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New stellar constraints on dark photons

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ABSTRACT

We consider the stellar production of vector states V within the minimal model of "dark photons". We show that when the Stückelberg mass of the dark vector becomes smaller than plasma frequency, the emission rate is dominated by the production of the longitudinal modes of V, and scales as $\kappa^2 m_V^2$, where κ and m_V are the mixing angle with the photon and the mass of the dark state. This is in contrast with widespread assertions in the literature that the emission rate decouples as the forth power of the mass. We derive ensuing constraints on the (κ, m_V) parameter space by calculating the cooling rates for the Sun and horizontal branch stars. We find that stellar bounds for $m_V < 10$ eV are significantly strengthened, to the extent that all current "light-shining-through-wall" experiments find themselves within deeply excluded regions.

1. Introduction

The Standard Model of particles and fields (SM) can be naturally extended by relatively light neutral states. Almost all possible ways of connecting such states to the SM have been explored, and several of such ways stand out as the most economical/natural. One of the most attractive possibilities is the so-called "hypercharge portal", or "kinetic mixing" portal that at low energy connects the electromagnetic current with another massive photon-like state [1]. This model has been under intense scrutiny in the last few years, both experimentally and observationally. The interest to this model is fueled by attractive (yet speculative) possibilities: the dark vector can be a promising mediator of the dark matter-SM interaction [2], or form super-weakly interacting dark matter itself [3,4]. Dark vectors were proposed as a possible solution to the muon $g-2$ discrepancy [5], and have been searched for (so far with negative results), both at high energy and in medium energy high-intensity particle physics experiments.

The region of small vector masses, $m_V <$ eV, can also be very interesting. On the theoretical side, there are speculations of dark photons contributing to dark matter (via an initial condensate-like state) [6] and dark radiation [7]. But perhaps more importantly, there are some hopes for the terrestrial detection of dark photons. So far, several avenues have been proposed: one can attempt observing a "visible-dark-visible" oscillation chain in "light-shining-through-wall" experiments (LSW) [8]. The quanta of dark photons emitted from the Sun can be searched for with "helioscopes" [9], neutrino [10] and dark matter experiments [11,12]. Some of these exciting possibilities have been summarized in the recent review [13]. We will refer to all proposals and experiments aimed at detection of dark vectors, produced astrophysically or in the laboratory, as direct searches.

At the same time, it is well-known that for many light ($m_V <$ keV) and weakly-coupled exotic particles the astrophysical constraints are often far stronger than direct laboratory constraints [14]. The astrophysical constraints are very important for the dark vectors as well, as they determine a surviving fraction of the parameter space that can be explored in direct searches. The most important limits to recon with are the constraints on the emission of dark vectors from solar luminosity, from the horizontal branch stars, neutron star and supernovae cooling rates.

To date, the only in-depth analysis of astrophysical bounds on sub-keV dark vectors was performed by Redondo in [9], where the production of longitudinal modes of the dark photon is treated incorrectly. We trace the mistake traced to a wrong use of the in-medium polarization effects for longitudinal modes. In this Letter we re-assess these bounds, provide correct calculations for the dark photon emission rates, and strengthen the astrophysical bounds in the LSW region by as much as ten orders of magnitude. Our findings significantly reduce the parameter space available for the direct searches and affect or completely change the conclusions

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of many papers written on this subject. In a separate forthcoming publication we will address new limits imposed by the most advanced WIMP detectors on the solar emission of dark vectors [15].

This Letter is organized as follows. The next section introduces the minimal model of the dark vector, and explains the main scaling of its production rate with m_V. Section 3 contains technicalities of the in-medium production of the dark vector. Section 4 contains practical formulae for the stellar emission rates, in application to the Sun and horizontal branch stars, and sets the constraints on the mass-mixing parameter space. We reach our conclusions in Section 5.

2. Dark photon production, in vacuum and in a medium

The minimal model of "dark vectors" extends the SM gauge group $SU(3)_c \times SU(2)_L \times U(1)_Y$ by an Abelian factor $U(1)_V$. Kinetic mixing of the hypercharge field strength $F^Y_{\mu\nu}$ with the field strength $V_{\mu\nu}$ of $U(1)_V$ links the SM to the new physics sector, while SM fields are assumed to be neutral under $U(1)_V$. We are interested in processes far below the electroweak energy scale, for which the relevant low-energy Lagrangian takes the form

$$\mathcal{L} = -\frac{1}{4}F^2_{\mu\nu} - \frac{1}{4}V^2_{\mu\nu} - \frac{\kappa}{2}F_{\mu\nu}V^{\mu\nu} + \frac{m_V^2}{2}V_\mu V^\mu + eJ^\mu_{\rm em}A_\mu. \quad (1)$$

Here $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$ is the photon field strength and V_μ is the "hidden photon" (also known as "dark vector", "secluded vector", "dark photon" etc.—an equivalent set of names). The coupling of A_μ and V_μ is regulated by the kinetic mixing parameter κ, redefined in an appropriate way to absorb the dependence on weak mixing angle. For all calculations in this Letter we use $\kappa \ll 1$, and consider only leading order terms in the mixing angle. Finally $J^\mu_{\rm em}$ is the usual electromagnetic current with electric charge $e < 0$.

It is important to comment on the origin of m_V in (1). The simplest possibility is that m_V is a Stückelberg-type mass. Because of the conservation of the Abelian vector current, m_V remains protected against sensitivity to UV scales, and such a model is technically natural even with very small m_V. An alternative generic possibility is a new scalar field(s) charged under $U(1)_V$ that develops a vacuum expectation value that Higgses the hidden group. This introduces a new interaction term of the physical hidden Higgs with vectors, $g' m_V h' V^2_\mu$, as well as h' self-interaction (see *e.g.* [16]). It is well understood that in the limit of m_V and $m_{h'}$ small compared to all energy scales in the problem, the production of dark sector states is dominated by the dark Higgsstrahlung [3,16], or equivalently, by the pair-production of the $U(1)_V$-charged Higgs scalar fields. Importantly, this process is insensitive to the actual mass of m_V in the small mass limit, and schematically

$$\text{Rate}_{SM\to V+h'} \propto \alpha' \kappa^2 (m_V)^0, \quad (2)$$

where we show only the dependence on dark sector parameters, leaving the SM part of the $V + h'$ production process completely general; $\alpha' = (g')^2/(4\pi)$ is the square of the coupling of dark Higgs to V_μ. For sub-keV dark vectors and Higgses, all previously derived constraints on "millicharged particles" apply [17], and limit the $\kappa g'$ combination to be below $\sim 10^{-13}$. The technical reason for not having any small m_V suppression of the rate (2) despite the interaction term $g' m_V h' V^2_\mu$ being proportional to m_V is of course tied to the production of the longitudinal modes of V in $V + h'$ final state.

The models with the hard (*i.e.* Stückelberg) mass m_V behave differently as the production rate of dark vectors *has to* decouple in the small m_V limit. The easiest way to see this is to restrict the interaction terms in (1) to on-shell V_μ, using $\partial_\mu V^\mu = 0$ and to leading order in κ, $\partial_\mu V^{\mu\nu} = -m_V^2 V^\nu$, so that

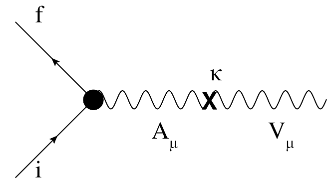

Fig. 1. Illustration of the dark photon emission process by the electromagnetic current.

$$\mathcal{L}_{\rm int} = -\frac{\kappa}{2}F_{\mu\nu}V^{\mu\nu} + eJ^\mu_{\rm em}A_\mu \xrightarrow{\text{on-shell V}}$$
$$\mathcal{L}_{\rm int} = -\kappa m_V^2 A_\mu V^\mu + eJ^\mu_{\rm em}A_\mu. \quad (3)$$

This expression is of course explicitly gauge invariant under $A_\mu \to A_\mu + \partial_\mu \chi$ due to the current conservation and on-shellness of V_μ conditions:

$$\partial_\mu J^\mu_{\rm em} = 0; \qquad \partial_\mu V^\mu = 0. \quad (4)$$

The appearance of m_V^2 in the coupling of V_μ and A_μ shows that two sectors are decoupled in $m_V = 0$ limit. The most important question in considering the production of V_μ states is the scaling of the production rate with m_V, in vacuum and inside a medium. The existing literature on the subject [9] and its subsequent follow-up papers claim that in-medium production decouples as $\text{Rate}_{SM\to V} \propto \kappa^2 m_V^4$ in the small m_V limit. This inference is wrong.

To demonstrate our point we consider a generic production process $i \to f + V$ due to (3), where i, f are any initial, final states of the SM particles. A schematic drawing of such a process is shown in Fig. 1. Without loss of generality we assume that V is emitted in z-direction, so that its four-momentum k_μ is given by $(\omega, 0, 0, |\vec{k}|)$, with $\omega^2 - \vec{k}^2 = m_V^2$. Moreover, we assume that the energy of the emitted V is much larger than its rest mass, $\omega \gg m_V$. Three polarization states can be emitted: two transverse states V_T with polarization vectors $\epsilon^T = (0, 1, 0, 0,)$ and $(0, 0, 1, 0)$, and one longitudinal mode V_L with polarization vector $\epsilon^L = m_V^{-1}(|\vec{k}|, 0, 0, \omega)$. In all cases $\epsilon^2_\mu = -1$ and $\epsilon_\mu k^\mu = 0$.

We include a boundary-free medium via some conducting plasma, characterized by the plasma frequency ω_p. We consider two regimes, [almost] vacuum: $\omega_p \ll m_V \ll \omega$, and in-medium: $m_V \ll \omega_p \ll \omega$. The choice of $|\vec{k}|, \omega \gg \omega_p$ is not essential, and we consider all ranges of ω in the next section. The matrix element for the production process induced by (3) is given by

$$\mathcal{M}_{i\to f+V_{T(L)}} = \kappa m_V^2 [eJ_{{\rm em}\,\mu}]_{fi} \langle A^\mu, A^\nu\rangle \epsilon_\nu^{T(L)}, \quad (5)$$

where $\langle A^\mu, A^\nu\rangle$ stands for the photon propagator with input momentum k_μ, and $[eJ^\mu_{\rm em}]_{fi}$ is the matrix element of the electromagnetic current. We disregard various m_V-independent phase factors and normalizations, as our goal in this section is to only consistently follow the powers of m_V.

For convenience, we fix the photon gauge to be Coulomb, $\nabla_i A_i = 0$, but the same results can be obtained in other gauge, of course. The photon propagator for the production of the transverse modes is given by (see, *e.g.* [18]):

$$\langle A_j, A_l\rangle = \frac{\delta^\perp_{jl}}{\omega^2 - |\vec{k}|^2 - \omega_p^2}$$
$$= \frac{\delta^\perp_{jl}}{m_V^2 - \omega_p^2} \longrightarrow \delta^\perp_{jl} \times \begin{cases} m_V^{-2} & \text{at } m_V \gg \omega_p, \\ -\omega_p^{-2} & \text{at } m_V \ll \omega_p, \end{cases} \quad (6)$$

where $\delta^\perp_{jl}$ is the projector onto transverse modes. Because of the existence of two different regimes for the transverse modes of the