Motivation for Weakly Interacting SubeV Particles

Ignatios Antoniadis[*1]

[1]Department of Physics, CERN - Theory Unit, CH-1211 Geneva 23, Switzerland

DOI: http://dx.doi.org/10.3204/DESY-PROC-2010-03/antoniadis_ignatios

We review sources and physics of WISPs (weakly interacting subeV particles) from theories beyond the Standard Model, such as string theory or models involving non-trivial anomaly cancellation. In particular, we discuss extra short range forces, axion-like particles (ALPs) and extra $U(1)$s.

WISPs appear in several theories beyond the Standard Model. In particular, they are generic in string compactifications but have varying properties in different classes of models, such as supersymmetric compactifications with high string scale or models with low string scale and large extra dimensions. Examples of WISPs are: light pseudoscalars, extra $U(1)$s, light scalars, as well as their possible superpartners, giving rise to axions or in general ALPs[1], and extra short range forces.

Indeed, string compactifications lead to scalar moduli whose vacuum expectation values (VEVs) parametrize the geometry of the internal compactified space, such as the size of cycles, shapes, the string coupling itself, etc. These VEVs are arbitrary because moduli have no potential (at least) at the classical level, which creates a serious problem since all low energy couplings are functions of moduli. In supersymmetric compactifications the moduli are complexified with pseudoscalars coming from internal components of higher-rank antisymmetric tensors present in the low energy spectrum. The moduli stabilization is a long outstanding problem, necessary to provide moduli masses (avoiding experimental conflict from long range forces and cosmology) and to fix their VEVs (allowing computation of the low energy couplings). A moduli potential can be generated either by non-perturbative effects or by turning on fluxes for the internal components of the higher-rank gauge potentials, generalizing ordinary magnetic fields.

Another source of light ALPs and extra $U(1)$s is related to non-trivial anomaly cancellation. In fact, theories in which fermions have chiral couplings with gauge fields are known to suffer from *anomalies* – a phenomenon of breaking of gauge symmetries of a classical theory at one-loop level. Anomalies make a theory inconsistent (in particular, its unitarity is lost). The only way to restore its consistency is to arrange for an exact *cancellation* of anomalies between the various chiral sectors. This happens, for example, in the Standard Model (SM), where the cancellation occurs between quarks and leptons within each generation [1]. Another well studied example is the Green-Schwarz anomaly cancellation mechanism [2] in string theory. In this case the cancellation arises between the anomalous contribution of chiral matter of the closed string sector with that of the open string.[2]

[*]On leave from CPHT (UMR CNRS 7644) Ecole Polytechnique, F-91128 Palaiseau.

[1]ALPs are axion like particles with no particular relation between their mass and decay constant ($m_a f_a = m_\pi f_\pi$ for ordinary axions).

[2]Formally, the Green-Schwarz anomaly cancellation occurs due to the anomalous Bianchi identity for the field

Particles involved in anomaly cancellation may have very different masses. For example, the mass of top quark in the SM is much higher than the masses of all other fermions. However, gauge invariance should pertain at all energies, including those which are smaller than the mass of some particles involved in anomaly cancellation. The usual logic of field theory is that interactions, mediated by heavy fermions running in loops, are suppressed by the masses of these fermions [4]. The case of anomaly cancellation presents a notable counterexample to this famous "decoupling theorem" – the contribution of *a priori* arbitrary heavy particles should remain unsuppressed at low energies. As it was pointed out in [5], this is possible because anomalous (i.e. gauge-variant) terms in the effective action have topological nature and therefore are scale independent. As a result, they are not suppressed even at energies much smaller than the masses of the particles producing these terms via loop effects. This gives a hope to see at low energies some signatures of new physics.

In the following, we first discuss sources of new short range forces and microgravity experiments, then axion like particles and finally effects of non-trivial anomaly cancellation involving extra $U(1)$s.

1 5th force and microgravity experiments

Theories with large extra dimensions predict modifications of gravitation in the sub-millimeter range, which can be tested in "table-top" experiments that measure gravity at short distances. There are three categories of such predictions:
(i) Deviations from the Newton's law $1/r^2$ behavior to $1/r^{2+n}$, which can be observable for $n = 2$ large extra dimensions of sub-millimeter size.
(ii) New scalar forces in the sub-millimeter range, related to the mechanism of supersymmetry breaking, and mediated by light scalar fields φ with masses [6, 7]:

$$m_\varphi \simeq \frac{m_{susy}^2}{M_P} \simeq 10^{-4} - 10^{-6} \text{ eV}\,, \tag{1}$$

for a supersymmetry breaking or string scale $m_{susy} \simeq 1 - 10$ TeV; they correspond to Compton wavelengths of 1 mm to 10 μm. A universal attractive scalar force is mediated by the so-called radion modulus field. Such a force can be tested in microgravity experiments and should be contrasted with the change of Newton's law due the presence of extra dimensions that is observable only for $n = 2$ [9, 10]. The resulting bounds from an analysis of the radion effects are [11]: $M_* \gtrsim 6\,\text{TeV}$.
(iii) Non universal repulsive forces much stronger than gravity, mediated by possible abelian gauge fields in the bulk [12, 13]. Such fields acquire tiny masses of the order of M_s^2/M_P, as in (1), due to brane localized anomalies [13]. Although their gauge coupling is infinitesimally small, $g_A \sim M_s/M_P \simeq 10^{-16}$, it is still bigger that the gravitational coupling E/M_P for typical energies $E \sim 1$ GeV, and the strength of the new force would be $10^6 - 10^8$ stronger than gravity.

In Fig. 1 we depict the actual information from previous, present and upcoming experiments [10, 8]. The solid lines indicate the present limits from the experiments indicated. The excluded regions lie above these solid lines. Measuring gravitational strength forces at short

strength of the closed 2-form. However, this modification of Bianchi identity arises from the 1-loop contribution of chiral fermions in the open string sector. A toy model, describing microscopically Green-Schwarz mechanism was studied e.g. in [3].

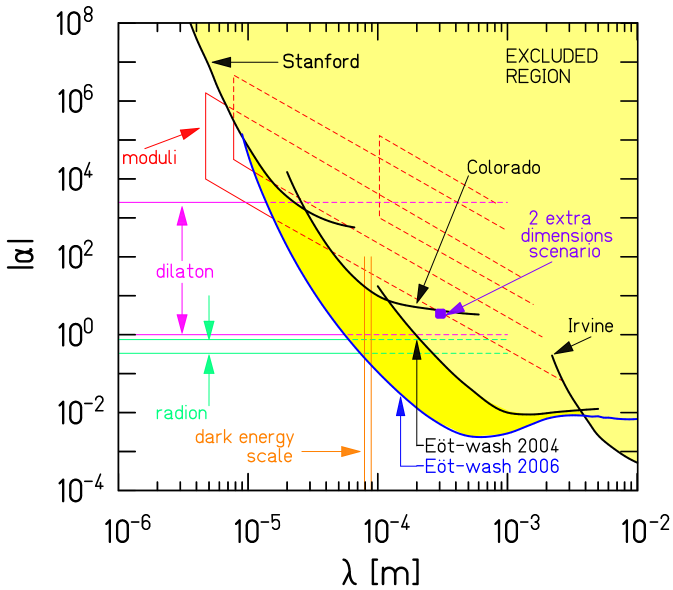

Figure 1: Present limits on new short-range forces (yellow regions), as a function of their range λ and their strength relative to gravity α.

distances is challenging. The horizontal lines correspond to theoretical predictions, in particular for the graviton in the case $n = 2$ and for the radion.

2 Axion Like Particles

As mentioned in the introduction pseudoscalar partners of moduli (Poincarée duals to two-index antisymmetric tensors in four dimensions) are often associated to perturbative shift symmetries and are thus candidates for axions or in general ALPs. Their common characteristic is the $\vec{E}\cdot\vec{B}$ type coupling to gauge fields, such as the photon:

$$\mathcal{L}_{a\gamma\gamma} = \frac{a}{4f_a}\epsilon_{\mu\nu\lambda\rho}F^{\mu\nu}F^{\lambda\rho} = -\frac{1}{2f_a}\epsilon^{\mu\nu\lambda\rho}\left(\partial_\mu a\right)A_\nu\,.F_{\lambda\rho} \tag{2}$$

Their mass m_a may be related to new strong interactions scales, or setup by the string scale, or suppressed by the compactification volume, or related to the supersymmetry breaking scale as in eq. (1). Similarly, their decay constant f_a may be related to a different physical scale independently of m_a but in general is strongly constrained from astrophysics to be $f_a \gtrsim 10^{10}$ GeV. Indeed the experimental bounds in the two-parameter space are shown in Fig. 2.

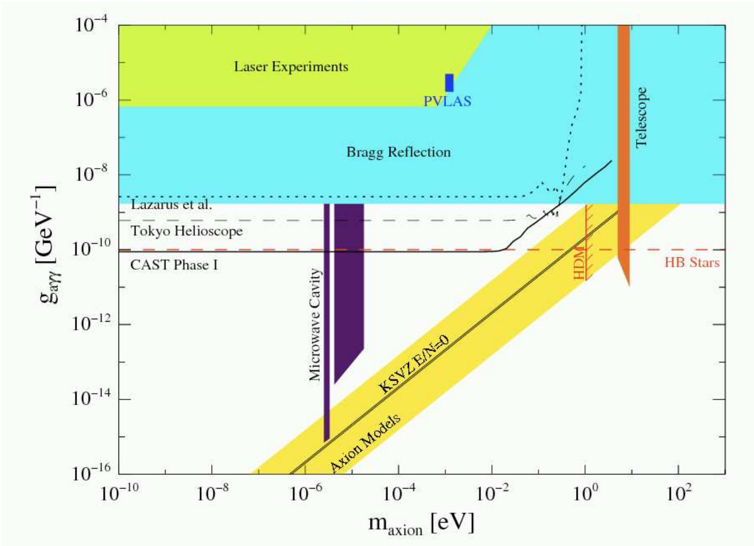

Figure 2: Present experimental bounds on ALPs.

3 Extra $U(1)$s and mixed anomaly involving photon

As discussed previously, non-trivial anomaly cancellation generically should involve at least one gauge field beyond the SM gauge sector. To reconcile this with existing experimental bounds, such an anomaly cancellation should take place between SM and "hidden" sector, with new particles appearing at relatively high energies. Here we concentrate on the case of one additional Abelian group. Extra $U(1)$ fields appear in many extensions of the Standard Model (see e.g. [14] and refs. therein). For example, additional $U(1)$s appear naturally in models in which $SU(2)$ and $SU(3)$ gauge factors of the SM arise as parts of unitary $U(2)$ and $U(3)$ groups (as e.g. in D-brane constructions of the SM [15, 16, 17, 18]). A common feature of these models is a non-trivial cancellation of anomalies between various sectors of the theory.

If the mixed anomaly cancellation between several groups of fermions involves the photon field A_μ, terms (often called *Generalized Chern-Simons*) can appear in the action

$$\mathcal{L}_{\rm CS} = -\frac{\kappa}{2}\epsilon^{\mu\nu\lambda\rho}X_\mu A_\nu F_{\lambda\rho} \tag{3}$$

where X_μ is an extra $U(1)$. Here κ is a dimensionless coupling constant. The Chern-Simons-like interaction (3) appears in various models (see e.g. [15, 18, 19, 20, 21, 22]). This term resembles an axion coupling to photon (2), under the identification of a derivatively coupled pseudo-scalar with the longitudinal part of the massive vector field X_μ, $\partial_\mu a \longrightarrow m_X X_\mu$, where m_X is the mass of this new vector boson. An analog of Peccei-Quinn scale f_a is played in this theory by the combination

$$f_a \leftrightarrow \frac{m_X}{\kappa} \tag{4}$$

Notice, that by making the coupling κ small, one can have $f_a \gg m_X$.

The simplest model involving the interaction term (3) is given by the effective action (with the masses generated via the Stückelberg mechanism):

$$S = \int d^4x \left(-\frac{1}{4}F_{\mu\nu}^2 - \frac{1}{4}X_{\mu\nu}^2 + \frac{m_X^2}{2}X_\mu^2 + \frac{m_\gamma^2}{2}A_\mu^2 + \frac{\kappa}{2}\epsilon^{\mu\nu\lambda\rho}X_\mu A_\nu F_{\lambda\rho} \right) \quad (5)$$

Here $X_{\mu\nu} = \partial_\mu X_\nu - \partial_\nu X_\mu$ is the field strength of the massive vector field X_μ. The Chern-Simons-like term is not gauge invariant under the electromagnetic gauge transformation $U(1)_{\text{QED}}$. To amend this drawback, we introduced a mass to the photon (m_γ), consistent with all existing restrictions (see e.g. [23] for current bounds on the photon mass). Alternatively, one can impose an additional constraint in the theory: $F_{\mu\nu}X_{\lambda\rho}\epsilon^{\mu\nu\lambda\rho} = 0$ [19].

For optical experiments the phenomenology of the theory (5) is similar to that of ALPs with (4) and ($m_X \leftrightarrow m_a$). However, at higher energies the phenomenology of the theory (5) can get significantly different. Indeed, if a massive vector field couples to a conserved current, all the processes involving the longitudinal degree of freedom are suppressed at energies much greater than its mass m_X. On the other hand, if the current is *not conserved*, for $E \gg m_X$ the longitudinal polarization behaves as a derivatively coupled scalar (the so called *Goldstone boson equivalence theorem* [24]).

This is what may happen in theory (5). Although the theory can be written in a formally gauge invariant form under the $U(1)_X$ gauge symmetry by introducing a Stückelberg field θ_X, the symmetry is realized by simultaneous gauge transformations of the X-field and θ_X-shifts. As a result, the field X_μ couples to a *non-conserved current* ($j_X^\mu \equiv \delta\mathcal{L}/\delta X_\mu$) and therefore its longitudinal polarization behaves as an axion (for $E \gg m_X$).

However, the theory (5) is an effective field theory, valid up to a certain energy scale Λ. This scale $\Lambda \lesssim \frac{m_X}{\kappa}$, as one can easily find by analyzing the unitarity bound in tree-level processes with outgoing longitudinally polarized X. It may naturally happen that for $E \gtrsim \Lambda$ the theory gets modified in such a way that the current j_X^μ becomes conserved. Then, all processes involving emission or absorption of the longitudinal polarization of X_μ get suppressed as $\left(\frac{m_X}{E}\right)^2$. As we are interested in the situation where the field X_μ can be produced at laboratory energies (e.g. in laser experiments), its mass should be $m_X \lesssim E_{lab} \sim$ eV.

The stringent constraints on ALPs come from stellar observations (see e.g. [25, 26]). The ALPs, created in stellar interior via the interactions (2), can significantly change burning cycles and life-times of stars (see [25]). To change the situation, as compared to a standard axion with energy-independent coupling, the scale of new physics Λ should be in the keV region $\Lambda \sim E_\star \sim$ keV. The conservation of the current j_X^μ implies a suppression of emission of longitudinal vector boson by *at least* $\sim (E_{lab}/E_\star)^2 \sim 10^{-6}$. Taking into account the astrophysical constraints, one finds $\kappa \lesssim 10^{-10}$ eV$/m_X$. Thus, the theory with Chern-Simons (CS) interaction (5) does not resemble the theory of ALPs. In particular, the production of X_μ is strongly suppressed by the small value of the dimensionless CS coupling κ.

To illustrate this idea, assume that there is an additional fermion with mass M_f, interacting with the fields of the theory (5) and giving rise to an effective action of the following (schematic) form:

$$\mathcal{L} = -\frac{1}{4}F_{\mu\nu}^2 - \frac{1}{4}X_{\mu\nu}^2 + \frac{m_X^2}{2}(D_\mu\theta_X)^2 + \frac{\kappa}{2}\epsilon^{\mu\nu\lambda\rho}X_\mu A_\nu F_{\lambda\rho} + \left(M_f^2\theta_X - \partial_\mu X^\mu\right)\frac{\kappa}{\Box + M_f^2}F\tilde{F} \quad (6)$$

where we introduced the notation $F\tilde{F} = \frac{1}{2}\epsilon^{\mu\nu\lambda\rho}F_{\mu\nu}F_{\lambda\rho}$. For simplicity of the presentation we work with the non-local action (6), but one can find an example of a renormalizable field theory which shares these properties in Ref. [20]. Recall that we add to this theory the constraint $\epsilon^{\mu\nu\lambda\rho}F_{\mu\nu}X_{\lambda\rho} = 0$ to make it gauge invariant with respect to $U(1)_A$ transformations.

Let us now demonstrate that this theory possesses the desired properties: At *low energies* (for $E < M_f$) one obtains the action (5) (formally taking $M_f \to \infty$). To analyze the theory at high energies ($E \gg M_f$), one can formally take $M_f \to 0$ and neglect the interaction term proportional to θ_X in the action (6). As a result at high energies, the field X_μ couples to the *conserved* current

$$j_X^\mu = \frac{\kappa}{2}\epsilon^{\mu\nu\lambda\rho}A_\nu F_{\lambda\rho} - \kappa\frac{\partial^\mu}{\Box}(F\tilde{F}) \ . \tag{7}$$

Therefore, at energies $E \gg M_f$ the production of the longitudinally polarized X_μ-field in theory (6) is suppressed. Of course, for $E > M_f$ the current (7) should be computed directly in the microscopic theory producing the non-local terms in (6), containing additional particles, rather than in the non-local effective theory. However, the conclusion remains the same.

The effect of decoupling of the longitudinal polarization of the vector boson at high energies, significantly changes the phenomenology. Most interestingly, it allows to reconcile the stellar constraints on ALPs (see e.g. [25, 26]) with a possible signal in the high precision optical experiments, outside the standard axion parameter space.

Notice, that such a model requires fermions with masses $E_{lab} \lesssim M_f < E_\star$, i.e. in the range from ~ 1 eV to ~ 1 keV. There are various restrictions on the charges q_f of such fermions. First, laboratory bounds, coming from the contribution to the Lamb shift and invisible orthopositronium decay [27] (based on the results of [28]) give $q_f < 10^{-4}$. A stronger bound on millicharged fermions ($q_f < 10^{-6}$) with sub-eV masses comes from the requirement that such fermions do not distort the CMB spectrum too much [29]. However, this restriction is not applicable in our case as the fermion masses are assumed to be above $M_f > E_{lab} \sim 1$ eV.

Finally, the strongest bound ($q_f < 10^{-14}$) on the charges of fermions with mass below $\lesssim 30$ keV comes from limiting the contribution of these particles to the energy transfer in stars [30] (see also [25]). To satisfy this bound, the vector field X_μ should be extremely light with $m_X \sim 10^{-10}$ eV and $\kappa \sim 10^{-28}$ [20] (which is a possibility). However, these bounds can be avoided in our model because the paraphoton field X_μ acquires a kinetic mixing with the photon due to the loop corrections coming from light fermions. Therefore, the mechanism of additional suppression of the coupling of fermions with the photon in stars, proposed in [31], is possible. The restriction then becomes $q_f \lesssim 10^{-14}\left(\frac{E_\star}{m_X}\right)^2$, i.e. the stellar bound of [30, 27] is weakened by *at least* six orders of magnitude, making the model compatible with existing observations (see [20] for details).

If a non-trivial anomaly cancellation *involves the electromagnetic* $U(1)$ *gauge group* observable effects may be present in optical experiments. Indeed, such high precision experiments (e.g. those measuring the change of polarization of light propagating in a strong magnetic field) could in principle see the anomalous terms, proportional to $\tilde{F}\cdot F = 4\vec{E}\cdot\vec{H} \neq 0$. There exists a significant experimental activity searching for such signals, as various ALPs are expected to couple to $\tilde{F}_{\mu\nu}F^{\mu\nu}$ and produce interesting signatures in parallel electric and magnetic fields. A different type of experiment using static fields, which may test effects caused by non-trivial anomaly cancellation in the electromagnetic sector, was suggested in [32].

Acknowledgements: Work supported in part by the European Commission under the ERC Advanced Grant "MassTeV" ERC-2008-AdG 20080228 and the ITN contract "UNILHC" PITN-GA-2009-237920.

References

[1] D. J. Gross, and R. Jackiw, *Phys. Rev.* **D6**, 477–493 (1972); C. Bouchiat, J. Iliopoulos, and P. Meyer, *Phys. Lett.* **B38**, 519–523 (1972); H. Georgi, and S. L. Glashow, *Phys. Rev.* **D6**, 429 (1972).

[2] M. B. Green, and J. H. Schwarz, *Phys. Lett.* **B149**, 117–122 (1984).

[3] A. Boyarsky, J. A. Harvey, and O. Ruchayskiy, *Annals Phys.* **301**, 1–21 (2002), `hep-th/0203154`.

[4] T. Appelquist, and J. Carazzone, *Phys. Rev.* **D11**, 2856 (1975).

[5] E. D'Hoker, and E. Farhi, *Nucl. Phys.* **B248**, 59 (1984); *ibid.* 77.

[6] N. Arkani-Hamed, S. Dimopoulos and G. R. Dvali, *Phys. Lett.* B **429** (1998) 263 [arXiv:hep-ph/9803315]; I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos and G. R. Dvali, *Phys. Lett.* B **436** (1998) 257 [arXiv:hep-ph/9804398].

[7] I. Antoniadis, S. Dimopoulos and G. Dvali, *Nucl. Phys.* B **516** (1998) 70; S. Ferrara, C. Kounnas and F. Zwirner, *Nucl. Phys.* B **429** (1994) 589.

[8] I. Antoniadis, K. Benakli, A. Laugier and T. Maillard, *Nucl. Phys.* B **662** (2003) 40 [arXiv:hep-ph/0211409].

[9] D. J. Kapner, T. S. Cook, E. G. Adelberger, J. H. Gundlach, B. R. Heckel, C. D. Hoyle and H. E. Swanson, Phys. Rev. Lett. **98** (2007) 021101.

[10] J. C. Long and J. C. Price, *Comptes Rendus Physique* **4** (2003) 337; R. S. Decca, D. Lopez, H. B. Chan, E. Fischbach, D. E. Krause and C. R. Jamell, *Phys. Rev. Lett.* **94** (2005) 240401; R. S. Decca et al., arXiv:0706.3283 [hep-ph]; S. J. Smullin, A. A. Geraci, D. M. Weld, J. Chiaverini, S. Holmes and A. Kapitulnik, arXiv:hep-ph/0508204; H. Abele, S. Haeßler and A. Westphal, in 271th WE-Heraeus-Seminar, Bad Honnef (2002).

[11] E. G. Adelberger, B. R. Heckel, S. Hoedl, C. D. Hoyle, D. J. Kapner and A. Upadhye, Phys. Rev. Lett. **98** (2007) 131104.

[12] N. Arkani-Hamed, S. Dimopoulos and G. Dvali, *Phys. Rev.* D **59** (1999) 086004.

[13] I. Antoniadis, E. Kiritsis and J. Rizos, *Nucl. Phys.* B **637** (2002) 92.

[14] J. L. Hewett, and T. G. Rizzo, *Phys. Rept.* **183**, 193 (1989).

[15] I. Antoniadis, E. Kiritsis, and T. N. Tomaras, *Phys. Lett.* **B486**, 186–193 (2000), `hep-ph/0004214`.

[16] L. E. Ibanez, F. Marchesano, and R. Rabadan, *JHEP* **11**, 002 (2001), `hep-th/0105155`.

[17] I. Antoniadis, E. Kiritsis, J. Rizos, and T. N. Tomaras, *Nucl. Phys. B* **660**, 81–115 (2003).

[18] P. Anastasopoulos, M. Bianchi, E. Dudas, and E. Kiritsis, *JHEP* **11**, 057 (2006a), `hep-th/0605225`; C. Coriano, N. Irges, and E. Kiritsis, *Nucl. Phys.* **B746**, 77–135 (2006), `hep-ph/0510332`; C. Coriano, M. Guzzi, and S. Morelli (2008a), `arXiv:0801.2949[hep-ph]`; C. Coriano, and M. Guzzi (2007), `0711.3424[hep-ph]`; `0709.2111`.

[19] I. Antoniadis, A. Boyarsky, and O. Ruchayskiy (2006), `hep-ph/0606306`.

[20] I. Antoniadis, A. Boyarsky, and O. Ruchayskiy, *Nucl. Phys.* **B793**, 246–259 (2008a), `0708.3001`.

[21] I. Antoniadis, A. Boyarsky, S. Espahbodi, O. Ruchayskiy, and J. D. Wells, *Nucl. Phys.* **B824**, 296–313 (2010), `0901.0639`.

[22] P. Anastasopoulos, et al., *Phys. Rev.* **D78**, 085014 (2008), `0804.1156`; J. A. Harvey, C. T. Hill, and R. J. Hill, *Phys. Rev.* **D77**, 085017 (2008), `0712.1230`; E. Dudas, Y. Mambrini, S. Pokorski, and A. Romagnoni (2009), `0904.1745`.

[23] C. Amsler, et al., *Phys. Lett.* **B667**, 1 (2008).

[24] J. M. Cornwall, D. N. Levin, and G. Tiktopoulos, *Phys. Rev.* **D10**, 1145 (1974).

[25] G. G. Raffelt, *Stars as laboratories for fundamental physics*, UofC Press, Chicago, USA, 1996.

[26] K. Zioutas, et al., *Phys. Rev. Lett.* **94**, 121301 (2005), `hep-ex/0411033`.

[27] S. Davidson, S. Hannestad, and G. Raffelt, *JHEP* **05**, 003 (2000a), `hep-ph/0001179`.

[28] T. Mitsui, et al., *Phys. Rev. Lett.* **70**, 2265–2268 (1993).

[29] A. Melchiorri, A. Polosa, and A. Strumia, *Phys. Lett.* **B650**, 416–420 (2007), `hep-ph/0703144`.

[30] S. Davidson, B. Campbell, and D. C. Bailey, *Phys. Rev.* **D43**, 2314–2321 (1991).

[31] E. Masso, and J. Redondo, *Phys. Rev. Lett.* **97**, 151802 (2006), `hep-ph/0606163`.

[32] A. Boyarsky, O. Ruchayskiy, and M. Shaposhnikov, *Phys. Lett.* **B626**, 184–194 (2005b).