Available online at www.sciencedirect.com

ScienceDirect

Nuclear Physics B Proceedings Supplement 00 (2014) 1–6

Nuclear Physics B Proceedings Supplement

Dilaton vs Higgs: Nearly Conformal Physics

G.A. Kozlov

Joint Institute for Nuclear Research,
Joliot Curie st., 6, Dubna, Moscow region, 141980 Russia

Abstract

We consider the model in which the conformal symmetry can be broken spontaneously, and a light scalar dilaton could emerge in the low-energy spectrum. The contribution of the dark photon production relevant to two photons decays of a Higgs boson/dilaton is discussed.

Keywords: Higgs, dilaton, conformal symmetry breaking, dark photon

1. Introduction

The interest to theories in which the Standard Model (SM) is strongly coupled to a conformal sector does not become weaker; on the contrary, it grows more and more every year (see, e.g., [1-4] and the references therein). In the classical gravity scheme at very high energies the scale invariance and its breaking is characterized by the scalar field, the dilaton $\bar{\sigma}$, which appears in the action

$$S = \int d^4x\,\sqrt{-\tilde{g}}\,\frac{1}{2}\left[\kappa r^2 + (\partial_\mu\bar{\sigma})^2 - \eta\bar{\sigma}^2\,r + ...\right]$$

with the scalar curvature r and positive dimensionless parameters κ and η. Once the conformal invariance is broken, the dilaton gets a vacuum expectation value (vev) f, and the last visible term in S yields the Einstein-Hilbert action at low energies. The spectrum of states would then contain the scalar resonance (the dilaton mode), associated with the pseudo-Goldstone boson due to the spontaneous breaking of conformal symmetry at the scale f. There is a class of models in which the breaking of conformal invariance at the scale $\Lambda_{CFT} = 4\pi f$ triggers electroweak symmetry breaking (EWSB) at the electroweak (EW) scale $\Lambda_{EW} = 4\pi v < \Lambda_{CFT}$, $v = 246$ GeV [2,3]. The dilaton itself exists in scale invariant world, and its production and decay in high energy experiments are model-dependent, however the dilaton can be seen as the Higgs-boson in the limit $f = v$. The mass of the dilaton is naturally light, $\sim \varepsilon \cdot f$, where the small parameter ε controls the deviations from the exact scale invariance. The dilaton becomes massless when the conformal invariance is recovered. One may have nearly conformal dynamics at a scale $\sim \Lambda_{EW}$ below which the scale symmetry is broken and one feeds into an EW sector.

We develop the theory in which the SM is coupled to scale invariant sector which is given by a dark matter (DM) in terms of the spin 1 dark photon (DP). We study the effects on the EW sector from the conformal sector, and the new bounds on DM physics are predicted. The coupling to the scalar dilaton sector is most important in our analysis.

The results of constraints on the dilaton mass m_σ and the conformal breaking scale f from the direct search at LEP and Tevatron have been carried out in [5]. In particular, based on the Tevatron data, there is a widely allowed range for a light dilaton below 200 GeV even for $f \sim O(1\,TeV)$.

The model is based on the extended group $SU_L(2) \times U(1)_Y \times U'_B(1)$, where the index B in $U'_B(1)$ is associated with an extra gauge boson B_μ that would be the DP $\gamma^\star$. The standard photon γ may oscillate into $\gamma^\star$ followed by the invisible decay to neutrino-antineutrino pair, $\gamma^\star \to \bar{\nu}\,\nu$, or by the electron-positron pair, $\gamma^\star \to e^-e^+$. As to the phenomenological realization we develop the model containing the hidden (dark) matter sector, which can be explored in collider physics experiments. The DP mixes with γ via the kinetic term $\sim \epsilon F_{\mu\nu} B^{\mu\nu}$ with ϵ being the kinetic mixing strength; $F_{\mu\nu}$ and $B_{\mu\nu}$ are the strength tensors of electromagnetic A_μ and DP B_μ fields, respectively. The coupling between A_μ and B_μ could be responsible for production of DP in the decay of the scalar particle $S \to \gamma\gamma^\star$, where S should either be the SM Higgs boson H or the dilaton σ.

No deviations with the expectations from the SM have been observed in two photon decays of the Higgs boson in the experiments ATLAS [6] and CMS [7] at the LHC. An enhanced rate of $S \to \gamma\gamma$, if observed, with respect to the SM prediction should be due to scale invariant breaking sector, where the contribution from the DP with the small mass $m_{\gamma^\star}$ to the branching ratio of the two photons decay $H \to \gamma\gamma$ of the Higgs-boson with the mass m_H would be significant

$$BR(S \to \gamma\gamma) \simeq (1 + a\,\epsilon^2\,\Omega) BR^{SM}(H \to \gamma\gamma), \quad (1)$$

where a is the positive constant, $\Omega \sim (1 - m^2_{\gamma^\star}/m^2_H)^b$, $b > 1$, BR^{SM} is the branching ratio of the decay $H \to \gamma\gamma$ due to SM calculations. The mass of the DP $m_{\gamma^\star}$ is unknown, however, the mixing factor ϵ in (1) is predicted in various models and to be in the range $10^{-12} - 10^{-2}$. On phenomenological grounds, however, DP masses in eV range are favored for the present work. If no excess events are found, the obtained results would give the bounds on ϵ as a function of $m_{\gamma^\star}$.

It is known that the characteristic feature of the gauge quantum field theory is to exhibit an infrared (IR) singularities which are incompatible with the positivity of the metric [8]. A relevant example is related to the dipole singularity. In particular, in Abelian Higgs model the breaking of the gauge symmetry requires such a type of singularity in two-point Wightman function (TPWF) of the dipole field, where, e.g., a scalar field σ satisfies the equation of 4th order $(\partial^2)^2\sigma(x) = 0$ [9].

In this paper DP is considered in the framework of the dipole field model [9-11], where the interacting dipole fields are local, relativistic quantum fields with a genuinely indefinite metric on the space of states generated from the vacuum. They converge asymptotically to free dipole fields. One of the crucial points of dipole quantum fields is that the massless dipole fields in four-dimensional space-time have a logarithmic increase at spatially separated arguments that features the confinement-like phenomenon [12].

The Higgs interpretation of a discovery at the LHC is not the only possibility. We suppose that the dilaton $\bar{\sigma} = f\sigma$ might be considered as the possible candidate for the scalar particle of the mass ~ 126 GeV observed at the LHC. It is supported by a) the pseudo-Goldstone nature of $\bar{\sigma}$-field with finite mass due to spontaneous breaking of the scale symmetry; b) the dilaton preserves a non-linear realization of the scale symmetry; c) the dilaton serves as a portal to the hidden sector, and it guarantees the renormalizability of the theory.

2. Couplings and constraints

The couplings of a dilaton to two gluons are crucial for collider phenomenology. These couplings can be significantly enhanced under very mild assumption about high scale physics [3]. At energies below $4\pi f$ the effective dilaton couplings to massless gauge bosons are provided by the SM quarks lighter than the dilaton: $\sigma[c_{EM}(F_{\mu\nu})^2 + c_s(G^a_{\mu\nu})^2]/(8\pi f)$. Here $G^a_{\mu\nu}$ is the gluon field strength tensor; $c_{EM} = -\alpha \cdot 17/9$ if $m_W < m_\sigma < m_t$, $c_{EM} = -\alpha \cdot 11/3$ if $m_\sigma > m_t$; $c_s = \alpha_s \cdot (11 - 2\,n_{light}/3)$; n_{light} is the number of quarks lighter than the dilaton; α and α_s are EM and strong coupling constants, respectively; m_W and m_t are masses of W-boson and the top-quark, respectively. The second term in the effective coupling above mentioned indicates a $(33/2 - n_{light})$-factor increase of the coupling strength compared to that of the SM Higgs boson. The upper limit of f is estimated on the level of O (6 TeV) for the light dilaton [13].

In general, the ultraviolet (UV) coupling of an operator O_{UV} of dimension d_{UV} to a SM operator O_{SM} of dimension d_{SM} at the UV scale M (UV messenger) is

$$\frac{1}{M^{d_{SM}-4}}\,O_{SM}\,\frac{1}{M^{d_{UV}}}\,O_{UV}. \quad (2)$$

No masses are allowed in the Lagrangian of the effective theory containing (2). All masses can be generated dynamically in IR. The hidden sector is formed when the dilaton field $\bar{\sigma} = f\sigma$ is coupled to a $U(1)$ gauge theory. Consider the coupling of $\bar{\sigma}(x)$ to DM sector

$$\frac{1}{M^{d_{UV}-2}}\,|\bar{\sigma}|^2\,O_{UV}, \quad (3)$$

which flows to coupling of the Higgs boson H to DM operator O_{IR} of dimension d_{IR} in IR

$$const\,\frac{\Lambda^{d_{UV}-d_{IR}}}{M^{d_{UV}-2}}\,|H|^2\,O_{IR}, \quad (4)$$

when the scale invariance is almost breaking. Here, Λ is the strong coupling scale. Once $H(x)$ acquires v, theory becomes nonconformal below the scale $\tilde{\Lambda}$, where

$$\tilde{\Lambda}^{4-d_{IR}} = \frac{\Lambda^{d_{UV}-d_{IR}}}{M^{d_{UV}-2}}\,v^2. \quad (5)$$

Below $\tilde{\Lambda}$ the DM sector becomes a standard particle sector. For a typical energy $Q = \sqrt{s}$ of a collider experiment $\tilde{\Lambda} < \sqrt{s} < \Lambda$, which leads to the constraint for the energy

$$s^{2-d_{IR}/2} > \left(\frac{\Lambda}{M}\right)^{d_{UV}-d_{IR}} M^{2-d_{IR}}\, v^2. \qquad (6)$$

Based on the operator form (2) the mixing strength ϵ in the kinetic term $\epsilon F_{\mu\nu} B^{\mu\nu}$ (as an observable) is

$$\epsilon = \left(\frac{\sqrt{s}}{M}\right)^{2(d_{SM}-4)} \left(\frac{\sqrt{s}}{\Lambda}\right)^{2d_{IR}} \left(\frac{\Lambda}{M}\right)^{2d_{UV}}. \qquad (7)$$

Then the effect of DM sector on observables has no the dependence on d_{UV}, d_{IR} and Λ, and is bounded by

$$\epsilon < \frac{s^{d_{SM}}}{(v^2\, M^{d_{SM}-2})^2}. \qquad (8)$$

It is clear from (7) that the signals of new physics containing DM increase with energy, the dimension of the SM operator d_{SM}, and would be seen if the parameter M of heavy messenger is not too high. Assuming that the deviation from the SM is detected at the level of order 3%, the DM would be visible at the LHC ($\sqrt{s}$ ∼ O(10 TeV)) as long as $M <$ 1000 TeV if $d_{SM} = 4$.

In the decay $S \to \gamma\gamma^\star$ the gauge invariant operator structure is $O_{SM}\, O_{IR} \sim \epsilon\bar{\psi}\, \gamma_\mu\, \psi\, H\, B^\mu\, M^{-1}$, where ψ is the operator of the quark (in the loop) with the mass m, and the relevant energy scale is $Q \sim m$. Since the SM is operated on the scale $\sim O$ (1 TeV), we expect $\epsilon < 10^{-5}$ in the case of the top-quark with $d_{SM} = 4$, or $\epsilon < 6\cdot 10^{-2}$ if the fourth generation quarks with the mass ∼ O(1 TeV) give the contribution to decay amplitude. At small energies s, the mixing ϵ is almost zero, and the only SM Higgs boson decay into two photons would be appropriate.

The result for upper limit (8) is shown in Fig.1 for $\sqrt{s}$ = 8-14 TeV and M = 800-1000 TeV at $d_{SM} = 4$. The LHC is the very good facility where the DM physics can be tested well.

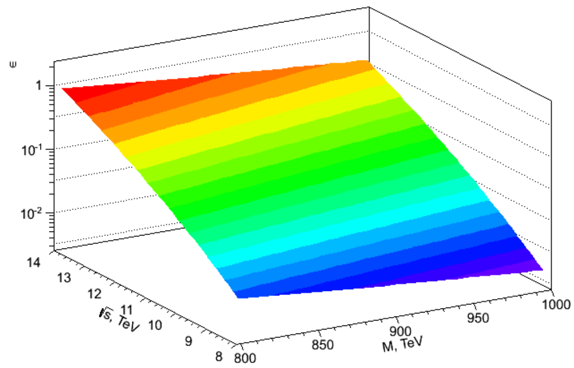

Fig. 1: *Upper limit of ϵ (8) as a function of $\sqrt{s}$ and M.*

where $c_i(\mu)$ is running coupling, the operator $O_i(x)$ has the scaling dimension d_i. Under the scale transformations $x^\mu \to e^\omega x^\mu$, one has $O_i(x) \to e^{\omega d_i}\, O_i(e^\omega x)$, $\mu \to e^{-\omega}\mu$. This gives for the dilatation current $S^\mu = T^{\mu\nu} x_\nu$

$$\partial_\mu S^\mu = T^\mu_\mu = \sum_i \left[c_i(\mu)(d_i - 4) O_i(x) + \beta_i(c) \frac{\partial}{\partial c_i} L \right],$$

where $T^{\mu\nu}$ and $\beta_i(c)$ are the energy-momentum tensor and the running β-function, respectively. At energies below f one has $c_i(\mu) \to (\sigma)^{4-d_i}\, c_i\,(\mu\,\sigma)$. The theory would be nearly scale invariant if $d_i = 4$ and $\beta(c) \to 0$. Thus, the breaking of chiral symmetry is triggered by the dynamics of nearly conformal sector.

The model is formulated in terms of the Lagrangian which features: the scalar CP even dilaton field $\sigma(x)$ as the local one arising as a generic pseudo-Goldstone boson from the breaking of conformal strong dynamics and from which the vector potential $A_\mu(x)$ is derived, the conformal field given by the operator O^μ_U and a set of the SM fields. The conformal invariance can be broken by the couplings with non-zero mass dimension effects. The Lagrangian density with a small explicit breaking of the conformal symmetry is $L = L_1 + L_2$, where

$$L_1 = -B\partial_\mu A^\mu + \frac{1}{2\xi} B^2 - \frac{1}{\Lambda_U^{d-3}} (A_\mu - \partial_\mu \sigma) O^\mu_U$$

$$+ \bar{\psi}(i\hat{\partial} - m + g\hat{A})\psi - \sigma \sum_\psi \left(m + \epsilon y_\psi v\right) \bar{\psi}\psi,$$

$$L_2 = \frac{1}{\Lambda_U^{d-1}} \sum_q \bar{\psi}(c_v\, \gamma^\mu - a_v\, \gamma^\mu \gamma_5)\psi\, O_{U_\mu}$$

$$+ \frac{1}{\Lambda_U^{d+1}} W^a_{\mu\alpha} W^{a\mu}_\beta \left(\partial^\alpha O^\beta_U + \partial^\beta O^\alpha_U\right).$$

3. Model

We start with partition function

$$Z = \int \mathbf{D}\bar{\sigma}_i\, \exp\left[-\int_0^\infty d\tau \int d^3x L(\tau, \vec{x})\right],$$

$$L(x) = \sum_i c_i(\mu)\, O_i(x),$$

The field B plays the role of the gauge-fixing Lagrangian multiplier, and it remains free. We assume that the real positive parameter $\xi \neq \infty$, otherwise the model becomes trivial. The scaling dimension d may appear as a non-integer number d of invisible particles [14]. The vector operator O^{μ}_{U} describes a scale-invariant hidden sector that possesses IR fixed point at a high scale Λ_U, presumably above the EW scale; c_v and a_v are vector and axial-vector couplings, respectively. A dilaton acquires a mass and its couplings to quarks can undergo variations from the standard form. In particular, since the scale symmetry is violated by operators involving quarks, the shifts in the dilaton Yukawa couplings to quarks can appear. This is given by $\varepsilon = m^2_{\sigma}/f^2$ which parametrizes the size of the deviation from exact scale invariance through m_{σ} as the measure of this deviation. The nine additional contributions to Yukawa couplings y_{ψ} are taken into account (y_{ψ} are 3×3 diagonal matrices in the flavor space); g is the coupling constant.

In the case of a Higgs-boson decays into two photons, the W-bosons contribute more significantly. In the model considered here, the only SM quarks contribution is dominated, because the W-boson loop contribution is suppressed by two more powers of Λ_U, and due to significantly large value of Λ_U one can ignore it.

The equations of motion are ($\nabla \equiv \partial/\partial x_{\mu}$)

$$\partial_{\mu}\sigma \simeq A_{\mu} - \frac{1}{\Lambda^2_U}\bar{\psi}(c_v\,\gamma_{\mu} - a_v\,\gamma_{\mu}\gamma_5)\psi, \quad \partial_{\mu} A^{\mu} = \xi^{-1} B,$$

$$\partial_{\mu} B = -J_{\mu} + \frac{1}{\Lambda^{d-3}_U} O_{U_{\mu}}, \quad J_{\mu} = g\,\bar{\psi}\,\gamma_{\mu}\,\psi,$$

$$\frac{1}{\Lambda^{d-3}_U}\,\partial_{\mu} O^{\mu}_U + \frac{1}{f}(m + \varepsilon\, y_{\psi}\, v)\bar{\psi}\,\psi = 0,$$

$$\left[i\,\hat{\partial} - m\left(1 + \frac{\sigma}{f}\right) + g\,\hat{A} - \frac{\sigma}{f}\,\varepsilon\, y_{\psi}\, v\right]\psi$$

$$+\left[\frac{1}{\Lambda^{d-1}_U}\, O^{\mu}_U (c_v\,\gamma_{\mu} - a_v\,\gamma_{\mu}\gamma_5)\right]\psi = 0.$$

In the nearly conformal sector (NCS) supported by the weakly changing operator O^{μ}_U in the space-time and the conservation of the current J_{μ}, the $\sigma(x)$-field looks like the dipole field obeying the equation of the 4th order

$$\lim_{m_{\sigma}\to 0}\left(\nabla^2 + m^2_{\sigma}\right)^2 \sigma(x) \simeq 0 \tag{9}$$

and the canonical commutation relation [15]

$$[\sigma(x), \sigma(y)] = \frac{1}{\xi}\int \frac{d^4 p}{(2\,\pi)^3}\, sgn(p^0)\,\delta'(p^2)\, e^{-i\,p\,(x-y)}$$

$$= \frac{1}{8\,\pi\, i\,\xi}\, sgn(z^0)\,\theta(z^2), \quad z = x - y,$$

where $sgn(p^0)\,\delta'(p^2)$ is well-defined as the odd homogeneous generalized function from the space $S'(\Re_4)$ of the temperate distributions on $\Re_4$.

4. The propagator and the effective potential

The propagator of the $\sigma(x)$-field in NCS in terms of TPWF $W(z)$ has the form [10]

$$W(z) = \langle \Omega, \sigma(x)\,\sigma(y)\,\Omega\rangle = \frac{1}{i\,\xi}\, E^{-}(z), \tag{10}$$

where $E^{-}(x)$ is the negative-frequency part of the generalized function $E(x) = E^{+}(x) + E^{-}(x) = (8\,\pi)^{-1}\,\theta(x^2)\, sgn(x^0)$, which is the only distribution among the solutions of the equation $\left(\nabla^2\right)^2 W(x) = 0$ obeying locality, Poincare covariance and the spectral conditions. However the metric is not positive defined. The vacuum Ω-vector satisfies the following conditions: $\sigma^{-}(x)|\Omega\rangle = 0$, $\langle\Omega, \Omega\rangle = 1$, where $[\sigma^{-}(x)]^{*} = \sigma^{+}(x)$ in the decomposition $\sigma(x) = \sigma^{-}(x) + \sigma^{+}(x)$. The solutions of (9) can be classified by TPWF's, where

$$E^{-}(x) = \int \frac{d^4 p}{(2\pi)^3}\,\theta(p^0)\,\delta'(p^2)\, e^{-ipx}$$

$$= \frac{-1}{(4\pi)^2}\ln(-\mu^2 x^2 + i\,\epsilon\, x^0),$$

where μ is the positive constant required for dimensioneless reason. The product of the generalized functions $\theta(p^0)\,\delta'(p^2)$ is defined uniquely only with the basic functions $u(p)$ which are equal to zero at $p = 0$. To separate the IR μ-dependence, $E^{-}(x)$ can also be given in the form $(-4\,\pi)^{-2}\left\{\ln|\mu^2\, x^2| + i\,\pi\, sgn(x^0)\,\theta(x^2)\right\}$.

The Fourier transform of $E^{-}(x)$ is $\tilde{E}^{-}(p) = 2\,\pi\,\theta(p^0)\,\delta^{(1)}(p^2, \tilde{\mu}^2)$, where $\tilde{\mu} = 2\, e^{-\gamma + 1/2}\,\mu$, $\gamma = -\Gamma'(1)$ beying the Euler's constant. The functional $\delta^{(1)}(p^2, \tilde{\mu}^2)$ is defined on the space $S(\Re_4)$ of the complex Schwartz test functions on $\Re_4$ as

$$\delta^{(1)}(p^2, \tilde{\mu}^2) = \frac{1}{16}\left(\frac{\partial^2}{\partial p^2}\right)^2\left[\theta(p^2)\,\ln\frac{p^2}{\tilde{\mu}^2}\right].$$

One can verify that $p^2\,\delta^{(1)}(p^2, \tilde{\mu}^2) = -\delta(p^2)$, $(p^2)^2\,\delta^{(1)}(p^2, \tilde{\mu}^2) = 0$. The presence of the parameter μ in $E^{-}(x)$ breaks its covariance under dilatation transformations $x_{\mu} \to \lambda\, x_{\mu}$ ($\lambda > 0$) and implies spontaneously symmetry breaking of the dilatation invariance of (9). This is one of the reasons for the special role of the dilaton field $\sigma(x)$ in what follows.

The Green's function in $\Re^4$ space-time is given by

$$G(z) = \langle\Omega, T[\sigma(x)\,\sigma(y)]\Omega\rangle = \frac{1}{i\,\xi}\,E_c(z),$$

where the causal function

$$E_c(x) = \theta(x^0)\,E^-(x) + \theta(-x^0)\,E^+(x)$$

$$= \frac{1}{i\,(4\,\pi)^2}\,\ln(-\mu^2\,x^2 + i\,\epsilon)$$

satisfies the following equations

$$\nabla^2 E_c(x) = \frac{i}{4\,\pi^2}\,\frac{1}{-x_\mu^2 + i\,\epsilon}, \quad (\nabla^2)^2 E_c(x) = \delta^4(x).$$

In $\Re_4$-momentum space the propagator of the dilaton field is given in terms of distributions

$$\tilde{G}(p) = \frac{-1}{(4\,\pi)^2\,\xi}\int d^4x\,e^{i\,p\,x}\,\ln(-\mu^2\,x^2 + i\,\epsilon).$$

Following [10] one can calculate $\tilde{G}(p)$ through the second order derivative $\tilde{G}(p) = (\partial^2/\partial p^2)H(p)$, where

$$H(p) = \frac{-1}{(4\,\pi)^2\,\xi}\int d^4x\,e^{i\,p\,x}\,\frac{\ln(-\mu^2\,x^2 + i\,\epsilon)}{-x_\mu^2 + i\epsilon}. \quad (11)$$

Finally, the result for (11) is

$$H(p) = \frac{i}{4\,\xi}\,\frac{\ln\left[e^{2\gamma}(-p^2 - i\,\epsilon)/(4\,\mu^2)\right]}{p^2 + i\,\epsilon},$$

which leads to

$$\tilde{G}(p) = \frac{1}{2\,i\,\xi}\,\frac{\partial}{\partial p^\mu}\left\{\frac{p^\mu\left[\ln\left(-p^2/\tilde{\mu}^2 - i\,\epsilon\right)\right]}{(p^2 + i\,\epsilon)^2}\right\}. \quad (12)$$

The following equation is strightforward: $(-p^2)^2\,i\,\xi\,\tilde{G}(p^2) = 1$. In the propagator (12) one finds an appearance of the parameter $\tilde{\mu}$ with the dimension of mass which otherwise would appear in the theory as a renormalization mass, and which distinguishes our model from the standard EW theory as conventionally formulated. The differentiation over p_μ in (12) with $\partial/\partial p^\mu$ being the weak derivative has to be understood in the sense of distribution where for any test function $u(p)$ we have

$$\int \tilde{G}(p)u(p)d^4p = \frac{i}{2\xi}\int d^4p\frac{\ln(-p^2/\tilde{\mu}^2 - i\epsilon)}{(p^2 + i\epsilon)^2}p^\mu\frac{\partial}{\partial p^\mu}u(p)$$

and the extra power of momentum p^μ explicitly eliminates the IR divergence at $p = 0$.

The lowest order energy (potential) of a static "charge" is given by the static Fourier transform ($|\vec{x}| \equiv r$)

$$E(r) \sim \int \frac{d^3\vec{p}}{(2\,\pi)^3}\,e^{i\vec{p}\vec{x}}\,D(p^0 = 0, \vec{p}; \tilde{M})$$

with $D(p, \tilde{M}) = \tilde{M}^2\,\tilde{G}(p)$, $\tilde{M}$ has the dimension one in mass units. Using the propagator $\tilde{G}(p)$ (12) in the form

$$\tilde{G}(p) = \frac{1}{4\,\xi\,i}\,\frac{\partial^2}{\partial p^2}\left\{\frac{\ln\left[e^{2\,\gamma}\,(-p^2/\tilde{\mu}^2 - i\,\epsilon)\right]}{-p^2 - i\,\epsilon}\right\}$$

one can find

$$E(r) \sim \frac{\tilde{M}^2}{8\,\pi\,\xi}\,r\,\left[a + b\ln(\tilde{\mu}r)\right],$$

where a and b are known constants. The energy of a dilaton in NCS is linearly rising as r. The result is stable both at short and large distances in any finite order of perturbation theory.

At small distances the dominant interactions between a heavy quark and antiquark with the mass m_q can be described by the effective potential in $SU(3)$

$$V_{eff}(r) \simeq -\frac{4}{3r}\alpha_s(m_q) - \frac{\lambda(m_q, \eta_{\sigma q})}{r}\,\exp(-m_\sigma r)$$

with

$$\lambda(m_q, \eta_{\sigma q}) = \frac{m_q^2}{4\,\pi\,f^2}\,\eta_{\sigma q}^2,$$

where $\eta_{\sigma q} = 1$ in SM, otherwise $\eta_{\sigma q} > 1$. The lower bound for m_q is

$$m_q > \frac{f}{\eta_{\sigma q}}\left(\frac{16}{3}\,\pi\,\alpha_s\right)^{1/2}$$

which can exceed the top quark mass even for $f \simeq v$. Thus, the mediator field in a (super)heavy quark sector is evident. Once the new force is discovered, there will be the first ever seen not related to a gauge symmetry.

The energy distribution of the emitted photon (γ) in the decay of the dilaton into γ and the DP in terms of the vector unparticles stuff U is

$$\frac{d\Gamma(\sigma \to \gamma U)}{dE_\gamma} = \frac{A_d}{(2\,\pi)^2}\,m_\sigma\,E_\gamma^3\left(P_U^2\right)^{d-2}\,|A(x_q, y_q)|^2,$$

where [14]

$$A_d = \frac{16\,\pi^{5/2}}{(2\,\pi)^{2d}}\,\frac{\Gamma(d + 1/2)}{\Gamma(d - 1)\,\Gamma(2d)},$$

$A(x_q, y_q)$ is the decay amplitude (see [12] for details), P_U is the momentum of DP in terms of unparticles..

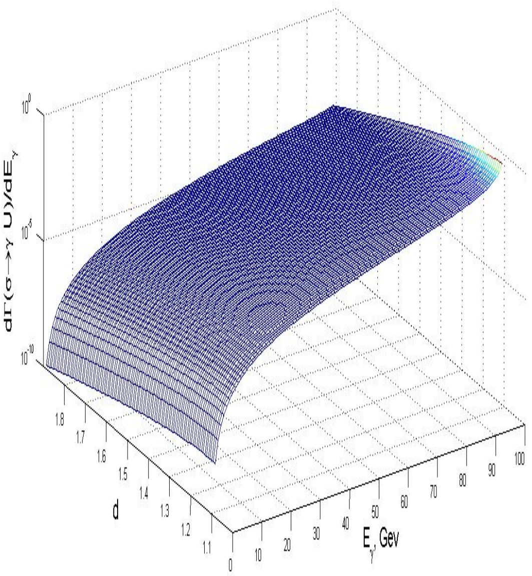

Fig. 2: *Energy spectrum of the photon in decay $\sigma \to \gamma U$ for various values of d.*

In Fig. 2, we show the energy spectrum of the emitted photon in decay $\sigma \to \gamma U$ for various values of d with the dilaton mass $m_\sigma \sim 0.2$ TeV, $c_v = 1$, $\Lambda_U = 1$ TeV, $f \simeq v$. The only top quarks in the loop are included for the calculations because of the negligible contributions from lighter quarks.

5. Conclusion

Once the scale symmetry is broken spontaneously, the dilaton could emerge in the low-energy spectrum. For a certain relation between couplings in NCS the field solutions are dictated by 4th order differential equation (9).

We have studied the decay of a dilaton into a real photon and the DP the nature of which is maintained through the hidden sector including the effects of scale invariance breaking. This mode is very useful to probe the nearly conformal sector containing the dilaton and the DM U- vector state having the continuous (small) mass. Unless the LHC can collect a very large samples of σ, the detection of DP through $\sigma \to \gamma U$ would be quite challenging, because the couplings of the dilaton are similar to those of the SM Higgs-boson, and the dilaton, if observed, could open the window to the conformal pattern of the strong sector. In this case the scale invariant sector is close to EW sector, that could provide the decay $\sigma \to \gamma\gamma$ to be compared with the LHC data in searching for new light scalar object with the mass close to 126 GeV.

The energy of the real photon $E_\gamma = (m_\sigma^2 - P_U^2)/(2\, m_\sigma)$ contains the information about the missing energy or the momentum P_U carried out by the DP. A nontrivial scale invariant sector of dimension d may give rise to peculiar missing energy distributions in $\sigma \to \gamma\, U$ that can be treated in the experiment. In particular, this energy distribution can discriminate d and estimate Λ_U. When combined with $gg \to \sigma$, the decay $\sigma \to \gamma\, U$ provides especially valuable information regarding possible loop contributions from new particles lighter than a dilaton.

The decay mode $\sigma \to \gamma\, U$ is particular useful for a light σ-dilaton ($m_\sigma \sim 100-200$ GeV) since it may have a nearly comparable decay rate with the $\gamma\gamma$ discovery mode of the Higgs-boson decay. Our result implies that the decay mode $\sigma \to \gamma\, U$ is near a border of the conformal invariance breaking.

The low rate $\Gamma(\sigma \to \gamma\, U)/\Gamma_{Higgs-total}$ is compensated by the enhancement of the order $(33/2 - n_{light})^2 \sim O(100)$ of the gluon fusion production cross-section compared to that of the SM Higgs-boson.

In addition to discovery, the implementation of our prediction in the LHC analyses should be straightforward and lead to more precise determination or limits of the dilaton and the DP couplings to gauge bosons, top quarks and the quarks of 4th generation.

References

[1] E. Gildener and S. Weinberg, Phys. Rev. D13 (1976) 3333.
[2] P. Fox, A. Rajaraman and Yu. Shirman, Phys. Rev. D76 (2007) 075004.
[3] W.D. Goldberger, B. Grinstein, and W. Skiba, Phys. Rev. Lett. 100 (2008) 111802.
[4] M. Holthausen, J. Kubo, K.S. Lim and M. Lindner, JHEP 12 (2013) 076.
[5] Y. Bai, M. Carena, and J. Lykken, Phys. Rev. Lett. 103 (2009) 261803.
[6] S. Laplace, Higgs to two photons with the ATLAS, Talk at ICHEP2014, July 3, 2014, Valencia, Spain.
[7] The CMS Collaboration, Observation of the diphoton decay of the Higgs boson and measurement of its properties, arXiv:1407.0558 [hep-ex].
[8] F. Strocchi, Phys. Rev. D17 (1978) 2010.
[9] R. Ferrari, IL Nuovo Cim. 19A (1974) 204.
[10] D. Zwanziger, Phys. Rev. D17 (1978) 457.
[11] E. d'Emilio and M. Mintchev, Phys. Lett. B89 (1980) 207.
[12] G. Kozlov, PoS (Confinement X) 086.
[13] G.A. Kozlov, I.N. Gorbunov, Int. J. Mod. Phys. A26 (2011) 3987.
[14] H. Georgi, Phys. Rev. Lett. 98 (2007) 221601; Phys. Lett. B650 (2007) 275.
[15] N.N. Bogolyubov, A.A. Logunov, A.I. Oksak, I.T. Todorov, General principles of quantum field theory, Springer (1989).