Composite magnetic dark matter and the 130 GeV line

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We propose an economical model to explain the apparent 130 GeV gamma ray peak, found in the Fermi/LAT data, in terms of dark matter (DM) annihilation through a dipole moment interaction. The annihilating dark matter particles represent a subdominant component, with mass density 7–17% of the total DM density; and they only annihilate into $\gamma\gamma$, $\gamma Z$, and $ZZ$, through a magnetic (or electric) dipole moment. Annihilation into other standard model particles is suppressed, due to a DM mass splitting in the magnetic dipole case, or to $p$-wave scattering in the electric dipole case. In either case, the observed signal requires a dipole moment of strength $\mu \sim 2/\text{TeV}$. We argue that composite models are the preferred means of generating such a large dipole moment, and that the magnetic case is more natural than the electric one. We present a simple model involving a scalar and fermionic techniquark of a confining SU(2) gauge symmetry. We point out some generic challenges for getting such a model to work. The new physics leading to a sufficiently large dipole moment is below the TeV scale, indicating that the magnetic moment is not a valid effective operator for LHC physics, and that production of the strongly interacting constituents, followed by technihadronization, is a more likely signature than monophoton events. In particular, 4-photon events from the decays of bound state pairs are predicted.

In the past few months a number of studies of publicly available Fermi Large Area Telescope [1] data have emerged that find evidence for a spectral peak of energy $\sim 130$ GeV from the galactic center [2–4], possibly accompanied by a second peak with lower energy $\sim 116$ GeV [5, 6]. Ref. [7] moreover finds corroborating evidence from outside the galaxy, originating from several galactic clusters, and ref. [8] sees marginal evidence for sources unassociated with known structures, though this has been questioned [9–11]. The morphology of the signal from the galactic center is consistent with dark matter annihilating into monoenergetic photons, but not with dark matter decays [7, 12]. In terms of the thermally averaged cross section required for the standard DM relic density, $\langle \sigma v \rangle_0 = 3 \times 10^{-26} \text{cm}^3/\text{s}$, the annihilating DM hypothesis is consistent with a cross section estimated to be $0.04 \langle \sigma v \rangle_0$ [3, 12] or $0.1 \langle \sigma v \rangle_0$ [4, 7]. Fermi/LAT itself places a limit of $\langle \sigma v \rangle_{\gamma\gamma} < 0.035 \langle \sigma v \rangle_0$ on this cross section (for an Einasto profile) [13, 14], marginally below the required value. Searches for the signal from dwarf galaxies yield the weaker limit $\langle \sigma v \rangle_{\gamma\gamma} < 1.3 \langle \sigma v \rangle_0$ [16]. If the 130 GeV line is accompanied by a gamma continuum due to additional annihilation channels into charged particles, the limit on the cross section for these processes is estimated by ref. [15] to be $\langle \sigma v \rangle_{\text{other}} < 3.5 \langle \sigma v \rangle_{\gamma\gamma}$ for the $b\bar{b}$ channel, while ref. [17] obtains the weaker limit $\langle \sigma v \rangle_{\text{other}} < 20 \langle \sigma v \rangle_{\gamma\gamma}$. (Weaker limits are found in ref. [18], and stronger ones in [19].) While the latter is only marginally in conflict with the DM getting its standard relic density from annihilating primarily into charged particles, the former would definitely forbid this scenario, presenting a challenge to many models.

Although some authors have expressed skepticism about the DM interpretation of the 130 GeV line [20, 21], there has been a great deal of interest in this possibility both from the experimental analysis perspective [22–26] and in theoretical model-building [27–36]. (Ref. [37], which preceded the observation of the 130 GeV line, is also a competing theory.) In the present paper, we consider a distinctive model similar to those discussed in refs. [38, 39], in which the dark matter interacts primarily through a large magnetic dipole moment (MDM). Direct detection of dipolar dark matter has been the focus of many recent papers, dealing either with elastic [42–46] or inelastic models [47–51] (for recent work on indirect detection, see [52]). Here we show that Majorana DM with a transition MDM and a relatively large mass splitting can economically explain the observed gamma ray line, in particular if we relax the usual assumption that the candidate particle must be the dominant form of dark matter. (As we will discuss, DM with an electric dipole moment (EDM) and a smaller mass splitting can also work, though we do not favor this scenario.) Moreover we remove the tension of too large a $\gamma$-ray continuum by arranging for the $2\gamma$ annihilation channel (along with the accompanying $\gamma Z$ and $ZZ$ channels) to be the dominant one. Majorana DM $\chi_1$ can have a transition magnetic moment to a heavier state $\chi_2$,

$$\frac{1}{2} \mu_{12} \chi_1 \sigma_{\mu \nu} \chi_2 F^{\mu \nu}. \quad (1)$$

It is natural to assume that the interaction (1) came originally from standard model hypercharge, and is therefore...
accompanied by a corresponding coupling to the \( Z \) boson field strength.

With the interaction (1), there exist both annihilations \( \chi_{1,2} \chi_{1,2} \rightarrow gg, Z, ZZ \), as well as coannihilations \( \chi_1 \chi_2 \rightarrow f \bar{f} \), where \( f \) is any charged standard model particle, or \( hZ \) in the case of the \( Z \) intermediate state. These are pictured in fig. 1. As was shown in ref. [39], the coannihilations suppress the relict density below its required value; but a mass splitting \( m_{\chi_2} - m_{\chi_1} \gtrsim 10 \) GeV suppresses the coannihilations enough to get the right relic density. Therefore a rather large mass splitting is a needed ingredient in the magnetic dipole case. Such a large mass splitting of course makes direct detection to be smaller (by a factor of \( 10 \)) than the standard relic density \( n_0 \). The total 130 GeV photon production rate scales as \( \langle \sigma v \rangle n^2 \). According to refs. [3, 4] the observed \( \gamma \) ray signal corresponds to

\[
\langle \sigma v \rangle n^2 = (0.04-0.1)\langle \sigma v \rangle n_0^2,
\]

where \( \langle \sigma v \rangle n_0 \) is the cross-section which would lead, via the usual freezeout calculation, to a relic density of \( n_0 \). Since the relic density is approximately inversely proportional to the annihilation cross-section,

\[
n \sim \frac{\langle \sigma v \rangle 0}{\langle \sigma v \rangle n_0} \, ,
\]

the photon producing cross-section should satisfy

\[
\frac{\langle \sigma \gamma \rangle}{\langle \sigma v \rangle 0} \sim \frac{\cos^8 \theta_w}{(0.04-0.1)} \, .
\]

Accounting for the dependence on the mass splittings given in appendix B, we find that the dipole moment should lie in the range

\[
1.6 f(r) \, \text{TeV}^{-1} < \mu_2 < 2.0 f(r) \, \text{TeV}^{-1}
\]

where \( r = m_{\chi_2}/m_{\chi_1} \) and \( f = \sqrt{(1 + r^2)/2r} \). The corresponding fractional relic abundance is

\[
0.07 \frac{f(r)}{r} < \frac{n}{n_0} < 0.17 \frac{f(r)}{r} \, .
\]

Notice that a larger dipole moment leads to more complete annihilation during freezeout, and hence to a smaller abundance and a smaller photon flux.

If another annihilation mechanism was active in the early Universe—such as coannihilation, which can occur if the \( \chi_1 \) to \( \chi_2 \) mass splitting is not too large [39]—then \( \langle \sigma \gamma \rangle \) and hence \( \mu_1 \) need to be smaller, and the relic density fraction is larger. In the extreme case that coannihilation controls the relic density and \( n = n_0 \), then \( \langle \sigma \gamma \rangle \) can be about 100 times smaller, and hence \( \mu_1 \) can be 10 times smaller. The coannihilation cross section is just large enough to give the right relic density (if the mass splitting is small) for this smaller value of \( \mu_1 \).

Determining \( \mu \) and relic abundance. We assume that the dominant annihilation channels are to \( \gamma \gamma, \gamma Z, ZZ \) via a dipole moment, as shown in fig. 1(a)—both in the context of current galactic dark matter signals, as well as that of the annihilations which establish the relic abundance.\(^2\) Let us call the cross-section to two photons \( \langle \sigma \gamma \rangle \); then using Eq. (2), the total cross-section is \( \langle \sigma v \rangle \sim \langle \sigma \gamma \rangle / \cos^4 \theta_w \) (neglecting corrections from the different kinematics of a \( Z \) final state, see appendix B). We also write the density of 130 GeV dark matter as \( n \), which will be smaller than the standard relic density \( n_0 \). The total 130 GeV photon production rate scales as \( \langle \sigma \gamma \rangle n^2 \). According to refs. [3, 4] the observed \( \gamma \) ray signal corresponds to

\[
\langle \sigma \gamma \rangle n^2 = (0.04-0.1)\langle \sigma v \rangle n_0^2,
\]

\(^1\) at first order in perturbation theory; integrating out the excited DM state of course leads to a small elastic transition at higher order [40].

\(^2\) This is a more detailed discussion of a comment made in [38] for DM with a Rayleigh interaction.
**Direct detection.** If the DM candidate $\chi$ was Dirac, an elastic EDM or MDM as large as in Eq. (6) would be ruled out by many orders of magnitude, by direct detection searches for elastic scattering on nuclei. In ref. [49] (see also [47]-[48],[50]-[51]) it was shown that by splitting the Dirac particle into two Majorana states with mass splitting of $\sim 150$ keV, a magnetic moment $\sim 10^{-2} \mu_N = 1.6$/TeV is compatible with direct detection limits (though somewhat too large to make the DAMA annual modulation signal compatible with these bounds). In the case of MDM, scatterings can be dominated either by dipole-charge interactions (which lack the $1/v^2$ enhancement present in the EDM case) or dipole-dipole interactions, depending on the nucleus.

The above determination was made assuming that the dipolar DM had the standard relic density; in our case the density is lower by about a factor of 10. Thus our DM candidate looks similar to DM with the full relic density of course with the charge of the proton, and is enhanced by 1 order magnitude over the EDM case) or dipole-dipole interactions, depending on the nucleus.

On the other hand, for EDM dark matter, the annihilations are $p$-wave suppressed even if the mass splitting is small, which suggests the possibility of combining the EDM explanation of the 130 GeV line with direct detectability. EDM scattering is dominated by interaction with the charge of the proton, and is enhanced by $1/v^2$ relative to dipole-charge scattering of an EDM (see the erratum of ref. [42]). For inelastic scattering with mass splitting $\delta M$, we find that this translates to an enhancement by the factor $m_\eta/\delta M$ in the cross section. Therefore one would have to reduce the EDM by a factor of $\sqrt{m_\eta/\delta M} \sim 100$ relative to the above estimate for it to be relevant for direct detection, which would then make it too small to explain the 130 GeV line. We find the same conclusion whether dipole-dipole or dipole-charge scattering dominates in the MDM interaction. (Notice that one cannot increase $\delta M$ much above 100 keV for direct detection since the minimum required DM velocity would then exceed the escape velocity of the galaxy.)

**Model requirements.** Consider first that the dark matter is a neutral fundamental field $\chi$. One way it can obtain a large dipole moment is through a Yukawa interaction $y\bar{\chi}\phi S$ to heavy (hyper)charged states $\psi, S$. If all the states have masses $\sim 100$ GeV, then an EDM of order $1$/TeV is only achieved by pushing $y$ to the perturbative limit $\sqrt{4\pi}$ and tuning the particles in the loop to be nearly on shell [53]. Large MDMs require similar tuning of parameters. Thus any perturbative origin requires stretching the limits of perturbation theory; probably the effective theory describing such states cannot be valid much higher than the particle mass scales. In fact, [54] give a perturbative construction of a large MDM, which requires large couplings, messenger masses not far above the DM mass, and dimension 7 Rayleigh operators to explain the Fermi signal.

On the other hand, large magnetic moments are known to arise for neutral composite particles. The neutron is a good example, with $\mu = -1.91(e/2m_n)$, which is approximately the sum of the magnetic moments of the constituent quarks (treating them to have constituent quark masses $\sim m_N/3$, see for example ref. [55]). If the DM is a bound state and analogously $\mu \sim e/2m_\chi$, this gives $\mu \sim 1.1$/TeV, which has the right order of magnitude. This could happen if the hidden sector has a confining gauge symmetry such as SU(3), analogous to QCD, that becomes strong at the 100 GeV scale. However we cannot push the analogy to QCD too far. Not only do we need the DM particle to be absolutely stable, we need it to be the lightest bound state, since otherwise it would have a large annihilation cross-section into lighter bound states. Therefore it cannot be a baryon of SU($N_c > 2$), since there will always be lighter mesonic states.

There will however be fermionic meson states in theories with both scalar and spinor matter, say a charged techniquark $S$ and an oppositely-charged anti-techniquark $\psi$ forming a spin-$1/2$ bound state $\eta = S\psi$. Such models automatically also have composite neutral bosonic mesons $\eta_S = S^*S$ and $\eta_\psi = \bar{\psi}\psi$. If the $\eta$ is lighter than the other mesons and baryons, it is safe from annihilating into them; and if the $S, \psi$ mass scales are below or comparable to the confinement scale, the glueballs are also heavier. The details of the strong dynamics may determine which meson is the lightest, but we argue in the next section that it can plausibly be the $\eta$.

A further requirement is to split the Dirac $\eta$ state into Majorana states at different mass, in order to have a transition dipole moment. In the theoretically preferred scenario of large magnetic moment, this splitting has to be sizable, $\gtrsim 10\%$, to get a sufficient relic density. A loop effect (such as that which gives the neutron a mass splitting in $R$-parity violating supersymmetry [56, 57]) will be too small. But there is a simple way to get a large mass splitting at tree level, if the dark matter is an admixture of an elementary fermion $\chi$ and the bound state $S\psi$, by having a bare mass term $\frac{1}{2}m_\chi \bar{\chi}\chi$ and a Yukawa coupling $y\bar{\chi}\psi S$. The latter becomes an off-diagonal mass term with mass of order $y\Lambda$ at scales below the confinement scale $\Lambda$. Large mass splittings result as long as $m_\chi$ is comparable to $m_\eta$, the (unmixed) bound state mass. The generation of the transition moment can be visualized through the diagrams of fig. 2.

Finally, the exotic charged states that appear as bound states (or within the loop for a perturbative origin of the dipole) must not be stable since there are very stringent constraints on charged relics. Thus any model introducing exotic heavy charged particles to induce the dipole moment must ensure that they can decay into charged particles within the standard model, while respecting the stability of the dark matter.
Explicit model. Based on the above considerations, we construct the simplest model of composite dipole dark matter\(^3\) that meets all the requirements. We take the new confining gauge group to be SU(2)\(\lambda\), and introduce a fundamental (SU(2)\(\lambda\), doublet) Dirac fermion \(\psi\) which carries electric charge \(n+\frac{1}{2}\) for integer \(n\), and a scalar \(S\) with the opposite charges. There is an elementary Majorana fermion \(\chi\) that mixes with the neutral \(S\) bound state to form the dark matter. The particle content and quantum numbers are listed in table I.

The relevant mass terms and interactions in the potential at the scale above \(\Lambda\) (the compositeness scale) are

\[
V = \frac{1}{2}m_\chi \bar{\chi}\chi + m_\psi \bar{\psi}\psi + m_S^2|S|^2 + \lambda|S|^4 + \bar{\chi} \chi S^a(y+i\gamma_5\gamma_5)\psi + y/\epsilon ab \epsilon_{ab} S^a \bar{\psi}\psi \left( \frac{\bar{\psi}\epsilon R}{\Lambda^2} \right)^n + \text{h.c.}
\]

The interactions are invariant under a discrete \(Z_1\) symmetry, shown in table I, which guarantees the stability of the dark matter. The last interaction involves electrons for simplicity, but in general any combination of right-handed charged leptons can appear.

The physical states of the system are the neutral mesonic bound states \(\eta = S^a \bar{\psi}_a\), \(\bar{\eta}_5 = S^a \bar{\psi}_a\), \(\bar{\eta}_5 = \bar{\psi}_a\), and the charged baryonic\(^4\) ones \(N^- = \epsilon_{ab} S^a \bar{\psi}_b\), \(\bar{N}^+ = \epsilon_{ab} S^a \psi_b \). The state \(\eta\) will be the lightest for some part of the parameter space, since \(\eta_5\) gets a positive mass correction from the repulsive quartic scalar interaction\(^5\), there is a positive spin-spin interaction energy for \(\bar{\eta}_5\), and we are still free to choose \(m_\chi, m_S\) – which however cannot be much heavier than the confinement scale so that \(m_\eta < m^{\text{glueball}}\). The charged states are heavier due to positive Coulombic energy shifts.

The bound state \(\eta\) mixes with \(\chi\) through the first Yukawa interaction in (9).

\[\frac{1}{2} m_\chi \bar{\chi}\chi + m_\psi \bar{\psi}\psi + m_S^2 S^a \bar{\psi}\psi \left( \frac{\bar{\psi}\epsilon R}{\Lambda^2} \right)^n + \text{h.c.}\]  

where \(m_\psi\) is of order \((y_5^2 + y_S^2)^{1/2} \Lambda\), \(\theta_5 = \tan^{-1}(y_5/y_\chi)\), and \(m_\eta \sim \Lambda\) is the mass of the \(\eta\) composite fermion. For simplicity, we impose parity to set \(\theta_5 = 0\) in the following, though the qualitative features of the model do not depend upon this choice. In the basis of Weyl components \((\eta^c, \eta, \chi)\) of a single handedness, this corresponds to a \(3 \times 3\) Majorana mass matrix of the form

\[
M = \begin{pmatrix}
0 & m_\eta & m_y \\
m_\eta & 0 & m_y \\
m_y & m_y & m_\chi
\end{pmatrix}.
\]

As long as \(m_y \neq 0\), the mass eigenstates are distinct Majorana particles, none of which can be paired up into a Dirac particle. This is important, since as noted above, a stable Dirac particle with as large a direct magnetic moment as \(\mu_\chi = 1/\text{TeV}\) is ruled out by direct detection constraints. Notice that in the same basis, the magnetic moment matrix takes the form

\[
\mu = \begin{pmatrix}
0 & \mu_\eta & 0 \\
-\mu_\eta & 0 & 0 \\
0 & 0 & 0
\end{pmatrix}.
\]

Thus if \(R^\dagger MR\) diagonalizes the mass matrix, then \(R^\dagger \mu R\) gives the transition moments in the mass eigenbasis.

In the parity-conserving case, the spectrum of DM states has a simple analytic form, which we give in appendix C. The combination \((\eta - \eta^c)/\sqrt{2}\) does not mix

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\(^3\) For another recent model of composite DM, without a dipole interaction however, see ref. [58].

\(^4\) For SU(2) there is no real distinction between mesons and baryons; so our choice to call these baryons is just a convention.

\(^5\) There is a potential danger that \(\Lambda\) runs to negative values before the confinement scale, due to a \(y_\chi^4\) term in its beta function \(\beta_\chi\). This is partly compensated by the \(-y_\chi^4\) term in \(\beta_\psi\). Here we will assume that parameters can be found where it does not go negative below the scale \(m_S\).
with $\chi$ and retains mass of exactly $m_\eta$. For small $m_\eta$, the spectrum can be written in the form

$$m_i = (m_\eta - \delta m, m_\eta, m_\chi + \delta m)$$  

(12)

where $\delta m \approx 2m_\eta^2/(m_\chi - m_\eta)$. We argue that to avoid suppressing $\mu_{12}$ by a small mixing angle, it is desirable to choose $m_\eta < m_\chi$ so that the lowest two states are mostly $\eta_1, \eta_2$. Then the mixing angle between the lowest and highest states is $\theta_{13} \equiv \sqrt{2m_\eta}/(m_\chi - m_\eta)$, and the middle state is purely composite; if $m_\chi < m_\eta$, the masses take the inverse hierarchy $m_3 < m_2 < m_1$. The transition moments in either case are $\mu_{12} \cong \cos \theta \mu_\eta \cong \mu_\eta$ and $\mu_{23} \cong \theta \mu_\eta$. By analogy to the nucleons of QCD, we can estimate $\mu_\eta \cong (n + \frac{1}{2})e/(2m_\eta/2) \approx (2n + 1) \times 1.2/\text{TeV}$. Based on the quantitative success of the quark model in reproducing the baryon magnetic moments [55], this can be considered as good to $\sim 10\%$, and not just an order of magnitude estimate. On the other hand, the process shown in fig. 1 involves a $\chi_2$ particle which is far off shell, so the hadronic effective picture is not reliable. There can be a nontrivial form factor, as well as other two-photon operators such as $\chi_i^3 \chi F^2 F$. It is best to consider the annihilation as one effective operator to two photons, and to adopt the dipole annihilation calculation, eq. (B1), as a rough estimate for the cross-section which probably receives $O(1)$ corrections. It would be difficult to improve this estimate analytically, but it is possible that a lattice calculation could clarify the actual annihilation rate. We will moreover see in the following section that it is necessary to take the mass ratio $r = m_\chi/m_\eta$ to be greater than 2, hence $f(r) > 1.2$ in eq. (6). Therefore to achieve a sufficient annihilation rate we may need to consider $n = 1$ (charge-3/2 techniquarks), leading to $\mu_\eta \cong 3.6/\text{TeV}$.

We thus arrive at the favored scenario in which $m_\chi > m_\eta > m_\eta$. We have computed the ratio of the transition moment $\mu_{12}$ to its maximum value $\mu_\eta$ as a function of $\delta m/m_\eta$, which is shown for several values of $m_\chi/m_\eta$ in fig. 3. The exact mass splittings and mixing angles are given in appendix C in this parity-symmetric case. (We leave the parity-violating case for future work; generally, all three states will mix, and none of the possible transition moments will vanish.)

The baryonic bound states carry electric charge, and must not be stable because of very stringent bounds on relic charged particles. The $y'$ Yukawa coupling ensures that they can decay into $2n+1$ charged leptons and a photon or dark matter. For the $N^-$ baryon, the $y'$ coupling in (9) leads directly to a decay term $y'(\Lambda/\Lambda^3 m) (\bar{e}_\nu N^-)(\bar{e}_\bar{e}_\nu)^n$. Because of the nontrivial $Z_4$ charges, the bosonic baryons must decay via $\bar{N}_i^- \rightarrow \eta + (2n+1)^+ \eta$ and $N^- \rightarrow \eta + (2n+1)^- \eta$ respectively, through the elementary processes $S \rightarrow \psi + (2n+1)^+ \eta$ or $\psi \rightarrow S + (2n+1)^- \eta$. Clearly, only a small mass splitting between the baryons and the DM is needed to make these kinematically possible. We remind the reader that $\epsilon_\eta$ (hence $\epsilon_\pm$) stands for any flavor of right-handed charged leptons.

**Challenges to the model.** Having invented a specific model, we need to check that no important new annihilation channels were introduced, that were not present when we assumed the dark matter interacted only through its magnetic moment coupling. Due to the strong interactions of the DM constituents, there are indeed new possible diagrams shown in fig. 4 that can be problematic.

The first one, fig. 4(a), illustrates the process $\chi_1\chi_1 \rightarrow \tilde{\eta}_i\gamma$ (where $i = \psi, S$) which involves a Yukawa interaction

$$\bar{y}_i\tilde{\eta}_i\chi_1\chi_2$$

(13)

in the effective theory below the confinement scale. This vertex only requires the rearrangement of the constituents into different bound states (with the annihilation of $SS^*$ or $\psi\bar{\psi}$), so there is no *a priori* reason for it to be suppressed, and the couplings $\tilde{y}_i$ might be relatively large. Thus this channel will dominate over that of fig. 1(a) if it is kinematically allowed, and one therefore needs to have $2m_\chi < m_\bar{\eta}_i$. This can be accomplished by judicious choices of the mass matrix elements in (10). For example, with $m_\psi/\Lambda = 0.55$ and $m_\chi/\Lambda = 1.5$, we find that $m_\chi_1 : m_\chi_2 : m_\chi_3 = 2.1 : 1 : 0.43$, which is sufficient. And the transition dipole moments are not much suppressed: $\mu_{21}/\mu_\eta = 0.6, \mu_{12}/\mu_\eta = 0.8$. The process with off-shell $\tilde{y}_i$ leads to $\gamma$-ray continuum $\chi_1\chi_1 \rightarrow 3\gamma$ annihilations, but they are suppressed by a power of $\alpha$ and should be small.

The second diagram, fig. 4(b), again involves a presumably large effective Yukawa coupling, and an off-shell s-channel $\tilde{\eta}_i$ particle. To appraise it, we need an estimate of the $\tilde{\eta}_i$ coupling to two photons. To compute this from first principles would require knowledge of the nonperturbative matrix elements, but by comparing to the two-photon decays of the $\eta$ and $y'$ mesons in the real world, we estimate that the interaction takes the form

$$e^2 \frac{\bar{y}_i \tilde{\eta}_i F_{\mu \nu} F^{\mu \nu}}{\Lambda}$$

(14)

with $e \cong 0.1$, where $q = (2n+1)/2$ is the charge of the constituents in our model. The suppression by $q$ is ac-
compounded by an additional $p$-wave suppression due to the scalar coupling (as opposed to pseudoscalar) in (13). The cross section, ignoring interference with fig. 1(a), is

$$\sigma = \frac{2\mu^2}{\pi} \frac{c^2 \gamma^2 S^4 \Lambda^4 m_N^4}{\Lambda^2 (4m_N^2 - m_N^2)^2}.$$  \hspace{1cm} (15)

Therefore this interaction turns out to be subdominant to that of fig. 1(a) unless it happens to be very close to resonance. Note however that the phenomenology of this process is the same as that of the magnetic moment interaction, in that it also produces photon pairs which would give a line feature. Therefore we might also consider the model where the magnetic moment annihilation rate is too small, but this process proceeds near threshold and actually dominates the two-photon production rate. In this case the model would be viable with the small ($n=0$) charge assignment for the $S, \psi$ techniquarks.

**Collider signatures.** Past studies of the collider signatures of magnetically interacting dark matter have focused on the production of the DM and its excited state $q\bar{q} \rightarrow \chi_1 \chi_2$, mediated by $s$-channel $\gamma$ or $Z$ and the transition moment $\mu_{12}$ of $\chi_1-\chi_2$. The subsequent decay $\chi_2 \rightarrow \chi_1 \gamma$ through the dipole moment can produce a hard monophoton that would pass experimental cuts [59, 60] if the $\chi_1-\chi_2$ mass splitting exceeds 125-150 GeV. Otherwise the pair-production of DM can be searched for in a more generic way, through missing energy and initial state radiation of a hadronic monojet [61]-[64].

However the previous considerations assume that the dipole interaction remains hard at LHC energies. We have shown that whether one invokes a loop effect or more plausibly compositeness to generate the large moment needed to explain the 130 GeV line, it arises from physics well below the TeV scale, and therefore the magnetic moment effective description is not valid at high energies. It will open up to reveal the constituent particles rather than behaving like a dipole moment. Equation (16) below estimates the cross section for dark matter production (though it is likely further suppressed due to cancellation between the oppositely charged DM constituents and slightly due to the DM mixing angles).

In the model we have proposed, the primary process will be to produce $\psi\bar{\psi}$ and $SS^*$ pairs via virtual photon and $Z$. These will then hadronize into the mesons and baryons listed in table I. The final states will include the $\chi_1$ and $\chi_2$ (mostly $\eta$) dark matter particles, producing monophotons and monojets. But other composite states will also be produced, and they can have very different signatures because of their visible decays. For example, the bosonic bound states $\tilde{\eta}_s$ and $\tilde{\eta}_\psi$ can decay into two photons by direct annihilation of their constituents, as shown in fig. 5. Thus a very clean signal of two diphoton pairs, each of which has the same invariant mass $\gtrsim 260$ GeV, is predicted. The production of these particles is suppressed near threshold because of the overall charge neutrality of the mesons, leading to e.g., destructive interference in the photon+$Z$ coupling to $\psi$ and to $\bar{\psi}$. But at center-of-mass energies well above the particle masses, the vector boson resolves the constituents and couples fully to both of them. The spin and color summed/averaged partonic cross section for $\psi\bar{\psi}$ production from up quarks, at leading order in $1/s$, is given by

$$\sigma \cong \frac{4\pi\alpha^2 Q_\psi^2}{9s \cos^2 \theta_W} \left( \frac{4}{9} + \frac{1}{36} \right) \sim 6 \text{ fb} \hspace{1cm} (16)$$

where for the numerical estimate we take $Q_\psi = 1/2$ and $\sqrt{s} = 1$ TeV. Of course a serious prediction requires

6 $4/9$ and $1/36$ are the squared hypercharges of the right and left components of the up quark.
weighting by parton distribution functions, which will reduce (16) by around an order of magnitude because of the need to get the antiquark out of the sea, as well as a better treatment of the hadronization process. But this estimate already shows that the production is at a potentially interesting level for discovery at the LHC; indeed, since the four photon signal is so striking and involves particles of such high energy, the existing LHC data may already have sufficient sensitivity to exclude or discover this model.

Even more interesting perhaps is the pair production of charged technibaryons, especially the $N^-$, which can decay via $N^- \rightarrow \left(2n+1\right)e + \gamma$, or $N^- \rightarrow \left(2n+1\right)e + \bar{\eta}_{\beta, \gamma}$ if the latter is kinematically allowed. Unlike the neutral pairs, here the production is not suppressed at threshold because both constituents have charges of the same sign. If both channels are open, the possible final states are $\left(2n+1\right)e + \left(2n+1\right)\bar{e}$ plus two, three or four photons, due to the subsequent decays of the $\bar{\eta}_{\beta, \gamma}$ mesons. This is illustrated in fig. 6. However it is quite possible that the $N^- \rightarrow \left(2n+1\right)e + \bar{\eta}_{\beta, \gamma}$ decays are kinematically blocked, in which case only the $\left(2n+1\right)e \left(2n+1\right)\bar{e} \gamma \gamma$ channel is present. Moreover, there is no reason to suppose that the $y'$ coupling in (9) is only to electrons, and it may be natural to assume that the dominant coupling is to $\tau$ leptons, which would be more difficult to identify in the final state.

Finally, there will be pair production of the charged bosonic technibaryons $N^+_{\beta}$ and $N^+_{\gamma}$, which decay into leptons and dark matter. This however has a large standard model background (if $n = 0$) since such events could be mistaken for $W$ boson decays.

**Conclusions.** If the 130 GeV gamma ray line and the tentative accompanying line at 114 GeV are indeed due to dark matter annihilation into $\gamma \gamma$ and $\gamma Z$, it is a challenge to explain its relatively strong annihilation into two monoenergetic photons, while keeping its annihilation into other particles that produce a continuum of secondary photons within bounds. We observed that a DM species comprising just 10-15% of the total DM mass density, possessing a transition magnetic moment of $\sim 2/\text{TeV}$ and a mass splitting $\gtrsim 10$ GeV can satisfy these requirements, with its relic density determined by the annihilations into photons and $Z$ bosons. (Although an electric dipole moment of the same strength and smaller mass splitting could do the same, and be relevant for direct detection, we find it theoretically difficult to generate such a large EDM without an accompanying MDM of the same size.) This scenario leaves open the possibility of direct detection of some other, dominant DM species.

We further observed that it is difficult to explain a dipole moment of the needed magnitude through a loop effect (see beginning of “Model requirements” section), which motivated us to study a composite DM particle. We presented a simple model that accomplishes this with an electrically charged “quark” and scalar “quark” of a new SU(2) gauge interaction that confines at the 100 GeV scale, and whose neutral mesonic bound state mixes with an elementary fermion to make the dark matter and its excited state. Remarkably, the magnetic moment is determined by the charge and mass of the constituents, leaving relatively little room for adjustment given that the DM has mass 130 GeV, yet coming out to approximately the desired value. Unfortunately, complicated confining dynamics occur at the scale of the annihilation momentum transfer, so the cross-section can have an $\mathcal{O}(1)$ difference from the pure dipole calculation.

Moreover the model makes an unambiguous prediction for hadron collider production of pairs of “unflavored” scalar-scalar or fermion-fermion bound states of mass $\gtrsim 200$ GeV, whose only decay channel is into two photons, and whose production cross section is estimated to be within reach of the LHC. This double diphoton pair signal comes in addition to the more discussed signature of DM pair production with missing energy and monophotons or monojets. And due to the pair production of charged bound states, further exotic final states can appear where two photons and one lepton (or three leptons in a related model) have the invariant mass of the charged parent. Although there is some freedom in varying the details that accommodate the decays of the charged bound states (necessary in order to avoid highly constrained charged relics), the two-photon decay of the neutral bound states is a robust prediction in any model that uses compositeness to generate large magnetic moments.

The dark matter interpretation of the 130 GeV line is expected to be tested definitively by upcoming gamma ray experiments [67], in particular by HESS-II which begins very soon. It is exciting to consider that complementary indirect evidence could come from the LHC on a similar time scale, if a new peak at mass $\gtrsim 260$ GeV (according to the discussion below eq. (13)), coming from pairs of diphoton events, should be discovered.

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**Appendix A:** $\chi_1 \chi_2 \rightarrow f\bar{f}, WW, hZ$ coannihilation

The annihilation cross section for $\chi_1 \chi_2 \rightarrow f\bar{f}$ for a standard model fermion pair at low velocity, via a magnetic moment interaction, can be expressed as

$$\sigma_{f\bar{f}v} = \alpha \mu^2 \sum_i N_i \tag{A1}$$

where $N_i = 1$ for a hypothetical particle with unit electric charge and no coupling to the $Z$. The sum is over all
kinematically accessible final states. For standard model fermions, and similarly parametrizing the contributions from \( W^+W^- \) and \( Zh \) final states, \( N_{\text{eff}} \) is

\[
N_i = \begin{cases} 
q_i^2 (1 - 2\nu_t l_W \xi + (\nu_t^2 + \alpha_t^2) l_W^2 \xi^2)), & f\bar{f} \\
\frac{m_{\nu_i}^2}{16 \epsilon_w} \xi^2 \psi_{ww}, & WW \\
\frac{m_{\nu_i}^2}{16 \epsilon_w} \xi^2 \psi_{hZ}, & hZ
\end{cases}
\]

The \( \bar{f}\bar{f} \) result is as given in ref. [38], where \( q_i \) is the electric charge, \( \nu_t, \alpha_t \) are the vector and axial-vector couplings respectively to the \( Z \), in units of \( q_i e \), and \( \xi = (1 - m_Z^2 / 4(m_W^2))^{-1} \). For \( WW \) we define \( \psi_{ww} = (1 + 4\epsilon_w - \frac{1}{4} \epsilon_w^2 - \frac{1}{4} \epsilon_w^3) (1 - \epsilon_w)^{1/2} \) with \( \epsilon_w = m_W^2 / m_Z^2 \).

We find that the contribution from neutrinos is \( \sum_i N_i \equiv 3 \xi^2 / (8 \epsilon_W^2) \), and for all charged leptons or quarks \( \nu_t = t_w (4s_w^2 q_i)^{-1} -1 \), \( a_t = t_w / (4s_w^2 q_i) \), giving \( \sum_i N_i = 10.4 \) from fermionic final states, assuming \( m_\chi = 130 \) GeV and \( s_W^2 = 0.23 \). Because of a strong cancellation between the virtual \( \gamma \) and \( Z \) contributions to the \( WW \) channel, it contributes only 0.20 to \( \sum_i N_i \). For \( hZ \), we define \( \psi_{hZ} = (1 - \frac{1}{4} (\epsilon_h - 5 \epsilon_Z) \epsilon_Z + \frac{1}{2} \epsilon_h - \epsilon_Z)^{1/2} \) contributing 0.22 to \( \sum_i N_i \). Hence the total is \( \sum_i N_i = 10.8 \).

For the \( f\bar{f} \) channels, the result from an electric dipole moment interaction is the same, except for an additional velocity-squared suppression factor of \( v^2 / 3 \).

**Appendix B: \( \chi_1\chi_1 \rightarrow \gamma\gamma, \gamma Z, ZZ \) annihilation**

Defining \( r = m_\chi^2 / m_\chi \), and \( \epsilon_Z = \frac{1}{2} (m_Z / m_\chi)^2 \) (superseding the definition of \( \epsilon_Z \) in appendix A), the cross sections (times relative velocity) for \( \chi_1\chi_1 \) to annihilate into \( \gamma\gamma, \gamma Z \) and \( ZZ \) are given by

\[
\langle \sigma v \rangle = \frac{\mu^4 m_\chi^2}{4\pi} \left\{ \begin{array}{ll} 
\frac{4r^2}{(1+r)^2(1-r^2)} & \gamma\gamma \\
2r^2 \frac{(1-r^2)^2(r+r^2)}{(1+r)^2 - r^2} & \gamma Z \\
2 \frac{r^4}{(1-r^2)^2(1+r)^2} & ZZ
\end{array} \right. \]

(B1)

The Dirac case corresponds to \( r = 1 \).

**Appendix C: Explicit \( 3 \times 3 \) DM mixing.**

The parity symmetric \( 3 \times 3 \) Majorana mass matrix \( M \) given in (10) can be diagonalized exactly, by converting it to a block diagonal \( 1 \times 1 \oplus 2 \times 2 \) matrix. The exact mass eigenvalues are \( -m_q \), which corresponds to an eigenvector entirely in the subspace of composite particles, and

\[
\frac{1}{2} \left[ m_\chi + m_\eta \pm \sqrt{(m_\chi - m_\eta)^2 + 8m_y^2} \right]
\]

(C1)

with eigenvectors mixing \( \chi \) and the composite states. The physical masses are then as in equation (12) with

\[
\delta m \equiv \frac{1}{2} (m_\chi - m_\eta) \left[ 1 - \sqrt{1 + 8m_y^2 / (m_\chi - m_\eta)^2} \right]
\]

(C2)

Notice that \( \delta m > 0 \) when \( m_\chi > m_\eta \), leading to level repulsion.

We can now write the three eigenvectors as \( \chi_1 = \cos \theta \eta_2 + \sin \theta \chi_1, \chi_2 = \eta_1, \), and \( \chi_3 = \cos \theta \chi_1 - \sin \theta \eta_2 \) in order of increasing mass and promote them to 4-component Majorana spinors. Then the transition dipole moment between the two composite states becomes

\[
\mu_\eta \eta_2 \sigma_\mu \sigma_\nu F^{\mu\nu} = \mu_\eta \chi_2 \sigma_\mu (\cos \theta \chi_1 - \sin \theta \chi_3) F^{\mu\nu}.
\]

(C3)

The transition moment between the two lightest states is therefore suppressed by the cosine of the mixing angle \( \theta \), and there is no transition moment between the lightest and heaviest states.

Finally, explicitly finding the eigenvectors yields a mixing angle

\[
\cos \theta = 2m_y / \sqrt{2\delta m^2 + 4m_y^2}.
\]

(C4)

In terms of \( m_\chi, m_\eta, \delta m \), this becomes

\[
\cos \theta = \frac{\delta m + m_\chi - m_\eta}{2\delta m + m_\chi - m_\eta}.
\]

(C5)

A straightforwards further algebraic rearrangement is necessary to write this in terms of \( m_\chi / m_\eta \) and the relative mass splitting \( \delta m / (m_\eta - \delta m) \) as in figure 3.

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