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Strong laser fields as a probe for fundamental physics

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Abstract. Upcoming high-intensity laser systems will be able to probe the quantum-induced nonlinear regime of electrodynamics. So far unobserved QED phenomena such as the discovery of a nonlinear response of the quantum vacuum to macroscopic electromagnetic fields can become accessible. In addition, such laser systems provide for a flexible tool for investigating fundamental physics. Primary goals consist in verifying so far unobserved QED phenomena. Moreover, strong-field experiments can search 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 photon experiments in strong electromagnetic fields. The interaction of particle-physics candidates with photons and external fields can be parameterized by low-energy effective actions and typically predict characteristic optical signatures. I perform first estimates of the accessible new-physics parameter space of high-intensity laser facilities such as POLARIS and ELI.

PACS. 12.20.-m Quantum electrodynamics – 42.81.Gs Birefringence, polarization – 14.80.-j Other particles (including hypothetical)

1 Introduction

The superposition principle of classical electrodynamics is violated on the quantum level. Heisenberg's uncertainty principle allows for fluctuations of electron-positron pairs on top of the vacuum. As these fluctuations separate charges on a time and length scale on the order of the Compton wavelength, electromagnetic fields can couple to the charge fluctuations. On average, these vacuum-mediated interactions induce nonlinearities among the electromagnetic field itself, giving rise to nonlinear corrections to Maxwell's theory [1–3].

On a microscopic level, fluctuating electrons and positrons and their interaction with electromagnetic fields is described by quantum electrodynamics (QED). QED has been tested to an extremely high precision in atomic physics and at accelerator experiments. By contrast, the predicted violation of the superposition principle for macroscopic classical fields has not been verified so far. Sizable deviations from linearity require field strengths on the order of the energy scale set by the mass of the fluctuating particle, namely the electron mass: $B_{\rm cr} = m^2/e \simeq 4.3 \times 10^9$ T or $E_{\rm cr} \equiv B_{\rm cr} = 1.3 \times 10^{18}$ V/m (we use units such that $\hbar = c = 1$). This exceeds standard laboratory electromagnetic field strengths by far. For such fields with $E, B \ll B_{\rm cr}$, the nonlinear interactions to lowest order are governed by the Heisenberg-Euler effective action, $\Gamma_{\rm HE} = \int d^4x\, \mathcal{L}_{\rm HE}$. The corresponding effective Lagrangian is given by [1–3]

$$\mathcal{L}_{\rm HE} = -\mathcal{F} + \frac{8}{45}\frac{\alpha^2}{m^4}\mathcal{F}^2 + \frac{14}{45}\frac{\alpha^2}{m^4}\mathcal{G}^2, \tag{1}$$

where $\mathcal{F} = \frac{1}{4}F_{\mu\nu}F^{\mu\nu} = \frac{1}{2}(\mathbf{E}^2 - \mathbf{B}^2)/2$, $\mathcal{G} = \frac{1}{4}F_{\mu\nu}\widetilde{F}^{\mu\nu} = \mathbf{E}\cdot\mathbf{B}$. In addition to the Maxwell part, $\mathcal{L}_{\rm M} = -\mathcal{F}$, the four-photon interaction terms $\sim \mathcal{F}^2, \mathcal{G}^2$ arise from the average over the electron-positron fluctuations, see [4,5] for reviews. Apart from the suppression by the electron mass scale, nonlinear phenomena are also suppressed by factors of $\alpha \simeq 1/137$.

In order to observe the nonlinearities for weak fields in the laboratory, long interaction lengths or times are required. This scheme has been pioneered by the BFRT experiment [6] and more recently by PVLAS [7]. Here, an interaction of optical laser photons with magnetic fields of order $\mathcal{O}(1-10\ \mathrm{T})$ is searched for using cavity techniques in order to increase the optical path length up to $\mathcal{O}(10\ \mathrm{km})$. Even though QED nonlinearities are still a few orders of magnitude below the current sensitivity scale, these experiments have already probed a new window of particle physics, see below.

By contrast, upcoming high-intensity laser systems have the potential to get a more direct access to strong-field nonlinearities. Whereas the field intensity parameter $\nu = B^2/B_{\rm cr}^2 \sim 10^{-18}$ for BFRT or PVLAS, current petawatt lasers such as POLARIS at Jena [8] will reach up to $\nu \sim 10^{-7}$. At planned facilities such as ELI [9], a maximum of $\nu \sim 10^{-3}$ is expected. Since this drastic enhancement of field strengths also requires

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a strong focusing of the laser pulse, the possible interaction region is only on the order of $\mathcal{O}(10-100\ \mu\mathrm{m})$. A discussion of the benefits and disadvantages of laser-based experiments will be given in this article.

The discovery of the quantum-induced nonlinearities of electrodynamics would verify our most successful theory QED in a parameter region which has been little explored so far. It would complete a quest which has begun in the thirties. But beyond this, there is another strong motivation to investigate strong-field nonlinearities, since these experiments have a discovery potential of new fundamental physics. This is because the source of fluctuation-induced nonlinear self-interactions of strong electromagnetic fields in vacuum is not restricted to electrons and positrons. Any quantum degree of freedom that couples to photons can lead to modifications of Maxwell's electrodynamics. Strong-field experiments therefore also investigate the general field content of fluctuating particles in the quantum vacuum. If so far unknown particles mediate apparent photon self-interactions in a manner similar to electron-positron fluctuations, they can generate nonlinear corrections analogous to the Heisenberg-Euler Lagrangian (1). Moreover, such hypothetical particles could even be created from strong fields. As high-intensity lasers such as POLARIS and ELI will substantially push the frontier of strong fields available in a laboratory, they have the potential to search directly and indirectly for new fundamental particles.

As the scale set, for instance, by the ELI peak field strength is expected to be of order $\mathcal{O}(100\ \mathrm{keV})$ in particle-mass/energy units, a typical strong-field experiment will be sensitive to particle masses up to this scale and particularly to much lower scales. On the other hand, the high field strength together with modern optical techniques provides for a strong handle on very weakly coupled particles. With this particular sensitivity to potentially light but weakly coupled degrees of freedom, strong-field experiments are complementary to accelerator searches for new particles [10].

Indeed a number of extensions of the *standard model of particle physics* predict the existence of weakly interacting sub-eV particles (WISPs) which couple to the electromagnetic sector. A popular candidate is the axion [11] which provides for a possible solution of the strong CP problem; more generally, we can think of axion-like particles (ALPs) as an uncharged scalar or pseudo-scalar degree of freedom with a coupling to two photons. Further candidates are mini-charged particles (MCPs), i.e., matter fields with charge ϵe and $\epsilon \ll 1$, which arise naturally in scenarios with gauge-kinetic mixing [12] or extra-dimensional scenarios [13]. More generally, many standard model extensions not only involve but often require – 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. Hence, a discovery of hidden-sector properties could decisively single out the relevant theoretical fundament.

In this article, we summarize both our well-founded QED expectations for nonlinear phenomena as well as well-motivated speculations of possible signatures for new physics in strong-field experiments. By performing first estimates of the parameter region accessible to high-intensity laser systems, we will argue that such systems can become an important source of information for fundamental physics both within and beyond the standard model of particle physics.

2 Low-energy effective actions

From a bottom-up viewpoint, QED as well as many extensions of the standard model of particle physics lead to similar consequences for low-energy laboratory experiments. These are parameterizable by effective nonlinear interactions such as equation (1) or effective couplings between photons and the new effective degrees of freedom. In the following, we summarize a set of generic low-energy effective actions as well as typical observables for optical experiments involving strong fields.

2.1 Heisenberg-Euler effective action

The first example is given by classic QED, effectively summarized by the Heisenberg-Euler effective action equation (1) governing the dynamics of macroscopic laboratory fields. This action gives rise to a number of non-classical phenomena, most prominently vacuum-electromagnetic birefringence [4,14,15]. The quantum-modified equations of motion resulting from equation (1) yield

$$0 = \partial_\mu \left(F^{\mu\nu} - \frac{4}{45}\frac{\alpha^2}{m^4} F^{\alpha\beta} F_{\alpha\beta} F^{\mu\nu} - \frac{7}{45}\frac{\alpha^2}{m^4} F^{\alpha\beta} F_{\alpha\beta} \widetilde{F}^{\mu\nu} \right). \tag{2}$$

In comparison to Maxwell's equation $\partial_\mu F^{\mu\nu} = 0$, this equation no longer admits plane wave solutions traveling at the speed of light. But for a probe beam with small amplitude propagating in a strong background field, we may linearize the field equation and solve for the dispersion relation of the probe beam. For instance for a strong background magnetic field B, the phase and group velocities of the probe field satisfy

$$v_\parallel \simeq 1 - \frac{14}{45}\frac{\alpha^2}{m^4} B^2 \sin^2\theta_B, \quad v_\perp \simeq 1 - \frac{8}{45}\frac{\alpha^2}{m^4} B^2 \sin^2\theta_B, \tag{3}$$

where θ_B is the angle between the propagation direction and the magnetic field. The indices $\parallel$ and $\perp$ distinguish the two polarization modes, where the polarization of the probe beam is in or perpendicular to the plane spanned by the propagation direction and the B field. For a probe field in a counter-propagating laser field, $B^2 \sin^2\theta_B$ has to be replaced by the laser intensity I [16].

The vacuum modified by a strong external field thus is birefringent in a manner similar to a uniaxial crystal with refractive indices $n_{\parallel,\perp} = 1/v_{\parallel,\perp}$. As an observable, an initially linearly polarized probe laser can pick up an *ellipticity* by traversing the strong beam. The ellipticity angle ψ is given by

$$\psi = \frac{\omega}{2} L \Delta n \sin 2\theta, \tag{4}$$

where θ is the angle between the probe polarization and the fast eigenmode's polarization, $\Delta n = n_\parallel - n_\perp$, and L is the path length inside the strong field.

Equation (4) can teach a lot about the specific advantages and disadvantages of the various possible setups. Classic strong-field experiments such as BFRT or PVLAS at comparatively low field strengths need to detect a tiny refractive-index difference on the order of $\Delta n \simeq 10^{-22}$. Using optical probe photons, e.g., $\omega \simeq 1\ \mu$m, high-finesse interferometry can increase the optical path length to up to $L \simeq \mathcal{O}(10\ \text{km})$. By contrast, high-intensity lasers such as POLARIS can reach up to $\Delta n \simeq 10^{-12}$, but are restricted to optical path lengths of order $L \simeq \mathcal{O}(10\ \mu\text{m})$ dictated by focusing as close to the diffraction limit as possible. Since cavity techniques are useless anyway for short-pulsed high-intensity fields, a higher probe frequency ω can be used in order to enhance the ellipticity signal, as proposed in [16].

For instance, for X-ray photons with $\omega \simeq 1$keV probing a POLARIS high-intensity field, the induced ellipticity angle is given by $\psi \simeq 6 \times 10^{-7}$. Higher frequencies such as $\omega \simeq 12$ keV even yield $\psi \simeq 7 \times 10^{-6}$. This should be compared to the sensitivity scale of an X-ray polarization measurement required for detecting ellipticity. Theoretical estimates predict that sensitivity bounds on the order of $\psi \simeq 3 \times 10^{-6}$ should be measurable with high-precision techniques for these frequencies [17]. Therefore, quantum-induced nonlinearities of electromagnetic fields may already be discovered at Peta-Watt lasers such as POLARIS. Indeed, even the required synchronized X-ray beam can be generated via laser-driven electron acceleration and subsequent Thomson backscattering with the same system.

Another optical observable can be important: any effect which modifies the amplitudes of the $\parallel$ or $\perp$ components in a polarization-dependent manner but leaves the phase relations invariant will induce 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. Further possibly rotation inducing effects such as photon splitting [15] or neutrino-pair production [18] in a strong field are severely suppressed for typical laboratory parameters. Therefore, a sizeable signal for vacuum rotation $\Delta\theta$ in a strong-field experiment would be a signature for new fundamental physics. (Note, however, that rotation can also be generated for pure kinematical reasons, depending on the details of the optical setup [19].)

Let us close this QED part by stressing, that we have concentrated on nonlinear QED phenomena with external fields as sole asymptotic states. Another very important and interesting nonlinear phenomenon is the spontaneous decay of the electromagnetized vacuum itself into electron-positron pairs. This Schwinger pair production is technically related to the imaginary part of the effective action $\mathrm{Im}\,\Gamma$ [1–3]. In addition to being nonlinear, it is also nonperturbative in the coupling to the, say, electric background field eE, see [5,20] for a review. An experimental verification of this phenomenon would therefore explore a particularly interesting and incompletely understood branch of quantum field theory. For subcritical fields, the pair production rate is unfortunately exponentially suppressed $\sim \exp(-\pi E_{\mathrm{cr}}/E)$ according to a constant-field approximation. Time-dependent fields can enhance the production rate significantly [5,20,21]. Recent results indeed indicate that the exponential suppression might be overcome with time-dependent tailored pulses even for ELI parameters [22].

2.2 Axion-like particle (ALP)

As a first example of a new particle candidate beyond those degrees of freedom of the standard model, we consider a new neutral scalar ϕ or pseudo-scalar degree of freedom ϕ^- such as an axion which is coupled to the photon by,

$$\mathcal{L}_{\mathrm{ALP}} = \left\{ -\frac{g}{4}\phi^{(-)} F^{\mu\nu} \overset{(\sim)}{F}_{\mu\nu} - \frac{1}{2}(\partial\phi^{(-)})^2 - \frac{1}{2}{m_\phi}^2 \phi^{(-)2} \right\}, \tag{5}$$

parameterized by the mass m_ϕ of this axion-like particle (ALP) and the dimensionful coupling g. In optical experiments in strong fields, ALPs can induce both ellipticity and rotation [23], since only one polarization mode couples to the ALP and the strong field. For instance, coherent photon-ALP conversion causes a depletion of one photon mode, implying rotation. In order to make contact with the literature, we approximate the strong field by a homogeneous magnetic field B as may be provided by a slowly beating standing wave formed from counter-propagating laser beams. We stress that detailed studies employing all relevant properties of the field provided by systems such as ELI still need to be performed. From the equations of motion for the photon-ALP system for the pseudo-scalar case, the induced ellipticity and rotation can be calculated:

$$\begin{aligned} \psi^- &= \frac{1}{2}\left(\frac{gB\omega}{m_\phi^2}\right)^2 \left(\frac{Lm_\phi^2}{2\omega} - \sin\left(\frac{Lm_\phi^2}{2\omega}\right)\right)\sin 2\theta, \\ \Delta\theta^- &= \left(\frac{gB\omega}{m_\phi^2}\right)^2 \sin^2\left(\frac{Lm_\phi^2}{4\omega}\right)\sin 2\theta, \end{aligned} \tag{6}$$

for single passes of a probe beam through a strong B field of length L. For the scalar ALP, we have $\Delta\theta = -\Delta\theta^-$, $\psi = -\psi^-$. Measuring ellipticity and rotation signals uniquely determines the two model parameters, ALP mass m_ϕ and ALP-photon coupling g. Measuring the signs of $\Delta\theta$ and ψ can even resolve the parity of the involved particle [24].

The effective interaction (5) is representative for various underlying particle scenarios. In the axion case, only the weak coupling to the photon is relevant and all other potential matter couplings are negligible. This facilitates the interesting experimental option to shine the ALP component through a wall which blocks all photons. Behind the wall, a second strong field can induce the reverse process and photons can be regenerated out of the ALP beam [25]. The regeneration rate is

$$n_{\rm out} = n_{\rm in}\,\frac{1}{16}\,(gBL\cos\theta)^4 \left[\sin\left(\frac{Lm_\phi^2}{4\omega}\right)\Big/\frac{Lm_\phi^2}{4\omega}\right]^4, \quad (7)$$

where $n_{\rm in}$ is the initial photon rate, and the fields B and its extension L are assumed to be identical on both sides of the wall.

Further models with ALPs have been proposed in cosmology. Cosmological scalar fields are discussed, e.g., in the context of dark energy, the cosmological coincidence problem and also dark-matter abundance. An interesting candidate also for optical experiments are those scalar fields with a chameleon mechanism which have been developed in the context of the fifth-force problem [26]. As an interesting property, a chameleonic ALP cannot penetrate the end caps of the vacuum chamber but gets reflected back. Whereas this has no influence on the formulas for ψ and $\Delta\theta$ in equation (6), a new detection mechanism arises: synchronizing a short laser probe pulse with the strong pulse, chameleons can be created inside the vacuum chamber and stored in a parallel cavity. By a synchronized second strong pulse, the chameleons can be re-converted into photons again inside the strong field; this would result in an afterglow phenomenon which is characteristic for a chameleonic ALP [27]. In the parameter range where $gB/m_\phi \ll 1$, the number of photons in the first afterglow pulse $n_{\rm out}$ is again given by equation (7) where this time $n_{\rm in}$ is the number of photons in the synchronized probe pulse initially generating the chameleons.

2.3 Minicharged particle (MCP)

In addition to neutral particles coupling to photons by dimensionful coupling constants such as ALPs, new particle candidates can also be charged. In order to have evaded detection so far, they have to be either very heavy or very weakly charged. In the latter case, they can also be very light and thus become ideal candidates for laser-based searches. These so-called minicharged particles (MCPs) arise naturally in scenarios with gauge-kinetic mixing [12] or extra-dimensional scenarios [13], and find a natural embedding in string-theory models with intermediate string scale [28]. The latter property makes optical searches for MCPs particularly attractive, because a possible optical signal at low energies could already help singling out the relevant class of microscopic highest-energy models.

From the bottom-up viewpoint of effective photon interactions and nonlinearities, the fluctuations of minicharged particles with mass m_ϵ and charge ϵe, $\epsilon \ll 1$ induce photon self-interactions in the same way as electrons do. However, the weak-field Heisenberg-Euler Lagrangian (1) is not sufficient to describe the physics of MCPs properly, as the expansion parameters $\epsilon eB/m_\epsilon^2$ and ω/m_ϵ in a strong B field varying with ω are not necessarily small. As an interesting consequence, the probe laser frequency can be above the pair-production threshold $\omega > 2m_\epsilon$ such that a rotation signal in addition to birefringence-induced ellipticity becomes possible [29].

All relevant information is encoded in the polarization tensor which is well known from QED [4,14,30]. Explicit results are available in asymptotic limits, e.g., for the rotation signal induced by a Dirac-fermionic MCP [29],

$$\Delta\theta \simeq \frac{1}{12}\frac{\pi}{\Gamma(\frac{1}{6})\Gamma(\frac{13}{6})}\left(\frac{2}{3}\right)^{\frac{1}{3}}\epsilon^2\alpha(m_\epsilon L)\left(\frac{m_\epsilon}{\omega}\right)^{\frac{1}{3}}\left(\frac{\epsilon eB}{{m_\epsilon}^2}\right)^{\frac{2}{3}},$$
$$\text{for } \frac{3}{2}\frac{\omega}{m_\epsilon}\frac{\epsilon eB}{m_\epsilon^2} \gg 1, \quad (8)$$

which is valid above threshold and for a high number of allowed MCP Landau levels. Similar formulas exist for ellipticity or the case of spin-0 MCPs [24]. Note that this rotation appears to become independent of m_ϵ in the small-mass limit. In practice, once the associated Compton wavelength $\sim 1/m_\epsilon$ becomes larger than the size of the strong field, the field size acts as a cutoff reducing the effect. Precise predictions then require computations of polarization tensors in inhomogeneous fields which is a challenge for standard methods and remains an interesting question for future research.

3 New-physics sensitivity scales of strong-field experiments

In order to put the capabilities of the particle-physics potential of optical experiments with high-intensity lasers into a greater context, let us draw a comparison with other currently performed optical experiments, such as PVLAS [7], BMV [31], ALPS [32], LIPSS [33], OSQAR [34], GammeV [35,36]. In all these experiments, optical probe lasers traverse a magnetic field of $\mathcal{O}(1{-}10$ T) and length $\mathcal{O}(1{-}10$ m). Whereas the reachable field strengths are comparatively small, e.g., if measured in units of the QED critical field strength of $B_{\rm cr} \simeq 4\times10^9$ T, the length of the interaction region is macroscopic. As already discussed above in the context of QED birefringence, the latter can even be enhanced by placing the field into a high-finesse cavity such that the signal is increased by a factor $N_{\rm pass}$ counting the number of passes of the probe laser inside the cavity.

By contrast, high-intensity laser systems provide for an interaction region only of the order of $\mathcal{O}(10{-}100\ \mu$m); also cavities are of no use, since pulse durations on the femtosecond scale are far too short compared to the time scale for a multiple pass. Nevertheless, the extreme intensity can compensate for these disadvantages. In the following, we base our estimates on a reference scenario with a peak intensity of $I = 10^{27}$ W/cm^2 and a laser focus spot size of $L \simeq 50\ \mu$m [37]. Although these are optimistic values even for ELI, the technological developments are rather rapid these days, such that these parameters may be reliably accessible by the time when ELI is operating in a stable manner. Note that recent ideas on higher-harmonic focusing by oscillating plasma mirrors or schemes based on relativistically flying mirrors might lead to even higher intensities [38].

Let us first concentrate on ALP scenarios, focusing on a parameter range satisfying $Lm_\phi^2/\omega \ll 1$ (sub-eV particle masses). Then the relevant combined dimensionless parameter is gBL. For instance for PVLAS, this parameter is $gBL|_{\rm PVLAS} \simeq 5 \times (g/\mathrm{GeV}^{-1})$; the other magnet-based experiments mentioned above lie in a similar ball park. As a result, typical bounds on the coupling g, (e.g., resulting from the nonobservation of photon regeneration behind a wall) are in the range of $g \lesssim 10^{-5} \dots 10^{-6}$ GeV^{-1} for sub-eV masses m_ϕ. The corresponding ELI parameters are expected to give

$$gBL = 3.3 \times 10^3 \, \frac{g}{[\mathrm{GeV}^{-1}]} \sqrt{\frac{I}{[10^{27}\ \frac{\mathrm{W}}{\mathrm{cm}^2}]}} \, \frac{L}{[50\,\mu\mathrm{m}]}, \tag{9}$$

yielding a prefactor which exceeds that of magnet-based optical experiments by up to 3 orders of magnitude.

However, this improvement does not directly translate into a comparable increase of sensitivity, due to the lack of cavity enhancements and the necessity of pulse-probe synchronization. As an example estimate, let us consider a regeneration or afterglow experiment (7), using a probe laser of the Peta-Watt class delivering $\sim 10^{21}$ photons per shot. Alternatively, a fraction of the ELI beam could be coupled out of the beam and successively be used as a probe beam. The latter setup would also be advantageous for issues of pulse synchronization.

Assuming single-photon detection per pulse behind the wall or in the afterglow, the sensitivity range for the ALP-photon coupling g yields,

$$\frac{g}{[\mathrm{GeV}^{-1}]} \gtrsim 3.4 \times 10^{-9} \sqrt{\frac{[10^{27}\ \frac{\mathrm{W}}{\mathrm{cm}^2}]}{I}} \, \frac{[50\,\mu\mathrm{m}]}{L} \times \sqrt[4]{\frac{10^{21}[\mathrm{pulse}^{-1}]}{n_{\rm in}}} \tag{10}$$

for ALP masses in the sub-eV range. This should be compared with the current best laboratory bounds excluding ALP couplings of $g \gtrsim 10^{-6}$ GeV^{-1} or chameleonic couplings $g \gtrsim 2.5 \times 10^{-7}$ GeV^{-1}. Also ALP rotation and ellipticity signals could be enhanced in comparison with standard optical experiments, but the potential improvement of ALP parameter bounds might not be as dramatic as from regeneration or afterglow experiments. In any case, we conclude that ELI has the potential to significantly improve existing laboratory bounds for standard model extensions involving ALPs.

Let us turn to the MCP case. From equation (8), we deduce that ELI may yield the following maximum rotation $\delta\theta$ of the polarization axis of a probe beam (at $\theta = \pi/2$):

$$\delta\theta = 4.1 \times 10^8 \epsilon^{8/3} \left(\frac{I}{[10^{27}\ \frac{\mathrm{W}}{\mathrm{cm}^2}]} \right)^{4/3} \left(\frac{\mathrm{eV}}{\omega} \right)^{1/3} \frac{L}{[50\,\mu\mathrm{m}]}. \tag{11}$$

Assuming a detection sensitivity of $\delta\theta|_{\rm sens} \simeq 10$ nrad, ELI will be sensitive to minicharge couplings down to $\epsilon \gtrsim \mathcal{O}(10^{-7})$ for optical probe lasers and sub-eV MCP masses. This is of the same order of magnitude as the current best laboratory bounds from PVLAS [7,39] and in the same ball park as cosmological observations [40].

4 Conclusions

The prospect of high-intensity laser systems being currently worldwide under intense development is a strong motivation for reconsidering aspects of fundamental physics in strong fields. This subject had early been initiated even before the full advent of quantum field theory. In particular, the long-standing prediction of quantum-induced nonlinear self-interactions of macroscopic magnetic fields by Heisenberg, Euler and others is still awaiting its experimental verification. High-intensity laser systems are good candidates for completing this quest.

In addition to confirming our expectations about nonlinearities induced by known degrees of freedom of the standard model of particle physics, strong-field experiments have recently proved very useful to explore new regions in the parameter space of hypothetical new-physics degrees of freedom. Strong fields in combination with optical probes have turned out to be particularly sensitive to weakly coupled particles with light masses in the sub-eV range. High-intensity lasers can add a new chapter to this story by giving experimental access to unprecedented field-strength values. In this work, we have argued that such laser systems may not only discover the QED-induced nonlinearities for macroscopic fields for the first time, but also search for unexpected optical signals that would point to new particle candidates such as axion-like particles (ALPs) or minicharged particles (MCPs). In particular for ALPs, optical probing of strong fields based on the light-shining-through-wall scheme or the afterglow mechanism can exceed current bounds by 2 or 3 orders of magnitude.

It is worthwhile to emphasize that optical experiments typically test a regime characterized by momentum transfers below the eV scale. This clearly distinguishes them from experiments looking for astrophysical bounds. Nevertheless, astrophysical observations in combination with energy-loss arguments impose strong constraints, e.g., on the ALP coupling $g \lesssim 10^{-10}$ GeV^{-1} for ALP masses in the eV range and below [41], or on MCP couplings $\epsilon \leq 2 \times 10^{-14}$ for m_ϵ below a few keV [42]. However, since the underlying solar physics involves keV momentum transfer scales, these bounds apply to laboratory transfer scales ($\sim\mu$eV) only if the coupling values are extrapolated over these many orders of magnitude [43]. It is precisely this assumption which has been put into question by various models [26,44–48] and which can be checked or falsified by a particle discovery at strong-field experiments such as POLARIS or ELI. Indeed, current strong-field laboratory experiments begin to enter the parameter regime which has previously been accessible only to cosmological and astrophysical considerations [39].

For the special case of scalar ALPs, it has been pointed out [49,50] that the scalar-ALP-photon coupling can be severely constrained by direct searches for non-Newtonian forces [51]. For instance, in simple models the bounds on the allowed parameter region can reach up to $g \lesssim 1.6 \times 10^{-17}$ GeV^{-1}. (The bounds on chameleonic ALP couplings are somewhat relaxed owing to the skin-depth effect [26], but a significant parameter range is excluded.) Since gravitational or fifth-force experiments measure ALP-matter interactions, the bounds rely on implicit assumptions on additional ALP-matter couplings. Typically these interactions are assumed to be generated by fluctuations from the ALP-photon coupling. However, if also microscopic ALP-matter couplings are present, these experiments only measure a sum of both microscopic and photon-induced ALP-matter interactions, the combination of which could be accidentally small. Optical measurements in high-intensity laser systems can therefore complement these bounds and exclude this loophole, as they are directly sensitive to the ALP-photon interaction.

In case of a positive signal, high-intensity lasers could not only discover a new particle but also contribute to the particle's identification. Whereas the field strength, length and frequency dependence can distinguish between ALPs or MCPs, the signs of ellipticity and rotation are characteristic for spin and parity [24]. Light-shining-through-wall or afterglow experiments are indicative for additional matter couplings. Further experiments have been suggested such as Schwinger-type MCP pair production [52] or hidden-photon searches [53] which may also become realizable at ELI.

It should nevertheless be stressed once more that all the above estimates are based on various approximations. In particular, the homogeneous-field assumption is questionable as the typical variation scale of the strong field can be of the same order of magnitude as the new particle's Compton wavelength. More detailed theoretical analyses are certainly required for precise estimates, and new unknown and surprising effects may arise from this interesting equality of scales.

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