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External fields as a probe for fundamental physics

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Abstract
Quantum vacuum experiments are becoming a flexible tool for investigating fundamental physics. They are particularly powerful in searching for new light but weakly interacting degrees of freedom and are thus complementary to accelerator-driven experiments. I review recent developments in this field, focusing on optical experiments in strong electromagnetic fields. In order to characterize potential optical signatures, I discuss various low-energy effective actions which parameterize the interaction of particle-physics candidates with optical photons and external electromagnetic fields. Experiments with an electromagnetized quantum vacuum and optical probes do not only have the potential to collect evidence for new physics, but special-purpose setups can also distinguish between different particle-physics scenarios and extract information about underlying microscopic properties.

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1. Introduction

With the advent of quantum field theory, our understanding of the vacuum has changed considerably from a literal 'nothing' to such a complex 'something' that its quantitative description requires us to know almost 'everything' about a given system.

Consider a closed quantum field theoretic system in a box with boundaries, where all matter density is already removed (pneumatic vacuum). Still, the walls of the system which are in contact with surrounding systems may have a temperature, releasing black-body radiation into the box. Charges and currents outside the box can create fields, exerting their influence on the box's inside. The box may furthermore be placed on a gravitationally curved manifold. Finally, the boundaries themselves do generally impose constraints on the fluctuating quantum fields inside.

A pure quantum vacuum which is as close to trivial as possible requires us to take the limit of vanishing parameters which quantify the influence on the quantum fluctuations, i.e.

temperature, fields $\to 0$ and boundaries $\to \infty$. Even then, the quantum vacuum may be thought of as an infinity of ubiquitous virtual processes—fluctuations of the quantum fields representing creations and annihilations of wave packets ('particles') in spacetime—which are compatible with Heisenberg's uncertainty principle.

Even if the ground state realizes the naive anticipation of vanishing field expectation values, we can probe the complex structure of the quantum vacuum by applying external fields, boundaries, etc, and measuring the response of the vacuum to a suitable probe. For instance, let us send a weak light beam into the box; it may interact with the virtual fluctuations and will have finally traveled through the box at just 'the speed of light'. If we switch on an external magnetic field, the charged quantum fluctuations in the box are affected and reordered by the Lorentz force. This has measurable consequences for the speed of the light probe which now interacts with the reordered quantum fluctuations. Thus, quantum field theory invalidates the superposition principle of Maxwell's theory. The quantum world creates nonlinearities and also nonlocalities [1–3].

Quantum vacuum physics inspires many research branches, ranging from mathematical physics studying field theory with boundaries and functional determinants to applied physics where the fluctuations may eventually be used as a building block to design dispersive forces in micro- and nanomachinery. Many quantum vacuum phenomena such as the Casimir effect are similarly fundamental in quantum field theory as the Lamb shift or $g-2$ experiments, and hence deserve to be investigated and measured with the same effort. Only a high-precision comparison between quantum vacuum theory and experiment can reveal whether we have comprehensively understood and properly computed the vacuum fluctuations.

In this paper, I will argue that, with such a comparison, one further step can be taken: a high-precision investigation can then also be used to look for systematic deviations as a hint for new physics phenomena. Similar to $g-2$, quantum vacuum experiments can be systematically used to explore new parameter ranges of particle-physics models beyond the Standard Model (BSM).

What are the scales of sensitivity which we can expect to probe? Consider a typical Casimir experiment: micro- or mesoscopic setups probe dispersive forces between bodies at a separation $a = \mathcal{O}(\text{nm—10}\,\mu\text{m})$. This separation a also sets the scale for the dominant quantum fluctuation wavelengths which are probed by the apparatus. The corresponding energy scales are of order $\mathcal{O}(10\text{ meV–100 eV})$. As another example, consider an optical laser propagating in a strong magnetic field of a few Tesla. Again, the involved energy scales allow us to probe quantum fluctuations below the $\mathcal{O}(10\,\text{eV})$ scale. Therefore, quantum vacuum experiments can probe new physics below the eV scale and hence are complementary to accelerators. Typical candidates are particles with masses in the meV range, i.e. *physics at the milli scale* [4].

The particular capability of these experiments is obviously not a sensitivity to heavy particles, but a sensitivity to light but potentially very weakly coupled particles. In the following, I will especially address optical experiments. Here, there are at least two lever arms for increasing the sensitivity toward weak coupling: consider a laser beam entering an interaction region, say a magnetized quantum vacuum; some photons may leave the region toward a detector. Let us assume that the setup is such that the Standard Model predicts zero photons in the detector; this implies that the observation of a single photon (which is technically possible) is already a signature for new physics. On the other hand, an incoming beam, for instance, from an optical 1000 Watt laser contains $\sim 10^{21}$ photons per second. It is this ratio of 10^{21}:1 which can be exploited for overcoming a weak-coupling suppression. Second, the interaction region does not have to be microscopic as in accelerator experiments, but can be of a laboratory size (meters) or can even increase to kilometers if, e.g., the laser light is stored in a high-finesse cavity.

Why should we care for the milli scale at all? First of all, exploring a new particle-physics landscape is worthwhile in itself; even if there is no discovery, it is better to *know* about non-existence than to *assume* it. Second, we already know about physics at the milli scale: neutrino mass differences and potentially also their absolute mass are of order $\mathcal{O}(1\text{–}100\,\text{meV})$; also, the cosmological constant can be expressed as $\Lambda \sim (2\,\text{meV})^4$. A more systematic search for further particle physics at the milli scale is hence certainly worthwhile and could perhaps lead to a coherent picture. Third, a large number of Standard-Model extensions not only involves but often requires—for reasons of consistency—a hidden sector, i.e. a set of so far unobserved degrees of freedom very weakly coupled to the Standard Model. A discovery of hidden-sector properties could much more decisively single out the relevant BSM extension than the discovery of new heavy partners of the Standard Model.

Optical quantum vacuum experiments can be very sensitive to new light particles which are weakly coupled to photons. From a bottom-up viewpoint, I will first discuss low-energy effective theories of the Standard Model and of BSM extensions which allow for a classification of possible phenomena and help relating optical observables to fundamental particle properties. Subsequently, current bounds on new-physics parameters are critically examined. In section 3, I briefly describe current and future experimental setups, and discuss recently published data. An emphasis is put on the question of how dedicated quantum vacuum experiments can distinguish between different particle-physics scenarios and extract information about the nature of the involved degrees of freedom. Section 4 gives a short account of underlying microscopic models that would be able to reconcile a large anomalous signal in the laboratory with astrophysical bounds. Conclusions are given in section 5.

2. Low-energy effective actions

A plethora of ideas for BSM extensions can couple new particle candidates to our photon. From a bottom-up viewpoint, many of these ideas lead to similar consequences for low-energy laboratory experiments, parameterizable by effective actions that describe the photon coupling to the new effective degrees of freedom. In the following, we list different effective actions that are currently often used for data analysis. This list is not unique nor complete.

2.1. QED and Heisenberg–Euler effective action

The first example is standard QED as a low-energy effective theory of the Standard Model: if there are no light particles coupling to the photon other than those of the Standard Model, present and near-future laboratory experiments will only be sensitive to pure QED degrees of freedom, photon and electron. If the variation of the involved fields as well as the field strength is well below the electron mass scale, the low-energy effective action is given by the lowest order Heisenberg–Euler effective action [1–3],

$$\Gamma_{\text{HE}} = \int_x \left\{ -\frac{1}{4} F_{\mu\nu} F^{\mu\nu} + \frac{8}{45} \frac{\alpha^2}{m^4} \left(\frac{1}{4} F_{\mu\nu} F^{\mu\nu} \right)^2 + \frac{14}{45} \frac{\alpha^2}{m^4} \left(\frac{1}{4} F_{\mu\nu} \widetilde{F}^{\mu\nu} \right)^2 + \mathcal{O}\left(\frac{F^6}{m^8}, \frac{\partial^2 F^2}{m^2} \right) \right\}, \tag{1}$$

which arises from integrating out the 'heavy' electron–positron degrees of freedom to one-loop order. In addition to the Maxwell term, the second and third term exemplifies the fluctuation-induced nonlinearities. The corresponding quantum equations of motion thus entail a photon self-coupling. As an example, let us consider the propagation of a laser beam with a weak amplitude in a strong magnetic field B. From the linearized equations of motion for the laser beam, we obtain a dispersion relation which can be expressed in terms of refractive indices for the magnetized quantum vacuum [2, 5, 6]:

$$n_{\parallel} \simeq 1 + \frac{14}{45}\frac{\alpha^2}{m^4}B^2\sin^2\theta_B, \qquad n_{\perp} \simeq 1 + \frac{8}{45}\frac{\alpha^2}{m^4}B^2\sin^2\theta_B, \tag{2}$$

where θ_B is the angle between the B field and the propagation direction. Most importantly, the refractive indices, corresponding to the inverse phase velocity of the beam, depend on the polarization direction $\parallel$ or $\perp$ of the laser with respect to the B field. The magnetized quantum vacuum is birefringent. As a corresponding observable, an initially linearly polarized laser beam which has nonzero components for both $\parallel$ and $\perp$ modes picks up an *ellipticity* by traversing a magnetic field: the phase relation between the polarization modes changes, but their amplitudes remain the same. The ellipticity angle ψ is given by $\psi = \frac{\omega}{2}\ell(n_\parallel - n_\perp)\sin 2\theta$, where θ is the angle between the polarization direction and the B field, and ℓ is the path length inside the magnetic field.

So far, a direct verification of QED vacuum magnetic birefringence has not been achieved; if measured it would be the first experimental proof that the superposition principle in vacuum is ultimately violated for macroscopic electromagnetic fields.

Another optical observable is important in this context: imagine that some effect modifies the amplitudes of $\parallel$ or $\perp$ components in a polarization-dependent manner, but leaves the phase relations invariant. By such an effect, a linearly polarized beam will then effectively change its polarization direction after a passage through a magnetic field by a *rotation* angle $\Delta\theta$. Since amplitude modifications involve an imaginary part for the index of refraction, rotation from a microscopic viewpoint is related to particle production or annihilation. In QED below threshold $\omega < 2m$, electron–positron pair production by an incident laser is excluded. Only photon splitting in a magnetic field would be an option [6]. However, for typical laboratory parameters, the mean free path exceeds the size of the universe by many orders of magnitude and hence is irrelevant[1]. We conclude that a sizeable signal for vacuum magnetic rotation $\Delta\theta$ in an optical experiment would be a signature for new fundamental physics.

2.2. *Axion-like particle (ALP)*

As a first BSM example, we consider a neutral scalar ϕ or pseudo-scalar degree of freedom ϕ^- which is coupled to the photon by a dimension-5 operator,

$$\Gamma_{\text{ALP}} = \int_x \left\{ -\frac{g}{4}\phi^{(-)}F^{\mu\nu}\overset{(\sim)}{F}_{\mu\nu} - \frac{1}{2}(\partial\phi^{(-)})^2 - \frac{1}{2}m_\phi{}^2\phi^{(-)2} \right\}. \tag{3}$$

This effective action is parameterized by the particle's mass m_ϕ and the dimensionful coupling g. For the pseudo-scalar case, this action is similar to axion models [8], where the two parameters are related, $m_\phi \sim g$. Here, we have a more general situation with free parameters in mind which we refer to as axion-like particles (ALP). In optical experiments in strong B fields, ALPs can induce both ellipticity and rotation [9], since only one photon polarization mode couples to the axion and the external field: the $\parallel$ mode in the pseudo-scalar case and the $\perp$ mode in the scalar case. For instance, coherent photon-axion conversion causes a depletion of the corresponding photon mode, implying rotation. Solving the equation of motion for the coupled photon-ALP system for the pseudo-scalar case yields a prediction for the induced

[1] It is amusing to observe that it is neutrino-pair production which could be the largest Standard-Model contribution to an optical rotation measurement in a strong electromagnetic field [7], but, of course, it is similarly irrelevant for current and near-future experiments.

ellipticity and rotation:

$$\Delta\theta^{-}=\left(\frac{gB\omega}{m_\phi^2}\right)^2 \sin^2\left(\frac{Lm_\phi^2}{4\omega}\right)\sin 2\theta, \qquad \psi^{-}=\frac{1}{2}\left(\frac{gB\omega}{m_\phi^2}\right)^2\left(\frac{Lm_\phi^2}{2\omega}-\sin\left(\frac{Lm_\phi^2}{2\omega}\right)\right)\sin 2\theta, \tag{4}$$

for single passes of the laser through a magnetic field of length L. For the scalar, we have $\Delta\theta = -\Delta\theta^{-}$, $\psi = -\psi^{-}$. This case is a clear example of how fundamental physics could be extracted from a quantum vacuum experiment: measuring ellipticity and rotation signals uniquely determines the two model parameters, ALP mass m_ϕ and ALP-photon coupling g respectively. Measuring the signs of $\Delta\theta$ and ψ can even resolve the parity of the involved particle.

Various microscopic particle scenarios lead to a low-energy effective action of the type (3). The classic case of the axion represents an example in which only the weak coupling to the photon is relevant and all other potential couplings to matter are negligible. In this case, the laser can be frequency-locked to a cavity such that both quantities are enhanced by a factor of N_{pass} accounting for the number of passes. For the generated ALP component, the cavity is transparent. This facilitates another interesting experimental option, namely to shine the ALP component through a wall which blocks all the photons. Behind the wall, a second magnetic field can induce the reverse process and photons can be regenerated out of the ALP beam [10]. The regeneration rate is

$$\dot{N}^{(-)}_{\gamma\text{reg}} = \dot{N}_0\left(\frac{N_{\text{pass}}+1}{2}\right)\frac{1}{16}(gBL\cos\theta)^4\left[\sin\left(\frac{Lm_\phi^2}{4\omega}\right)\Big/\frac{Lm_\phi^2}{4\omega}\right]^4, \tag{5}$$

where $\dot{N}_0$ is the initial photon rate, and the magnetic fields are assumed to be identical.

In other models, such as those with a chameleon mechanism [11], the ALP cannot penetrate into the cavity mirrors but gets reflected back into the cavity. Whereas this has no influence on the single-pass formulae for ψ and $\Delta\theta$ in equation (4), the use of cavities and further experimental extensions can be used to distinguish between various microscopic models; see in the following.

2.3. Minicharged particle (MCP)

In addition to the example of a neutral particle, optical experiments can also search for charged particles. If their mass is at the milli scale, these experiments can even look for very weak coupling, i.e. minicharged particles (MCPs) [12], the charge of which is smaller by a factor of ϵ in comparison with the electron charge. If the MCP is, for instance, a Dirac spinor ψ_ϵ, the corresponding action is

$$\Gamma_{\text{MCP}} = -\bar{\psi}(\mathrm{i}\partial\!\!\!/ + \epsilon e A\!\!\!/)\psi + m_\epsilon\bar{\psi}\psi, \tag{6}$$

where we again encounter two parameters, ϵ and the MCP mass m_ϵ. At a first glance, the system looks very similar to QED. However, since the particle mass m_ϵ can be at the milli scale or even lighter, the weak-field expansion of the Heisenberg–Euler effective action for slowly varying fields (1) is no longer justified. Both field strength and laser frequency can exceed the electron mass scale with various consequences [13]: the laser frequency can be above the pair-production threshold $\omega > 2m_\epsilon$ such that a rotation signal becomes possible. Second, there is no perturbative ordering anymore as far as the coupling to the B field is concerned; hence the MCP fluctuations have to be treated to all orders with respect to B. All relevant information is encoded in the polarization tensor corresponding to an MCP loop with two photon legs and an infinite number of couplings to the B field which is well known from the QED literature [2, 3, 14]. Explicit results are available in certain asymptotic limits, for instance, for the rotation,

$$\Delta\theta \simeq \frac{1}{12}\frac{\pi}{\Gamma\left(\frac{1}{6}\right)\Gamma\left(\frac{13}{6}\right)}\left(\frac{2}{3}\right)^{\frac{2}{3}}\epsilon^2\alpha(m_\epsilon\ell)\left(\frac{m_\epsilon}{\omega}\right)^{\frac{1}{3}}\left(\frac{\epsilon eB}{{m_\epsilon}^2}\right)^{\frac{2}{3}}, \qquad \text{for} \quad \frac{3}{2}\frac{\omega}{m_\epsilon}\frac{\epsilon eB}{m_\epsilon^2} \gg 1, \tag{7}$$

which is valid above threshold and for a high number of allowed MCP Landau levels. Similar formulae exist for ellipticity or the case of spin-0 MCPs [15]. Note that this rotation becomes independent of m_ϵ in the small-mass limit, such that equation (7) apparently implies a sensitivity to arbitrarily small masses. In practice, this sensitivity is limited for other reasons: for instance, once the associated Compton wavelength $\sim 1/m_\epsilon$ becomes larger than the size of the magnetic field, the constant-field assumption, which is often used for calculating the polarization tensor, is no longer valid. The rotation depends on the size of the magnetic field, the scale of which acts as a cutoff for the sensitivity toward smaller masses; e.g. for $\ell \simeq 1$ m, the MCP mass should satisfy $m_\epsilon \gg 0.2\,\mu$eV. Let me stress that the computation of polarization tensors in inhomogeneous fields is a challenge for standard methods and remains an interesting question for future research. Progress may come from modern worldline techniques [16, 17].

2.4. Paraphotons

In addition to neutral scalars or weakly charged particles, we may also consider additional (hidden) gauge fields which interact weakly with the photon. A special coupling is provided by gauge-kinetic mixing which occurs only between Abelian gauge fields, hence involving a second photon, i.e. a paraphoton, γ' [12],

$$\Gamma_{\gamma\gamma'} = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} - \frac{1}{4}F'_{\mu\nu}F'^{\mu\nu} - \frac{1}{2}\chi F^{\mu\nu}F'_{\mu\nu} - \frac{1}{2}\mu^2 A'_\mu A'^\mu, \tag{8}$$

with a mixing parameter χ and a paraphoton mass term μ. Without the mass term, the kinetic terms could be diagonalized by a non-unitary shift $A'_\mu \to \hat{A}'_\mu - \chi A_\mu$ which would decouple the fields at the expense of an unobservable charge renormalization. The mass term does not remain diagonal by this shift, such that observable $\gamma\gamma'$ oscillations arise from mass mixing in this basis. The pure paraphoton theory is special in the sense that $\gamma\gamma'$ conversion is possible without an external field and is not sensitive to polarizations. For instance, the conversion rate after a distance L is given by $\mathcal{P}_{\gamma\to\gamma'} = 4\chi^2\sin^2\frac{\mu^2 L}{4\omega}$. Therefore, paraphotons can be searched for in future light-shining-through-walls experiments [18]. Below, we discuss microscopic scenarios in which paraphotons and MCPs naturally occur simultaneously.

2.5. Bounds on low-energy effective parameters

Many different observations seem to constrain the parameters in the effective theories listed above. The strongest constraints typically come from astrophysical observations usually in combination with energy-loss arguments. Consider, for instance, the ALP low-energy effective action (3). Assuming that it holds for various scales of momentum transfer, we may apply it to solar physics. Thermal fluctuations of electromagnetic fields in the solar plasma, giving rise to non-vanishing $F^{\mu\nu}\overset{(\sim)}{F}_{\mu\nu}$, act as a source for $\phi^{(-)}$ ALPs. In the absence of other sizeable interactions, ALPs escape the solar interior immediately and contribute to stellar cooling. A similar argument for the helium-burning lifetime of HB stars leads to a limit $g \lesssim 10^{-10}\,\text{GeV}^{-1}$ for ALP masses in the eV range and below [19]. Monitoring actively a potential axion flux from the sun as done by the CAST experiment even leads to a slightly better constraint for ALP masses <0.02 eV [20].

Astrophysical energy-loss arguments also constrain MCPs [21]: for instance, significant constraints on ϵ come from helium ignition and helium-burning lifetime of HB stars, resulting in $\epsilon \leqslant 2 \times 10^{-14}$ for m_ϵ below a few keV.

Without going into detail, let us stress that all these bounds on the effective-action parameters depend on the implicit assumption that the effective actions hold equally well at solar scales as well as in the laboratory. But whereas solar processes typically involve momentum transfers on the keV scale, laboratory quantum vacuum experiments operate with much lower momentum transfers, a typical scale being μeV. In other words, the above bounds can only be applied to laboratory experiments, if one accepts an extrapolation of the underlying model over nine orders of magnitude. In fact, it has been shown quantitatively how the above-mentioned bounds have to be relaxed, once a possible dependence of these effective-action parameters, e.g., on momentum transfer, temperature, density or other ambient-medium parameters, is taken into account [22]. This observation indeed provides for another strong imperative to perform well-controlled laboratory experiments.

Previous laboratory experiments have also produced more direct constraints on the effective action parameters. For instance, the best laboratory bounds on MCPs previously came from limits on the branching fraction of ortho-positronium decay or the Lamb shift [21, 23, 24], resulting in $\epsilon \lesssim 10^{-4}$. Similarly, pure laboratory bounds on ALP parameters used to be much weaker than those from astrophysical arguments.

3. From optical experiments to fundamental particle properties

A variety of quantum vacuum experiments is devoted to a study of optical properties of modified quantum vacua. The BFRT experiment [25] has pioneered this field by providing upper bounds on vacuum-magnetically induced ellipticity, rotation as well as photon regeneration. Improved bounds for ellipticity and rotation have recently been published by the PVLAS collaboration [26][2]. Further polarization experiments such as Q&A [28] and BMV [29] have also already taken and published data.

The PVLAS experiment uses an optical laser ($\lambda = 1064$ nm and 532 nm) which is locked to a high-finesse Fabry–Perot cavity ($N = \mathcal{O}(10^5)$) and traverses an $L = 1$ m long magnetic field of up to $B = 5.5$ Tesla. Owing to the high finesse, the optical path length ℓ inside the magnet effectively increases up to several tens of kilometers.

The improved PVLAS bounds for ellipticity and rotation can directly be translated into bounds on the refractive-index and absorption-coefficient differences, $\Delta n = n_\parallel - n_\perp$,

$$|\Delta n(B = 2.3\ \mathrm{T})| \leqslant 1.1 \times 10^{-19}/\mathrm{pass}, \qquad |\Delta\kappa(B = 5.5\ \mathrm{T})| \leqslant 5.4 \times 10^{-15}\ \mathrm{cm}^{-1}. \tag{9}$$

As an illustration, an absorption coefficient of this order of magnitude would correspond to a photon mean free path in the magnetic field of the order of a hundred times the distance from earth to sun, demonstrating the quality of these laboratory experiments.

These bounds imply new constraints, e.g., for the ALP parameters, $g \simeq 4 \times 10^{-7}$ GeV^{-1} for $m_\phi < 1$ meV. More importantly, for MCPs, we find $\epsilon \lesssim 3 \times 10^{-7}$ for $m_\epsilon < 30$ meV. This bound is indeed of a similar size as a cosmological MCP bound which has recently been derived from a conservative estimate of the distortion of the energy spectrum of the cosmic microwave background [30]. Hence, laboratory experiments begin to enter the parameter regime which has previously been accessible only to cosmological and astrophysical considerations.

[2] The new data are no longer compatible with the PVLAS rotation signal reported earlier [27]. Nevertheless, this artifact deserves the merit of having triggered the physics-wise well-justified rapid evolution of the field which we are currently witnessing.

Imagine that an anomalously large signal, say, for ellipticity ψ and rotation $\Delta\theta$ is observed by such a polarization experiment, thereby providing evidence for vacuum-magnetic birefringence and dichroism. How could we extract information about the nature of the underlying particle-physics degree of freedom? The two data points for ψ and $\Delta\theta$ can be translated into parameter pairs g and m_ϕ for ALPs, ϵ and m_ϵ for MCPs, etc., leaving open many possibilities. As already mentioned above, a characteristic feature is the sign of ψ and $\Delta\theta$; that is, identifying the polarization modes $\parallel$ or $\perp$ as fast or slow modes reveals information about the microscopic properties [15]: e.g., a pseudo-scalar ALP goes along with $\Delta\kappa$, $\Delta n > 0$, whereas a scalar ALP requires $\Delta\kappa$, $\Delta n < 0$. A mixed combination, say $\Delta\kappa > 0$, $\Delta n < 0$, would completely rule out an ALP, leaving a spinor MCP as an option.

Another test would be provided by varying the experimental parameters such as length or strength of the magnetic field, or the laser frequency [15]. For instance, an ALP-induced rotation exhibits a simple B^2 dependence and the nonperturbative nature of MCP-induced rotation results in a $B^{2/3}$ law, cf equations (4), (7).

The underlying degree of freedom can more directly be identified by special-purpose experiments that probe a specific property of particle candidate. The light-shining-through-walls experiment is an example for such a setup. A magnetically induced photon regeneration signal in such an experiment would clearly point to a weakly interacting ALP degree of freedom; the outgoing photon polarization would distinguish between scalar ($\perp$ mode) and pseudo-scalar ($\parallel$ mode) ALPs. For this reason, a number of light-shining-through-walls experiments is currently being built or already taking data: ALPS at DESY [31], LIPSS at JLab [32], OSQAR at CERN[3] and GammeV at Fermilab[4]. PVLAS will shortly be upgraded accordingly, and BMV and GammeV have already published first results, yielding new bounds: $g \lesssim 1.3 \times 10^{-6}$ GeV^{-1} for $m_\phi \lesssim 2$ meV (BMV [29]) and $g \lesssim 3.2 \times 10^{-7}$ GeV^{-1} for $m_\phi \lesssim 1$ meV (GammeV [33]).

MCPs do not contribute to a photon regeneration signal, since pair-produced MCPs inside the magnet are unlikely to recombine behind the wall and produce a photon[5]. A special-purpose quantum vacuum experiment for MCP production and detection has been suggested in [34]: a strong electric field, e.g., inside an RF cavity can produce an MCP dark current by means of the nonperturbative Schwinger mechanism [1, 17]. A first signature could be provided by an anomalous fall-off of the cavity quality factor (the achievable high-quality factor of TESLA cavities already implies the bound $\epsilon \lesssim 10^{-6}$ [34]). Owing to the weak interaction, the MCP current can pass through a wall where a dark current detector can actively look for a signal.

In the case of a strongly interacting ALP, photon regeneration behind a wall would not happen either, since the wall would block both photons and generated ALPs. A special example is given by chameleon models which have been developed in the context of cosmological scalar fields and the fifth-force problem [11]. Somewhat simplified, chameleons can be viewed as ALPs with a varying mass that increases with the ambient matter density. As a result, low-energy chameleons which are initially produced in vacuo by photon conversion in a magnetic field cannot penetrate the end caps of a vacuum chamber and are reflected back into the chamber. After an initial laser pulse, the chameleons can again be reconverted into photons inside the magnetized vacuum; this would result in an afterglow phenomenon which is characteristic for a chameleonic ALP [35]. First estimates indicate that the chameleon parameter range accessible to available laboratory technology is comparable to scales familiar from astrophysical stellar energy-loss arguments, i.e. up to $g^{-1} \sim 10^{10}$ GeV for $m_\phi \lesssim 1$ meV. Afterglow measurements are already planned at ALPS [31] and GammeV [33].

[3] P Pugnat *et al* CERN-SPSC-2006-035, see http://graybook.cern.ch/programmes/experiments/OSQAR.html.

[4] See http://gammev.fnal.gov/.

[5] Photon regeneration can still be a decisive signal for models with both MCPs and paraphotons [18].

In the near future, quantum vacuum experiments could also be realized that involve strong fields generated by a high-intensity laser; for a concrete proposal aiming at vacuum birefringence, see [36] and also [37]. As major differences, laser-driven setups can generate field strengths that exceed conventional laboratory fields by several orders of magnitude. The price to be paid is that the spatial extent of these high fields is limited to a few microns. We expect that laser-driven experiments can significantly contribute to MCP searches in the intermediate-mass range whereas ALP and paraphoton searches, which are based on a coherence phenomenon, typically require a spatially sizeable field.

Both spatially extended and strong fields are indeed available in the vicinity of certain compact astrophysical objects. Also cosmic magnetic fields though weak may be useful due to their extreme spatial extent. For suggestions how to exploit these fields as a probe for fundamental physics, see, e.g., [38–40].

4. Microscopic models

So far, we argued that quantum vacuum experiments do not only serve as a probe for fundamental physics and BSM extensions, but also are required to provide for model-independent information about potential weakly coupled light degrees of freedom. Nevertheless, in the case of a positive anomalous experimental signal a puzzle of how to reconcile this signal with astrophysical bounds would persist on the basis of the low-energy effective actions discussed above. A resolution of this puzzle has to come from the underlying microscopic theory that interconnects solar scales with laboratory scales.

A number of ideas have come up to separate solar physics from laboratory physics; for a selection of examples, see [11, 41–45]. A general feature of many ideas is to suppress the coupling between photons and the new particle candidates at solar scales by a parameter of the solar environment such as temperature, energy or momentum transfer, or ambient matter density. A somewhat delicate alternative is provided by new particle candidates that are strongly interacting in the solar interior, resulting in a small mean free path (similar or smaller than that of the photons!), such that they do not contribute to the solar energy flux [43].

A paradigmatic example for a parametrical coupling suppression is given by the Masso–Redondo model [41] which, in addition to resolving the above puzzle, finds a natural embedding in string-theory models [46]. As a prerequisite, let us consider the paraphoton model of equation (8) and include a hidden-sector parafermion h which couples only to the paraphoton A' with charge e_{h} and interaction $e_{\mathrm{h}}\bar{h}\not{A}'h$. After the shift $A'_\mu \to \hat{A}'_\mu - \chi A_\mu$ which diagonalizes the kinetic terms, the parafermion acquires a coupling to our photon: $-\chi e_{\mathrm{h}}\bar{h}\not{A}h$. Since χ is expected to be small, we may identify $-\chi e_{\mathrm{h}} = \epsilon e$. As a result, the hidden-sector fermion appears as minicharged with respect to our photon. The bottom line is that a hidden sector with further $U(1)$ fields and correspondingly charged particles automatically appears as MCPs for our photon if these further $U(1)$'s mix weakly with our $U(1)$. However, if the paraphoton is massive the coupling of on-shell photons to parafermions is suppressed by this mass μ, since the on-shell condition cannot be met by the massive paraphoton.

The Masso–Redondo model now involves two paraphotons, one massless and one massive, with opposite charge assignments for the parafermions. The latter charge assignment indeed cancels the parafermion-to-photon coupling at high virtuality (as, e.g., for the photon plasma modes in the solar interior), implying that solar physics remains unaffected. At low virtualities such as in the laboratory, the massive paraphoton decouples which removes the cancellations between the two $U(1)$'s. A photon-paraphoton system is left over in which the parafermions indeed appear as MCPs with respect to electromagnetism. In this manner, the astrophysical bounds remain satisfied, but laboratory experiments could discover unexpectedly large anomalous signatures.

In fact, hidden sectors also involving further $U(1)$'s and correspondingly charged matter as required for the Masso–Redondo mechanism can not only be embedded naturally in more fundamental models, but also are often unavoidable in model building for reasons of consistency.

5. Conclusions

Quantum vacuum experiments such as those involving strong external fields can indeed probe fundamental physics. In particular, optical experiments can reach a high precision and thereby constitute an ideal tool for searching for the hidden sector of BSM extensions containing weakly interacting and potentially light degrees of freedom at the milli scale. A great deal of current experimental activity will soon provide for a substantial amount of new data which will complement particle-physics information obtained from accelerators.

From a theoretical viewpoint, many open problems require a better understanding of fluctuations of light degrees of freedom, the small mass of which often inhibits conventional perturbative ordering schemes. Modern quantum-field theory techniques for external-field problems such as the worldline approach [16, 17, 47] will have to be used and developed further hand in hand with experimental progress in probing the quantum vacuum.

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