

# Dark photons and resonant monophoton signatures in Higgs boson decays at the LHC

Emidio Gabrielli,<sup>1,2</sup> Matti Heikinheimo,<sup>1</sup> Barbara Mele,<sup>3</sup> and Martti Raidal<sup>1,4</sup><sup>1</sup>*NICPB, Ravala 10, 10143 Tallinn, Estonia*<sup>2</sup>*Dipartimento di Fisica Teorica, Università di Trieste, Strada Costiera 11, I-34151 Trieste, Italy*  
*and INFN, Sezione di Trieste, Via Valerio 2, I-34127 Trieste, Italy*<sup>3</sup>*INFN, Sezione di Roma, P. le A. Moro 2, I-00185 Rome, Italy*<sup>4</sup>*Institute of Physics, University of Tartu, Ravila 14c, 50411 Tartu, Estonia*

(Received 22 May 2014; published 29 September 2014)

Motivated by dark-photon  $\tilde{\gamma}$  scenarios extensively considered in the literature, we explore experimentally allowed models where the Higgs boson coupling to photon and dark photon  $H\gamma\tilde{\gamma}$  can be enhanced. Correspondingly, large rates for the  $H \rightarrow \gamma\tilde{\gamma}$  decay become plausible, giving rise to one monochromatic photon with  $E_\gamma \simeq m_H/2$  (i.e., more than twice the photon energy in the rare standard-model decay  $H \rightarrow \gamma Z \rightarrow \gamma\nu\nu$ ), and a similar amount of missing energy. We perform a model-independent study of this exotic resonant monophoton signature at the LHC, featuring a distinctive  $E_T^\gamma$  peak around 60 GeV, and  $\gamma + E_T$  transverse invariant mass ruled by  $m_H$ . At parton level, we find a  $5\sigma$  sensitivity of the present LHC data set for a  $H \rightarrow \gamma\tilde{\gamma}$  branching fraction of 0.5%. Such large branching fractions can be naturally obtained in dark  $U(1)_F$  models explaining the origin and hierarchy of the standard model Yukawa couplings. We urge the LHC experiments to search for this new exotic resonance in the present data set and in future LHC runs.

DOI: [10.1103/PhysRevD.90.055032](https://doi.org/10.1103/PhysRevD.90.055032)

PACS numbers: 14.80.-j, 14.80.Bn, 13.85.Rm

## I. INTRODUCTION

Although dark matter (DM) is five times more abundant in the Universe than ordinary baryonic matter [1], its properties are yet unknown. It is plausible that the dark sector, which is weakly coupled to the standard model (SM), possesses rich internal structure and interactions. Among the most popular scenarios is the idea that the dark sector contains light or massless gauge bosons [2] that mediate long-range forces between dark particles. In cosmology, the dark photons may solve the small-scale structure formation problems [3,4] and, for massless dark photons [5], predict dark discs of galaxies [6]. In astroparticle physics, dark photons may induce Sommerfeld enhancement of the DM annihilation cross section needed to explain the PAMELA-Fermi-AMS2 positron anomaly [7], may assist light DM annihilations to reach the phenomenologically required magnitude, and make asymmetric DM scenarios phenomenologically viable [8]. Dark/hidden photon scenarios have also been extensively considered in beyond-the-SM frameworks in particle physics [9–15].

Recently, a new paradigm has been proposed for generating exponentially spread SM Yukawa couplings from unbroken  $U(1)_F$  quantum numbers in the dark sector [16]. In this approach, nonperturbative flavor- and chiral-symmetry breaking is transferred from the dark to visible sector via heavy scalar messenger fields [16,17] that might give distinctive new physics (NP) signals at the LHC. For massless dark photons [5], the  $U(1)_F$  kinetic mixing

with  $U(1)_Y$  can be tuned away [9] on shell, in agreement with all existing constraints [2], while off-shell contributions give rise to higher-dimensional contact operators strongly suppressed by the scale of the heavy messengers' mass. Therefore, in this scenario direct tests of dark photons may require new ideas. On the other hand, the photon kinetic mixing can induce millicharge couplings of dark fermions with ordinary photons, that can already be probed at the LHC [18]. This could allow one to constrain some regions of the model parameter space.

In this work we show that, in the unbroken dark  $U(1)$  scenarios, the Higgs-boson two-body decay  $H \rightarrow \gamma\tilde{\gamma}$  to one photon  $\gamma$  and one dark photon  $\tilde{\gamma}$  can be enhanced despite existing constraints, providing a very distinctive NP signature of a single photon plus missing energy at the Higgs resonance. If this signature will be discovered at the LHC,  $CP$  invariance will imply the spin-1 nature of the missing energy, excluding axions or other ultralight scalar particles.

Monophoton plus  $E_T$  signatures have been used by the LHC experiments to search for NP scenarios such as extra dimensions, supersymmetry, DM pair production [19], and SM continuous  $Z\gamma$  production [20]. In those cases the photon and  $E_T$  distributions are mostly monotonic and not much structured, corresponding to the nonresonant production of different invisible particles that carry away broadly distributed missing energy. A resonant monophoton plus  $E_T$  signature occurs in the SM rare Higgs decays  $H \rightarrow Z\gamma \rightarrow \bar{\nu}\nu\gamma$  with a  $\gamma$  energy of about 30 GeV, which is much lower than the  $m_H/2$  photon energy in  $H \rightarrow \gamma\tilde{\gamma}$ . To our knowledge this exotic Higgs decay

channel, giving rise to a striking experimental signature, has not been considered so far (for a review of exotic Higgs signatures, see [21,22]). The aim of this work is to show that the corresponding  $\gamma\bar{\gamma}$  resonance can be realistically detected at the LHC, providing a nontrivial test of dark-photon scenarios at the LHC.

Inspired by the model in [16], we present a more general model-independent framework that can predict enhanced  $H \rightarrow \gamma\bar{\gamma}$  decay rates. We perform a parton-level Monte Carlo study of this process versus relevant SM backgrounds, and show that, for a significant part of the model parameter space, this process could be observed at the LHC. Detailed detector-level studies of the proposed signature will be needed to find the actual LHC sensitivity to massless dark-photon scenarios.

## II. THEORETICAL FRAMEWORK

The aim of the model in [16] is to explain the observed hierarchies in fermion masses, i.e., in the SM Yukawa couplings, by exponential hierarchies due to quantum numbers of an exact new  $U(1)_F$  gauge symmetry in the dark sector. In this model the hidden sector consists of dark fermions charged under  $U(1)_F$ . As previously noted in [23], spontaneous chiral symmetry breaking can be triggered by the presence of a higher derivative kinetic term in the gauge sector, suppressed by a scale  $\Lambda$ , which can be interpreted as the mass scale of the associated Lee-Wick ghost [24] of the  $U(1)_F$  gauge theory. The dark fermion masses  $M_i$  can be dynamically generated via a nonperturbative mechanism *à la* Nambu–Jona-Lasinio [25] as a nontrivial solution of the (finite) mass-gap equation.

The SM Yukawa couplings  $Y_i$  are dynamically generated at one loop by the messenger fields that carry the SM quantum numbers of squarks and sleptons of supersymmetric models. In the approximation of a universal average mass  $\bar{m}$  for the messenger fields, we get

$$Y_i = Y_0(M_i/\bar{m}) \exp\left(-\frac{2\pi}{3\bar{\alpha}q_i^2}\right), \quad (1)$$

where  $\bar{\alpha}$  is the  $U(1)_F$  fine structure constant and  $q_i$  are dark fermion  $U(1)_F$  charges. The loop function  $Y_0(M_i/\bar{m})$  (see [16]) has a weak dependence on  $M_i/\bar{m}$  and is proportional to  $Y_0 \sim \langle S \rangle \Lambda / \bar{m}^2$ , where  $\langle S \rangle$  is the vacuum expectation value (VEV) of the singlet scalar field  $S$  required to break the  $H \rightarrow -H$  parity. Equation (1) implies that the origin of flavor in the SM Yukawa couplings resides in the non-universality of the  $U(1)_F$  charges in the dark sector. Vacuum stability bounds and Eq. (1) require the average mass of colored messengers to be above 50 TeV [16]. The dark fermions are the lightest dark particles which, due to  $U(1)_F$ , are all stable and can potentially contribute to the dark matter density of the universe. Because of the long-range  $U(1)_F$  interaction with nonuniversal charges, the cosmology of the dark sector is nontrivial, and constraints

apply on the masses and couplings of the dark fermions [5]. We will not discuss the DM phenomenology further in this work.

An analogous model has been recently proposed in [17], although there the dynamics responsible for generating the hierarchy in the dark fermion spectrum is missing. We then compute  $\text{BR}(H \rightarrow \gamma\bar{\gamma})$  in a model-independent way to extend our results to all models of this type.

## III. HIGGS DECAYS TO $H \rightarrow \gamma\bar{\gamma}$

Consider a generic messenger sector like in [16,17], consisting of left doublet and right singlet scalars  $S_L^i, S_R^i$ , with a flavor universal mass term. The latter carry squark and slepton quantum numbers under the SM gauge group, and additional  $U(1)_F$  charges to couple to dark fermions. Their couplings to the Higgs boson are (omitting the flavor indices)

$$\mathcal{L}_{\text{MS}}^I = \lambda_S S(\tilde{H}^\dagger S_L^U S_R^U + H^\dagger S_L^D S_R^D) + \text{H.c.} \quad (2)$$

After the singlet  $S$  scalar gets a VEV, a  $H \rightarrow \gamma\bar{\gamma}$  decay rate proportional to  $\mu_S = \lambda_S \langle S \rangle$  is induced at one loop. After electroweak symmetry breaking, the Lagrangian for generic  $S_{L,R}$  is

$$\mathcal{L}_S^0 = \partial_\mu \hat{S}^\dagger \partial^\mu \hat{S} - \hat{S}^\dagger M_S^2 \hat{S}, \quad (3)$$

where  $\hat{S} = (S_L, S_R)$ , and the mass term is given by

$$M_S^2 = \begin{pmatrix} m_L^2 & \Delta \\ \Delta & m_R^2 \end{pmatrix}, \quad (4)$$

where  $\Delta = \mu_S v$  parametrizes the left-right mixing of scalars, and  $v$  is the SM Higgs VEV. Then, if  $\epsilon_1^\mu(k_1)$  and  $\epsilon_2^\mu(k_2)$  are the photon and dark-photon polarization vectors, respectively, we express the  $H \rightarrow \gamma\bar{\gamma}$  amplitude as

$$M_{\gamma\bar{\gamma}} = \frac{1}{\Lambda_{\gamma\bar{\gamma}}} T_{\mu\nu}(k_1, k_2) \epsilon_1^\mu(k_1) \epsilon_2^\nu(k_2), \quad (5)$$

where  $\Lambda_{\gamma\bar{\gamma}}$  parametrizes the effective scale associated to the NP, and  $T^{\mu\nu}(k_1, k_2) \equiv g^{\mu\nu} k_1 \cdot k_2 - k_2^\mu k_1^\nu$ . The total width is then

$$\Gamma(H \rightarrow \gamma\bar{\gamma}) = m_H^3 / (32\pi \Lambda_{\gamma\bar{\gamma}}^2). \quad (6)$$

If we neglect the Higgs boson mass with respect to the messenger masses  $m_{L,R}$  in the loop, we obtain

$$\frac{1}{\Lambda_{\gamma\bar{\gamma}}} = \frac{\mu_S \sqrt{\alpha\bar{\alpha}} R}{12\pi} \left( \frac{\sqrt{(m_L^2 - m_R^2)^2 + 4\Delta^2}}{m_L^2 m_R^2 - \Delta^2} \right) \sin 2\theta, \quad (7)$$

where  $R = N_c \sum_{i=1}^3 (e_U q_{U_i} + e_D q_{D_i})$ , with  $q_{U_i}, q_{D_i}$  the  $U(1)_F$  charges in the up and down sectors, and  $e_U = \frac{2}{3}$ ,

$e_D = -\frac{1}{3}$  the corresponding EM charges;  $\alpha$  is the EM fine structure constant,  $N_c = 3$  is the number of colors, and  $\theta$  is the mixing angle diagonalizing Eq. (4). The above result can be easily generalized to include the contributions of messengers in the leptonic sector, in this case  $N_c = 1$ ,  $e_U = 0$  and  $e_D = -1$ . We assume mass universality for  $S_L$  and  $S_R$ ,  $m_L \simeq m_R \equiv \bar{m}$ , giving  $\theta = \pi/4$ . Defining  $\xi = \Delta/\bar{m}^2$ , the eigenvalues of Eq. (4) become  $m_{\pm}^2 = \bar{m}^2(1 \pm \xi)$ , and the  $\Lambda_{\gamma\bar{\gamma}}$  scale simplifies to

$$\Lambda_{\gamma\bar{\gamma}} = \frac{6\pi v}{R\sqrt{\alpha\bar{\alpha}}} \frac{1 - \xi^2}{\xi^2}. \quad (8)$$

To avoid tachyons, one needs  $0 \leq \xi \leq 1$ .

The messengers induce new contributions also to the Higgs decays  $H \rightarrow \gamma\gamma$  and  $H \rightarrow \bar{\gamma}\bar{\gamma}$ . The corresponding amplitudes have the same structure as (5), and we obtain

$$\Lambda_{\gamma\gamma} = \Lambda_{\gamma\bar{\gamma}} \frac{R}{R_0} \sqrt{\frac{\bar{\alpha}}{\alpha}}, \quad \Lambda_{\bar{\gamma}\bar{\gamma}} = \Lambda_{\gamma\bar{\gamma}} \sqrt{\frac{\bar{\alpha}}{\alpha}} \frac{R}{R_1}, \quad (9)$$

where  $R_0 = 3N_c(e_U^2 + e_D^2)$ , and  $R_1 = N_c \sum_{i=1}^3 (q_{U_i}^2 + q_{D_i}^2)$ .

A model-independent parametrization for the branching ratios (BRs) of the decays  $H \rightarrow \gamma\gamma$ ,  $H \rightarrow \gamma\bar{\gamma}$ , and  $H \rightarrow \bar{\gamma}\bar{\gamma}$  can be expressed as follows:

$$\text{BR}_{\gamma\gamma} = N(1 \pm \sqrt{r_{\gamma\gamma}})^2, \quad \text{BR}_{AB} = N r_{AB}, \quad (10)$$

where  $AB \equiv \{\gamma\bar{\gamma}, \bar{\gamma}\bar{\gamma}\}$ ,  $N = \text{BR}_{\gamma\gamma}^{\text{SM}} / (1 + r_{\bar{\gamma}\bar{\gamma}} \text{BR}_{\gamma\gamma}^{\text{SM}})$ , and the ratios  $r_{AB}$  are given by

$$r_{\gamma\gamma} = 2r_{\gamma\bar{\gamma}} \frac{R^2}{R_0^2} \left(\frac{\bar{\alpha}}{\alpha}\right), \quad r_{\bar{\gamma}\bar{\gamma}} = r_{\gamma\bar{\gamma}} \frac{R_1^2}{R_0^2} \left(\frac{\bar{\alpha}}{\alpha}\right)^2, \quad (11)$$

where  $r_{\gamma\gamma} \equiv \Gamma_{\gamma\gamma}^{\text{NP}} / \Gamma_{\gamma\gamma}^{\text{SM}}$ . Here  $\Gamma_{\gamma\gamma}^{\text{NP}}$  and  $\Gamma_{\gamma\gamma}^{\text{SM}}$  correspond to the  $H \rightarrow \gamma\gamma$  decay widths, mediated by new particles and SM ones, respectively. The  $\pm$  signs in Eq. (10) correspond to the constructive or destructive interference with the SM amplitude. In the scenario [16], the sign in  $\text{BR}_{\gamma\gamma}$  is predicted to be positive, while the corresponding value for  $r_{\gamma\gamma}$  is given by

$$r_{\gamma\gamma} = \left( \frac{R_0 \xi^2}{3F(1 - \xi^2)} \right)^2, \quad (12)$$

where  $F$  is the SM contribution, given by  $F = F_W(\beta_W) + \sum_f N_c Q_f^2 F_f(\beta_f)$ , with  $\beta_W = 4M_W^2/m_H^2$ ,  $\beta_f = 4m_f^2/m_H^2$ , and  $F_W(x)$  and  $F_f(x)$  can be found in [26]. Once the corresponding Higgs BRs are measured, the  $U(1)_F$  charges  $q_i$  can be derived from the Yukawa couplings by Eq. (1).

To quantify predictions of this scenario, in Fig. 1 we plot  $\text{BR}(H \rightarrow \gamma\bar{\gamma})$  as a function of  $\bar{\alpha}$ , assuming that there is only one messenger contributing, with a charge  $e = q = 1$ . The curves are evaluated for  $r_{\gamma\gamma} = 0.1, 0.2, 0.5, 1$ . The red

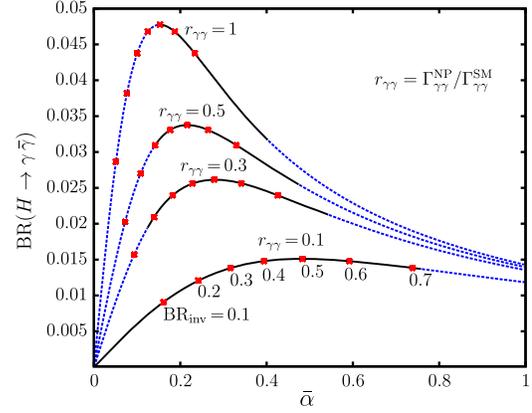


FIG. 1 (color online). Predictions for  $\text{BR}(H \rightarrow \gamma\bar{\gamma})$  as functions of  $\bar{\alpha}$  for different  $\text{BR}_{\text{inv}}$  and  $r_{\gamma\gamma}$  in the minimal model.

dot bullets correspond to different  $\text{BR}_{\bar{\gamma}\bar{\gamma}}$  values (or Higgs invisible branching ratios  $\text{BR}_{\text{inv}}$ ), as shown in the plot (in the experimentally allowed range [27]). The full lines correspond to the interval  $\text{BR}_{\gamma\gamma}^{\text{SM}}/2 \leq \text{BR}_{\gamma\gamma} \leq 2\text{BR}_{\gamma\gamma}^{\text{SM}}$ , where  $\text{BR}_{\gamma\gamma}^{\text{SM}} = 2.28 \times 10^{-3}$ , while the dashed lines correspond to predictions outside that range. We find that the signal  $\text{BR}(H \rightarrow \gamma\bar{\gamma})$  can be as large as 5% (that is more than 1 order of magnitude larger than  $\text{BR}_{\gamma\gamma}^{\text{SM}}$ ), consistently with all model parameters and the LHC constraints.

We stress that large values of the messenger mixing-mass parameter  $\xi$  are natural in the present scenario, in order to generate a large top-quark Yukawa coupling radiatively, and all EW precision tests can be satisfied due to the heavy and flavor universal messenger sector [16]. In addition, large values of  $\bar{\alpha} \gg \alpha$  are naturally expected in this scenario from Eq. (1), provided the splitting among the  $q_i$  charges is not too small. Consequently, the relatively large  $\text{BR}(H \rightarrow \gamma\bar{\gamma})$  shown in Fig. 1 can be considered a generic prediction of the present theoretical framework.<sup>1</sup>

#### IV. MODEL-INDEPENDENT ANALYSIS OF $H \rightarrow \gamma\bar{\gamma}$ AT THE LHC

The process  $pp \rightarrow H \rightarrow \gamma\bar{\gamma}$  gives rise to the signal  $\gamma + \cancel{E}_T$ , where  $E_\gamma = m_H/2$  in the Higgs rest frame. In the lab frame, one can define the variable  $M_T$ , that is the transverse invariant mass of the  $\gamma + \cancel{E}_T$  system, as

$$M_T = \sqrt{2p_T^\gamma \cancel{E}_T (1 - \cos \Delta\phi)}, \quad (13)$$

where  $p_T^\gamma$  is the photon transverse momentum, and  $\Delta\phi$  is the azimuthal distance between the photon momentum and the missing transverse momentum  $\cancel{E}_T$ .

<sup>1</sup>Large values of the mixing parameter  $\xi$  can be safely generated from the purely EW messenger sector, since the latter does not affect the Higgs production cross section in gluon fusion.

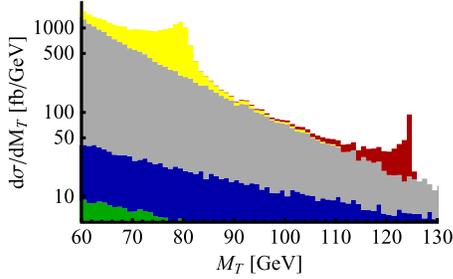


FIG. 2 (color online). The  $\gamma + E_T$  transverse invariant mass distribution (in fb/GeV) of the signal (red), and the main backgrounds  $\gamma j$  (grey),  $\gamma Z$  (blue),  $jZ$  (green), and  $W$  (yellow). For illustration, we show the signal for  $\text{BR}(H \rightarrow \gamma\bar{\gamma}) = 5\%$ .

Like in the  $W \rightarrow e\nu$  production, the  $M_T$  observable features a narrow peak at the mass of the original massive particle (that is  $m_H$ ; see Fig. 2). Also the  $p_T^\gamma$  distribution will exhibit a similar structure around  $m_H/2$ . These features allow for a very efficient cut-based search strategy, looking for events with a single photon and missing energy, with no jets or leptons, and cutting around the expected maximum of the  $M_T$  and  $p_T^\gamma$  distributions. These peaks could be relatively easy to pinpoint on top of the continuous relevant backgrounds, for sufficiently large  $H \rightarrow \gamma\bar{\gamma}$  decay rates. Thus we formulate the criteria for event selection as follows:

- (i) One isolated photon with  $50 \text{ GeV} < p_T^\gamma < 63 \text{ GeV}$  and  $|\eta^\gamma| < 1.44$ .
- (ii) Missing transverse momentum with  $E_T > 50 \text{ GeV}$ .
- (iii) Transverse mass in  $100 \text{ GeV} < M_T < 126 \text{ GeV}$ .
- (iv) No isolated jets or leptons.

The most relevant backgrounds for the above selection criteria are, in order of importance,

- (1)  $pp \rightarrow \gamma j$ , where large apparent  $E_T$  is created by a combination of real  $E_T$  from neutrinos in heavy quark decays and mismeasured jet energy;
- (2)  $pp \rightarrow \gamma Z \rightarrow \gamma\nu\bar{\nu}$  (irreducible background);
- (3)  $pp \rightarrow jZ \rightarrow j\nu\bar{\nu}$ , where the jet is misidentified as a photon;
- (4)  $pp \rightarrow W \rightarrow e\nu$ , where the electron (positron) is misidentified as a photon;
- (5)  $pp \rightarrow \gamma W \rightarrow \gamma\ell\nu$ , where the lepton is missed;
- (6)  $pp \rightarrow \gamma\gamma$ , where one of the photons is missed.

The  $pp \rightarrow \gamma j$  background is expected to be dominant for the  $E_T$  range relevant here, and also the most difficult to estimate without detailed information about the detector performance [28]. We have evaluated this background by simulating events with one photon and one jet, treating jets with  $|\eta| > 4.0$  as missing energy, following [29] (a more detailed investigation of the  $pp \rightarrow \gamma j$  background, although crucial for assessing the actual experiment potential, is beyond the scope of this work). All the other backgrounds have also been estimated through a parton-level simulation, expected to be relatively accurate for electroweak processes (applying a probability  $10^{-3}$  and

TABLE I. The cross section times acceptance (in fb) for the signal and background processes at 8 TeV for the selections ( $A_1$ )  $50 \text{ GeV} < p_T^\gamma < 63 \text{ GeV}$ ; ( $A_2$ )  $60 \text{ GeV} < p_T^\gamma < 63 \text{ GeV}$ . In all cases  $|\eta^\gamma| < 1.44$ , and  $S/\sqrt{S+B}$  is for  $20 \text{ fb}^{-1}$ . The significance improves with tighter cuts, but this is subject to experimental resolution and radiative corrections.

	$\sigma \times A_1$	$\sigma \times A_2$
Signal $\text{BR}_{H \rightarrow \gamma\bar{\gamma}} = 1\%$	65	34
$\gamma j$	715	65
$\gamma Z \rightarrow \gamma\nu\bar{\nu}$	157	27
$jZ \rightarrow j\nu\bar{\nu}$	63	11
$W \rightarrow e\nu$	22	0
Total background	957	103
$S/\sqrt{S+B}$ ( $\text{BR}_{H \rightarrow \gamma\bar{\gamma}} = 1\%$ )	9.1	13.0
$S/\sqrt{S+B}$ ( $\text{BR}_{H \rightarrow \gamma\bar{\gamma}} = 0.5\%$ )	4.6	6.9

1/200 to misidentify a jet and an electron, respectively, as a photon). We will neglect the subdominant backgrounds from processes 5 and 6 (the  $H \rightarrow \gamma\gamma$  background is also negligible). The contribution of relevant backgrounds passing the cuts is shown in Table I, and the scaling of the different components with the transverse mass is shown in Fig. 2. Although our leading-order parton-level analysis, after applying a cut on  $p_T^\gamma$  is not much affected by a further cut on the  $M_T$  variable, we expect the latter to be very effective in selecting our structured signal over the continuous reducible QCD background [28].

With the existing data set of  $20 \text{ fb}^{-1}$ , for  $\text{BR}(H \rightarrow \gamma\bar{\gamma}) = 1\%$ , we get a significance  $S/\sqrt{S+B}$  of 9 standard deviations ( $9\sigma$ ), with  $S(B)$  the number of signal (background) events passing the cuts. The sensitivity limit for a  $5\sigma$  discovery is then estimated to be  $\text{BR}(H \rightarrow \gamma\bar{\gamma}) \sim 0.5\%$  with the existing data set.

## V. CONCLUSIONS

Motivated by possible cosmological and particle physics hints for the existence of massless dark photon  $\bar{\gamma}$ , we have performed a model-independent study of the exotic  $H \rightarrow \gamma\bar{\gamma}$  decay. At the LHC this results in a single photon plus  $E_T$  signature, with both energies peaked at  $m_H/2$ . At parton level, we estimate that a  $5\sigma$  discovery can be reached with the existing 8 TeV LHC data sets if  $\text{BR}(H \rightarrow \gamma\bar{\gamma}) \sim 0.5\%$ . Such a large branching ratio can be easily obtained in dark  $U(1)_F$  models explaining the origin and hierarchy of the SM Yukawa couplings. The proposed experimental signature is new, and requires detailed detector-level studies to draw realistic conclusions on the LHC sensitivity to dark photons.

## ACKNOWLEDGMENTS

We thank S. Chauhan, J. P. Chou and J. Alcaraz Maestre for communications, and C. Spethmann for collaboration in the early stages of the project. This work was supported by Grants No. MTT60 and No. IUT23-6, CERN+, and by EU through the ERDF CoE program.

- [1] P. Ade *et al.* (Planck Collaboration), arXiv:1303.5076 [Astron. Astrophys. (to be published)].
- [2] For a review and references, see R. Essig *et al.*, arXiv:1311.0029.
- [3] D. N. Spergel and P. J. Steinhardt, *Phys. Rev. Lett.* **84**, 3760 (2000); M. Vogelsberger, J. Zavala, and A. Loeb, *Mon. Not. R. Astron. Soc.* **423**, 3740 (2012).
- [4] L. G. van den Aarssen, T. Bringmann, and C. Pfrommer, *Phys. Rev. Lett.* **109**, 231301 (2012); S. Tulin, H.-B. Yu, and K. M. Zurek, *Phys. Rev. D* **87**, 115007 (2013).
- [5] L. Ackerman, M. R. Buckley, S. M. Carroll, and M. Kamionkowski, *Phys. Rev. D* **79**, 023519 (2009).
- [6] J. Fan, A. Katz, L. Randall, and M. Reece, *Phys. Rev. Lett.* **110**, 211302 (2013).
- [7] N. Arkani-Hamed, D. P. Finkbeiner, T. R. Slatyer, and N. Weiner, *Phys. Rev. D* **79**, 015014 (2009).
- [8] K. M. Zurek, *Phys. Rep.* **537**, 91 (2014).
- [9] B. Holdom, *Phys. Lett.* **166B**, 196 (1986).
- [10] K. S. Babu, C. F. Kolda, and J. March-Russell, *Phys. Rev. D* **57**, 6788 (1998).
- [11] D. Feldman, Z. Liu, and P. Nath, *Phys. Rev. D* **75**, 115001 (2007).
- [12] S. A. Abel, M. D. Goodsell, J. Jaeckel, V. V. Khoze, and A. Ringwald, *J. High Energy Phys.* **07** (2008) 124.
- [13] M. Pospelov, *Phys. Rev. D* **80**, 095002 (2009).
- [14] E. J. Chun, J. C. Park, and S. Scopel, *J. High Energy Phys.* **02** (2011) 100.
- [15] S. Y. Choi, C. Englert, and P. M. Zerwas, *Eur. Phys. J. C* **73**, 2643 (2013).
- [16] E. Gabrielli and M. Raidal, *Phys. Rev. D* **89**, 015008 (2014).
- [17] E. Ma, *Phys. Rev. Lett.* **112**, 091801 (2014); S. Fraser and E. Ma, arXiv:1402.6415.
- [18] J. Jaeckel, M. Jankowiak, and M. Spannowsky, *Phys. Dark Univ.* **2**, 111 (2013).
- [19] S. Chatrchyan *et al.* (CMS Collaboration), *Phys. Rev. Lett.* **108**, 261803 (2012); G. Aad *et al.* (ATLAS Collaboration), *Phys. Rev. Lett.* **110**, 011802 (2013).
- [20] G. Aad *et al.* (ATLAS Collaboration), *Phys. Rev. D* **87**, 112003 (2013); S. Chatrchyan *et al.* (CMS Collaboration), *J. High Energy Phys.* **10** (2013) 164.
- [21] D. Curtin *et al.*, arXiv:1312.4992; J. F. Kamenik and C. Smith, *Phys. Rev. D* **85**, 093017 (2012); H. Davoudiasl, H.-S. Lee, I. Lewis, and W. J. Marciano, *Phys. Rev. D* **88**, 015022 (2013); Y. Sun and D.-N. Gao, *Phys. Rev. D* **89**, 017301 (2014).
- [22] A. Falkowski and R. Vega-Morales, arXiv:1405.1095.
- [23] E. Gabrielli, *Phys. Rev. D* **77**, 055020 (2008).
- [24] T. D. Lee and G. C. Wick, *Phys. Rev. D* **3**, 1046 (1971).
- [25] Y. Nambu and G. Jona-Lasinio, *Phys. Rev.* **122**, 345 (1961); **124**, 246 (1961).
- [26] W. J. Marciano, C. Zhang, and S. Willenbrock, *Phys. Rev. D* **85**, 013002 (2012).
- [27] S. Chatrchyan *et al.* (CMS Collaboration), Report No. CMS-PAS-HIG-13-013; G. Aad *et al.* (ATLAS Collaboration), *Phys. Rev. Lett.* **112**, 201802 (2014); S. Chatrchyan *et al.* (CMS Collaboration), arXiv:1404.1344.
- [28] S. Chatrchyan *et al.* (CMS Collaboration), Report No. CMS-PAS-JME-12-002; G. Aad *et al.* (ATLAS Collaboration), *Eur. Phys. J. C* **72**, 1844 (2012).
- [29] C. Petersson, A. Romagnoni, and R. Torre, *J. High Energy Phys.* **10** (2012) 016.