Based on tentative evidence for a peak in the Fermi gamma-ray spectrum originating from near the center of the galaxy, it has been suggested that dark matter of mass \( \sim 130 \text{ GeV} \) is annihilating directly into photons with a cross section \( \sim 24 \) times smaller than that needed for the thermal relic density. We propose a simple particle physics model in which the DM is a scalar \( X \), with a coupling \( \lambda_X X^2 |S|^2 \) to a scalar multiplet \( S \) carrying electric charge, which allows for \( XX \rightarrow \gamma \gamma \) at one loop due to the virtual \( S \). We predict a second monochromatic photon peak at 114 GeV due to \( XX \rightarrow \gamma Z \). The \( S \) is colored under a hidden sector SU(N) or QCD to help boost the \( XX \rightarrow \gamma \gamma \) cross section. The analogous coupling \( \lambda_h h^2 |S|^2 \) to the Higgs boson can naturally increase the partial width for \( h \rightarrow \gamma \gamma \) by an amount comparable to its standard model value, as suggested by recent measurements from CMS. Due to the hidden sector SU(N) (or QCD), \( S \) binds to its antiparticle to form \( S \)-mesons, which will be pair-produced in colliders and then decay predominantly to \( XX \), \( hh \), or to glueballs of the SU(N) which subsequently decay to photons. The cross section for \( X \) on nucleons is close to the Xenon100 upper limit, suggesting that it should be discovered soon by direct detection.

Dark matter (here denoted by \( X \)) should couple only weakly to photons, if at all, at tree-level \([9, 10]\). One way to insure the “darkness” of the DM is for it to couple to photons only via loops. At one loop, the DM should couple directly to charged particles \( S \). To make a renormalizable coupling of this type, both \( X \) and \( S \) must be bosons, since the stability of \( X \) and the conservation of charge require \( X^2 \) and \( |S|^2 \). This leads us to consider the interactions

\[
\mathcal{L}_{\text{int}} = \frac{\lambda_X}{2} X^2 |S|^2 + \lambda_h |H|^2 |S|^2 + \frac{\lambda_h}{2} |H|^2 X^2 \tag{1}
\]

between \( X \), the Higgs doublet \( H \), and \( S \). The second coupling is not necessary, but neither is there is any reason to forbid it, and in fact we will show that it can naturally give rise to an interesting enhancement in the \( h \rightarrow \gamma \gamma \) branching ratio, for the same values of the \( S \) mass and charge as needed to explain the Fermi line. The third coupling is useful for achieving the correct relic density of \( X \) \([11]\), as we will discuss. The stability of \( X \) is insured by the \( \Sigma_2 \) symmetry \( X \rightarrow -X \).

**Decays of \( S \).** It is necessary to make \( S \) unstable in order to avoid charged relics, on whose abundance there are very stringent bounds from terrestrial searches for anomalous heavy isotopes \([12, 13]\) and from their effects on big bang nucleosynthesis \([14, 15]\). We will also find it useful to let \( S \) transform under QCD or a hidden SU(N) gauge symmetry, in order to boost the cross section for \( XX \rightarrow \gamma \gamma \). Suppose \( S \) is in the fundamental representation of SU(N) for definiteness. If SU(N) is QCD and \( S \) has charge \( 4/3 \), it can decay into right-handed up-type quarks through the renormalizable operator \( \epsilon_{abc} S \bar{u}_{R, a} \bar{u}_{R, b} u_{R, c} \). If the SU(N) is exotic, then \( S \) could decay into a lighter, neutral fundamental representation field \( T \) and two charged right-handed fermions through a dimension 5 operator. For example, if \( S \) has charge \( q_S = 2 \), the decay into \( T + e^+ + e^+ \) occurs via the
which could arise from a renormalizable theory by integrating out a heavy colored fermion $N_a$ carrying charge 1. For $q_S = 1$, there is an analogous operator to (2) involving left-handed lepton doublets, $T^\alpha_\alpha S^a_e \tilde{L}_e \sigma^2_L e_e$, that mediate the decay $S \rightarrow T e^+ \tilde{\nu}$. We will focus on the example (2) because the larger charge $q_S = 2$ helps to increase the $XX \rightarrow \gamma \gamma$ cross section, and the decays into electrons can lead to interesting collider signatures. Of course, higher-generation leptons can also appear in addition to electrons through analogous operators, as well as lepton flavor-violating versions, whose contribution to the rare process $\mu \rightarrow e \gamma$ is suppressed by two loops and two powers of $M$.

Obviously the lifetime of $S$ depends upon the scale $M$ of the heavy $N_a$ particle. Estimating the decay rate for $S \rightarrow T e e$ as $m_S^2/(16\pi M^2)$ and demanding the lifetime to be less than $10^3$ s [14], we find the limit $M < 10^{15}$ GeV. We will show that the relic neutral particle $T_a$ binds into stable “baryons” that make a small contribution to the total dark matter population.

For simplicity we have assumed that additional couplings of $T$ such as $\lambda_{XT} X^2 |T|^2$ and $\lambda_{XY} |S|^2 |T|^2$ are small. The former could provide a significant annihilation channel $XX \rightarrow TT^*$ if $m_T < m_X/2$ (the factor of 1/2 coming from the fact that each $T$ must hadronize into $TT^*$ bound states) and if $\lambda_{XT}$ is sufficiently large, while the latter has no particular impact on the points that follow.

**Annihilation to two photons.** The model parameters relevant for the Fermi line are $\lambda_X$, the mass $m_X$, the charge $q_S$ (in units of $e$), the mass $m_S$, and the number of colors $N_c$ of QCD or the hidden SU(N) gauge group. The annihilation cross section corresponding to the diagrams of fig. 1 is given by

\[ \langle \sigma v \rangle = \frac{\sum |M|^2}{64 \pi m_X^2} \]  

where the squared matrix element, summed over photon polarizations, is

\[ \sum |M|^2 = \frac{\alpha^2}{2 \pi^4} q_S^2 \lambda_X^2 N^2 \tau^2 A_0^2(\tau) \]  

with $\tau = m_X^2/m_S^2$ and

\[ \tau A_0(\tau) = 1 - \tau^{-1} \arcsin(\sqrt{\tau})^2 \]  

for $\tau \leq 1$, which we presume to be the case. Eqs. (4,5) can be deduced by comparing to the well-known result for $h \rightarrow \gamma \gamma$ from fig. 2, in which $\tau \rightarrow m_h^2/4m_S^2$; see for example ref. [16].

Ref. [2] determined that $\langle \sigma v \rangle$ should be approximately 0.042 in units of the thermal relic density value $(\sigma v)_0 = 1$ pb-c, in order to explain the Fermi gamma-ray line. (Version 1 of ref. [5] found a larger value, comparable to $(\sigma v)_0$, but this was due to an error that has now been corrected.) Taking this value for $(\sigma v)$ and $m_X = 130$ GeV, we can find the relation between $q_S \sqrt{\lambda_X N e}$ and $m_S$, shown in fig. 3.

From fig. 3 we see that even if the coupling is rather large, $\lambda_X \sim 3$, and $N_c = 3$, the $S$ charge is typically greater than 1 (in units of $e$), and only reaches 1 for $m_S$ close to $m_X$. At the other extreme of $q_S \approx 6$, the corresponding interaction strength $q_S^2 \alpha = 0.26$ would still be under perturbative control, indicating that models with $m_S$ up to at least 400 GeV are viable. On the other hand, the rate goes like $q_S^2 \lambda_X^2$, so with $q_S = 6$ one could alternatively lower $\lambda_X$ from 3 to 0.08 by taking $m_S \approx m_X$. Thus it is not necessary to invoke a very large value of $\lambda_X$.

**Electroweak Precision Constraints.** $S$ must inherit its electric charge from weak hypercharge, and therefore it couples also to the $Z$ boson with strength $q_S e \tan \theta_W$, where $\theta_W$ is the Weinberg angle. Such a particle, if neutral under SU(2)$_L$, is unconstrained by precision electroweak constraints since its contribution to the $\rho$ or $T$ parameters vanishes identically [17], and it

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2 This limiting case might be of interest because it allows for the possibility of embedding $S$ into an SU(2)$_L$ doublet with the standard hypercharge assignment for extra Higgs doublets (see arxiv version 1 of this paper for more details). However this is not necessary, and we will focus on the case where $S$ is a singlet of SU(2)$_L$ and carries only weak hypercharge.
contributes only to the $Y$ parameter [18] which is weakly constrained.

**Annihilation to Z bosons and photons.** As noted above, $S$ necessarily couples to the $Z$ boson as well as to photons. Comparing the seagull vertices for $|S|^2 \gamma\gamma$ and $|S|^2 Z Z$, we can deduce that the cross section for $XX \to Z\gamma$ is related to that for $XX \to \gamma\gamma$ by the factor

$$\frac{\langle \sigma v \rangle_{XX \to Z\gamma}}{\langle \sigma v \rangle_{XX \to \gamma\gamma}} = 2 \tan^2 \theta_W \left( 1 - \frac{m_Z^2}{4m_X^2} \right)^{1/2} = 0.56 \quad (6)$$

taking into account the reduced phase space for identical particles in the case of $XX \to \gamma\gamma$. Since the former process produces only one photon, its intensity will be 0.28 times that of the 2-photon line. The energy of this single photon is given by

$$E_\gamma = m_X - \frac{m_Z^2}{4m_X} = 114 \text{ GeV} \quad (7)$$

We therefore predict that the spectral feature will resolve into two peaks separated by 16 GeV in energy. The relative strength of the peaks could be modified by giving a different SU(2)$_L$ assignment to $S$ so that both components of the doublet become electrically charged (see footnote 2). The current Fermi/LAT limit on $\langle \sigma v \rangle_{XX \to Z\gamma}$ is $2 \times 10^{-27} \text{cm}^3\text{s}^{-1}$ [3] is compatible with our value.

**Relic density.** The cross section for $XX \to \gamma\gamma$ is well below that which is needed to obtain the right relic density, but this can still be achieved using the $XX \to hh, WW, ZZ$ channels mediated by the interactions $\frac{1}{4} \lambda_{hh} h^2 X^2, \frac{1}{4} \lambda_{hh} v h X^2, \lambda v h^3$ (from the standard model Higgs potential) and the gauge couplings of the Higgs, with intermediate Higgs $h$ in the s channel [11]. The cross section for $XX \to hh$ is

$$\langle \sigma v \rangle_{hh} = \frac{\lambda_{hh}^2}{64 \pi m_X^4} \left( \frac{r_h}{r_h - 2} \frac{\lambda_{hh}}{\lambda} + 1 + \frac{r_h/2}{1 - r_h/4} \right)^2 \sqrt{1 - r_h} \quad (8)$$

where $r_h = m_h^2/m_X^2 \approx 0.94$ and $\lambda = 0.13$ assuming $m_h = 126$ GeV. For the $WW$ final state, we find

$$\langle \sigma v \rangle_{WW} = \frac{\lambda_{hh}^2}{8 \pi m_X^4} \left( \frac{r^2_w}{4 - r^2_w} \right)^2 \left( 2 + \left( 1 - \frac{2}{r_w} \right)^2 \right) \sqrt{1 - r_w} \quad (9)$$

where $r_w = m_W^2/m_X^2$. The cross section $\langle \sigma v \rangle_{ ZZ}$ is the same as $\langle \sigma v \rangle_{ WW}$ with the replacement $r_w \rightarrow r_Z \equiv m_Z^2/m_X^2$ and a factor of 1/2 for identical particles in the final state. Demanding that $\langle \sigma v \rangle_{tot} = \langle \sigma v \rangle_0$ implies that $|\lambda_{hh}| = 0.051$. If there are other significant annihilation channels, this contribution must be correspondingly reduced so that in general $|\lambda_{hh}| \leq 0.05$. For example if $S$ carries QCD color, then $XX \to gluons$ resulting in hadrons can be significant. There is also the possibility of $XX \to TT^*$ if the coupling $\lambda_{XX} T^2 |T|^2$ is sufficiently large and if $m_T < m_X/2$.

**Direct detection.** The $\frac{1}{2} \lambda_{hh} h^2 X^2$ vertex gives rise to the trilinear interaction $\lambda_{hh} vh X^2$ from electroweak symmetry breaking. (Note that the neutral component of the Higgs doublet is $H^0 = \frac{1}{\sqrt{2}} (h + v)$). The Higgs can therefore mediate scattering of $X$ on nucleons $N$. The cross section for $XN \to XN$ elastic scattering is [19]

$$\sigma = \frac{f^2 \lambda_{hh}^4 m_X^4}{4 \pi m_X^2} \quad (10)$$

where $f m_N/v$ is the Higgs-nucleon coupling, with $f = \sum_q u_{d,s} f_{Tq}^p + (2/9) f_{TG}^p = 0.319$ [20]. Using the constraint $\lambda_{hh} \lesssim 0.05$ from the preceding relic density determination, we can evaluate (10) to find $\sigma \lesssim 1.3 \times 10^{-45} \text{ cm}^2$. This is an order of magnitude lower than the 2011 Xenon100 90\% c.l. limit of $1.2 \times 10^{-44} \text{ cm}^2$ at $m_X = 130$ GeV [21], but only 3 times lower than the new limit of $3 \times 10^{-45} \text{ cm}^2$ that was recently announced [22]. Thus the model could be confirmed or ruled out in the near future by anticipated improvements [23] in the direct detection limit.

**Implications for Higgs decays.** Because of the close similarity between the diagrams of figs. 1 and 2, there is a simple relation between $\langle \sigma v \rangle$ and the extra contribution to $h \to \gamma\gamma$, which is especially interesting in light of the recent observation by CMS of an upward fluctuation in that branching ratio, relative to the standard model expectation [24], assuming of course that the indications of discovery of the Higgs boson with $m_h \equiv 126$ GeV are borne out [25–27]. Specifically, the squared matrix element for $h \to \gamma\gamma$ is related to (4) by replacing $\lambda_X \rightarrow \lambda_h v$ (where $v$ is the Higgs VEV) and $\tau \rightarrow \tau = m_h^2/4m_X^2$. The contribution of the charged scalar interferes constructively with that of the SM if $\lambda_h > 0$ [16]. To give an idea its relative size, we can express the extra contribution to the partial width of $h \to \gamma\gamma$ (here ignoring interference effects) in terms of its ratio to the SM contribution (see eq. (5),

$$\frac{\Gamma_{new}}{\Gamma_{SM}} = \frac{q^4 S_h^2 N_e^2}{A_{SM} G_F m_h^4} \tilde{s}^2 A_h^2(\tilde{s}) \quad (11)$$

where $m_h$ is the Higgs boson mass and $q$ is the Higgs-nucleon coupling. The factor $S_h$ is defined in the figure.
where the SM amplitude is given by $A_{SM} = -6.52$ for $m_h = 126$ GeV. For given values of $\lambda_h/\lambda_X$, $N_\pi$ and $m_S$ (and assuming that $m_X = 130$ GeV) the ratio (11) is fixed by the Fermi-LAT cross section, and is $O(1)$ for a wide range of $m_S$ if $\lambda_h \sim 0.5\lambda_X$. We plot it as a function of $m_S$ in fig. 4 for several values of $\lambda_h/\lambda_X$.

**Confinement of S particles.** Considering the case where $q_s = 2$ and $S$ decays are mediated by a dimension 5 operator (2), generically one would expect the $S$ lifetime to be much larger than the hadronization time, $1/\Lambda$ for a gauge theory with confinement mass scale $\Lambda$. Thus any $S$ particles produced in a collider will have time to form “mesonic” or “baryonic” bound states with other $S$ or $T$ particles. We first discuss the $SS^*$, $ST^*$ and $TT^*$ mesons, which we denote by $\phi_S$, $\pi_S$ and $\phi_T$ respectively.

Through the interaction (2), $\pi_S$ can decay into two electrons with a relatively long lifetime. But $\phi_S$ decays much faster by its constituents annihilating into photons, which we denote by $\Omega$. ($\Omega$ is unstable to decay into photons in the theory, these gluons hadronize into glueballs, SU(N) theory. Since there are no lighter colored particles produced in a collider will have time to form “mesonic” or “baryonic” bound states with other $S$ or $T$ particles. We first discuss the $SS^*$, $ST^*$ and $TT^*$ mesons, which we denote by $\phi_S$, $\pi_S$ and $\phi_T$ respectively.

The partial decay widths can be estimated using $\Gamma \sim n(\sigma v)$ where the density $n$ is the square of the $S$-meson wave function at the origin, $n \sim (\alpha' m_S)^2$, $\alpha'$ is the strength of the SU(N) gauge interaction at the scale $m_S$, and $\sigma$ is the cross section for $SS^*$ scattering into the desired final state. In this way we find the partial widths

$$\Gamma_{\phi_S} \approx \frac{\alpha'^3 m_S}{8} \left\{ \begin{array}{l} \frac{1}{16\pi} \frac{\lambda_X^2}{m_X^2/m_S^2}, \phi_S \rightarrow XX \\ \frac{1}{16\pi} \frac{\lambda_h^2}{m_h^2/m_S^2} \left( 1 + \frac{\lambda_h v^2}{m_h^2 - 2m_S^2} \right)^2 \times \frac{1}{1 - m_h^2/m_S^2}, \phi_S \rightarrow hh \\ \alpha^2, \phi_S \rightarrow \gamma\gamma, ZZ \\ \alpha^2, \phi_S \rightarrow \Omega\Omega \end{array} \right. \right. \tag{12}$$

Because of the large coupling $\lambda_X$, the invisible width for decays into $XX$ is typically the most important; even if the mass splitting $m_S - m_X$ is small, for example 3 GeV, $\Gamma_{XX} = 170\Gamma_{\gamma\gamma}$ when $\lambda_X = 3$. However, the branching ratio into glueballs $\Omega$ can also be large and even dominant, depending upon the unknown value of $\alpha'$. The glueballs can decay into photons via the Euler-Heisenberg interaction $F\bar{G}G/m_S^2$ induced by an $S$ loop. If the SU(N) confinement scale $\Lambda$ is below the weak scale, then photons will be the only kinematically available final states for the glueballs to decay into. The partial width for $\phi_S \rightarrow hh$ can be significant if $\lambda_h \sim \lambda_X$.

The $\phi_T$ meson can decay into all the same final states as $\phi_S$ by first going into two gluons of the SU(N) theory, which turn into two other particles by going through a loop of $S$. These amplitudes thus occur at two loops. They can be dominated by a one-loop contribution due to the possible interaction $\lambda|S|^2|T|^2$. Thus we expect the $\phi_T$ partial widths to be proportional to those of $\phi_S$, with an extra loop and coupling constant suppression.

The $S$ and $T$ particles can also bind together into baryons of the SU(N). For example if $N = 3$, we have the $SS^*$, $ST^*$ and $TT^*$ baryons, which can decay to states with fewer $S$'s via $S \rightarrow T + 2e$ until reaching the stable, neutral $TT^*$ state. The former would be stable charged relics in the absence of this decay mechanism. The latter is a dark matter candidate, but its annihilation cross section is too large for it to provide a significant contribution to the total DM density. The $TT^*$ annihilation cross section is of order $\alpha'^2/(3m_S^2)$, which for $\alpha' = 0.1$ (similar to the strength of QCD at these energy scales) is some 2 orders of magnitude larger than needed to get the correct thermal relic density.

**Collider Signatures.** The $S$-mesons would be produced at LHC mainly through intermediate s-channel photons and $Z$ bosons in the case where their color permutes to an exotic SU(N) gauge symmetry, as shown in fig. 5. Because of confinement the initially produced $S$-$S^*$ pair must hadronize to form the $S$-meson bound states $\phi_S$, $S^*$ or $\pi_S$, $ST^*$, or $TT^*$, respectively. If $m_{\phi_S} \lesssim m_S$ (but not $\lesssim m_{\pi_S}$), then $\phi_S$ and $\pi_S$ pairs will form with roughly equal probability. In the case of $\pi_S$ production, if the $S$ lifetime is short enough so that it decays within the detector, there will be a distinctive signal of pairs of charged leptons and antileptons, where each like-sign pair has an invariant mass equal to that of the $\pi_S$.

If on the other hand a $\phi_S$ pair is produced, the typical decay products will be dark matter pairs, $XX$, or glueballs, $\Omega\Omega$, as shown in fig. 5. Each glueball decays into two photons because of the one-loop Euler-Heisenberg interaction. Thus pairs of photons with invariant mass $m_{\phi_S}$ will be produced, and pairs of pairs will reconstruct to invariant mass $m_{\phi_S} \approx 2m_S$. In the example shown, this will be accompanied by missing energy $m_{\phi_S}$. A detailed study should be done to see if existing searches of this nature [28] already exclude our model. Another likely signature would be four pairs of photons due to decay of both mesons into glueballs.

**Conclusions.** We have shown that scalar dark matter $X$ with mass 130 GeV could produce a gamma ray spectral feature tentatively identified in Fermi-LAT data, with the addition of just one scalar multiplet $S$ transforming as $(Y_3,1,3)$ under $U(1)_Y \times SU(2)_L \times SU(3)$, where the SU(3) might be the gauge group of QCD, or else some new hidden sector interaction. The coupling $\lambda_X X^2|S|^2$ between $S$ and $X$ need not be very large unless $q_S \sim 1$ and $m_S/m_X$ is greater than a few. The strong interactions of $S$ serve two purposes: they confine the stable charged relic component of $S$, and the number of colors helps to increase the $XX \rightarrow \gamma\gamma$ cross section while keeping $\lambda_X$ reasonably small.

Because $S$ has similar quantum numbers to right-handed squarks, it is tempting to make this identification, but such light squarks are ruled out by LHC with the possible exception of the third generation. Squarks have a smaller electric charge than $S$ in our preferred examples, and all three generations would need to con-
FIG. 5: Production and decay of $S$-mesons at hadron collider. Attributed to compensate for the resulting decrease in the $XX \rightarrow \gamma\gamma$ cross section.

It is interesting that we rely upon the $h^2X^2$ coupling between $X$ and the Higgs boson to get the thermal relic density of dark matter, and that the same coupling leads to a cross section for $X$ scattering on nucleons that is just a factor of 3 below the 2012 Xenon100 direct detection limit. The model has additional links to Higgs physics: the possibility of increasing the $h \rightarrow \gamma\gamma$ branching ratio by a factor of a few, and the existence of bound states of $S$ and $S^*$ ($\phi_S$), which could have a large branching ratio to decay into Higgs bosons, though more generally their decay products are dominated by glueballs $\Omega$ of the exotic SU(N) or $X$ bosons. We suggest that LHC might discover the $S$-mesons, whose mass should be $> 260$ GeV, by observation of two photon pairs each with invariant mass of $m_\Omega$, accompanied by the same amount of missing energy, or four photon pairs each of mass $m_\Omega$. In addition, lighter charged $ST^*$ mesons ($\pi_S$) should be produced, decaying into like-sign lepton pairs, which might or might not occur within the detector. A more detailed study of LHC signals is contemplated [30].

After releasing version 1 of this work, we were informed of a similar model in [29], resembling ours in the case where $m_X \cong m_{\phi_S}$, using the process $XX \rightarrow \phi_S\phi_S$ followed by $\phi_S \rightarrow \gamma\gamma$, which becomes the dominant decay channel for $\phi_S$ if $\lambda_\phi \ll 1$ and no hadronic channels are available (as in the case of a hidden SU(N) with glueball mass greater than $m_{\phi_S}/2$). This seems to be another viable region of parameter space for our model, relying only upon tree-level amplitudes. In this case the DM mass would be 260 rather than 130 GeV. We acknowledge A. Ibarra for pointing this out.

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