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AXIONS AND DILATONS: THE SEARCH FOR
VERY LIGHT SCALAR PARTICLES*

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ABSTRACT

Axions, dilatons and other very light scalar particles are possible low energy manifestations of new interactions beyond the Standard Model. We discuss the connection between axion (a) and dilaton (σ) in supersymmetric theories and derive a prediction for their mass ratio, $m_a/m_\sigma \approx 10^{-4} - 10^{-5}$. We also review three recent proposals to search for light scalar particles in Bragg scattering, in a laser experiment and via the Mößbauer effect.

1. Light scalar particles and large mass scales

The success of the Standard Model as theory of strong and electroweak interactions [1] has reached an almost worrisome extent. Unless new, unexpected phenomena are found at the current generation of accelerators the conclusion seems unavoidable that the mass scale f of new interactions beyond the Standard Model is much larger than the Fermi scale of weak interactions, i.e., $f \gg G_F^{-1/2} \sim 300$ GeV.

How can such a new mass scale manifest itself at currently accessible energies ? One possibility are new "superweak" interactions which are suppressed by inverse powers of f. Their form is restricted by requiring invariance under the Standard Model gauge group. To order 1/f there is only one interaction, which yields Majorana mass terms for neutrinos:

$$L^{(5)}_{eff} = \frac{1}{f}\,(l_L\varphi)^T C(l_L\varphi) = \frac{v^2}{f}\,\nu_L^T C\nu_L + \dots \,, \qquad l = \begin{pmatrix}\nu\\ e^-\end{pmatrix} \;; \tag{1}$$

here v = 174 GeV is the vacuum expectation value of the Higgs doublet φ. To order $1/f^2$ there are many interaction terms, in particular four-fermion operators with various flavour and Lorentz structures,

$$L^{(6)}_{eff} = \frac{1}{f^2}\,C^{AB}_{ijkl}\;\bar\psi_i\Gamma^A\psi_j\bar\psi_k\Gamma^B\psi_l \quad , \tag{2}$$

which give rise to deviations from weak interactions, flavour changing neutral currents, proton decay etc. The weakest lower bound on f stems from quark-lepton universality of the charged current weak interactions, which yields f > 5 TeV [2]. Other processes, such as flavour changing neutral currents or CP violating operators yield much larger bounds on f which almost reach the Planck mass $m_{PL} \sim 10^{18}$ GeV if baryon number violating interactions are allowed.

Another possibility is the existence of pseudo-Goldstone bosons, i.e., new, very light scalar particles whose mass is of order 1/f, e.g.,

$$m \sim \alpha_{S.M.}\frac{\Lambda^2_{S.M.}}{f} < \begin{cases} 100\ \text{MeV},\ \Lambda_{S.M.} \sim G_f^{-1/2} \sim 300\ \text{GeV} \\ 1\ \text{keV},\ \Lambda_{S.M.} \sim \Lambda_{QCD} \sim 300\ \text{MeV}, \end{cases} \tag{3}$$
$$\text{for } f > 10\ \text{TeV}.$$

Such light scalar particles have to exist if the mass scale f is associated with the spontaneous breaking of some global symmetry of the more fundamental theory which contains the Standard Model as its low energy limit. As in the case of the pion, the only known pseudo-Goldstone boson, these (pseudo)scalar particles are expected to have a 2-photon coupling, which may allow their experimental detection.

Many candidates [3] for light pseudo-Goldstone bosons have been advocated by theorists. From the point of view of the Standard Model, the axion [4] has the best motivation, since it is an unavoidable consequence of the Peccei-Quinn mechanism which has survived as the only viable solution of the strong CP-problem. Also interesting is the dilaton, a Brans-Dicke type scalar [5], which arises in theories with spontaneously

*Lecture given at the Workshop on Higgs Particles, Erice, July 1989 broken scale invariance [6] and which has recently been discussed in connection with the cosmological constant problem [7-13]. As we will see in the following section axion and dilaton occur together in supersymmetric theories.

2. Dilatons and axions: symmetries, couplings and masses

One of the most remarkable features of the Standard Model is that all particle masses are generated through spontaneous symmetry breaking. As a consequence the lagrangian

$$L = L_{gauge} + L_{scalar} + L_{fermion} , \tag{4a}$$

where

$$L_{gauge} = -\frac{1}{4} G^A_{\mu\nu} G^{A\mu\nu} - \frac{1}{4} W^I_{\mu\nu} W^{I\mu\nu} - \frac{1}{4} B_{\mu\nu} B^{\mu\nu}, \tag{4b}$$

$$L_{scalar} = -(D_\mu \varphi)^\dagger (D^\mu \varphi) - V_o(\varphi^\dagger \varphi) , \tag{4c}$$

$$V_o = \frac{\lambda}{2} (\varphi^\dagger \varphi + \frac{\mu^2}{2\lambda})^2 , \tag{4d}$$

contains only one parameter of dimension mass, μ, which sets the Fermi scale of weak interactions through the vacuum expectation value of the Higgs doublet φ. D_μ, $G^A_{\mu\nu}$, $W^I_{\mu\nu}$ and $B_{\mu\nu}$ are the gauge covariant derivative and the SU(3), SU(2) and U(1) field strengths respectively.

What is the origin of μ ? If the Fermi scale itself is dynamically generated through the vacuum expectation value f of some field Ψ, the Standard Model is part of a more fundamental theory whose effective lagrangian at distances larger than 1/f is obtained from eq. (4) by substituting

$$\mu \to \mu\, e^{\sigma/f} = g\, \Psi , \qquad \text{with}$$

$$\Psi = f\, e^{\sigma/f}, \quad g = \frac{\mu}{f} . \tag{5}$$

Here the dilaton field σ corresponds to the fluctuations of Ψ around the vacuum expectation value f.

The new effective action is well known to be invariant under

dilatations [14]:

$$\delta\sigma = f + x^\mu \partial_\mu \sigma, \qquad \delta\varphi = (1 + x^\mu \partial_\mu)\, \varphi, \ldots \tag{6}$$

This symmetry of the action, $\delta\Gamma = 0$, leads to a conserved dilatation current:

$$\partial^\mu D_\mu = 0 . \tag{7}$$

Associated with this invariance under scale and conformal transformations is the existence of a flat direction of the scalar potential which, as we shall see, is lifted by quantum corrections which also generate a small mass for the dilaton. The kinetic term for the dilaton can be choosen as

$$e^{2\sigma/f} (\partial_\mu \sigma)^2 \quad \text{or} \quad (\partial_\mu \sigma)^2 , \tag{8}$$

depending on whether scale invariance is treated as fundamental symmetry of the theory or just as accidental symmetry of a certain part of the lagrangian.

It is interesting that in this latter case, if gravity is encluded by adding the Einstein-Hilbert action, one obtains precisely the Brans-Dicke theory of gravity with the ordinary Standard Model as matter sector [12]. In order to see this let us use the field variables σ and $\phi = e^{-\sigma/f}\varphi$ in terms of which the lagrangian (4) reads:

$$L_{S.M.}(\hat{g}_{\mu\nu}, \phi, \ldots) = -\sqrt{\hat{g}} \left[\hat{g}^{\mu\lambda} \hat{g}^{\nu\tau} \frac{1}{4} G^A_{\mu\nu} G^A_{\lambda\tau} + \ldots + \hat{g}^{\mu\nu} (D_\mu \phi)^\dagger (D_\nu \phi) + \frac{\lambda}{2} (\phi^\dagger \phi + \frac{\mu}{2\lambda})^2 \right] , \tag{9}$$

where

$$\hat{g}_{\mu\nu} = \eta_{\mu\nu}\, e^{2\sigma/f}, \qquad \hat{g}^{\mu\nu} \hat{g}_{\nu\lambda} = \delta^\mu_{\ \lambda} . \tag{10}$$

This is the Standard Model lagrangian in a conformally flat background metric, where the conformal factor is given by the dilaton field. Let us now add the gravitational field in the standard manner, i.e.,

$$L = -\sqrt{g} \left(\frac{\kappa}{2} R + \frac{1}{2} g^{\mu\nu} \partial_\mu \sigma \partial_\nu \sigma \right) + L_{S.M.}(g_{\mu\nu} e^{2\sigma/f}, \phi, \ldots), \tag{11}$$

$$\kappa = 1/8\pi G_N ,$$

and perform a Weyl transformation

$$\bar{g}_{\mu\nu} = g_{\mu\nu}\, e^{2\sigma/f} \;, \qquad (12)$$

$$R = e^{2\sigma/f} \left[\bar{R} - 6 \frac{1}{\sqrt{\bar{g}}}\, \partial_\mu \left(\bar{g}^{\mu\nu} \sqrt{\bar{g}}\, \partial_\nu \frac{\sigma}{f} \right) + 6\, \bar{g}^{\mu\nu} \partial_\mu \frac{\sigma}{f}\, \partial_\nu \frac{\sigma}{f} \right] . \qquad (13)$$

We then obtain

$$L = - \sqrt{\bar{g}} \left\{ \chi \bar{R} + \frac{\omega}{\chi}\, \bar{g}^{\mu\nu} \partial_\mu \chi\, \partial_\nu \chi - L_{S.M.}(\bar{g}_{\mu\nu}, \phi, \dots) \right\} , \qquad (14)$$

$$\chi = \frac{\kappa}{2}\, e^{-\sigma/f} \;, \quad \omega = \frac{f^2 - 6\kappa}{4\kappa} \;,$$

which is precisely the Brans-Dicke lagrangian. We note that the strong experimental bounds on the parameter ω only apply if the dilaton mass vanishes, which is usually assumed, but which is not the case for the theory defined by the lagrangian (9), where in general quantum corrections lead to a nonvanishing mass.

Let us briefly digress and discuss the possible implication of the dilaton field for the cosmological constant problem. The field equations of the Brans-Dicke theory (14) read:

$$R = - \frac{1}{\kappa} \left(g^{\mu\nu} \partial_\mu \sigma\, \partial_\nu \sigma \right) + T^\mu{}_\mu \;, \qquad (15a)$$

$$f\, \sigma = T^\mu{}_\mu \;, \qquad (15b)$$

and imply that the vacuum state, which is characterized by constant fields $\sigma = \sigma_o$ = const. and $\varphi = \varphi_o$= const., has vanishing curvature, i.e., $R_o = 0$ [12]. Unfortunately, this cannot be regarded as solution of the cosmological constant problem, since the existence of the constant solution of eq. (15b) requires fine-tuning of a constant in the scalar potential. Non-linearly realized scale invariance is certainly not sufficient [9,11,12] to obtain flat space-time as uniquely determined ground state, but the presence of the dilaton field does have some intriguing aspects which may turn out to be essential for understanding

the observed flatness of space-time. We note that the Brans-Dicke theory emerges from one particular definition of "dilatons" in curved space-time. Other options are discussed in refs. [9] and [12].

Classical scale invariance is not preserved in the full quantum theory where one has:

$$\delta\Gamma = M \frac{\partial}{\partial M} \Gamma = \int d^4x\, A_D + O(\hbar^2) \;, \qquad (16a)$$

$$\partial^\mu D_\mu = A_D \;. \qquad (16b)$$

Here Γ and M are effective action and renormalization mass, and A_D is the well known conformal anomaly which is given by the various operators appearing in the lagrangian multiplied by the appropriate β-functions [15]:

$$A_D = \frac{\beta(g_s)}{2g}\, G^A_{\mu\nu} G^{A\mu\nu} + \quad \dots \quad - \frac{\beta(\lambda)}{\lambda} \frac{\lambda}{2} (\varphi^\dagger \varphi)^2 \;. \qquad (17)$$

The presence of the dilaton, the Goldstone boson of spontaneously broken scale invariance, allows to restore scale invariance by adding a Wess-Zumino term [16], satisfying

$$\delta\Gamma = \delta\Gamma_{WZ} + O(\hbar^2) \;. \qquad (18)$$

In the case of dilatons one has

$$\Gamma_{WZ} = \int d^4x\, A_D \frac{\sigma}{f} \;. \qquad (19)$$

If the theory is renormalized such that the conformal anomaly (16) is kept, the dilaton aquires a mass which is determined by the Wess-Zumino term (19) [8]:

$$m_\sigma = - \frac{4}{f^2} \langle A_D \rangle_o = \frac{1}{8\pi^2 f^2}\, \mathrm{Str}\, M^4 = \frac{1}{8\pi^2 f^2} \left(\sum_{bosons} m_b^4 - \sum_{fermions} m_f^4 \right) . \qquad (20)$$

The ground state of the theory is now uniquely determined.

The coupling of the dilaton to the conformal anomaly is reminiscent of the coupling of the axion to the chiral anomaly of a Peccei-Quinn current. A possible connection between "dilaton" and "axion" in superstring theories has been particularly emphasized by Nilles [17]. In supersymmetric theories the analogue of conformal invariance is superconformal invariance, and in the nonlinearly realized version of this symmetry [18] dilaton- and axion-like degrees of freedom occur together.

In these supersymmetric theories dilaton and axion form together a complex scalar field,

$$\chi = \sigma + ia \quad , \tag{21}$$

which is part of a chiral superfield Σ:

$$\Sigma(x,\Theta,\bar{\Theta}) = \chi(y) + \sqrt{2}\Theta\psi(y) + \Theta\Theta F(y) \quad , \quad y^\mu = x^\mu + i\Theta\sigma^\mu\bar{\Theta} \quad . \tag{22}$$

Instead of the scalar potential with dilaton one now has the superpotential

$$W = e^{3\Sigma/f}\Big[g_U QH_1U_c + g_D QH_2D_c + g_E LH_2E_c + \lambda_1 S\,(H_1H_2 + \mu_1{}^2) + \frac{1}{2}\,\mu_2 S^2 + \frac{1}{3}\,\lambda_2 S^3 \Big] \quad , \tag{23}$$

where Q, U_c, D_c, L and E_c are the familiar quark and lepton superfields, and H_1, H_2 and S are Higgs superfields (cf. ref. [19]). The corresponding action is invariant under superconformal transformations [20] which include dilatations and chiral R-transformations which play the role of the Peccei-Quinn symmetry. The kinetic term may be chosen as

$$\frac{1}{64}\, D^2\bar{D}^2\, \bar{\Sigma}\, \Sigma + \text{c.c.} \tag{24a}$$

$$\text{or} \quad \frac{f^2}{64}\, D^2\bar{D}^2 \exp\Big(\frac{\Sigma + \bar{\Sigma}}{f} \Big) + \text{c.c.} \tag{24b}$$

The second option is easily identified as the superconformally invariant lagrangian first obtained by Kobayashi and Uematsu who studied nonlinear realizations of the superconformal group [18]:

$$\frac{f^2}{64}\, D^2\bar{D}^2 \exp\Big(\frac{\Sigma + \bar{\Sigma}}{f} \Big) + \text{c.c.}$$

$$= - e^{2\sigma/f} \Big(\frac{1}{2}\,(\partial_\mu\sigma)^2 + \frac{1}{2}\,(\partial_\mu a)^2 - \frac{1}{2}\,(\bar{F} - \frac{1}{2f}\,\overline{\psi\psi})(F - \frac{1}{2f}\,\psi\psi) + \frac{i}{2f}\,\psi\sigma^\mu\partial_\mu\bar{\psi} + \frac{1}{2f}\,\psi\sigma^\mu\bar{\psi}\partial_\mu a \Big) \quad . \tag{25}$$

Quantum corrections do not respect this classical invariance. The variation of the effective action with respect to R-transformations yields the superconformal anomaly [20,21]:

$$\delta^R\Gamma = \int d^4x\; A + O(\hbar^2) \quad , \tag{26a}$$

$$A = \partial^\mu J_\mu = \frac{i}{2}\,(\bar{D}^2\bar{S} - D^2 S) \quad . \tag{26b}$$

Here $J_\mu = (R_\mu, Q_{\mu\alpha}, \bar{Q}_{\mu\dot{\alpha}}, \Theta_{\mu\nu})$ is the supercurrent which contains R-current, supersymmetry currents and energy-momentum tensor, and the chiral superfield S is again given by the various operators appearing in the lagrangian multiplied by the appropriate β-functions:

$$S \sim \begin{cases} \dfrac{\beta(\lambda)}{\lambda}\, \bar{D}^2\bar{\Phi}\Phi \;, & \text{chiral models} \\ \dfrac{\beta(g)}{g}\, \mathrm{tr}\big[\lambda^\alpha\lambda_\alpha\big] \;, & \text{gauge theories} \;, \end{cases} \tag{27}$$

$$\lambda_\alpha = - \frac{1}{4}\, D_\alpha\, \bar{D}^2\, e^{-V}\, D_\alpha\, e^{V} \quad .$$

In the supersymmetric case, instead of eq.(16b) the following trace identities hold:

$$\bar{D}^\alpha J_{\alpha\dot{\alpha}} = D_\alpha S \quad , \qquad J_{\alpha\dot{\alpha}} = \frac{1}{2}\,\sigma^\mu_{\alpha\dot{\alpha}}\, J_\mu \quad , \tag{28}$$

which include eq. (26b). The supersymmetric extension of the Wess-Zumino term (19) reads ($\delta^R\Sigma = i\frac{2}{3}f$):

$$\delta^R\Gamma = \delta^R\Gamma_{WZ} + O(\hbar^2) \quad ,$$

$$\Gamma_{WZ} = \frac{3}{4f} \int d^4x \Big[\bar{D}^2(\bar{S}\,\bar{\Sigma}) + D^2(S\,\Sigma) \Big] \quad . \tag{29}$$

Γ_{WZ} contains couplings of dilaton and axion to the conformal anomaly and the chiral anomaly of the R-current, respectively, and also interaction terms of the "dilatino" ψ.

What is left of this connection between axion and dilaton after supersymmetry breaking ? This question can be answered in a model independent way by means of "secret supersymmetry", i.e., nonlinear realizations of local supersymmetry which were systematically studied by

Samuel and Wess [22]. The starting point is the goldstino λ_α, the Goldstone fermion of spontaneously broken supersymmetry, which transforms as

$$\delta_\xi \lambda_\alpha = \kappa \xi_\alpha - \frac{2i}{\kappa} \lambda \sigma^\mu \bar{\xi} \partial_\mu \lambda_\alpha \quad . \tag{30}$$

Here δ_ξ and κ denote supersymmetry transformation and supersymmetry breaking scale, respectively. From λ_α one then obtains a superfield Λ_α which allows to construct supersymmetric lagrangians for fields which have no superpartners. Physically, they correspond to low energy effective lagrangians which are valid at distances larger than $1/\kappa$. In the "unitary gauge" the goldstino is "eaten" and one obtains a massive gravitino. In this way one can construct a supersymmetric Standard Model without scalar quarks and leptons, gauginos and higgsinos. The only remnant of supersymmetry at low energies is an enlarged Higgs sector with complex scalars H_1, H_2 and S and some relations among coupling constants which are determined by the superpotential of the original lagrangian.

In a theory with nonlinearly realized superconformal invariance one finds, that the low energy effective lagrangian contains dilaton <u>and</u> axion together. Their couplings to quarks and leptons follow from eq. (23):

$$\begin{aligned} L_P = &- \frac{3\sigma}{f} \left(m_{q_i} \bar{q}_i q_i + m_{l_i} \bar{l}_i l_i \right) \\ &- \frac{3a}{f} \left(m_{q_i} \bar{q}_i i\gamma_5 q_i + m_{l_i} \bar{l}_i i\gamma_5 l_i \right) + O\left(\frac{1}{f^2}\right) \quad . \end{aligned} \tag{31}$$

The Wess-Zumino term yields (cf. eqs. (17), (29)):

$$L_{WZ} = \frac{1}{f} \left(\sigma A_D + a A_R \right) \quad , \tag{32a}$$

$$A_D = \partial^\mu D_\mu \sim \frac{\beta(g_s)}{2g_s} \; G^A_{\mu\nu} \, G^{A\mu\nu} + \ldots \quad , \tag{32b}$$

$$A_R = \partial^\mu R_\mu \sim \frac{\beta(g_s)}{2g_s} \; G^A_{\mu\nu} \, \tilde{G}^{A\mu\nu} \quad + \quad \ldots \quad . \tag{32c}$$

In eq. (29), which leads to (32a), we have assumed that the superfield Σ couples to the complete superconformal anomaly S (cf. (28)). This is plausible but not necessary. The breaking of superconformal invariance through the anomaly would also be realized if Σ would couple to different

contributions to S, which are indicated in (27), with different strength.

Given the axion and dilaton couplings (32a) it is straightforward to compute axion and dilaton masses (cf. [23], [12]). We find:

axion:
$$m_a^{\,2} = \frac{16}{9} N_g^2 \, m_\pi^{\,2} \left(\frac{f\pi}{f} \right)^2 \frac{m_u m_d}{(m_u + m_d)^2} \quad , \tag{33a}$$

$$m_a = 3 \text{ keV} \left(\frac{10 \text{ TeV}}{f} \right) , \quad N_g = 3 \quad ; \tag{33b}$$

dilaton:
$$m_\sigma^{\,2} = \frac{1}{8\pi^2 f^2} \left(6\, m_W^4 + 3\, m_Z^4 + \sum_i m_{H_i}^4 - 12\, m_t^4 \right) , \tag{34a}$$

$$m_\sigma \approx 250 \text{ MeV} \left(\frac{10 \text{ TeV}}{f} \right) \quad . \tag{34b}$$

The numerical factor in (33a) is due to the R-charges of quarks; in (34b) we have assumed that the contributions from Higgs particles and t-quark to the dilaton mass approximately cancel. Furthermore we have assumed that the strengths of QCD and electroweak contributions to the superconformal anomaly are determined by the particle content of the Standard Model. Eqs. (33) and (34) are based on the one-loop contributions to the anomaly. To this order the different currents which are used in connection with the superconformal anomaly [21] all give the same result. From eqs. (33) and (34) we conclude:

$$\frac{m_a}{m_\sigma} \approx 10^{-4} - 10^{-5} \quad . \tag{35}$$

This result is a generic feature of theories in which superconformal invariance is broken by the superconformal anomaly and does not depend on the value of the symmetry breaking scale f.

How large is f ? On phenomenological grounds, as discussed in section 1, we expect $f > 10$ TeV. A simple, theoretically motivated guess is that f has the same order of magnitude as the supersymmetry breaking scale κ which is a function of the gravitino mass $m_{3/2}$ and the Planck mass $m_{PL} \sim 10^{18}$ GeV, i.e., $f \sim \kappa = (m_{3/2} m_{PL})^{1/2}$ [22]. If $m_{3/2}$ should be related to the Fermi scale, e.g., $m_{3/2} \sim G_F^{-1/2} \sim 300$ GeV one would have $f \sim \kappa \sim 10^{10}$ GeV and, correspondingly, $m_\sigma \sim 200$ eV and $m_a \sim 10^{-3}$ eV. This is the range of masses and couplings discussed for "invisible" axions [3].

3. The search for very light scalar particles

In this section we will discuss three recent suggestions how to produce and detect very light scalar particles in laboratory experiments. The first two, Bragg scattering and a laser experiment, are based on the two-photon coupling of light scalars which, like the two-gluon coupling, is generated by the superconformal anomaly (cf. eqs. (29), (32a)). The third experiment, which makes use of the Mößbauer effect, is based on the coupling of light scalars to nucleons.

(1) Bragg scattering [24]

In order to be specific we will restrict ourselves to the pseudoscalar axion in the following. The result which we will finally obtain holds also for the scalar dilaton. The effective lagrangian for axions and photons reads

$$L = -\frac{1}{4} F_{\mu\nu}F^{\mu\nu} + \frac{1}{2}(\partial_\mu a)^2 - \frac{m_a^2}{2} a^2 - \frac{1}{4M} F_{\mu\nu}\tilde{F}^{\mu\nu} a \quad , \qquad (36)$$

and the corresponding field equations are

$$(\Box + m_a^2) a = -\frac{1}{M} \vec{E}\cdot\vec{B} \quad , \qquad (37a)$$

$$\vec{\nabla}\times\vec{B} - \frac{\partial}{\partial t}\vec{E} = -\frac{1}{M}\left(\vec{B}\frac{\partial}{\partial t} a - \vec{E}\times\vec{\nabla} a\right) \quad , \qquad (37b)$$

$$\vec{\nabla}\cdot\vec{E} = \frac{1}{M}\vec{B}\cdot\vec{\nabla} a \ . \qquad (37c)$$

They show that in an external electromagnetic field electromagnetic waves generate axion waves and vice versa. This "Primakoff-effect" (cf. Fig. 1) is the basis of Sikivie's methods [25] for the detection of axions in strong external magnetic fields.

Fig. 1. Primakoff process: photon - axion (dilaton) conversion in an external electromagnetic field

In crystals X-rays penetrate strong electric fields. For instance, the strength of the electric field at a distance d = 1 Å away from a nucleous with charge Z = 10 is

$$|\vec{E}| = \frac{Ze}{4\pi d^2} \approx 10^6\ eV^2 \approx 10^4\ \text{Tesla} \quad . \qquad (38)$$

This suggests that Bragg scattering may be an efficient way to produce axions. An experiment for production and detection of axions is shown in Fig. 2: After the first Bragg scattering the reflected beam contains photons and axions; only the axions penetrate the absorber and produce in a second Bragg reflection, again via the Primakoff effect, photons which are detected.

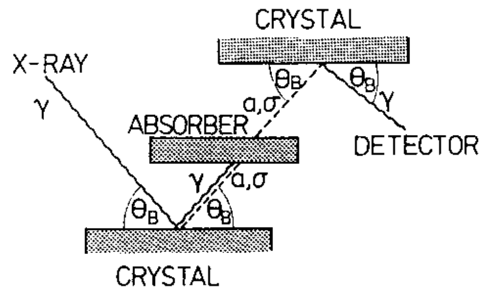

Fig. 2. Experimental setup to search for light scalar particles in Bragg scattering (from ref. [24])

In order to calculate the intensity of the final outgoing electromagnetic wave one first has to calculate the scattering amplitudes for photon-axion and axion-photon conversion. From the field equations (37) one finds that incoming plane waves

$$\vec{B}(t,\vec{x}) = \vec{B}^{(o)} e^{i(\omega t - \vec{k}\vec{x})} \quad , \qquad (39a)$$

$$\dot{a}(t,\vec{x}) \equiv \frac{\partial}{\partial t} a(t,\vec{x}) = \dot{a}^{(o)} e^{i(\omega t - \vec{k}\vec{x})} \quad , \quad \dot{a}^{(o)} = i\omega a^{(o)} \quad , \qquad (39b)$$

generate the outgoing spherical waves

$$\dot{a}(t,\vec{x}) = \frac{F_a(2\Theta)}{4\pi M} \vec{e}_r \cdot \vec{B}^{(o)} \frac{1}{r} e^{i(\omega t - kr)} \quad , \qquad (40a)$$

$$\vec{B}(t,\vec{x}) = -\frac{F_a(2\Theta)}{4\pi M}\ \vec{e}_r \times (\vec{e}_r \times \hat{\vec{k}})\dot{a}^{(o)}\ \frac{1}{r}\ e^{i(\omega t - kr)} \quad , \tag{40b}$$

where

$$F_a(2\Theta) = k^2 \int d^3x\ \phi(\vec{x})\ e^{i\vec{q}\vec{x}} \quad , \qquad \vec{E} = -\vec{\nabla}\phi \quad ,$$

$$\vec{q} = \vec{k}' - \vec{k}, \quad \vec{k}' = k\vec{e}_r, \quad \vec{e}_r = \frac{\vec{x}}{r}, \quad 2\Theta = \sphericalangle(\vec{k},\vec{k}'). \tag{40c}$$

Here we have neglected the axion mass for simplicity. Comparison with eq. (37a) shows that the formfactor $F_a(2\Theta)$ corresponds to an average electric field seen by the X-rays,

$$\bar{E}(k,2\Theta) = \frac{1}{kd^3} F_a(2\Theta) \quad , \tag{41}$$

which, as we will see, determines the probability for photon-axion conversion in a crystal.

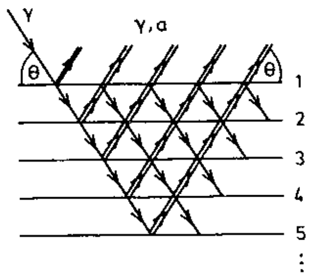

Fig. 3. Coherent production of axions in a crystal

Given the scattering amplitudes for photon-axion transitions and elastic photon (Thomson) scattering by a single atom one can calculate intensity and opening angle of the outgoing beam in a Bragg reflection on a crystal by means of Darwins dynamical theory [26]. Multiple scattering inside the crystal (cf. Fig. 3) leads to transmitted and reflected photon and axion plane waves, which at the n-th layer of atoms satisfy the linear system of equations (cf. Fig. 4):

$$A_{n+1} = (1+i\rho_o)A_n + i\rho\ e^{2in\phi}\ B_{n+1} + i\xi\ e^{2in\phi} D_{n+1} \quad , \tag{42a}$$

$$B_n = i\rho e^{-2in\phi}\ A_n + (1+i\rho_o)B_{n+1} + i\xi\ e^{-2in\phi}\ C_n \quad , \tag{42b}$$

$$C_{n+1} = i\xi\ e^{2in\phi}\ B_{n+1} + C_n \quad , \tag{42c}$$

$$D_n = i\xi\ e^{-2in\phi}\ A_n + D_{n+1}, \tag{42d}$$

where

$$\phi = k\ d\ \sin\Theta \ , \tag{43}$$

$$\rho = \rho(2\Theta) = \frac{\alpha F_\gamma(2\Theta) N_s \lambda}{m \sin\Theta} \quad , \qquad \rho_o = \rho F_\gamma(0)/\ F_\gamma(2\Theta) \ , \tag{44}$$

$$\xi = \xi(2\Theta) = \frac{F_a(2\Theta) N_s \lambda}{4\pi M \sin\Theta}\ \sin 2\Theta \ . \tag{45}$$

Here α, m, λ, F_γ, N_s and d are fine structure constant, electron mass, photon wave length, atomic structure factor (cf. [26]), density of scattering centers per unit area and lattice spacing, respectively. For $\phi = \pi$ the Bragg condition is fullfilled and one has constructive interference. In the case ξ=o eqs. (42) reduce to a system of equations familiar from Thomson scattering [26].

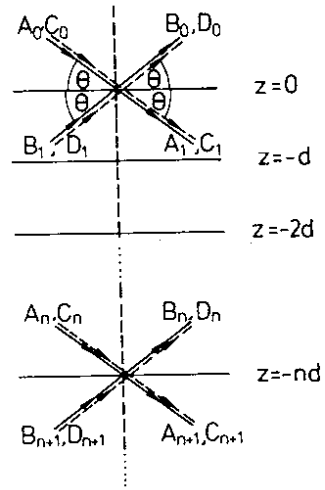

Fig. 4. Beams of photons and scalar particles in the crystal: A(B) = transmitted (reflected) photon beam, C(D) = transmitted (reflected) axion (dilaton) beam (from ref. [24])

Since $\xi \ll \rho$ eqs. (42) can be solved as for Thomson scattering and obtains for the axion-photon and photon-axion transition propabilities up to terms of order $\rho\xi$, ξ^2:

$$P(a\to\gamma) = \left|\frac{D_o}{A_o}\right|^2 = \left|\frac{B_o}{C_o}\right|^2$$

$$= P(\gamma\to a) = \left(\frac{\xi}{\rho}\right)^2 = \left(\frac{\bar{E}}{2}\frac{1}{M}\sin 2\Theta\right)^2 \quad , \qquad (46)$$

where

$$l = \frac{md}{\alpha N_S \lambda F_\gamma(2\Theta)} \qquad (47)$$

is the penetration depth of the X-ray into the crystal. Eq. (46) is very similar to the transition probability for photon-axion conversion in an external magnetic field.

In a realistic experiment [24,27], which makes use of the very intense source of X-rays with a brightness of $\Phi \sim 10^{18}$/sec(0.1% BW) soon available at the European Synchrotron Radiation Facility the following lower bound on the mass scale M can be achieved:

$$M > 1\cdot 10^3 \text{ GeV} \left(\frac{\bar{E}}{1\text{keV}^2}\frac{l}{1\mu\text{m}}\frac{\sin 2\Theta}{0.34}\right)\left(\frac{\Phi}{10^{18}/\text{s } 0.1\%\text{BW}}\frac{T}{100\text{d}}\frac{10}{N^{obs}}\right)^{1/4} . \qquad (48)$$

Here we have normalized the running time to T = 100 days and the number of observed photons to $N^{obs} = 10$. This lower bound will be slightly decreased if finite detection efficiency and temperature effects, i.e. the Debye-Waller factor, are taken into account. In the case of nonvanishing axion mass the axions are emitted under an angle $\bar{\Theta}_B < \Theta_B$. In principle the experiment is sensitive to masses $m_a < \omega \sin\Theta_B$, where ω is the photon energy.

The lower bound on M is linear in the penetration depth l which is only 1 μm for a typical Bragg scattering process. In the case of Laue scattering (cf. Fig. 5) a much larger penetration depth, up to 1 cm for 100 keV photons and scattering angle $\Theta \sim 1^\circ$, can be achieved [27]. If eq. (46) for the transition probability $P(a\to\gamma) = P(\gamma\to a)$ also holds for Laue scattering, which remains to be seen, the lower bound on M given in (48) could be improved by three orders of magnitude up to $\sim 10^6$ GeV.

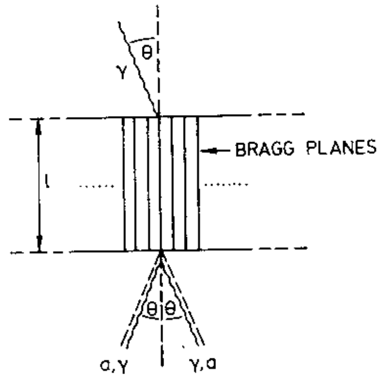

Fig. 5. Axion production in Laue scattering

(2) Laser experiment [28]

The experiment described in the last subsection is based on the idea to discover axions by "shining light through walls", which was first put forward in connection with the laser experiment shown in Fig. 6 [28]. Here the photon-axion transitions take again place via the Primakoff effect, now with an external magnetic field. The transition probability in the homogeneous magnetic Field B of a dipole magnet with length L can be computed from eqs. (37). For zero mass axions one finds [25, 28]:

$$P(a\to\gamma) = P(\gamma\to a) = \left(\frac{BL}{2M}\right)^2 , \qquad (49)$$

which is the analog of the transition probability (46) in an external electric field.

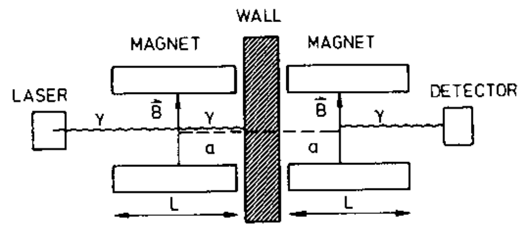

Fig. 6. "Shining light through walls": photon-axion and axion-photon conversion in an external magnetic field (cf. [28])

In a realistic experiment a rather large lower bound on the mass scale M can be achieved (cf. [28]):

$$M > 4\cdot 10^{8}\ \text{GeV}\left(\frac{B}{10\text{TeV}}\,\frac{L}{10\text{m}}\right)\left(\frac{\Phi}{2.5\cdot 10^{21}\text{s}^{-1}}\,\frac{T}{100\text{d}}\,\frac{10}{N^{obs}}\right)^{1/4} . \qquad (50)$$

In comparison with Bragg scattering one looses 3 orders of magnitude in the field strength but one gains 7 orders in the penetration length L. Furthermore the photon flux is larger: $\Phi = 2.5\cdot 10^{21}\text{s}^{-1}$ corresponds to a 1 kW laser with frequency $\omega = 2.5$ eV. However, due to this small energy the laser experiment is sensitive only to axions with masses below ~ 1 eV, which is 4 - 5 orders of magnitude below the sensitivity of X-ray experiments.

An alternative laser experiment, where the change in polarization due to the photon-axion coupling is measured, has been suggested by Maiani, Petronzio and Zavattini [29] and is currently carried out at Brookhaven. First results yield the lower bound $M > 10^{5}$ GeV for $m_a < 10^{-3}$ eV [30].

(3) Mößbauer effect [31]

Another variant of the "shining light through walls"-idea is the recent suggestion by de Rújula and Zioutas to search for axions by means of the Mößbauer effect [31]. The experimental set-up is sketched in Fig. 7: An intense Mößbauer source, which emits M1 photons in the process $N^{*} \to N + \gamma$, is used as source for axions through the transition $N^{*} \to N + a$; the source is shielded; only the axions penetrate the wall and resonantly excite the nuclei of the absorber, i.e., $N + a \to N^{*}$; the deexcitation of these nuclei by emission of photons or internal conversion is the axion signal.

A realistic experiment[1] can reach the following upper bound on the axion-nucleon coupling [29]:

[1]The required specifications are [31]: 200 mCi of ^{119m}Sn (produced from 0.1 gr of ^{118}Sn exposed for 200 days to a flux of $3\cdot 10^{14}$neutrons/cm^{2}s) as source; 10 layers of 100 mg/cm^2 ^{119}Sn as resonant absorber; the frequency of the γ-signal is ω = 24 keV.

$$\frac{\alpha_a}{\alpha_{em}} < 5\cdot 10^{-8} \quad , \quad \alpha_a = \frac{g^2_{aNN}}{4\pi} \quad . \qquad (51)$$

With $g_{aNN} \sim \frac{m_N}{f} \sim 1/f[\text{GeV}]$ [32] the corresponding lower bound on the mass scale f is

$$f > 10^{4}\ \text{GeV} \quad . \qquad (52)$$

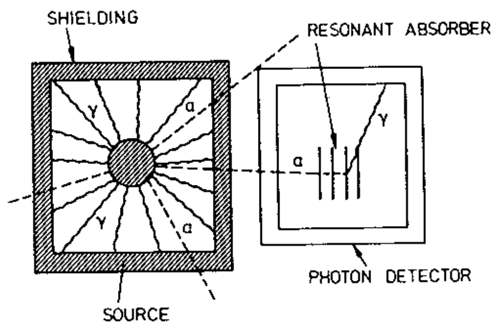

Fig. 7. Sketch of experimental setup to produce and detect axions via the Mössbauer effect

Comparing this bound on f with the bounds (48) and (50) on M, one has to take into account that in models where the two-photon coupling is generated by the anomaly, the connection between f and M is given by $f \sim 10^{-2}$ M. Hence the bound on f from the Mößbauer effect appears more stringent than the bound achievable in Bragg scattering, less stringent than the bound obtainable with the laser experiment and possibly comparable to Laue scattering.

The lower bounds on the axion - two photon coupling, which can be obtained in the three experiments discussed in this section, have to be compared to the astrophysical bounds [33]. Fig. 7, which is partly taken from ref. [28], shows that the range accessible by the simplest version of the laser experiment is already excluded by bounds derived from the allowed energy loss of the sun. The advantage of the X-ray and Mößbauer experiments is that the sensitivity with respect to the axion mass extends to about 10 keV, i.e., about five orders of magnitude beyond the sensitivity of the laser experiment. However, if bounds from lifetimes of

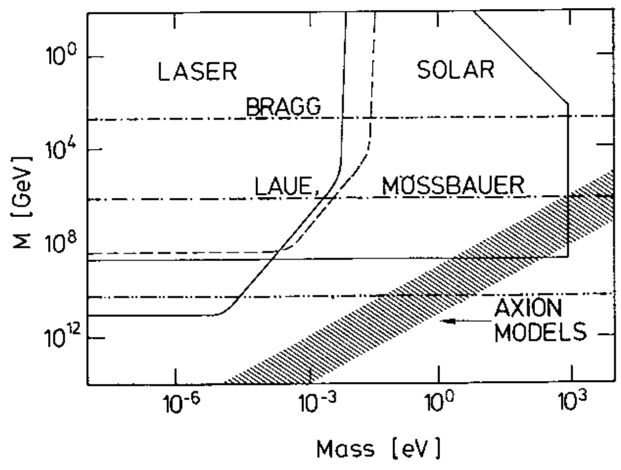

Fig. 8. Comparison of the sensitivity of the three laboratory experiments with astrophysical bounds. The solar and laser bounds are taken from ref. [28]. The dashed double-dotted line is obtained from helium burning red giants [34].

helium burning red giants [34] are included, no "window" in the $M - m_a$ plane remains for the laboratory experiments, which is not already excluded by astrophysical considerations.

4. Summary

New interactions beyond the Standard Model, which are characterized by a mass scale f much larger than the Fermi scale of weak interactions, can manifest themselves at low energies through pseudo-Goldstone bosons, i.e., light, weakly interacting (pseudo)scalar particles.

Symmetry arguments based on the Standard Model single out two spin-0 particles, dilaton and axion, which are related to spontaneously broken symmetries - scale invariance and a chiral, Peccei-Quinn type symmetry. In supersymmetric extensions of the Standard Model, dilaton and axion belong to the same supermultiplet. If their masses are generated by the

superconformal anomaly, their mass ratio is $m_a/m_\sigma \approx 10^{-4} - 10^{-5}$, independent of the mass scale f.

Several novel experiments have been suggested which can significantly extend the bounds on masses and couplings of light scalars obtained from current laboratory experiments. So far, however, none of these new experiments can improve the present astrophysical bounds. We leave it to the reader to decide whether such experiments may nevertheless be sufficiently interesting to be carried out.

The proposed experiments to search for axions and dilatons all require production <u>and</u> detection of the scalar particles, i.e., the rate of observed events is proportional to $(1/f)^4$. The sensitivity on the mass scale f might be improved if some new indirect signature for the production of light scalars could be found, since the corresponding signal would then only be proportional to $(1/f)^2$.

On the theoretical side it appears most important to obtain a prediction of, or restrictions on the mass scale f. This may be possible in theories where f and the Fermi scale $G_F^{-1/2}$ are related, as it is the case in some supergravity models.

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References

[1] P. Langacker, in Proc. of the XXIV Int. Conf. on High Energy Physics, eds. R. Kotthaus and J.H. Kühn (Munich, 1988) p. 190

[2] W. Buchmüller and D. Wyler, Nucl. Phys. B268 (1986) 621

[3] For a review, see J.E. Kim, Phys. Rep. 150 (1987) 1

[4] R.D. Peccei and H. Quinn, Phys. Rev. Lett. 38 (1977) 1440; S. Weinberg, Phys. Rev. Lett. 40 (1978) 223; F. Wilczek, Phys. Rev. Lett. 40 (1978) 279

[5] P. Jordan, Z. Phys. 157 (1959) 112; C. Brans and R.H. Dicke, Phys. Rev. 124 (1961) 925

[6] G. Mack, Nucl. Phys. B5 (1968) 499; P.G.O. Freund and Y. Nambu, Phys. Rev. 174 (1968) 1741

[7] R.D. Peccei, J. Solà and C. Wetterich, Phys. Lett. B 195 (1987) 183

[8] W. Buchmüller and N. Dragon, Phys. Lett. B 195 (1987) 417

[9] C. Wetterich, Nucl. Phys. B302 (1988) 668

[10] S. Weinberg, Rev. Mod. Phys. 61 (1989) 1

[11] G.D. Coughlan, I. Kani, G.G. Ross and G. Segrè, Nucl. Phys. B316 (1989) 469

[12] W. Buchmüller and N. Dragon, Nucl. Phys. B321 (1989) 207

[13] E.T. Tomboulis, preprint UCLA/88/TEP/42 (1988)

[14] For a review, see S. Coleman, Aspects of Symmetry (Cambridge University Press, 1985) p. 67

[15] S.L. Adler, J.C. Collins and A. Duncan, Phys. Rev. D15 (1977) 1712; J.C. Collins, A. Duncan and S.D. Joglekar, Phys. Rev. D16 (1977) 438

[16] J. Wess and B. Zumino, Phys. Lett. B37 (1971) 95

[17] H.P. Nilles, in Proc. of the 9th Johns Hopkins Workshop, Florence, 1985; J.P. Derendinger, L.E. Ibañez and H.P. Nilles, Nucl. Phys. B267 (1986) 365

[18] K. Kobayashi and T. Uematsu, Nucl. Phys. B263 (1986) 309

[19] H.P. Nilles, Phys. Rep. 110 (1984) 1

[20] For a review, see O. Piguet and K. Sibold, Renormalized Supersymmetry, Birkhäuser Inc. (Boston 1986)

[21] O. Piguet and K. Sibold, Int. J. Mod. Phys. A (1986) 913; M.T. Grisaru and P.C. West, Nucl. Phys. B254 (1985) 249; M.A. Shifman and A.I. Vainshtein, Nucl. Phys. B277 (1986) 456

[22] S. Samuel and J. Wess, Nucl. Phys. B233 (1984) 488

[23] W.A. Bardeen, R.D. Peccei and T. Yanagida, Nucl. Phys. B279 (1987) 401

[24] W. Buchmüller and F. Hoogeveen, Hannover preprint ITP-UH 9/89 (1989)

[25] P. Sikivie, Phys. Rev. Lett. 51 (1983) 1415; 52 (1984) 695 (E); Phys. Rev. D32 (1985) 1988

[26] See, for instance, B.E. Warren, X-Ray Diffraction, Addison-Wesley Comp. (Reading, 1969)

[27] G. Materlik, private communication

[28] K. van Bibber et al., Phys. Rev. Lett. 59 (1987) 759

[29] L. Maiani, R. Petronzio and G. Zavattini, Phys. Lett. B 175 (1986) 359

[30] A. C. Melissinos, private communication

[31] A. de Rujula and K. Zioutas, Phys. Lett. B217 (1989) 354

[32] T.W. Donnelly et al., Phys. Rev. D18 (1987) 1607

[33] M. Yoshimura, in Proc. of the XXIII Int. Conf. on High Energy Physics, Berkeley, ed. S.C. Loken (1986) p. 189; G.G. Raffelt, in Proc. of the XIV Int. Conf. on High Energy Physics, Munich, eds. R. Kotthaus and J.H. Kühn (1988) p. 1519

[34] G.G. Raffelt, private communication; G.G. Raffelt and D.S.P. Dearborn, Phys. Rev. D36 (1987) 2211