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Using a bottom up phenomenological approach we constructed a simple doubly charged vector lepton $E^{\pm\pm}$ model for the possible 750 GeV diphoton resonance Φ at the LHC assuming it to be a scalar particle. Since no stable doubly charged leptons are seen, to facilitate their decays we complete the model by adding a charged SM electroweak scalar $S^{\pm}$. Φ is a SM singlet and can be either an inert scalar or a Higgs field. In the inert case more than one vector lepton are required to account for the photon fusion production of the resonance if the model is to remain perturbative. For a Higgssed case $S^{\pm}$ can assist the production mechanism without using more than one such lepton. We also found that precision measurements constrain the couplings of $E^{\pm\pm}$ and $S^{\pm}$ to SM particles to be small. This raises the possibility that they can be fairly long lived and can give rise to displaced vertices if produced at the LHC.

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A Standard Model (SM) singlet scalar is an important ingredient in the popular Higgs portal [1][2] scenario for dark matter models. In the simplest case the only SM field Φ interacts with is the Higgs field H via interactions such as $\Phi^{\dagger}\Phi H^{\dagger}H$ and possibly $\Phi H^{\dagger}H$ depending on whether an extra symmetry is invoked. This makes Φ very difficult to detect both at high energy collider experiments and low energy precision measurements. This is because effects of Φ can only arise through mixing with the Higgs boson and this mixing is known to be small < 0.04 [3]. Hence, it is important to explore ways to induce/enhance couplings of Φ to other SM fields. In doing so the detection probability Φ will be increased as there are more channels to explore if it is light enough to be produce at the LHC or a future circular collider. A simple possibility is to add vector like fermions that are charged under the SM gauge symmetries and they can couple to Φ. These fermions are vectorlike due to anomaly considerations. The simplest case is a $SU(2)$ singlet vector lepton with $U(1)$ hypercharge $Y \neq 0$. For $Y = 1$ such a lepton will mix with the righthanded SM leptons and thus lead to fine tuning of parameters of such model. Moreover, for $Y \geq 2$ there will be no tree level mixing with SM leptons and this greatly simplifies the analysis of such models. In this paper we will examine the phenomenology of adding to the SM a $Y = 2$ vector lepton which carries two units of electric charge. It will be denoted by $E^{\pm\pm}$. An immediate observation will be that Φ can now decay into 2γ and $Z\gamma$ via 1-loop effects of the vector lepton and thus making the Higgs portal particle directly observable at the LHC if it is sufficiently light and the parameters of the model are favorable.

In addition to the ability of enhancing the detectability of a portal Higgs scalar, $E^{\pm\pm}$ will also naturally lead to lepton flavor violating processes. We will explain this assertion later. Since the SM offers no understanding of why there are three generations of chiral fermions with masses apparently generated by electroweak symmetry breaking it is important to explore other avenues in flavor physics. Doubly charged vector leptons is one such venue which as far as we know has not being fully explored.

Recently the ATLAS and CMS collaborations have reported the observation of excess of events in the diphoton mass distribution around 750 GeV [4],[5] in the $\sqrt{s} = 13$ TeV Run II data recorded with $3.2\,\text{fb}^{-1}$ of pp collisions. This unexpected development has understandably generated a great deal of interest among theorists. It is common to interpret this in terms of a new spin-0 resonance, although a spin-2 particle is not ruled out . It is also noted the same is not seen in the dijet mass spectrum. On the other hand, the event rate and lack of signal at lower energies appear to favor a scalar or pseudo scalar resonance being produced and decays predominantly into two photons. A plausible explanation would be that this resonance couples predominantly to two photons and the couplings to gluons and other colored objects are suppressed or not exist at all. This was first discussed in [6],[7]. In this case, the resonance is produced via two photon fusion and both exclusive and inclusive processes can take place with the latter being more important. It can also be shown that other fusion mechanism such as photon-Z and two Z bosons are less important. More re-

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cent LHC Run II data from both CMS [8] and Atlas [9] do not support the initial data. With $12.9\,\text{fb}^{-1}$ pp collions collected no new signal above background was recorded. This reduces the local significance of the initial excess from $\sim 3.6\sigma$ to $\sim 2.3\sigma$.

In this paper we will take the data at face value and interpret the combined results of 2015 and 2016 data as an upper limit on the the photon fusion production of a spin 0 resonance of mass 750 GeV. This is to be taken as example of the limits on the parameters of the model we will discuss later. As reported by Atlas two of the 15 original excess events are consistent with background; and no new excess is found in the 2016 run thus far we estimate this cross section to be $\lesssim 0.84$ fb. [6]. Focusing on the case of a scalar ϕ the effective Lagrangian for the above process is given by

$$\mathcal{L} = \frac{1}{f_\gamma}\phi(F_{\mu\nu})^2 \tag{1}$$

and $f_\gamma \gtrsim 8.6$ TeV [6]. Since Eq.(1) can only come from a 1-loop effect this will set a limit on the coupling of ϕ to the particles in the loop .

In this paper we take ϕ to be a Standard Model (SM) singlet scalar field and identify it as the real part of a complex Higgs portal field discussed previously. We study the possibility that it is a bridge to new vector leptons E of hypercharge $Y = 2$ or higher which are also $SU(2)$ singlets. We do not extend the gauge symmetry of the SM and the vector nature of these leptons will not lead to anomalies. The number of such particles is not known. The physics we wish to explore is well illustrated by considering just one such particle. Extending to more is straightforward. Similarly higher charged particles can also be considered. For simplicity we take $Y = 2$ and comment on other possibilities when appropriate. Since there are no stable doubly charged leptons, it is mandatory that E has decay channels. Charge and angular momentum considerations dictate the decay to be $E^{--} \to \ell + \ell' + \nu^c$ where $\ell, \ell' = e, \mu, \tau$ and ν is an active SM neutrino. The flavor of ν is different from ℓ and/or ℓ'.[1] An explicit example will be $E^{--} \to e^-\mu^-\nu_\tau^c$. If we assign unit lepton number to E^{--} the decay conserves global lepton number; however,lepton flavor will be violated. The scale of lepton flavor violation will be given by the mass of the mediating particle that gives rise to the above decay. No SM fields can lead to the above decay. The simplest solution is to introduce a $Y = 1$ $SU(2)$ singlet scalar S. If S is lighter than E then the decay will be sequential : $E \to S + \ell$ followed by $S \to \ell' + \nu^c$. On the other hand, if E is lighter than S, the decay will be a 3-body mode similar to that of ordinary muon decays.

The quantum numbers of the new particles together with the relevant SM fields are given in Table (1) below

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[1] $E^{--} \to \ell + W^-$ is allowed by charge and angular momentum considerations but forbidden by SM gauge symmetries.

TABLE I: Quantum numbers of the SM Higgs H, leptons L, ℓ and E, S, ϕ

Field	$SU(2)$	$U(1)_Y$
H	**2**	$\frac{1}{2}$
L	**2**	$-\frac{1}{2}$
ℓ_R	**1**	-1
E	**1**	-2
S	**1**	1
Φ	**1**	0

where standard notations are used.

In addition to the SM Lagrangian that involving the new fields is given by

$$\begin{aligned}
\mathcal{L}' & \\
\overline{E}i\gamma^\mu(\partial_\mu - 2ig_1B_\mu)E & + [(\partial^\mu + ig_1B^\mu)S]^\dagger(\partial_\mu + ig_1B_\mu)S \\
- \big[f_{e\mu}(\overline{\nu_e^c}\mu_L - \overline{\nu_\mu^c}e_L) & + f_{e\tau}(\overline{\nu_e^c}\tau_L - \overline{\nu_\tau^c}e_L) \\
+ f_{\mu\tau}(\overline{\nu_\mu^c}\tau_L - \overline{\nu_\tau^c}\mu_L)\big] & S - y_E\overline{E}E\Phi - M_E\overline{E}E \\
- \sum_a^{e,\mu,\tau} y_a\overline{E}\ell_{Ra}S^\dagger & - V(H,S,\Phi) + h.c.
\end{aligned} \tag{2}$$

The scalar potential $V(H, S, \Phi)$ is

$$\begin{aligned}
V = & -\mu^2 H^\dagger H + \lambda(H^\dagger H)^2 + M_S^2 S^\dagger S + \lambda_S(S^\dagger S)^2 \\
& + \lambda_{SH}S^\dagger SH^\dagger H + \lambda_\phi(\Phi^\dagger\Phi)^2 + M_\phi^2\Phi^\dagger\Phi + \lambda_{\phi h}\Phi^\dagger\Phi H^\dagger H \\
& + \lambda_{\phi S}\Phi^\dagger\Phi S^\dagger S + \alpha\Phi + \beta\Phi^\dagger\Phi\Phi + \kappa_H\Phi H^\dagger H \\
& + \kappa_S\Phi S^\dagger S
\end{aligned} \tag{3}$$

In general, Φ can be complex. For simplicity we take Φ to be real. The imaginary part plays no role in what we are studying here since we take ϕ in Eq.(1) to be a scalar. The usual portal Higgs potential can be obtained from Eq.(3) by deleting the S field. How Φ connects to the dark is model depend and will not be pursued here. Note that lepton number is conserved in this model and neutrinos remain massless. In order to give masses to active neutrinos one can implement Type I seesaw by adding heavy singlet neutrinos or radiatively generate them by adding a second scalar doublet as in the Zee model [10]. Though interesting we shall not pursue this further here.

Taking the U-gauge for the Higgs doublet, we parameterize H and Φ by

$$H = \begin{pmatrix} 0 \\ \frac{v+h}{\sqrt{2}} \end{pmatrix}, \quad \Phi = \frac{w+\phi}{\sqrt{2}}. \tag{4}$$

where v, w are the respective VeV of H and Φ fields. The

stationary conditions for H, Φ are

$$v\left(-\mu^2 + \lambda v^2 + \frac{\lambda_{\phi h} w^2}{2} + \frac{\kappa_H w}{\sqrt{2}}\right) = 0,$$

$$w\left(M_\phi^2 + \lambda_\phi w^2 + \frac{\lambda_{\phi h} v^2}{2} + \frac{\alpha}{\sqrt{2}w} + \frac{3\beta w}{2\sqrt{2}} + \frac{\kappa_H v^2}{2\sqrt{2}w}\right) = 0.$$

(5)

If $w = 0$ then Φ is not in the Higgs phase. However, this requires $\alpha + \kappa_H v^2/2 = 0$.

A second possibility is $w \neq 0$ and Φ is also a Higgs field. The stationary condition can easily be satisfied for $M_\phi^2 < 0$ although not the only possibility. The trilinear terms ϕhh and $\phi S^+ S^-$ will be present whether Φ is in the Higgs phase or not. Furthermore, h and ϕ will in general mix.

If $w = 0$ then their mass square matrix of (h, ϕ) is expressed in

$$\frac{1}{2}\begin{pmatrix} h & \phi \end{pmatrix}\begin{pmatrix} 2v^2\lambda & \frac{v\kappa_H}{\sqrt{2}} \\ \frac{v\kappa_H}{\sqrt{2}} & \bar{M}_\phi^2 \end{pmatrix}\begin{pmatrix} h \\ \phi \end{pmatrix},$$

(6)

where $\bar{M}_\phi^2 = M_\phi^2 + \lambda_{\phi h} v^2/2$. (h, ϕ) is related to the mass eigenstates (h', ϕ') by the usual 2×2 rotation matrix define by the mixing angle θ which is given by

$$\tan 2\theta = \frac{\sqrt{2}v\kappa_H}{\bar{M}_\phi^2 - v^2\lambda}.$$

(7)

If we identify ϕ' as the 750 GeV resonance and h' as the SM like Higgs with mass 125 GeV it is natural to assume $M_\phi > v$. In the limit $\kappa_H \to 0$ the two fields will decouple from each other. The observed Higgs boson is SM-like and the mixing with another scalar is limited by $\sin^2\theta \lesssim 0.04$ [3] from an analysis of Higgs coupling strength data from LHC run-I. We also note that another analysis [12] gives a larger value of $\sin^2\theta \lesssim 0.33$. Using the more stringent constrain we estimate that $\frac{\kappa_H}{v} \lesssim 5$. it is interesting that current data allow κ_H to be $O(\text{TeV})$. It can be much smaller if the data on SM Higgs couplings become more stringent.[2] On the other hand κ_S remains unconstrained.

For notational simplicity we shall drop the prime in the mass eigenstates.

For $w \neq 0$ the neutral scalar mass square matrix is more complicated. It can be obtained from Eq.(6) by the following substitutions: $\bar{M}_\phi^2 \to \lambda_\phi w^2 - \frac{\alpha}{\sqrt{8}w} + \frac{3\beta}{\sqrt{8}} - \frac{\kappa_H v^2}{\sqrt{32}w}$ and $\frac{\kappa_H v}{\sqrt{2}} \to \lambda_{\phi h} vw + \frac{\kappa_H v}{\sqrt{2}}$. The mixing is given by Eq.(7) with the above substitutions.

The mass parameters M_E, M_S, M_ϕ in Eq.(3) are all free parameters. The physical masses depends on

whether Φ is in the Higgs phase or not. The mass of E is just M_E since there is no mixing with the SM leptons. For $w = 0$, the mass S is given by $M_S^2 + \lambda_{SH} v^2/2$. Similarly the mass of ϕ is approximately given by $M_\phi^2 + \lambda_{\phi h} v^2/2$. The relative size of these masses cannot be determined. The case of Higgsed Φ is similar with more complicated formulas. For definiteness we take E to be heavier than S and the physical mass of S is greater than 80 GeV from LEPII searches [11].

With the Lagrangian in place the effective Lagrangian Eq.(1) is obtained from calculating Figs.(1)
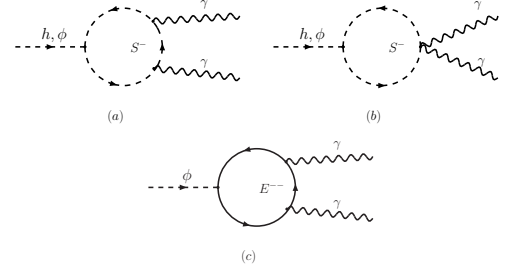

FIG. 1: (a,b) ϕ to two photons via charged S scalar loop for Φ in the Higgs phase. They also contribute to SM Higgs decays. (c) The E loop contribution

All three diagrams will contribute to ϕ decays whether it is Higgsed or not. For the same SM Higgs decays and neglecting the small $\phi - h$ mixing the S loop will contribute but the E loop will not. The calculation of the above diagrams give [13]

$$f_\gamma^{-1} =$$
$$\frac{\alpha}{4\pi M_\phi}\left(Q^2 N y_E \sqrt{\tau_E} F_{\frac{1}{2}}(\tau_E) + \frac{2(\lambda_{\phi S} w + \kappa_S)}{M_S}\sqrt{\tau_s} F_0(\tau_s)\right).$$

(8)

We have used the convention of [14] and define $\tau_i = M_\phi^2/(4M_i^2)$ and the 1-loop functions are

$$F_0(\tau) = -[\tau - f(\tau)]\tau^{-2} \tag{9}$$
$$F_{\frac{1}{2}}(\tau) = 2[\tau + (\tau - 1)f(\tau)]\tau^{-2} \tag{10}$$

with

$$f(\tau) = \begin{cases} \arcsin^2\sqrt{\tau} & \tau \leq 1 \\ -\frac{1}{4}\left[\log\frac{1+\sqrt{1-\tau^{-1}}}{1-\sqrt{1-\tau^{-1}}} - i\pi\right]^2 & \tau > 1. \end{cases} \tag{11}$$

From the event rates given one can deduce that $f_\gamma \sim 8-9$ TeV [6]. Eq.(8) gives strong constraints on the model parameters since the F functions are known. We plot them in Fig.(2).

First we consider Φ is not a Higgs scalar i.e. $w = 0$. Since there is no constrain on the value of κ_S we consider two limits

1. We first take it to be small, i.e. $\kappa_S << v$. Then the E-loop will have to account for the observed

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[2] Since $\phi \to hh$ is allowed and we will require that it does not dominate over the diphoton mode. We then obtain the constrain $\frac{\kappa_H}{M_\phi} \lesssim O(\frac{\alpha_{EM} y_E}{4\pi})$. This will become clear later.
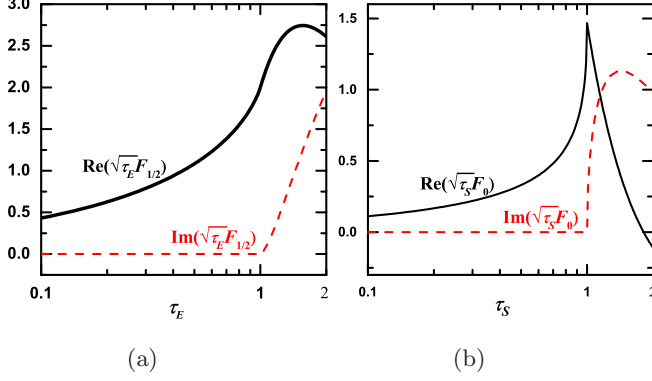
FIG. 2: Form factors for (a) spin-$\frac{1}{2}$ and (b) spin-0 particle contributions to $\gamma\gamma$ couplings to a scalar as a function of $\tau_i = M_\phi^2/4M_i^2$ with M_i mass of the loop particle.

events. For the simplest case of one E we have $Q = 2, N = 1$. The function $\sqrt{\tau}F_{\frac{1}{2}}(\tau)$ has the value of 2.0 at $\tau = 1$ and falls rapidly for $\tau < 1$. As a benchmark point we take $f_\gamma = 9$ TeV and obtain the constraint

$$y_E \ N \simeq 16.8. \tag{12}$$

Clearly for $N = 1$ the Yukawa coupling is so large as to invalidate perturbative calculations. In order for y_E to reside in the perturbative regime, i.e. $< 4\pi$, would require $N > 2$. In the region of $\tau_E < 1$, ϕ will not decay into $E\bar{E}$ pairs and hence will not lower the two photon branching ratio.

For the region $1 < \tau_E < 2$, we have $Re(\sqrt{\tau}(F_{\frac{1}{2}}))$ falling slowly and the imaginary part rising very fast by comparison. Taking the peak value of $\sqrt{\tau}F_{\frac{1}{2}} = 2.5$ will still require $y_E \sim O(1)$ unless $N > 10$. Barring the very high multiplicity case, $y_E \sim O(1)$ will imply that the dominant decay of ϕ is into $E\bar{E}$ pairs instead of the two photons mode.

An alternative to adding more doubly charged leptons will be to add a higher charged vector lepton. For example, we can add a $E^{\pm\pm\pm}$. This will have the same effect as adding two $E^{\pm\pm}$ assuming that their Yukawa couplings are not too different. We note that the triply charged E is also unstable and can decay via $E^{---} \to E^{--} + S^-$.

2. Next we take $\kappa_S \gtrsim v$. We have seen that $\kappa_H \sim$ TeV is allowed by the mixing data perhaps this not unreasonable domain for κ_S to be in. The scalar loop contribution is given by F_0, see Eq.(8). It is smaller than $F_{\frac{1}{2}}$ by a factor of 2 for the same value of the arguments in the region $\tau < 1$. This can be compensated by adjusting κ_S/M_S and we set $w = 0$ for now. The constraint is given by

$$N y_E \sqrt{\tau_E} F_{\frac{1}{2}}(\tau_E) + \frac{\kappa_S}{2M_S}\sqrt{\tau_s}F_0(\tau_s) \lesssim 34 \tag{13}$$

To have a significant effect, we require $\kappa_S \gtrsim 10$TeV. For κ_S in the $\gtrsim 20$ TeV range the scalar loop dominates and all couplings can remain perturbative even for $N = 1$.

The same considerations can be applied to the case of a Higgssed Φ. The main difference is we require $\lambda_{\phi S}w + \kappa_S$ to be in the 10 TeV range and the scalar loop will dominate and all couplings can remain perturbative. To give an example, let $w(\kappa_S) = 10(15)$TeV and $M_S = 400$ Gev and $M_E = 500$ GeV we obtain $y_E = 2.6$ for $\lambda_{\phi S} = 1$ and $N = 1$. Moreover, for large values of w a mild fine tuning of λ_ϕ is required to get a 750 GeV scalar. There is a further fine tuning problem with this solution. The physical mass of $S^\pm$ is given by

$$M_S^2(phy) = M_S^2 + \frac{\lambda_{SH}v^2}{2} + \frac{\lambda_{\phi S}w^2}{2} + \frac{\kappa_S w}{\sqrt{2}} \tag{14}$$

The sum of the first two terms has to be negative and large in order to provide cancelation to the large κ_S and w contributions, i.e. $O(10)$ TeV, so that $M_S(phy) \sim 0.4$ TeV . It is easy to see from Eq.(3) that the extremum condition on S can be written as $M_S^2(phy) + 2\lambda_S S^\dagger S$ at the electroweak and singlet minimum and is positive definite. Hence, there is no charge breaking vacuum here.

In passing we note that the region of $\tau_s > 1$ is ruled out since $\phi \to S^+S^-$ will be the dominant decay.

In either phase of the Φ field, both E and S have to be heavier than $M_\phi/2$ in order for the model to accommodate the current data for the 750 resonance. However, which one is heavier cannot be determined yet.

The model we constructed with new fields carrying only $U(1)_Y$ quantum numbers will lead to the prediction that the ratios of widths into SM gauge bosons are [15]

$$\Gamma_{\gamma\gamma} : \Gamma_{\gamma Z} : \Gamma_{ZZ} = 1 : \frac{2s_w^2}{c_w^2} : \frac{s_w^4}{c_w^4}$$
$$\approx 1 : 0.54 : 0.07 \tag{15}$$
$$\Gamma_{WW} = 0.$$

where $s_w(c_w)$ is the sine(cosine) of the weak mixing angle and in the limit that κ_H is small.

An important consideration of introducing new heavy charged states is to examine the constraints of low energy precision measurements put on their masses and couplings. If $E^{\pm\pm}$ and $S^\pm$ were to play roles in the diphoton resonance as we have discussed above, then both masses have to be > 375GeV. This is higher than the constraints imposed from the model independent bound from LEP II since no doubly charged leptons nor charged scalars were seen [16] . Closer examination of Fig.(2) reveals that if M_S stays close to 375 GeV it offers the most impact while M_E can be larger due to the slower fall off of $Re(\sqrt{\tau}F_{\frac{1}{2}})$. As stated before we take the vector lepton to be the heavier one and use $M_S = 400$ GeV as a benchmark. For the opposite case of $M_S > M_E$ there is no qualitative difference. Quantitatively since F_0 falls off very fast as M_s increases then a larger κ_S will be required.

Next we examine the constraints from low energy physics. The exchange of S in muon decays will modify the Fermi coupling G_F as measured by muon lifetime. With new physics in the leptonic sector we assume instead unitarity of quark mixing and extract the G_F from nuclear, kaon and B-meson decays [17]. This gives $G_F = 1.166309(350) \times 10^{-5} \text{GeV}^{-2}$. The effective Lagrangian due to S exchange yields

$$\mathcal{L} = \frac{i f_{e\mu}^2}{2M_S^2} \left(\overline{\nu_\mu} \gamma^\alpha \hat{L} \nu_e \right) \left(\bar{e} \gamma_\alpha \hat{L} \mu \right) \qquad (16)$$

whereas the SM has $-\frac{ig^2}{2M_W^2}$ in front of the 4-fermi operator. Here $\hat{L} = (1 - \gamma_5)/2$. Thus, we obtain

$$f_{e\mu} \leq 1.502 \times 10^{-1} \left(\frac{M_S}{400\text{GeV}} \right). \qquad (17)$$

Similarly using the leptonic τ decays ratio into μ, e we get

$$\frac{\Gamma(\tau \to \mu + \nu's)}{\Gamma(\tau \to e + \nu's)} = \frac{\left(1 - \frac{f_{\mu\tau}^2 M_W^2}{g^2 M_S^2} \right)^2 + \cdots}{\left(1 - \frac{f_{e\tau}^2 M_W^2}{g^2 M_S^2} \right)^2 + \cdots} \qquad (18)$$

$$\simeq 1 + 2(f_{e\tau}^2 - f_{\mu\tau}^2) \left(\frac{M_W^2}{g^2 M_S^2} \right)$$

where $\cdots$ denotes terms such as $f_{\mu e}^2 f_{\tau e}^2$ which come from diagrams that interfere incoherently with the SM ones. They are of order f^4 which we neglect. Thus,

$$f_{e\tau}^2 - f_{\mu\tau}^2 \leq \pm 2.25 \times 10^{-2} \left(\frac{M_S}{400\text{GeV}} \right)^2. \qquad (19)$$

where the experimental value of $\frac{\Gamma(\tau \to \mu^- \bar{\nu}_\mu \nu_\tau)}{\Gamma(\tau \to e^- \bar{\nu}_e \nu_\tau)} = 0.979 \pm 0.004$ has been used [18].

Next we consider the rare decay of $\mu \to e\gamma$. The Feynman diagrams to calculate are given by Fig.(3). We have
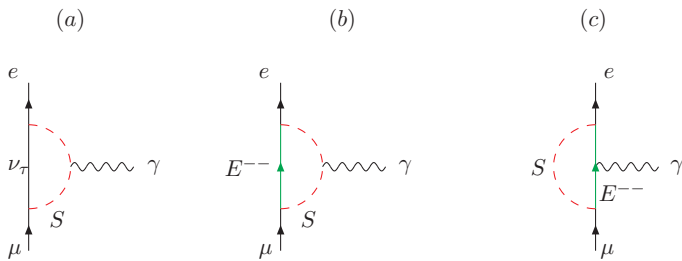

FIG. 3: Diagrams leading to $\mu \to e\gamma$. Wavefunction renormalization graphs are not shown

set $m_e = 0$ and note that diagram (a) has a different chiral structure than (b),(c) and they add incoherently. Using the experimental bound of $BR(\mu \to e\gamma) < 5.7 \times 10^{-13}$

[19] we get the strong constraint

$$f_{e\tau}^2 f_{\mu\tau}^2 + \left(\frac{y_e y_\mu x^2}{(1-x)^4} \right)^2 \left[(-4 + 9x - 5x^3) + 6x(2x - 1) \ln x \right]^2$$

$$\leq 2.235 \times 10^{-12} \left(\frac{M_s}{400\text{GeV}} \right)^4, \qquad (20)$$

where $x = \frac{M_S^2}{M_E^2}$. Assuming that the two terms on the left hand side are of the same order then $f_{e\tau} f_{\mu\tau} \lesssim 10^{-3}$. For $x < 1$ the terms multiplying $y_e y_\mu$ can give a large coefficient, e.g. for $x = 0.64$ this factor is ~ 18.18. This implies that $y_e y_\mu$ is about the same order of magnitude as that of the f's. Certainly it can be smaller in which case the $f_{e\mu} f_{\mu\tau}$ will saturate the bound.

Similar diagrams with the final electron state replaced by a muon will contribute to muon anomalous moment a_μ. This contribution is

$$\Delta a_\mu = \frac{m_\mu^2}{96\pi^2 M_S^2} \left((f_{\mu\tau}^2 + f_{e\mu}^2) - \frac{y_\mu^2 x}{(1-x)^4} \left[4 - 9x + 5x^3 - 6x(2x - 1) \ln x \right] \right) \qquad (21)$$

Putting in the numbers we get

$$f_{\mu\tau}^2 + f_{e\mu}^2 - \frac{y_\mu^2 x}{(1-x)^4} \left[4 - 9x + 5x^3 - 6x(2x - 1) \ln x \right]$$

$$\leq 39.08 \left(\frac{M_S}{400} \right)^2 \qquad (22)$$

where we have used $a_\mu^{exp} - a_\mu^{SM} = 2.88 \times 10^{-9}$ [20]. The above considerations constraint most of the parameters of Eq.(2).

On the other hand, there are fewer limits on the parameters of the $V(H, S, \Phi)$. Since the potential preserves custodial symmetry of the SM, most electroweak precision measurements takes the SM values. Moreover, from Fig.(1) we see the scalar loop can modify the $h \to \gamma\gamma$ signal comparing to the SM. In contrast the vector lepton makes no contribution at this level. Experimentally the signal strength of $h \to \gamma\gamma$ is very close to the SM value we can use it to constrain λ_{SH}. Define $R \equiv \Gamma^{new}/\Gamma^{SM}$ we obtain

$$R = \left| 1 + \frac{\lambda_{SH} v^2}{2M_S^2} \frac{F_0(\tau')}{F_1(\tau_w) + \frac{4}{3} F_{\frac{1}{2}}(\tau_t)} \right|^2 \qquad (23)$$

where $\tau' = M_h^2/4M_S^2$ and $F_1(\tau) = -[2\tau^2 + 3\tau + 3(2\tau - 1)f(\tau)]\tau^{-2}$. The current bound on R is 1.17 ± 0.27 [21]. This yields the constraint: $|\lambda_{SH}| < 8.1$.

The above can be used to estimate the lifetimes of E and S. Since we are interested in the case of E being

heavier than S which is the favored region from the perturbative view point. The main decays will be $S \to \ell + \nu$ and $E \to S + \ell$. The results are

$$\Gamma_S = \frac{M_S}{8\pi} \sum_{\ell' \neq \ell} \sum_{\ell} |f_{\ell\ell'}|^2, \qquad (24)$$

and

$$\Gamma_E = \frac{M_E}{32\pi} \left(1 - \frac{M_S^2}{M_E^2}\right)^2 \sum_{\ell} |y_l|^2. \qquad (25)$$

where the light lepton masses are all neglected. From Eqs.(17,19,20) we expect $f_{\ell\ell'} \lesssim 10^{-3}$ [see discussions following Eq.(20)]. For a 400 GeV S its lifetime is 1.4×10^{-20}sec. with the assumption $f_{e\mu} \sim f_{e\tau} \sim f_{\mu\tau} \sim 10^{-3}$. For smaller f's, e.g. $\sim 10^{-5}$ the lifetime is long enough to give displaced vertices which can aid in its detection when they are produced at the LHC. The lifetime for a 500 GeV vector lepton decaying to a 400 GeV S is estimated to be $\sim 3.4 \times 10^{-19}$ sec. These lifetime estimates are important for finding signatures for the production of these new particles. At the LHC production of S proceeds via quark anti-quark annihilation at the parton level

$$q + \bar{q} \to S^+ S^- \to \ell^+ \nu \ell' \nu^c \qquad (26)$$

The signature is two lepton pairs which need not have the same flavor and missing transverse energy, $\not{E}_T$ with no associated jets. However, the background from W boson pair production will be severe. On the other hand, at an $e^+ e^-$ collider the signals for $S^+ S^-$ will be much easier to unravel.

Similarly the production of E can be searched for by the sequence of reactions

$$q + \bar{q} \to E + \bar{E} \to S^- + \ell_a^- + S^+ + \ell_b^+ \to \ell_a^- + \ell_c^- + \ell_b^+ + \ell_d^+ + \not{E}_T \qquad (27)$$

where a, b, c, d denote the flavors of the charged leptons. Here the signal is four leptons plus $\not{E}_T$ with no associated jets. Furthermore, the charged leptons do not form invariant mass peaks. Interestingly if the couplings f's and y's are very small i.e. $< O(10^{-6})$ we will have displaced vertices as discussed before. As an example we take $y_e \simeq y_\mu \simeq y_\tau \sim 5 \times 10^{-6}$, which are values near the experimental limits, then the production of $E\bar{E}$ pairs will leave two $\sim$ 2mm tracks from the collision point. Each will subsequently lead to two more tracks depending on the decay modes. For smaller values of the y's longer tracks will be expected. For a discussion of displaced vertices search, see for example [22]. We have here the unusual case in which the LHC can cover very small couplings which precision measurements cannot reach in the foreseeable future.

Similar to the case of S high energy $e^+ e^-$ colliders offer cleaner signatures. Aside from the ratio of total cross sections to muon pair production R which gives 1 for $S^+ S^-$ and 4 for $E^{++} E^{--}$ respectively. If their production is way above threshold they can give spectacular signatures with a pair of same sign leptons going in one direction and a pair of anti-leptons going in the opposite direction. Looking for similar signatures at the LHC is more complicated. At the parton level, see Eq.(27), in the parton center of mass frame, the charged leptons pair and anti-leptons pair will emerge in opposite directions. However, since the quark and anti-quarks have different parton distribution functions, in the laboratory frame they will be boosted differently. Nevertheless, we can expect lepton and anti-leptons are still separated in rapidity. We will defer a detail study to a later investigation. A recent study of the Drell-Yan production of new particles related to the diphoton resonance is given in [23].

In conclusion, we employed a totally phenomenological approach to construct and study a simple model of doubly charged vector leptons $E^{\pm\pm}$ that may enhance the photon fusion production of singlet portal Higgs at the LHC. Using a 750 GeV diphoton scalar resonance as an example to evaluate the parameters of the model that can lead to its detectability. The coupling y_E will have to be $O(1)$. Furthermore, these vector leptons are phenomenologically interesting in their own rights and have been discussed in the context of excited leptons, see e.g. [24]. Moreover, they have to be unstable and thus must decay. A simple electroweak singlet charged scalar S is utilized to complete the model. It is found that if Φ is not in the Higgs phase one would require two or more vector leptons in order for the model to be amendable to perturbative treatments and account for the data if the κ_S term is small. On the other hand if κ_S is large then S can assist the lepton loop in giving a large enough effective coupling to accommodate the observed signal strength with only a single vector lepton required. The Yukawa couplings of Φ to these particles are of order 1 for the reference kinematic point we use. Given the preliminary nature of the data we did not pursue a detail parameters scan. This mechanism can be carried over to the case of a Higgssed Φ. Moreover, there is a price to be paid here in that fine tuning of parameters in the scalar potential is needed in order to keep $S^\pm$ relatively light and the singlet VeV in the tens of TeV range.

We also studied the low energy constraints on the model and found that couplings of the these new states to the SM matter fields must be small, i.e. $\lesssim O(10^{-3})$. Interestingly if theses couplings are $< 10^{-6}$ they can lead to displaced vertices of multileptons signals for the production and decays of $E^{\pm\pm}$ at the LHC for reference masses of $M_E = 500$ GeV and $M_S = 400$ GeV.

Other studies of doubly charged lepton contribution to a scalar diphoton resonance with different emphasis and context can be found in [25]. For an early summary of other approaches to the diphoton resonance see e.g.[26]

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