\, l_a\, v_d\, s + \text{h.c.}. \tag{10}$$

So, small admixture of sgoldstino to the lightest Higgs boson generates flavor violating couplings of the latter

$$Y^{\tilde{h}}_{\mu\tau(\tau\mu)} = \frac{v\, A_{\mu\tau(\tau\mu)} \cos\beta \sin\theta}{\sqrt{2}F}. \tag{11}$$

[1] Logarithmic divergences are also present in triangle diagrams contributing to $\tau \to 3\mu$ decay but here they are suppressed compared to the leading tree-level diagrams, see [15].

[2] Here parameters β,$M_{1,2}$,μ are conventional MSSM parameters (see [17] for details), $v = 174$ GeV, $g_{1,2}$ are $U(1)$ and $SU(2)$ coupling constants respectively.

The corresponding decay width reads

$$\Gamma(\tilde{h} \to \mu\tau) \equiv \Gamma(\tilde{h} \to \bar{\tau}\mu) + \Gamma(\tilde{h} \to \bar{\mu}\tau) = \frac{m_{\tilde{h}}}{8\pi}\left(\left|Y^{\tilde{h}}_{\mu\tau}\right|^2 + \left|Y^{\tilde{h}}_{\tau\mu}\right|^2\right). \quad (12)$$

3 Model analysis and results

We perform scan over the following parameter space (see Table 1). Recall that the consistency of the effective field theory approach to the model (3)–(6) requires that the parameters which become soft terms after the spontaneous supersymmetry breaking should be smaller than $\sqrt{F}$. In what follows we fix this value to 8 TeV. Values of $\sqrt{F}$ smaller than about 8 TeV turn out to be disfavored the

$\tan\beta$	$1.5\ldots50.5$
$\lvert\mu\rvert$	$100\ldots2000$ GeV
M_1	$100\ldots2000$ GeV
M_2	$200\ldots2000$ GeV
M_3	$1.5\ldots4.0$ TeV
$A_{\mu\tau}, A_{\tau\mu}$	$0.1\ \sqrt{F}\ldots\ \sqrt{F}$

Table 1. Parameter space used in the analysis

results of direct searches. Namely, at smaller $\sqrt{F}$ the coupling constants of sgoldstino to the SM particles increase and such sgoldstino is phenomenologically unacceptable. In this case, very light sgoldstino, which decays mostly to $b\bar{b}$ due to large mixing with the Higgs boson, becomes excluded by the LEP [22] and TeVatron [23] results. Heavier sgoldstino with $\sqrt{F}$ smaller than about 8 TeV is excluded by the results of the ATLAS and CMS searches for diboson resonances. If we enlarge our parameter space by increasing, in particular, the upper bound on μ in the Table 1, we expect that somewhat lower values of $\sqrt{F}$ and larger values of the sgoldstino mass will be allowed. We find that for the most interesting cases the mass parameter of the scalar sgoldstino m_s should not be very heavy or very small. In the case of heavy sgoldstino, the mixing angle (9) is small and, as a consequence, the width of $\tilde{h} \to \mu\tau$ decay is suppressed. On the other hand, very light sgoldstinos with large

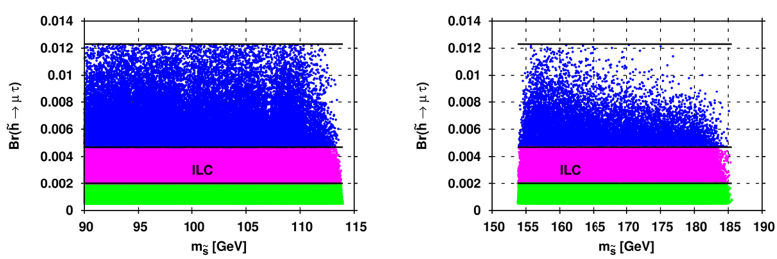

Figure 1. Scatter plots in plane (m_s, Br($h \to \tau\mu$)) for $\sqrt{F}$ = 8 TeV and sgoldstinos lighter (left) and heavier (right) than Higgs.

Higgs boson admixture are phenomenologically unacceptable due to results from the LEP [22] and

Tevatron [23] experiments. So, we limit ourselves to the regimes in which the scalar sgoldstino mass parameter is somewhat smaller (90–114 GeV) or larger (150–200 GeV) than the Higgs boson mass. Some parameters, which are not of primary importance for the analysis, were fixed to reasonable benchmark values. For example, soft trilinear constant of b-quark A_{bb} is fixed to $A_{bb} = 0.5\sqrt{F}$ and the mass of the pseudoscalar sgoldstino is fixed to $m_p = 200$ GeV. For each point in the mentioned parameter space we calculate signal strength

$$\mu_f = \frac{\sigma(pp \to \tilde{h}) \times \mathrm{Br}(\tilde{h} \to f)}{\sigma(pp \to h^{SM}) \times \mathrm{Br}(h^{SM} \to f)}, \tag{13}$$

where final state f stands for $b\bar{b}, W^+W^-, ZZ, \gamma\gamma, \tau\bar{\tau}$. We compare it with experimental constraints by CMS: [24–30] and ATLAS: [31–35]. Next, we check whether the scalar sgoldstino-like resonance is

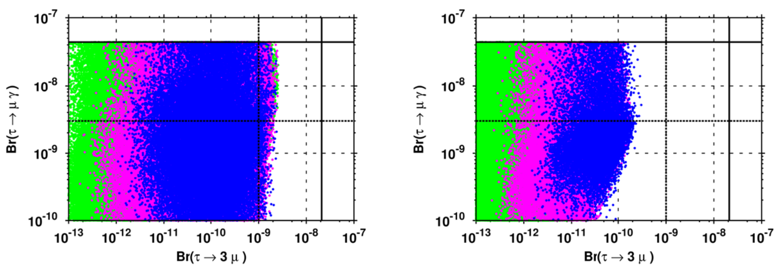

Figure 2. Scatter plots in plane $(\mathrm{Br}(\tau \to \mu\gamma), \mathrm{Br}(\tau \to 3\mu))$ for $\sqrt{F} = 8$ TeV and sgoldstinos lighter (left) and heavier (right) than Higgs. Solid lines correspond to present limits on branching fraction of both decays, whereas dashed line represents expected SuperKEKB sensitivity [46].

allowed by existing experimental results. For the case of the LHC searches for diboson resonances, we use the observables $\sigma(pp \to \tilde{s}) \times \mathrm{Br}(\tilde{s} \to \mathrm{f})$ where f stands for a pair of photons [36], W [37, 38] or Z [39] since these final states are the most constraining for sgoldstinos with discussed masses. Due to large tree level coupling to the gluons M_3, gluon-gluon fusion will be the dominating mechanism for sgoldstino production with cross-section given by the following expression

$$\sigma_{\tilde{s}} = \frac{\pi^2}{8}\frac{\Gamma(\tilde{s} \to gg)}{s m_{\tilde{s}}} \int_{m_{\tilde{s}}^2/s}^{1} \frac{dx}{x} f_{p/g}(x, m_{\tilde{s}^2}) f_{p,\bar{p}/g}\left(\frac{m_{\tilde{s}}^2}{xs}, m_{\tilde{s}}^2\right), \tag{14}$$

where $\Gamma(\tilde{s} \to gg)$ is a partial width of sgoldstino-like state decaying into two gluons, s is a center-of-mass energy squared and $f_{p/g}(x, Q^2)$ are the parton density functions defined at factorisation scale Q. We use CTEQ6L [40] parametrization for the parton density functions. Finally we make use of constraints on $\mathrm{Br}(\tau \to \mu\gamma)$ [41] and $\mathrm{Br}(\tau \to 3\mu)$ [42] decays. For more detailed discussion on calculation of $\mathrm{Br}(\tau \to \mu\gamma)$ in our model see arguments in section 2.2 and in [15]. The result of the scan over parameter space are presented on Fig.1-5. For illustrating purposes, we present only the models with sufficiently large branching fraction $\mathrm{Br}(\tilde{\mathrm{h}} \to \mu\tau) > 5 \cdot 10^{-4}$. By blue color we mark the models which are capable of explaining the CMS excess, by purple color to we mark points which lie somewhat below the CMS excess but still have significant (more than 0.2%). According to the latest studies this level of branching fraction will be reachable at experiments on HL-LHC and ILC; see refs. [43–45]. We also make predictions for flavour-violating decays of τ-lepton (Fig.2),signal strength μ_f for different final states f and Higgs production channels (Fig.3,4) and make predictions for observable $\sigma(pp \to \tilde{s}) \times \mathrm{Br}(\tilde{s} \to \mathrm{W^+W^-}$ or ZZ) for the case of 13 TeV (Figs.5,6). For more results and detailed discussion see [15].

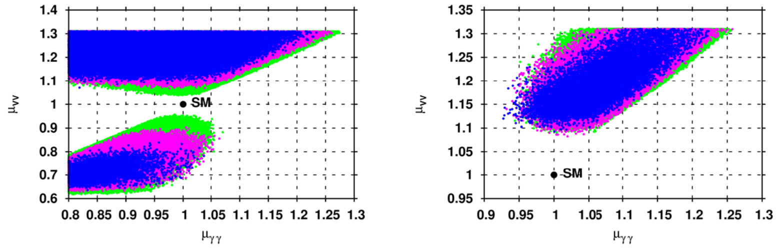

Figure 3. Scatter plots in plane $\left(\mu_{\gamma\gamma}, \mu_{VV}\right)$ for $\sqrt{F}$ = 8 TeV and sgoldstinos lighter (left) and heavier (right) than Higgs.

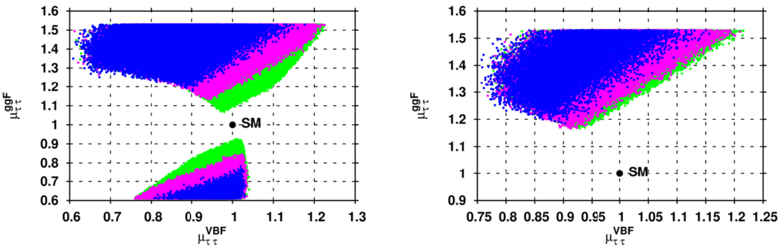

Figure 4. Scatter plots in plane $\left(\mu_{\tau\tau}^{\mathbf{VBF}}, \mu_{\tau\tau}^{\mathbf{ggF}}\right)$ for $\sqrt{F}$ = 8 TeV and sgoldstinos lighter (left) and heavier (right) than Higgs.

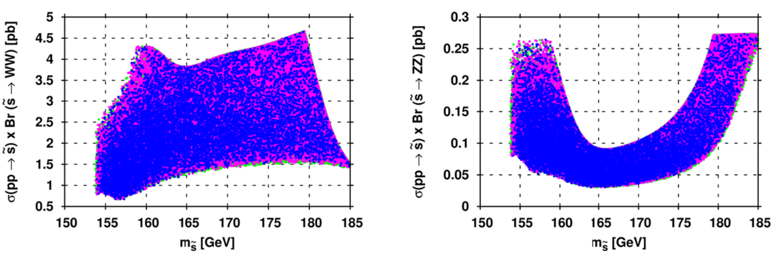

Figure 5. Scatter plots in $\sigma\,(pp \to \tilde{s}) \times \mathrm{Br}\,(\tilde{s} \to W^+W^-$ or $ZZ)$-plane for heavy sgoldstino and $\sqrt{s}$ = 13 TeV.

4 Conclusions

In this contribution we showed that in models with low scale supersymmetry breaking, the Higgs boson can have considerable branching ratio of $h \to \mu\tau$ decay due to mixing with sgoldstino and demonstrated that the CMS excess in $h \to \mu\tau$ decay can be explained in this framework. We stress, that the features we pointed out in [15] are common in the class of models in question.

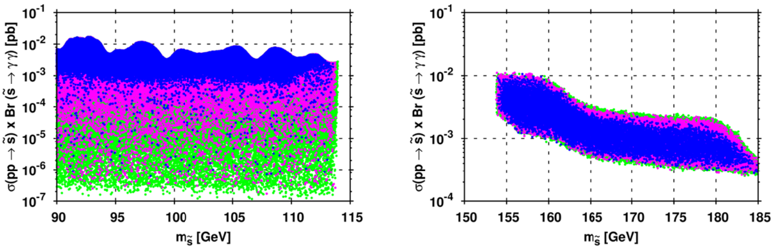

Figure 6. Scatter plots in $\sigma(pp \to \tilde{s}) \times \mathrm{Br}(\tilde{s} \to \gamma\gamma)$-plane for light (left panel) and heavy (right panel) sgoldstino and $\sqrt{s} = 13$ TeV.

Acknowledgements

This work is supported by the RSCF grant 14-22-00161. Numerical calculations have been performed on the Computational cluster of the Theoretical division of INR RAS.

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