Minicharged particles in light-shining-through-wall-experiments and the photon polarization tensor

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A first theoretical feasibility study for a novel "light-shining-through-wall" scenario in the presence of an external magnetic field is performed. In contrast to standard scenarios, the barrier is not traversed by means of weakly interacting on-shell particles, but by virtual minicharged particle-antiparticle states. The study of this process heavily relies on the knowledge of the photon polarization tensor in the non-perturbative regime, and in particular requires its full momentum dependence, thereby rendering conventional approximations inapplicable. A first study and its results are presented and discussed in this contribution.

1 LSW with virtual particles - a motivation

"Light-shining-through-wall" (LSW) experiments are a versatile means in the search for numerous theoretically well-motivated "weakly interacting slim particles" (WISPs) that are proposed to exist beyond the standard model of particle physics. In standard LSW scenarios the barrier is 'tunneled' by *real*, i.e., on-shell WISPs. Their paradigm is the LSW scenario with axions or, more generally, axion-like particles (ALPs). LSW with axion-like particles is possible if the laser probe photons are converted into real ALPs in front of the wall and reconverted into photons behind that wall.

Intriguingly, LSW setups, although originally aimed at the detection of axions, are also sensitive to other WISPs, in particular to hidden photons and minicharged particles (MCPs); see [1, 2] and references therein for an overview and recent experimental results. For example, if MCPs and hidden photons exist, "light-shining-through-wall" is also possible by means of a real hidden photon: Photons can be converted into hidden photons through an intermediate MCP loop within an external magnetic field. Similar to the LSW scenario with ALPs, the hidden photons are then assumed to traverse the barrier unhindered and can thereafter be reconverted into photons [3]. Note however that this LSW scenario can only provide combined bounds on the fractional charge of the minicharged particles and the hidden-photon coupling.

A different LSW scenario for MCPs proposed recently in [5], is the 'tunneling' of a barrier through generally *virtual* particle-antiparticle intermediate states, cf. Fig. 1. This scenario, also referred to as "tunneling of the third kind", is interesting from an experimental point of view since it could provide for *direct* bounds on the fractional charge of the MCPs, without any reference to the coupling strength of hidden photons. It has however been shown, that

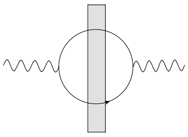

Figure 1: LSW scenario with virtual minicharged particles, also referred to as 'tunneling of the 3rd kind', cf. [5]. A spontaneous oscillation into a minicharged particle-antiparticle pair which traverses a light blocking barrier freely, enables the photon to effectively "shine through a wall".

- in the absence of an external field - bounds derived from this LSW scenario are typically less restrictive than current laboratory limits for minicharged particles [5]. In this note, we briefly discuss the same scenario in the presence of a constant external magnetic field. We limit the discussion to fermionic minicharged particles. ¿From the viewpoint of theory, tackling this problem in an external field is far more involved than the corresponding problem in the absence of a field. In particular, it turns out that new insights into the polarization tensor in external magnetic fields, which enters the computation of the photon-to-photon transition probability, are necessary. More detailed considerations and results are forthcoming [4].

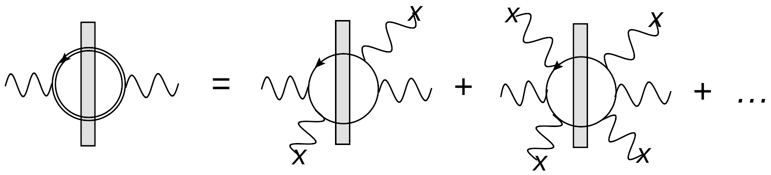

Figure 2: Same scenario as in Fig. 1, but now in the presence of an external magnetic field. The dressed propagator of the minicharged particles, accounting for all possible insertions of the external field, is represented by the double solid line. As argued in the main text, to make experimentally relevant predictions, a computational restriction to the lowest order perturbative corrections in the external field is insufficient.

Let us briefly introduce the basic equations which are needed to compute the transition probability for the tunneling scenario depicted in Fig. 2. The wall is assumed to be perpendicular to the propagation direction of the photons. A detailed presentation for the zero-field case can be found in [5]. The essential point to note here, is that the dressed particle-antiparticle loop traversing the wall in Fig. 2 corresponds to the one-loop photon polarization tensor in an external magnetic field, $\Pi^{\mu\nu}(x,x'|B)$, with the quantum loop run by minicharged particles. Hence, we start with the effective field theory describing photon propagation in a constant external magnetic field of strength $B=|\vec{B}|$, given by the following Lagrangian,

$$\mathcal{L}[A] = -\frac{1}{4}F_{\mu\nu}(x)F^{\mu\nu}(x) - \frac{1}{2}\int_{x'} A_\mu(x)\Pi^{\mu\nu}(x,x'|B)A_\nu(x')\,, \tag{1}$$

where $F_{\mu\nu}$ denotes the field strength tensor of the classical, macroscopic photon field A_μ. As the minicharged particles traverse the wall unhindered, translational invariance is maintained on the level of the polarization tensor at one-loop order, as long as the B field is homogeneous in the relevant space-time region, implying $\Pi^{\mu\nu}(x,x'|B) = \Pi^{\mu\nu}(x-x'|B)$.

Hence, upon a variation of Eq. (1) and a transformation to momentum space, the following equation of motion is obtained

$$\left(k^2 g^{\mu\nu} - k^\mu k^\nu + \Pi^{\mu\nu}(k)\right) A_\nu(k) = 0\ . \tag{2}$$

We use a metric with signature $(-,+,+,+)$, i.e., $k^2 = \vec{k}^2 - \omega^2$.

In a next step, we impose reflecting boundary conditions at the wall for the incoming photons and determine the fluctuation induced current behind the wall with absorbing boundary conditions. The detector is assumed to be positioned asymptotically far from the back side of the wall. The photon-to-photon transition probability depends on the polarization mode of the photons. In the presence of an external field there are three independent polarization modes, henceforth labeled by an index $p = 1, 2, 3$. As the vacuum speed of light in external fields deviates from its zero-field value, and the vacuum exhibits medium-like properties, the occurrence of three (instead of two in the absence of an external field) independent polarization modes is not surprising.

Defining projectors $P_p^{\mu\nu}$ onto these modes, as done explicitly in Sect. 2 below, the photon-to-photon transition probability for photons polarized in the mode p, is given by

$$P_p^{\gamma\to\gamma} = \left| \int_d^\infty \mathrm{d}x' \, \frac{e^{-i\omega x'}}{2\omega} \int_{-\infty}^0 \mathrm{d}x'' \, P_p^{\mu\nu} \Pi_{\nu\mu}(x' - x'') \sin(\omega x'') \right|^2 , \tag{3}$$

with ω denoting the photon frequency, and d the thickness of the wall.

2 The photon polarization tensor and its approximations

Let us now briefly introduce the photon polarization tensor in the presence of an external field. We stick to its representation in the propertime formalism [6] at one-loop level. The corresponding expression is known for arbitrary homogeneous, externally set electromagnetic field configurations in terms of a double parameter integral [7, 8]. Here we limit ourselves to the special case of a purely magnetic field [9]. This naturally suggests a decomposition of the photon four-momentum k^μ in components parallel and perpendicular to the magnetic field vector $\vec{B}$. Without loss of generality $\vec{B}$ is assumed to point in $\vec{e}_1$-direction, and the following decomposition is adopted,

$$k^\mu = k_\parallel^\mu + k_\perp^\mu , \qquad k_\parallel^\mu = (\omega, k^1, 0, 0) , \qquad k_\perp^\mu = (0, 0, k^2, k^3) . \tag{4}$$

In the same manner, tensors can be decomposed, e.g., $g^{\mu\nu} = g_\parallel^{\mu\nu} + g_\perp^{\mu\nu}$. It is then convenient to introduce the following projection operators onto photon polarization modes,

$$P_1^{\mu\nu} = g_\parallel^{\mu\nu} - \frac{k_\parallel^\mu k_\parallel^\nu}{k_\parallel^2} , \qquad \text{and} \qquad P_2^{\mu\nu} = g_\perp^{\mu\nu} - \frac{k_\perp^\mu k_\perp^\nu}{k_\perp^2} . \tag{5}$$

Defining a third projector as follows,

$$P_3^{\mu\nu} = g^{\mu\nu} - \frac{k^\mu k^\nu}{k^2} - P_1^{\mu\nu} - P_2^{\mu\nu} , \tag{6}$$

the three projectors $P_p^{\mu\nu}$ obviously span the transverse subspace. Whereas two of these polarization modes can be continuously related to the photon polarization modes in the absence of an external magnetic field, the third mode manifests itself in the presence of the external field only.

Note that $P_1^{\mu\nu}$ and $P_2^{\mu\nu}$ have an intuitive interpretation, given that $\vec{k} \nparallel \vec{B}$. Namely, they project onto photon modes polarized parallel and perpendicular to the plane spanned by the two vectors, $\vec{k}$ and $\vec{B}$. For $\vec{k} \nparallel \vec{B}$ these are the polarization modes, that can be continuously related to those in the zero-field limit. Remarkably, for the special alignment of $\vec{k} \parallel \vec{B}$, the situation is different. Here, the modes 2 and 3 can be continuously related to the two zero-field polarization modes.

With the help of Eqs. (5) and (6), the photon polarization tensor in a purely magnetic field can be decomposed as follows [7],

$$\Pi^{\mu\nu}(k) = \Pi_1(k)\, P_1^{\mu\nu} + \Pi_2(k) P_2^{\mu\nu} + \Pi_3(k) P_3^{\mu\nu}\,, \tag{7}$$

where the scalar functions Π_p are the components of the polarization tensor in the respective subspaces. The coupling of the MCPs to photons is ϵe, with ϵ referring to a dimensionless fractional coupling, and m denotes the mass of the MCPs. Specifying $\epsilon \equiv 1$ and identifying m with the electron mass, the photon polarization tensor of standard quantum electrodynamics is retained. The general structure of the scalar components is

$$\Pi_p(k) = (\epsilon e)^2 \int_0^\infty \frac{\mathrm{d}s}{s} \int_{-1}^{+1} \frac{\mathrm{d}\nu}{2}\, e^{-i\Phi_0(k,\nu,z)s} f_p(k,\nu,z)\,, \tag{8}$$

with the dependence on the magnetic field encoded in the variable $z \equiv \epsilon e B s$. Here, the parameter s denotes the propertime, and the parameter ν governs the momentum distribution within the loop. The explicit expressions for the functions f_p can be found in [7], and the phase factor in the argument of the exponential function reads

$$\Phi_0 = m^2 + \frac{1-\nu^2}{4} k_\parallel^2 + \frac{\cos\nu z - \cos z}{2z\sin z} k_\perp^2\,. \tag{9}$$

In particular due to the complicated functional dependence on s in Eq. (9), the propertime integral in general cannot be performed analytically, and is also hard to tackle numerically. Hence, basically all explicit insights into the photon polarization tensor in the presence of a constant magnetic field can be traced back to three well-established approximations:

- a perturbative expansion in the number of external field insertions (cf. Fig. 2), which can be associated with the limit $\frac{\epsilon e B}{m^2} \ll 1$,
- a quasi-classical approximation [10] developed in the seminal works of Tsai and Erber [11, 12], derived "on-the-light-cone", i.e., for $k^2 = 0$, and restricted to $\frac{k_\perp^2}{\epsilon e B} \gg 1$ only [4], and
- the restriction to the lowest Landau level, or equivalently a "large-z" expansion [13, 14], valid in the limit where $\frac{\epsilon e B}{m^2} \gg 1$, and commonly utilized below pair-creation threshold, $\omega^2 < 4m^2$.

Concerning MCPs, neither their fractional charge ϵe, nor their mass m is restricted *a priori*, which means that in principle arbitrary values for the ratio $\frac{\epsilon e B}{m^2}$ are possible. Given that the zero-field bounds, and therewith also these for $\frac{\epsilon e B}{m^2} \ll 1$, fall into a parameter regime in the ϵ-m plane already excluded by means of other laboratory experiments, we subsequently aim at gaining insights into the parameter regime where $\frac{\epsilon e B}{m^2} \gg 1$. As the full momentum dependence is essential for the virtual tunneling process, none of the approximations listed above is applicable here.

However, reconsidering Eqs. (8) and (9), a significant simplification can be expected in the special situation where $\vec{k} \parallel \vec{B}$, which implies that $k_\perp^2 \equiv 0$. In this limit the z-dependence in Eq. (9) drops out, and the propertime integration can even be performed analytically [4]; see also [15, 16]. Combining the "large-z" expansion with an analytic continuation in B, we recently devised a strategy to surpass the pair creation threshold [4], indicating that the 1-component of the polarization tensor in the limit $\vec{k} \parallel \vec{B}$ indeed results in the maximum transition probability, achievable in the tunneling process with virtual MCPs in the regime where $\frac{\epsilon e B}{m^2} \gg 1$. This can be traced back to the fact that the lowest Landau level is unscreened only for this component.

Note, that the situation is somewhat subtle, as in the strict limit $\vec{k} \parallel \vec{B}$, exactly this component corresponds to the polarization mode that cannot be continuously related to any of the two polarization modes in the absence of an external field. The question if and how this particular mode can be excited is disputed in the literature [15, 16]. However, as discussed above, it turns out that for any non-vanishing angle between $\vec{k}$ and $\vec{B}$, the 1-mode can be continuously related to one of the photon polarization modes in the absence of an external field, but as a function of the angle between $\vec{k}$ and $\vec{B}$ receives an exponential suppression as compared to its maximum in the asymptotic limit $\sphericalangle(\vec{k}, \vec{B}) \to 0$.

3 Transition probability for LSW with virtual MCPs

Hence, in order to estimate the maximum achievable photon to photon transition probability in the regime where $\frac{\epsilon e B}{m^2} \gg 1$, we subsequently exclusively focus on the 1-component of the photon polarization tensor and limit ourselves to the special alignment $\vec{k} \parallel \vec{B}$. Let us emphasize that this is an essential difference to conventional LSW setups, where the magnetic field is applied perpendicular to the propagation direction of the photons.

We then find that the maximum achievable photon transition probability (cf. Eq. (3)) in the limits $md \ll 1$ and $m \ll \omega$ is well approximated by the following simple formula,

$$P_{\rm max}^{\gamma\to\gamma} \simeq \left(\frac{\epsilon^2 e^2}{24\pi^2}\right)^2 \left(\frac{\epsilon e B}{m^2}\right)^2 . \tag{10}$$

To understand the elementary dependencies of this transition probability, it is useful to consider the fluctuations of the MCPs in position space where their associated space-time trajectories have an intrinsic length scale of the order of the Compton wavelength $\sim 1/m$. For $md \ll 1$, this size of the fluctuations exceeds the thickness of the wall. Thus, the transition probability becomes d-independent. In the limit $m/\omega \ll 1$, also the dependence on the probe photon wavelength drops out as the large Compton wavelength of the virtual fluctuations dominates all other length scales. Note, however that in an actual experiment, the smallest testable minicharged mass is limited by the extent and the scale of homogeneity of the external magnetic field.

As outlined above, the transition rate as given in Eq. (10) is the maximal transition rate, which can only asymptotically be reached in an experiment. In practice, it is therefore necessary to choose a finite but preferably very small angle between the propagation direction of the photons and the external magnetic field. This introduces an exponential suppression factor in Eq. (10) (cf. Sect. 2), which reduces the transition probability as a function of the photon incidence angle $\sphericalangle(\vec{k}, \vec{B})$ and the ratio $\frac{\epsilon e B}{m^2}$. However, for experimentally feasible photon incidence angles, this additional suppression does not result in any severe reduction of the testable parameter space for minicharged particles [4].

4 Towards a new test of MCPs

We have performed a first theoretical feasibility study for a novel "light-shining-through-wall" scenario in the presence of an external magnetic field. In this scenario, the barrier is traversed by means of virtual particle-antiparticle states, rather than by on-shell particles. We argued that the evaluation of the corresponding photon-to-photon transition probability requires profound knowledge of the photon polarization tensor in the presence of external fields, and in particular requires to keep its full momentum dependence. As established approximations for the photon polarization tensor are inappropriate here, we focused on a special alignment of the photon propagation direction and the external magnetic field. Based on the insights obtained in this limit, we devised a strategy that allows for a solid estimate of the transition probability beyond this special alignment also. Further details of this investigation are forthcoming [4].

Building on the in-depth experience with LSW setups in the experimental community, we hope that these considerations can soon lead to a novel experiment that puts the existence of minicharged particles to the test.

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References

[1] J. Redondo, A. Ringwald, Contemporary Physics **52**, 3 (2011) [arXiv:1011.3741 [hep-ph]].

[2] K. Ehret *et al.*, Phys. Lett. B **689**, 149 (2010) [arXiv:1004.1313 [hep-ex]].

[3] M. Ahlers, H. Gies, J. Jaeckel *et al.*, Phys. Rev. **D76**, 115005 (2007) [arXiv:0706.2836 [hep-ph]].

[4] B. Döbrich, H. Gies, N. Neitz and F. Karbstein, in preparation (2011)

[5] H. Gies and J. Jaeckel, JHEP **0908**, 063 (2009) [arXiv:0904.0609 [hep-ph]].

[6] J. S. Schwinger, Phys. Rev. **82**, 664 (1951).

[7] W. Dittrich and H. Gies, Springer Tracts Mod. Phys. **166**, 1 (2000).

[8] I. A. Batalin and A. E. Shabad, Sov. Phys. JETP **33**, 483 (1971).

[9] L. F. Urrutia, Phys. Rev. D **17**, 1977 (1978).

[10] V. N. Baier, V. M. Katkov, Phys. Lett. **A374**, 2201-2206 (2010). [arXiv:0912.5250 [hep-ph]].

[11] W. y. Tsai and T. Erber, Phys. Rev. D **10**, 492 (1974).

[12] W. y. Tsai and T. Erber, Phys. Rev. D **12**, 1132 (1975).

[13] A. E. Shabad, Annals Phys. **90**, 166 (1975).

[14] A. E. Shabad, arXiv:hep-th/0307214.

[15] R. A. Cover, G. Kalman, Phys. Rev. Lett. **33**, 1113-1116 (1974).

[16] W. Y. Tsai and T. Erber, Acta Phys. Austriaca **45**, 245 (1976).