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Multiphotons and photon jets from new heavy vector bosons

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We discuss an extension of the Standard Model with a new vector boson decaying predominantly into a multiphoton final state through intermediate light degrees of freedom. The model has a distinctive phase in which the photons are collimated. As such, they would fail the isolation requirements of standard multiphoton searches, but group naturally into a novel object, the "photon-jet". Once defined, the photon-jet object facilitates more inclusive searches for similar phenomena. We present a concrete model, discuss photon jets more generally, and outline some strategies that may prove useful when searching for such objects.

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New heavy vector bosons have long been discussed in the literature and searched for in colliders [1]. The production of the vector boson in high-energy collisions is usually searched for by looking for a heavy resonance in a difermion final state. In this paper we discuss the possibility of a vector boson decaying predominantly into a multiphoton final state and the associated phenomenology. A case of particular interest is when the vector boson decays into highly boosted light resonances, which in turn decay to collimated diphotons. Such diphoton pairs fail standard isolation criteria, instead giving rise to a novel object that is likely recognizable at colliders, dubbed a "photon jet."

Theories with new massive vector bosons usually involve a broken gauge symmetry. If the symmetry is broken by charged scalars that spontaneously break the Abelian gauge symmetry associated with the vector boson, then the vector boson necessarily couples to the uneaten components of these scalars and can decay into them if kinematically allowed. This decay dominates if no other light fields are charged under the broken gauge symmetry. The observation of the heavy vector boson then depends on the decay of the scalars, as is common in hidden valley models [2]. If these scalars are light and do not have any renormalizable interaction with the Standard Model (SM), then they may decay only into photons through terms of the form $\phi F_{\mu\nu}F^{\mu\nu}$ where $F_{\mu\nu}$ is the electromagnetic field strength. Such scalars are only very loosely constrained when their masses $\gtrsim 100$ MeV as we will consider in this paper [3]. Through relativistic kinematics, the mass spectrum of the theory dictates the degree of collimation of these photons. If the scalars are comparable in mass to the vector boson then the photons will be fairly well isolated and result in multiphoton signature. On the other hand, if the scalars are much lighter than the vector boson, they are ultrarelativistic so the photons resulting from their decay are tightly collimated (as is commonly seen in π^0 decay). Such "photon jets" would fail the isolation criteria present in general photon searches. In extreme cases, when the scalars are very light, each of the scalars may be reconstructed as a single photon. A collimated photon bunch of this sort was previously considered in Ref. [4] where the authors examined the possibility of the Higgs boson decay into two light pseudoscalars. Amusingly, a massive vector boson decaying into such narrow photon jets would be reconstructed as a diphoton decay of a spin-1 particle, in apparent violation the Landau-Yang theorem which forbids such transitions [5].

A concrete model realizing the above phenomenology includes a new gauge symmetry $U(1)'$ and its carrier, Z', two complex scalars, $S_{1,2}$, charged under this symmetry, and a set of heavy fermions ψ, ψ^c and χ, χ^c that are charged both under the new group as well as hypercharge as $(\pm 2, \pm 1)$ and $(\mp 1, \mp 1)$, respectively. The high-energy Lagrangian is

$$\mathcal{L} = \mathcal{L}_{U(1)'} + \mathcal{L}_{\text{fermions}} - \frac{\epsilon}{2} Z'_{\mu\nu} B^{\mu\nu}$$
$$\mathcal{L}_{U(1)'} = -\frac{1}{4} Z'_{\mu\nu} Z'^{\mu\nu} + \sum_{i=1,2} |D_\mu S_i|^2 + V(S_1, S_2) \quad (1)$$
$$\mathcal{L}_{\text{fermions}} = \text{kinetic terms} + y_1 S_1 \psi\chi + y_2 S_2 \psi^c \chi^c.$$

The fermions will generically drive the scalars to develop a vacuum expectation value and break the $U(1)'$. This will render the vector boson massive, with mass $M^2_{Z'} = g^2_{z'}(\langle S_1\rangle^2 + \langle S_2\rangle^2)$ where $\langle S_i\rangle$ is the expectation value of $\sqrt{2}S_i$. The fermions are also rendered massive with masses $y_1\langle S_1\rangle$ and $y_2\langle S_2\rangle$ for $\psi\chi$ and $\psi^c\chi^c$, respectively. The separate global phase symmetries on S_1 and S_2 result in two Goldstone bosons at low energies [6]. One combination becomes the longitudinal component of the $U(1)'$ through the Higgs mechanism. The remaining physical pseudoscalar, which we denote by a', may be naturally light with a mass sensitive only to terms that break the separate phase symmetry of the two complex scalars. The masses of the rest of the scalars are determined by the detailed quartic couplings of the scalar potential $V(S_1, S_2)$. We note that the matter content above is anomaly free, but the sum of $U(1)' \times U_Y(1)$ charges does not vanish, giving rise to a logarithmic running of kinetic mixing. This can be canceled by the introduction of a heavy Dirac fermion of charge $(-3, 1)$ with an arbitrary mass M. The kinetic mixing term [7] in this model is then given by $\epsilon = \epsilon_{UV} + (3g_{z'}g_Y/16\pi^2)\sum_i \log(y_i\langle S_i\rangle/M)$, which we treat as a free parameter. This mixing results in a Z' which interpolates between coupling dominantly to the hypercharge current when $M_{Z'} \gg M_{Z^0}$ to coupling dominantly to the electromagnetic current when $M_{Z'} \ll M_{Z^0}$ [8,9].

The theory as it is has no predictive power over the mass spectrum of the fermions and scalars. But there is an interesting phase of the theory where the broad phenomenology becomes independent of the detailed mass spectrum. This happens when the fermions are heavier than the vector boson, which is in turn much heavier than the scalars, $M_\Psi > M_{Z'} > m_s$. In this case, we can integrate out the fermions and generate an effective coupling between the scalars and hypercharge [10],

$$\mathcal{L}_{sBB} = \frac{1}{\Lambda_{h'}} h' B_{\mu\nu}B^{\mu\nu} + \frac{1}{\Lambda_{a'}} a' B_{\mu\nu}\tilde{B}^{\mu\nu}, \qquad (2)$$

where h' stands for either the heavy or the light CP-even scalar, a' is the pseudoscalar, and $\tilde{B}^{\mu\nu} = \epsilon^{\mu\nu\alpha\beta}B_{\alpha\beta}$ is the dual hypercharge field strength. The dimensionful coefficients are given by

$$\Lambda^{-1}_{h',a'} = \frac{\alpha_Y}{8\pi}\sum_f \left(\frac{y_{h',a'}}{M_f}\right) q_f^2 A\left(\frac{m^2_{h',a'}}{4M_f^2}\right), \qquad (3)$$

where the sum runs over the heavy fermions. Here, $\alpha_Y = \alpha/\cos^2\theta_W$ is the hypercharge coupling, $y_{h',a'}$ is the Yukawa coupling of the respective scalar to the fermions, and the function $A(m^2_{h',a'}/4M_f^2) \to 4/3$ and 2 in the limit $M_f \gg m_{h',a'}$ for the scalar and pseudoscalar, respectively (see Ref. [11] for the precise form of $A(x)$). Similar terms involving the Z' as well as mixed hypercharge-Z' are also generated with similar coefficients.

The collider phenomenology of this phase is dominated by the production of the Z' through the kinetic mixing, and its subsequent decay into the light scalars to which it couples directly. The scalars' subsequent decay into multiple photons therefore forms the distinctive signature of this type of model, and we are now set to explore the detailed structure of such events. The total production cross section is given approximately by (see e.g., Ref. [12])

$$\sigma_{\rm tot} \approx 15\ {\rm fb}\left(\frac{\epsilon}{0.1}\right)^2\left(\frac{1\ {\rm TeV}}{M_{Z'}}\right)^5, \qquad (4)$$

where the scaling with the mass is approximate. There are strong generic bounds on kinetically mixed Z' models, but for $M_{Z'} \gtrsim$ a few hundred GeV the mixing parameter can be rather large $\epsilon \gtrsim 0.1$ [13]. When more than one scalar is lighter than the Z' the decay $Z' \to h'a'$ would dominate over any decay into SM fermions since it is not suppressed by ϵ, the kinetic mixing parameter. We concentrate on this possibility for the reminder of this paper since it offers unique phenomenology which is not already covered by existing searches. However, before proceeding we briefly digress to comment on other possibilities that can arise when the direct decay into two scalars is not allowed.

If only the pseudoscalar, a', is lighter than the vector boson then the Z' may decay into a' + photon through the dimension-5 mixed operator $a'\tilde{Z}'_{\mu\nu}F^{\mu\nu}$. However, this decay mode is not parametrically much larger than the decay back to hypercharged SM fermions through kinetic mixing. These are independent operators, and whether one decay dominates over the other depends on the details of the high-energy theory. We thus refrain from elaborating on this possibility further except to note three things: first, for the kinetic mixing contribution to the width of the Z' to be smaller than the dimension-5 mixed operator contribution would require a kinetic mixing parameter of $\epsilon \lesssim 0.01$. As is clear from Eq. (4), this will result in too small a rate unless the Z' is in the several hundreds GeV range and below; second, the decay mode $Z' \to a'$ + photon is best searched for in 3 photon searches if the pseudoscalar mass is not much smaller than the Z' mass. Alternatively, when the pseudoscalar is very light, the diphoton resulting from its decay would form a photon jet and result in the interesting new topology of photon + photon jet. Third, when the decay into SM fermions through kinetic mixing dominates, the phenomenology of this model is similar to the well-studied phenomenology of kinetically mixed Z's [8,9,13]. We emphasize that these other decay possibilities are all suppressed by $\mathcal{O}(\epsilon^2)$ compared to the direct decay $Z' \to h'a'$ when that is present.

Returning to the main concern of this paper, the decay into two scalars ($Z' \to h'a'$), the scalars thus produced are unstable and will decay into diphotons. The interactions (2) result in a decay of the scalars to two photons with relatively long lifetime. The boosted lifetime of the pseudoscalar for instance is given by

$$\gamma c\tau_{a'} = 1\ {\rm mm}\left(\frac{137^{-1}}{\alpha'}\right)\left(\frac{10\ {\rm GeV}}{m_{a'}}\right)^4\left(\frac{M_{Z'}}{\rm TeV}\right)^3. \qquad (5)$$

More importantly for the structure of such decays, the differential width for the decay of a scalar s, where s stands for any of the scalars, with boost γ, velocity $\beta = \sqrt{1-\gamma^{-2}}$, and energy $E_s = \gamma m_s$ in the lab frame is

$$\frac{d^2\Gamma(s \to {\rm di\ photons})}{d\varepsilon d\phi} = \frac{\Gamma_s}{2\pi\gamma\beta}, \qquad (6)$$

where Γ_s is the total decay width, ϕ is the angle of the production plane with respect to the scalar's momentum, $\varepsilon = 2E_2/m_s$ is the dimensionless energy of the less energetic photon, and $\gamma(1-\beta) \le \varepsilon \le \gamma$. For the purpose of collider phenomenology, we will be especially interested in the behavior of the decay rate as a function of the separation $\Delta R = \sqrt{\Delta\phi^2 + \Delta\eta^2}$ between the two photons. This dependence can be worked out precisely, but the associated formulas are particularly simple in the large boost limit. For instance, the separation between the two photons is given by

$$\Delta R = \frac{2\cosh\eta}{\sqrt{\varepsilon\varepsilon'}} + \mathcal{O}(\gamma^{-1}), \tag{7}$$

where η is the rapidity of the scalar in the lab frame and $\varepsilon' = 2\gamma - \varepsilon$ is the energy of the more energetic photon. Thus the separation of the two photons is greater than a nominal minimum separation, $\Delta R_{\min}$, when

$$0 < \varepsilon < \frac{2\cosh^2\eta}{\gamma\Delta R^2_{\min}}. \tag{8}$$

Different experimental analyses use different choices for the minimum separation $\Delta R_{\min}$ and we therefore leave it as a free parameter and explore its effect on the phenomenology. Since decays are evenly distributed in the photon's energy, ϵ, the probability of the scalar decaying into two photons with $\Delta R > \Delta R_{\min}$ is simply given by

$$\frac{\Gamma(\Delta R > \Delta R_{\min})}{\Gamma_s} = \frac{2\cosh^2\eta}{\gamma^2\Delta R^2_{\min}}. \tag{9}$$

A similar approximate formula can be obtained for the case when a minimum transverse energy cut is placed on the photons. The approximation above breaks down when

$$\gamma\Delta R_{\min}\cosh^{-1}\eta \lesssim 1 \tag{10}$$

but is otherwise very good. Importantly, in this boosted limit, Eq. (9) is independent of the rapidity of the Z' in the lab frame, but depends only on the rapidity of the scalars in the rest frame since $\gamma^{-1}\cosh\eta = \hat{\gamma}^{-1}\cosh\hat{\eta}$, where $\hat{\gamma} = M_{Z'}/2m_s$ and $\hat{\eta}$ are the scalar's boost and rapidity in the Z' rest frame.

Computing more detailed observables requires accounting for geometric acceptance requirements and integrating over the full Z' phase space. Since the diphotons are fairly collimated, we can approximate the effect of a maximum detectable rapidity for photons, $y_{\max}$, by treating $y_{\max}$ as an upper limit on the scalars' rapidities. In the narrow-width approximation the leading-order [12] double-differential cross section for Z' production depends on the Z' lab frame rapidity, y, and the scalars rapidity in the Z' rest frame, $\hat{\eta}$, as

$$\frac{d^2\sigma}{dy\,d\hat{\eta}} = \rho_{q,\bar{q}}(x_1, x_2) \times \Theta(\hat{\eta})\sigma_0\,\mathrm{BR}(Z' \to s_i a'), \tag{11}$$

where the parton luminosity function is $\rho_{q,\bar{q}}(x_1, x_2)$ with $x_{1,2} = (M_{Z'}/\sqrt{s})e^{\pm y}$, $\sqrt{s}$ is the center of mass energy,

$$\sigma_0 = \frac{4\pi^2\alpha\epsilon^2}{3M^2_{Z'}\cos^2\theta_W}\left(\frac{Y_L^2 + Y_R^2}{2}\right), \tag{12}$$

where Y_L and Y_R represent the left- and right-handed coupling of the quarks to the Z' [8] and $\Theta(\hat{\eta})$ embodies the angular distribution of the Z' decay products,

$$\Theta(\hat{\eta}) = \frac{3}{4}\left(\frac{1 - \tanh^2\hat{\eta}}{\cosh^2\hat{\eta}}\right). \tag{13}$$

Here, the limits of integration on the lab frame rapidity are $|y| < \log(\sqrt{s}/M_{Z'})$. It is possible to integrate the full differential distribution including the decay of the scalar by applying the narrow-width approximation to the scalars. Specializing to the case $Z' \to h'a'$, the fraction of events with four isolated photons can be computed exactly, yielding

$$\frac{\sigma(\Delta R > \Delta R_{\min})}{\sigma_{\mathrm{tot}}} = \frac{6\langle y_{\max} - y\rangle}{\hat{\gamma}^2_{h'}\hat{\gamma}^2_{a'}\Delta R^4_{\min}}, \tag{14}$$

with

$$\langle y\rangle = \frac{\sum_q \int dy\, y\sigma_0\rho_{q,\bar{q}}(x_1, x_2)}{\sum_q \int dy\sigma_0\rho_{q,\bar{q}}(x_1, x_2)}, \tag{15}$$

where the limits of integration on the lab-frame rapidity were quoted above. Here we assumed $E^T_{\min} = 0$ for simplicity, but an analytic expression can be obtained also in the case of $E^T_{\min} \neq 0$. We note that the number of events with four isolated photons drops extremely rapidly with the rest-frame boost $\hat{\gamma} = M_{Z'}/2m_s$. Using the above expressions, we can also arrive at a simple approximate formula for the fraction of events where one diphoton pair is separated by less than $\Delta R_{\min}$ whereas the other is separated by more than $\Delta R_{\min}$ ($\Delta R_1 < \Delta R_{\min}$ and $\Delta R_2 > \Delta R_{\min}$):

$$\frac{\sigma(\Delta R_{1,2} \lessgtr \Delta R_{\min})}{\sigma_{\mathrm{tot}}} = \frac{6\langle\tanh(y_{\max} - y)\rangle}{\hat{\gamma}^2_*\Delta R^2_{\min}}, \tag{16}$$

where $\hat{\gamma}^2_* = 2/(\hat{\gamma}^{-2}_{h'} + \hat{\gamma}^{-2}_{a'})$ is the harmonic mean of the two scalars' squared boosts. In Figs. 1 and 2 we plot the approximate analytic formulas (16) and (14) for the

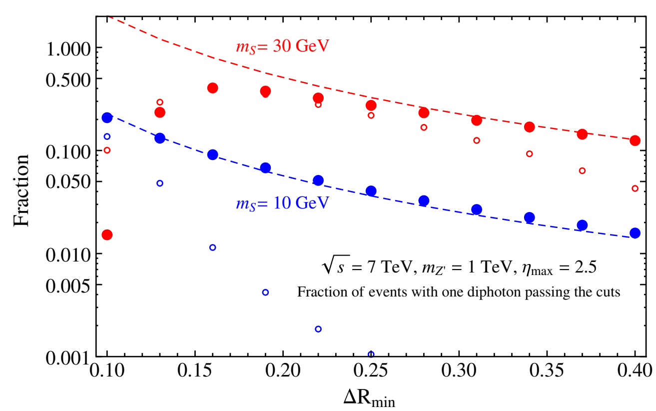

FIG. 1 (color online). The fraction of events with only one diphoton satisfying $\Delta R > \Delta R_{\min}$ for $M_{Z'} = 1$ TeV and two choices of the scalar mass. The dashed lines represent the approximate formula, Eq. (16) whereas the filled data points are from a Monte Carlo simulation of the full event at the LHC with $\sqrt{s} = 7$ TeV. The hollow data points are simulation including a transverse energy cut on the photons of 10 GeV.

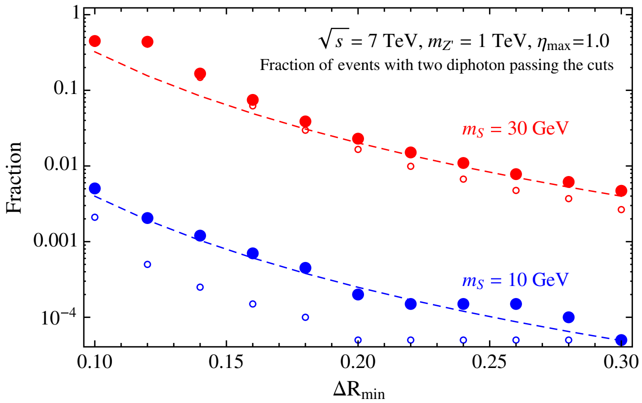

FIG. 2 (color online). The fraction of events with two diphotons both satisfying $\Delta R > \Delta R_{\rm min}$ for the same choice of parameters as in Fig. 1. We note the expected deterioration of the approximation, Eq. (14), as $\Delta R_{\rm min}$ diminishes as well as when the scalar mass increases resulting in a smaller boost. The loss of accuracy at higher $\Delta R_{\rm min}$ is only due to limited statistics.

fraction of events with one (two) diphotons satisfying $\Delta R > \Delta R_{\rm min}$, together with results of Monte Carlo simulation of the full event using the MadGraph4 package [14]. In each case we have simplified to the case of equal scalar masses for simplicity, i.e., $m_{h'} = m_{a'} = m_S$.

The above formulas make precise what is intuitively clear, namely that when the ratio of the Z' mass to that of the scalar is large, $M_{Z'}/m_s \gg 1$, most of the events will consist of two diphoton pairs, i.e., of two photon jets. The above formulas are useful for quick and reasonable estimates for the fraction of phase space where photon jets are important in a variety of circumstances. Depending on the detailed mass spectrum of the scalar sector, events may contain even more photon jets. This occurs for example when the CP-odd scalar is the only light particle. In that case, events where the Z' decays into $h'a'$ (or $H'a'$) will result in three photon jets as the CP-even scalars decay into two CP-odd ones ($h' \to 2a'$ or $H' \to 2a'$). Such events are particularly interesting since they contain several distinct resonances. The angular distribution of the scalars in the event also carries important information, as is well known (e.g., Ref. [15]).

It remains to consider the experimental reconstruction of photon jets, in particular whether they would pass standard photon identification requirements, and how they may be discriminated from both individual photons and isolated neutral pions. As a careful consideration requires full detector simulation, we restrict this discussion to several qualitative observations. The experimental signature of a photon jet in the electromagnetic calorimeter (ECAL) may not be resolvable as either one or two standard isolated photon objects. Less-boosted photon jets produce photons that are physically separated at the ECAL by several Moliere radii, so that their showers will not overlap but the resulting photon objects will generically fail isolation requirements. More-boosted photon jets result in overlapping showers from the two photons, producing a statistically broader $\eta - \phi$ profile in the electromagnetic calorimeter than normal photons. Even in this case, such objects will likely fail tight photon definitions, which are designed in part to reject the similar signal from decays of neutral pions and η's. At ATLAS for example, the first layer of the electromagnetic calorimeter is very finely spaced, $\Delta\eta = 0.003$, which allows rejection of extremely boosted pions. The above remarks apply to unconverted photons, but the conversion probability of photons in the tracker is significant. The double conversion of both photons may allow for a reconstruction of the photon pair invariant mass and serve as a discriminant against neutral pion decay.

Thus, if the scalars are lighter than the Z', the model discussed above would generically not give conventional multiphoton signatures with well-isolated photons [16]. One exception already noted is when the two photons in a photon jet merge into a single reconstructed photon that passes quality cuts, giving rise to an apparent diphoton resonance of spin 1. However, over most of the parameter space it is likely that a dedicated reconstruction algorithm would yield significantly higher efficiency. High-mass diphoton events will also arise in events where one photon in each photon jet is soft and/or wide angle enough that its harder companion passes isolation requirements. As can be seen from Figs. 1 and 2, this happens with very low efficiency, but the subset of events with two isolated photons would appear as a broadly peaked signal in a diphoton. Present limits on Randall-Sundrum diphoton resonances at TeV masses, on the scale of 3–10 fb [17], only mildly constrain the parameter space of interest once the inefficiency of photon isolation for these signals is accounted for.

We close by mentioning several novel search possibilities. In the simplest case of intermediate mass scalars, where no strong collimation is expected, the particular model presented in this Paper motivates searches for a new Z' resonance decaying into multiple photons. On the other hand, when the scalars, or more naturally only the pseudscalar, are much lighter than the Z' mass, the majority of the diphoton pairs will be highly collimated and fail typical photon identification requirements. This phase of the theory prompts the consideration of a photon jet as a distinct type of object in analogy with the recently proposed lepton jets [18]. We stress that such Z' need not be extremely heavy and searches for intermediate mass resonances are well motivated. Finally, searches for doubly converted photons may be particularly powerful if the light resonance mass can be accurately reconstructed from the charged tracks and used as a discriminant against background.

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