SM*A*S*H

Andreas Ringwald,

Deutsches Elektronen-Synchrotron (DESY), Hamburg, Germany

DOI: http://dx.doi.org/10.3204/DESY-PROC-2016-03/Ringwald_Andreas

We present a minimal model for particle physics and cosmology. The Standard Model (SM) particle content is extended by three right-handed SM-singlet neutrinos N_i and a vector-like quark Q, all of them being charged under a global lepton number and Peccei-Quinn (PQ) $U(1)$ symmetry which is spontaneously broken by the vacuum expectation value $v_\sigma \sim 10^{11}$ GeV of a SM-singlet complex scalar field σ. Five fundamental problems – neutrino oscillations, baryogenesis, dark matter, inflation, strong CP problem – are solved at one stroke in this model, dubbed "SM*A*S*H" (Standard Model*Axion*Seesaw*Higgs portal inflation). It can be probed decisively by upcoming cosmic microwave background and axion dark matter experiments.

1 The quest for a minimal model of particle cosmology

The discovery of the Higgs boson has marked the completion of the SM particle content. However, observations in particle physics, astrophysics, and cosmology point to the existence of particles and interactions beyond the SM. In fact, the SM lacks an explanation of *i)* neutrino oscillations, *ii)* the baryon asymmetry of the Universe, *iii)* dark matter, *iv)* inflation, and *v)* the non-observation of strong CP violation.

Remarkably, problems *1)-3)* are solved in the Neutrino Minimal SM (νMSM) [1, 2]: a minimal extension of the SM by three right-handed singlet neutrinos N_i, having Dirac masses $m_D = Fv/\sqrt{2}$ arising from Yukawa couplings F with the Higgs (H) and lepton (L_i) doublets, as well as explicit Majorana masses M,

$$\mathcal{L} \supset -\left[F_{ij} L_i \epsilon H N_j + \frac{1}{2} M_{ij} N_i N_j\right], \tag{1}$$

where we have exploited a Weyl spinor notation. In the seesaw limit, $M \gg m_D$, the neutrino mass spectrum splits into a light set given by the eigenvalues $m_1 < m_2 < m_3$ of the matrix $m_\nu = -m_D M^{-1} m_D^T$, with the eigenstates corresponding mainly to mixings of the active left-handed neutrinos ν_α, and a heavy set given by the eigenvalues $M_1 < M_2 < M_3$ of the matrix M, with the eigenstates corresponding to mixings of the sterile right-handed neutrinos N_i. Problem *1)* is thus solved by the usual seesaw type-I mechanism. Intriguingly, problems *2)* and *3)* can be solved simultaneously if $M_1 \sim$ keV and $M_2 \sim M_3 \sim$ GeV. In fact, in this case $N_{2,3}$ create flavored lepton asymmetries from CP-violating oscillations in the early Universe which are crucial for the generation of the baryon asymmetry of the Universe via flavored leptogenesis and of the lightest sterile neutrino N_1 – the dark matter candidate of the νMSM – by the MSW effect. Moreover, it was argued in Ref. [3] that also problem *4)* can be solved in the νMSM by allowing a non-minimal coupling of the Higgs field to the Ricci scalar, $S \supset -\int d^4x\sqrt{-g}\,\xi_H\, H^\dagger H R$, which promotes the Higgs field to an inflaton candidate.

However, the success of the νMSM as a minimal model of particle cosmology is threatened by several facts. First of all, recent findings in astrophysics have seriously constrained the parameter space for N_1 as a dark matter candidate [4, 5]. Secondly, the large value of the non-minimal coupling $\xi_H \sim 10^5\sqrt{\lambda_H}$, where λ_H is the Higgs self-coupling, required to fit the amplitude of the scalar perturbations inferred from the cosmic microwave background (CMB) temperature fluctuations, imply that perturbative unitarity breaks down at the scale $M_P/\xi_H \sim 10^{14}\,\text{GeV}$, well below the scale of inflation, $M_P/\sqrt{\xi_H} \sim 10^{16}\,\text{GeV}$, making the inflationary predictions unreliable [6, 7]. Thirdly, Higgs inflation cannot be realised at all if the Higgs quartic coupling λ_H runs negative at large (Planckian) field values due to the corrections from top quark loops. Although, given the current experimental uncertainties, a definite conclusion cannot yet be drawn, see e.g. [8, 9], the presently favoured central values of the strong gauge coupling and the Higgs and top quark masses imply that λ_H becomes negative at a field value corresponding to an energy scale $\Lambda_I \sim 10^{11}$ GeV, much lower than what is required for Higgs inflation, and is thus inconsistent with it.

These three obstacles of the νMSM are circumvented in SMASH - an extension of the SM which features the Axion, the type-I Seesaw and Higgs portal inflation [10, 11] - as we will review in these proceedings.

2 The SMASH model

The SM particle content is extended not only by three right-handed singlet neutrinos N_i, but also by a vector-like color-triplet quark Q, as in the KSVZ [12, 13] model. The SM quarks and leptons as well as the N_i and Q are assumed to be charged under a global lepton number and PQ $U(1)$ symmetry [14] which is spontaneously broken by the vacuum expectation value $v_\sigma \sim 10^{11}$ GeV of a SM-singlet complex scalar field σ. The most general scalar potential reads

$$V(H,\sigma) = \lambda_H\left(H^\dagger H - \frac{v^2}{2}\right)^2 + \lambda_\sigma\left(|\sigma|^2 - \frac{v_\sigma^2}{2}\right)^2 + 2\lambda_{H\sigma}\left(H^\dagger H - \frac{v^2}{2}\right)\left(|\sigma|^2 - \frac{v_\sigma^2}{2}\right),$$

while the most general Yukawa couplings of the new fields are given by

$$\mathcal{L} \supset -\left[F_{ij}L_i\epsilon H N_j + \frac{1}{2}Y_{ij}\sigma N_i N_j + y\,\tilde{Q}\sigma Q + y_{Q_d\,i}\sigma Q d_i + h.c.\right].$$

After $U(1)$ symmetry breaking the sterile neutrinos N_i, the particle excitation ρ of the modulus of the σ field, and the exotic quark Q get large masses $\propto v_\sigma \gg v = 246\,\text{GeV}$: $M_{ij} = \frac{Y_{ij}}{\sqrt{2}}v_\sigma + \mathcal{O}\left(\frac{v}{v_\sigma}\right)$, $m_\rho = \sqrt{2\,\lambda_\sigma}\,v_\sigma + \mathcal{O}\left(\frac{v}{v_\sigma}\right)$, and $m_Q = \frac{y}{\sqrt{2}}\,v_\sigma + \mathcal{O}\left(\frac{v}{v_\sigma}\right)$. Therefore, as far as physics around the electroweak scale or below is concerned, these heavy particles can be integrated out (unless one considers tiny Yukawa and self couplings). The corresponding low-energy Lagrangian of SMASH is identical to that of the SM, augmented by seesaw-generated neutrino masses, $m_\nu = 0.04\,\text{eV}\left(\frac{10^{11}\,\text{GeV}}{v_\sigma}\right)\left(\frac{-F\,Y^{-1}\,F^T}{10^{-4}}\right)$, and mixing (thus solving problem *1)*), plus one new particle: the particle excitation A of the angular degree of freedom of the complex σ field – the Nambu-Goldstone boson of the spontaneous symmetry breaking of the $U(1)$, which is dubbed "axion" in the literature dealing with the PQ solution of the strong CP

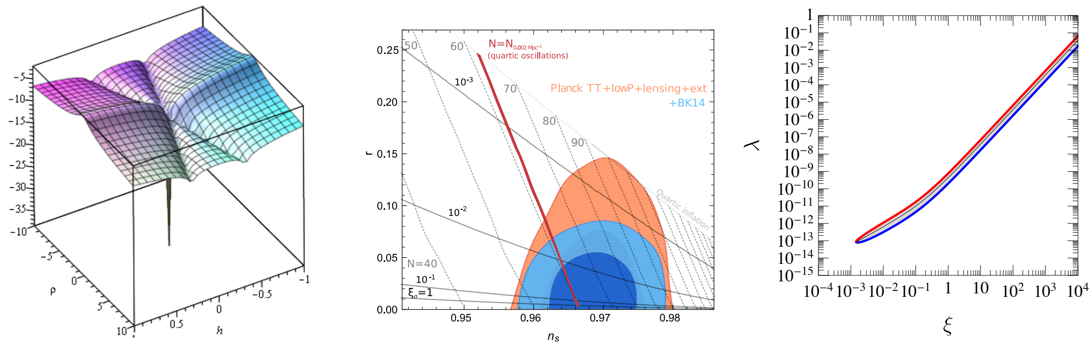

Figure 1: Left: Decadic log of the SMASH scalar potential in the Einstein frame, as a function of h and ρ, all in units of M_P, for $\kappa_H < 0$, $\kappa_\sigma < 0$, supporting mixed Higgs-Hidden-Scalar Inflation along one of the valleys. Middle: Bounds on r vs. n_s [18], compared to the predictions from (H)HSI in SMASH for fixed values of the non-minimal coupling ξ_σ and the number of e-folds N, respectively. Right: Required self-coupling versus non-minimal coupling to reproduce CMB results on inflation. All figures from [11].

problem [15, 16] and "majoron" in the literature dealing with the spontaneous breaking of a global lepton symmetry. Integrating out the exotic quark induces an anomalous coupling of the axion field to the topological charge density in QCD, $\mathcal{L} \supset -\frac{\alpha_s}{8\pi}\frac{A}{f_A}\, G^c_{\mu\nu}\tilde{G}^{c,\mu\nu}$, promoting the axion field to a dynamical theta parameter, $\theta(x) = A(x)/f_A$, which relaxes to zero in the vacuum, $\langle\theta\rangle = 0$, thereby solving problem *5)*. While the strong CP problem is solved for any value of the axion decay constant $f_A = v_\sigma$, the dark matter will be comprised by axions only if f_A is around 10^{11} GeV, as we will see later. In this case, the axion mass is predicted to be around $m_A = 57.0(7)\left(\frac{10^{11}\mathrm{GeV}}{f_A}\right)\mu\mathrm{eV}$ [15, 17].

3 Inflation

The non-minimal couplings in SMASH, $S \supset -\int d^4x\sqrt{-g}\,\left[\xi_H\, H^\dagger H + \xi_\sigma\,\sigma^*\sigma\right] R$, stretch the scalar potential in the Einstein frame, which makes it convex and asymptotically flat at large field values. Depending on the signs of the parameters $\kappa_H \equiv \lambda_{H\sigma}\xi_H - \lambda_H\xi_\sigma$ and $\kappa_\sigma \equiv \lambda_{H\sigma}\xi_\sigma - \lambda_\sigma\xi_H$, it can support Higgs Inflation (HI), Hidden Scalar Inflation (HSI), or even mixed Higgs-Hidden Scalar Inflation (HHSI) (cf. Fig. 1 (left)). For $\xi \simeq 10^5\sqrt{\lambda} \gtrsim 10^{-3}$, where

$$\xi \equiv \begin{cases} \xi_H, & \text{for HI,} \\ \xi_\sigma, & \text{for HSI,} \\ \xi_\sigma, & \text{for HHSI,} \end{cases} \qquad \lambda \equiv \begin{cases} \lambda_H, & \text{for HI,} \\ \lambda_\sigma, & \text{for HSI,} \\ \lambda_\sigma\left(1 - \frac{\lambda^2_{H\sigma}}{\lambda_\sigma\lambda_H}\right), & \text{for HHSI,} \end{cases} \tag{2}$$

the predicted values of the CMB observables such as the amplitude of scalar perturbations A_s, the spectral index n_s, and the tensor-to-scalar ratio r are in perfect consistency with the current observations, see Fig. 1 (middle). Importantly, for (H)HSI, the effective self-coupling λ is a free parameter and therefore can be chosen small, $\lambda \sim 10^{-10}$, such that the required non-minimal coupling to fit the amplitude of primordial scalar perturbations is of order unity,

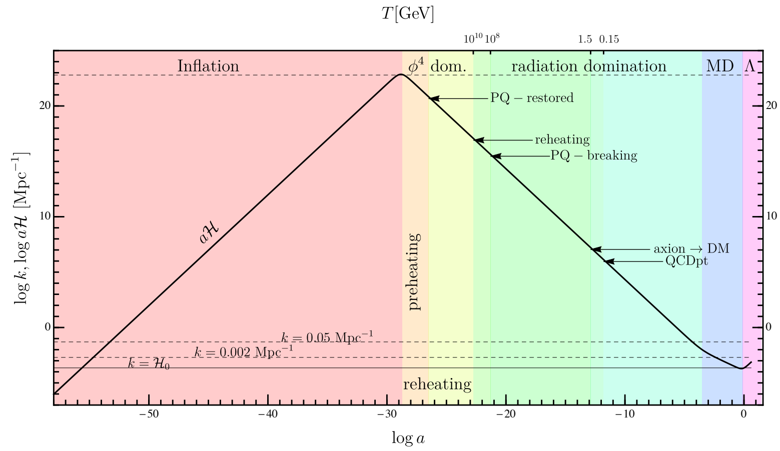

Figure 2: The history of the Universe in SMASH HHSI, emphasising the transition from inflation to radiation-domination-like Universe expansion $a\mathcal{H} \propto 1/a$ before standard matter and cosmological constant domination epochs [11].

$\xi_\sigma \sim 1$, cf. Fig. 1 (right). In this region of parameter space, the perturbative predictivity of SMASH is guaranteed and superior to HI, which necessarily operates at large ξ_H, since λ_H is not small. Remarkably, the requirement of predictive inflation, free of unitarity problems, demands $r \gtrsim 0.01$, which will be probed by CMB experiments such as LiteBIRD and PRISM.

4 Stability

Self-consistency of inflation in SMASH requires a positive scalar potential all the way up to the Planck scale. Importantly, the Higgs portal term $\propto \lambda_{H\sigma}$ in the scalar potential helps to ensure absolute stability in the Higgs direction via the threshold stabilisation mechanism pointed out in [19, 20]. We have found that stability can be achieved if the threshold parameter $\delta = \lambda_{H\sigma}^2/\lambda_\sigma$ is between 10^{-3} and 10^{-1}. Instabilities could also originate in the σ direction, due to quantum corrections from the right-handed neutrinos N_i and the exotic quark Q. Stability in the σ direction then restricts their Yukawas to $\sum Y_{ii}^4 + 6y^4 \lesssim 16\pi^2\lambda_\sigma/\log\left(30M_P/\sqrt{2\lambda_\sigma}v_\sigma\right)$.

5 Reheating

Both in (H)HSI, slow-roll inflation ends at a value of $\rho \sim \mathcal{O}(M_P)$, where the effect of $\xi_\sigma \sim 1$ is negligible and the inflaton starts to undergo Hubble-damped oscillations in a quartic potential, with the Universe expanding as in a radiation-dominated era, which lasts until reheating, cf. Fig. 2. After the latter, radiation domination continues, though driven by a bath of relativistic particles. This fixes the thick red line in Fig. 1 (middle) as the prediction for r, n_s and N in SMASH. The fluctuations of σ grow fast due to parametric resonance while the inflaton background oscillates in its quartic potential, leading to a rapid restoration of the PQ symmetry after about 14 oscillations. The following reheating stage differs considerably for HSI and HHSI.

In the former, the large induced particle masses quench inflaton decays or annihilations into SM particles, resulting in a low reheating temperature, $T_R \sim 10^7$ GeV, such that the produced relativistic axions are never thermalized. Correspondingly, HSI predicts a significant amount of cosmic axion background radiation (CAB): an increase $\Delta N_\nu^{\rm eff} = \mathcal{O}(1)$ of the effective number of relativistic neutrino species beyond the SM value $N_\nu^{\rm eff}({\rm SM}) = 3.046$ [21]. This disvafors HSI, since the current results from CMB and baryon acoustic oscillations yield $N_\nu^{\rm eff} = 3.04 \pm 0.18$ at 68% CL and thus do not allow an additional contribution of order one [18]. For this reason, inflation in SMASH must be of HHSI type and therefore the inflaton contains a (small) Higgs component. The latter allows for efficient reheating of the Universe by the production of SM gauge bosons. The reheating temperature in this case is predicted to be around $T_R \sim 10^{10}$ GeV. Such temperature ensures a thermal restoration of the PQ symmetry for the relevant region of parameter space, since the critical temperature T_c of the PQ phase transition goes as $T_c/v_\sigma \simeq 2\sqrt{6\lambda_\sigma}/\sqrt{8(\lambda_\sigma + \lambda_{H\sigma}) + \sum_i Y_{ii}^2 + 6y^2}$. A thermal background of axions is produced at this stage which later decouples and results in a moderate CAB corresponding to $\triangle N_\nu^{\rm eff} \simeq 0.03$, a prediction which may be checked in a future CMB polarisation experiment.

6 Dark Matter

Dark matter is produced in SMASH by the re-alignment mechanism [22, 23, 24] and the decay of topological defects (axion strings and domain walls) [25]. In order to account for all of the cold dark matter in the Universe, the PQ symmetry breaking scale is predicted to be in the range $3 \times 10^{10}\,{\rm GeV} \lesssim f_A \lesssim 1.2 \times 10^{11}\,{\rm GeV}$, corresponding to an axion mass in the range $50\,\mu{\rm eV} \lesssim m_A \lesssim 200\,\mu{\rm eV}$ [11, 17]. Here, the uncertainty originates mainly from the difficulty in predicting the relative importance of the two main production mechanisms of axionic dark matter, i.e. re-alignment and topolical defect decay. Fortunately, the axion dark matter mass window will be probed in the upcoming decade by axion dark matter direct detection experiments such as CULTASK, MADMAX, and ORPHEUS.

7 Baryogenesis

The origin of the baryon asymmetry of the Universe is explained in SMASH by thermal leptogenesis [26]. In HHSI, after reheating and thermal PQ restoration, the RH neutrinos become massive and at least the lightest RH neutrino N_1 will retain an equilibrium abundance. However the stability bound on $M_1 \lesssim 10^8\,(\lambda/10^{-10})^{1/4}(v_\sigma/10^{11}{\rm GeV})\,{\rm GeV}$, for a hierarchical N_i spectrum ($M_3 = M_2 = 3M_1$), is just borderline compatible with vanilla leptogenesis from the decays of N_1, which demands $M_1 \gtrsim 5 \times 10^8$ GeV [27, 28]. Nevertheless, leptogenesis can occur with a mild resonant enhancement [29] for a less hierarchical RH neutrino spectrum, which relaxes the stability bound and ensures that all the RH neutrinos remain in equilibrium after the phase transition.

8 Acknowledgments

Many thanks to Guillermo Ballesteros, Javier Redondo and Carlos Tamarit for the great collaboration.

References

[1] T. Asaka, S. Blanchet and M. Shaposhnikov, "The nuMSM, dark matter and neutrino masses," Phys. Lett. B **631** (2005) 151 [hep-ph/0503065].

[2] T. Asaka and M. Shaposhnikov, "The nuMSM, dark matter and baryon asymmetry of the universe," Phys. Lett. B **620** (2005) 17 [hep-ph/0505013].

[3] F. L. Bezrukov and M. Shaposhnikov, "The Standard Model Higgs boson as the inflaton," Phys. Lett. B **659** (2008) 703 [arXiv:0710.3755 [hep-th]].

[4] A. Schneider, "Astrophysical constraints on resonantly produced sterile neutrino dark matter," JCAP **1604** (2016) no.04, 059 [arXiv:1601.07553 [astro-ph.CO]].

[5] K. Perez, K. C. Y. Ng, J. F. Beacom, C. Hersh, S. Horiuchi and R. Krivonos, "(Almost) Closing the Sterile Neutrino Dark Matter Window with NuSTAR," arXiv:1609.00667 [astro-ph.HE].

[6] J. L. F. Barbon and J. R. Espinosa, "On the Naturalness of Higgs Inflation," Phys. Rev. D **79** (2009) 081302 [arXiv:0903.0355 [hep-ph]].

[7] C. P. Burgess, H. M. Lee and M. Trott, "Power-counting and the Validity of the Classical Approximation During Inflation," JHEP **0909** (2009) 103 [arXiv:0902.4465 [hep-ph]].

[8] D. Buttazzo, G. Degrassi, P. P. Giardino, G. F. Giudice, F. Sala, A. Salvio and A. Strumia, "Investigating the near-criticality of the Higgs boson," JHEP **1312** (2013) 089 [arXiv:1307.3536 [hep-ph]].

[9] A. V. Bednyakov *et al.*, "Stability of the Electroweak Vacuum: Gauge Independence and Advanced Precision," Phys. Rev. Lett. **115** (2015) 20, 201802 [arXiv:1507.08833 [hep-ph]].

[10] G. Ballesteros, J. Redondo, A. Ringwald and C. Tamarit, "Unifying inflation with the axion, dark matter, baryogenesis and the seesaw mechanism," arXiv:1608.05414 [hep-ph].

[11] G. Ballesteros *et al.*, "Standard Model-Axion-Seesaw-Higgs Portal Inflation. Five problems of particle physics and cosmology solved in one stroke," arXiv:1610.01639 [hep-ph].

[12] J. E. Kim, "Weak Interaction Singlet and Strong CP Invariance," Phys. Rev. Lett. **43** (1979) 103.

[13] M. A. Shifman, A. I. Vainshtein and V. I. Zakharov, "Can Confinement Ensure Natural CP Invariance of Strong Interactions?," Nucl. Phys. B **166** (1980) 493.

[14] R. D. Peccei and H. R. Quinn, "CP Conservation in the Presence of Instantons," Phys. Rev. Lett. **38** (1977) 1440.

[15] S. Weinberg, "A New Light Boson?," Phys. Rev. Lett. **40** (1978) 223.

[16] F. Wilczek, "Problem of Strong p and t Invariance in the Presence of Instantons," Phys. Rev. Lett. **40** (1978) 279.

[17] S. Borsanyi *et al.*, "Lattice QCD for Cosmology," arXiv:1606.07494 [hep-lat].

[18] P. A. R. Ade *et al.* [Planck Collaboration], "Planck 2015 results. XIII. Cosmological parameters," arXiv:1502.01589 [astro-ph.CO].

[19] O. Lebedev, "On Stability of the Electroweak Vacuum and the Higgs Portal," Eur. Phys. J. C **72** (2012) 2058 [arXiv:1203.0156 [hep-ph]].

[20] J. Elias-Miro, J. R. Espinosa, G. F. Giudice, H. M. Lee and A. Strumia, "Stabilization of the Electroweak Vacuum by a Scalar Threshold Effect," JHEP **1206** (2012) 031 [arXiv:1203.0237 [hep-ph]].

[21] G. Mangano, G. Miele, S. Pastor and M. Peloso, "A Precision calculation of the effective number of cosmological neutrinos," Phys. Lett. B **534** (2002) 8 [astro-ph/0111408].

[22] J. Preskill, M. B. Wise and F. Wilczek, "Cosmology of the Invisible Axion," Phys. Lett. B **120** (1983) 127.

[23] L. F. Abbott and P. Sikivie, "A Cosmological Bound on the Invisible Axion," Phys. Lett. B **120** (1983) 133.

[24] M. Dine and W. Fischler, "The Not So Harmless Axion," Phys. Lett. B **120** (1983) 137.

[25] M. Kawasaki, K. Saikawa and T. Sekiguchi, "Axion dark matter from topological defects," Phys. Rev. D **91** (2015) 6, 065014 [arXiv:1412.0789 [hep-ph]].

[26] M. Fukugita and T. Yanagida, "Baryogenesis Without Grand Unification," Phys. Lett. B **174** (1986) 45.

[27] G. F. Giudice, A. Notari, M. Raidal, A. Riotto and A. Strumia, "Towards a complete theory of thermal leptogenesis in the SM and MSSM," Nucl. Phys. B **685** (2004) 89 [hep-ph/0310123].

[28] W. Buchmuller *et al.*, "Leptogenesis for pedestrians," Annals Phys. **315** (2005) 305 [hep-ph/0401240].

[29] A. Pilaftsis and T. E. J. Underwood, "Resonant leptogenesis," Nucl. Phys. B **692** (2004) 303.