Gamma ray lines from TeV dark matter

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Abstract

We calculate, using unitarity, a lower bound on the branching ratio \( \chi \chi \rightarrow \gamma \gamma \) and \( \chi \chi \rightarrow \gamma Z \), where \( \chi \) is any halo dark matter particle that has \( W^+W^- \) as one of the major annihilation modes. Examples of such particles are supersymmetric particles with a dominant Higgsino component, or heavy triplet neutrinos. A substantial branching ratio is found for the \( \gamma \gamma \) and \( \gamma Z \) modes. We estimate the strength of the monoenergetic \( \gamma \) ray lines that result from such annihilations in the Galactic or LMC halos.
The question of the extent and nature of the dark matter in the Universe continues to be one of the most pressing problems of contemporary physics and astrophysics. Large scale structure and peculiar velocity observations seem to favour a value of the energy density much larger than what baryons can contribute due to nucleosynthesis constraints (for a recent review, see [1]). The recent candidate detections of microlensing events in the halo [2] are, although very interesting, not conclusive as concerns the baryonic content of the dark halo of our galaxy (see, e.g., [3]). Even if a large population of dark compact objects (such as brown dwarfs) were formed at an early stage of galactic evolution, the overdensity of baryonic matter thus formed would accrete surrounding cold dark matter creating a particle dark matter halo density at least comparable in magnitude to the baryonic one.

One of the favoured particle dark matter candidates is the lightest supersymmetric particle $\chi$, assumed to be a neutralino, i.e. a mixture of the supersymmetric partners of the photon, the $Z^0$ and the two neutral $CP$-even Higgs bosons present in the minimal extension of the supersymmetric standard model (see, e.g. [4]). The attractiveness of this candidate stems from the fact that its generic couplings and mass range naturally gives a relic density close to the critical one. Besides, its motivation from particle physics has recently become stronger due to the apparent need for 100 GeV - 10 TeV scale supersymmetry to achieve unification of the gauge couplings in view of recent LEP results [5].

When it comes to detecting cold dark matter particles, it seems that the technology is not yet advanced enough for the ultrasensitive detectors developed to register the nuclear recoil and/or ionization to put interesting bounds on supersymmetric dark matter. On the other hand, indirect detection methods look quite promising. With large neutrino telescopes like DUMAND [6] and AMANDA [7] now being deployed, there is a fair chance to detect energetic neutrinos from the center of the Sun or Earth coming from annihilations of captured neutralinos, if they constitute the dark matter halo [8]. In fact, some bounds have already been obtained from the much smaller Kamiokande detector [9].

The other way to detect particle dark matter in the halo is through indirect detection of positrons, antiprotons and $\gamma$ rays generated through the continuous annihilation of dark matter particles in the halo [10]. Observations of $\gamma$ rays have the advantage of giving a sensitive map of the galactic halo (since the annihilation rate depends on the square of the local dark matter density). In particular, if there is an enhancement of the density at the center of our galaxy [11, 12, 13] or the Large Magellanic Cloud [14], the $\gamma$ ray flux could stand out well above background.

A particularly interesting annihilation process in the halo is $\chi \chi \rightarrow \gamma \gamma$ or $\chi \chi \rightarrow Z^0 \gamma$. Since these are two-body final states and the annihilating massive particles move with non-relativistic speed in the halo (typically $v/c \sim 10^{-3}$) the produced $\gamma$s will be nearly monoenergetic, meaning a $\gamma$ ray line signature [15]. The calculations of $\chi \chi \rightarrow \gamma \gamma$ for $m_\chi < m_W$ have been made in quite some detail.
The result is that the signal is only marginally detectable with present-day space detectors [17]. On the other hand, it has recently been pointed out that a very massive $\chi$ in the TeV region could give a detectable signal in Air Cherenkov Telescopes (ACT) on the ground [18]. These detectors have a very large effective area (on the order of 20 000 m$^2$) and a good proton rejection and energy resolution ($\sim 10\%$) can be obtained with modern techniques.

In [18] the $\gamma\gamma$ process was estimated for the case of a nearly pure Bino, based on earlier calculations. However, the rates turn out to be very small [18, 19]. There remains to do the more difficult calculation for a Higgsino, both for the $\gamma\gamma$ and the $Z^0\gamma$ final states. The latter, which to our knowledge is considered for the first time here, will in fact turn out to be the most important one for Higgsinos (and for any other dark matter candidate that has a strong coupling to the $W^+W^-$ final state in its annihilations). Although a full calculation remains to be done for, say, the minimal supersymmetric extension of the standard model, we will use unitarity to put a strict lower bound on the $\gamma$ line signal as well as an estimate of the full rate.

We consider first the Higgsino annihilation $\chi^0\chi^0 \rightarrow W^+W^-$, mediated by chargino exchange in the $t$ and $u$ channel (there are also $Z$ and $H$ exchange contributions in general, but these vanish in the $v \rightarrow 0$ limit). The calculation is identical for any Majorana particle that annihilates to $W$ pairs through the exchange of a charged fermion. We write the $\chi^0\chi^\pm W^\mp$ coupling (see [20, 4] for conventions; the index 0 here identifies the lighter of the two charginos)

$$\Gamma^\mu = \frac{ig}{2}\gamma^\mu(f_0P_L + e_0P_R),$$ \hspace{1cm} (1)

where $P_{L,R} = (1 \pm \gamma^5)/2$.

In the limit where the lightest neutralino is a pure heavy Higgsino, the $\chi\chi \rightarrow WW$ amplitude is dominated by the exchange of the lightest chargino, nearly degenerate with the neutralino, and

$$|e_0| = |f_0| = 1$$ \hspace{1cm} (2)

Projecting out the $S$ wave part of the initial amplitude by the projector [21]

$$\mathcal{O}_{Ps} = -\frac{m_\chi}{\sqrt{2}}\gamma^5(1 - \frac{p_\chi}{m_\chi}),$$ \hspace{1cm} (3)

we find the effective $\chi\chi W^+W^-$ vertex for polarization indices $\mu$ and $\nu$ of the $W$ bosons with four-momenta $p_+$ and $p_-$:

$$V_{\chi\chi W^+W^-} = \left(\frac{g^2(e_0^2 + f_0^2)}{2\sqrt{2}(m_{\chi^0}^2 + m_{\chi^+}^2 - m_W^2)}\right)\epsilon[\mu\nu p_-p_+],$$ \hspace{1cm} (4)
where $\epsilon_{\mu\nu\rho\sigma}$ is the completely antisymmetric constant tensor, and we use the abbreviation $\epsilon[\mu\nu p_+] \equiv \epsilon_{\mu\nu\rho\sigma} p^\rho p^\sigma$. This gives the annihilation rate

$$\sigma v(\chi\chi \to W^+W^-) = \frac{g^4}{128\pi m_\chi m_W} \frac{(\omega - 1)^{3/2}(e_0^2 + f_0^2)^2}{(1 - \omega - \kappa)^2},$$  

(5)

where $\omega = (m_\chi^0/m_W)^2$, $\kappa = (m_\chi^+/m_W)^2$. This result coincides with the $v \to 0$ limit of the expression given in [20].

We are now prepared for the one-loop calculation of $\chi\chi \to \gamma\gamma$. We assume $m_{\chi^+} \gtrsim m_{\chi^0} >> m_W$. It has been shown in a similar type of calculation [22] that it is a very good approximation, even when $m_{\chi^+} \sim m_{\chi^0}$, to take the effective vertex (4) as pointlike:

$$V_{\chi\chi W W} \sim C \epsilon[\mu\nu k(k - q_1 - q_2)],$$  

(6)

with

$$C = \frac{g^2(e_0^2 + f_0^2)}{2\sqrt{2}m_{\chi^+}^2}$$  

(7)

(see Fig. 1). For the contribution of the imaginary part to the branching ratio which we will extract to obtain a lower bound on the cross section, expression (4) is even exact because the denominator of the chargino propagator is constant. One can easily convince oneself that the direct $\chi\chi W W \gamma$ vertex (Fig. 1(b)) that by gauge invariance has to be present due to the derivative coupling in the effective vertex (4) does not contribute to our process. In addition, the ordinary contact $W W \gamma\gamma$ term (Fig. 1(c)) contributes zero, since the initial low-velocity $\chi\chi$ state only projects out the part of the amplitude antisymmetric in the $W$ momenta, and the contact term is symmetric in its indices and thus vanishes upon contraction.

There remains to calculate the triangle graphs with $W$s and in the customary linear $R_\xi$ gauge also Goldstone bosons circulating in the loop (Fig. 1(a)). Considerable simplification is obtained by choosing a non-linear gauge condition as originally suggested by Fujikawa [23]. This has been used in various applications such as calculating Higgs decays into two photons [24, 25], the charge radius of the neutrino [26] or radiative neutralino decay [27]. The idea is to replace the usual gauge fixing term (we set $\xi = 1$ for simplicity)

$$\mathcal{L}_{g.f.} = -\frac{1}{2}(\partial_\mu A^\mu)^2 - \frac{1}{2}(\partial_\mu Z^\mu + m_Z G^0)^2 - |\partial_\mu W^{+\mu} + im_W G^+|^2,$$

(8)

where $G^{\pm,0}$ are the Goldstone bosons eaten by $W^\pm$ and $Z^0$, by

$$\mathcal{L}^{n.f.}_{g.f.} = -\frac{1}{2}(\partial_\mu A^\mu)^2 - \frac{1}{2}(\partial_\mu Z^\mu + m_Z G^0)^2 - |(\partial_\mu + igW^3_\mu)W^{+\mu} + im_W G^+|^2,$$

(9)

where $W^3_\mu$ is the neutral component of the $SU(2)_L$ gauge triplet ($W^3_\mu = Z^0_\mu \cos \theta_w - A^\mu_\mu \sin \theta_w$). The main advantage of this gauge is that the new contribution from
(9) to the trilinear $\gamma W^\pm G^\mp$ part of the Lagrangian actually cancels a corresponding piece of the original trilinear sector of the Lagrangian giving a vanishing total $\gamma W^\pm G^\mp$ coupling. In addition, the new terms entering the $\gamma W^\pm W^\mp$ and $Z^0 W^\pm W^\mp$ vertices, although making them superficially look more complicated (and less symmetric in the momenta of the three bosons) in reality make these vertices much simpler when at least one of the bosons is on mass shell. The diagram in Fig. 1(a) (and the one obtained by crossing the photon lines which actually gives an identical contribution) were calculated in the 't Hooft-Feynman version ($\xi = 1$) of the nonlinear gauge (in fact, we also checked our results using the more cumbersome linear gauge).

We find the amplitude for $\chi \chi \rightarrow \gamma(q_1)\gamma(q_2)$ with photon polarization four vectors $\epsilon^\mu(q_1)$ and $\epsilon^\nu(q_2)$ to be

$$A^\mu\nu = \frac{-8ie^2C}{(2\pi)^4} \int d^4k \frac{m^2_W[k\mu\nu(q_1 + q_2)] + k_\mu\epsilon[q\nu_1q_2] - k_\nu\epsilon[q\mu_1q_2]}{(k^2 - m^2_W)(k^2 - m^2_W)((k - q_1)^2 - m^2_W)((k - q_1 - q_2)^2 - m^2_W)},$$

where the ordinary prescription $m^2_W \rightarrow m^2_W - i\epsilon$ is understood in the denominator.

Using standard techniques, this integral can be transferred to a symmetric integral in $D = 4 - 2\epsilon$ dimensions after combining the factors in the denominator using Feynman parametrization. The $D$-dimensional integrals are of two types, one convergent and one logarithmically divergent as $D \rightarrow 4$ (signalled by an $1/\epsilon$ divergence as $D \rightarrow 4$). The (infinite) renormalization constant is real and does not affect the imaginary part, which can be calculated in terms of elementary functions since the 2-dimensional Feynman parameter integral becomes 1-dimensional after using the $\delta$-function coming from the $i\epsilon$ term in the denominator. (This corresponds to using the Cutkosky rules to extract the imaginary part of the amplitude.) Writing

$$A_{\gamma\gamma} = \frac{C e^2}{4\pi} \epsilon[q\nu_1q_2] (M_{re} + iM_{im}),$$

we find

$$M_{im} = \beta^2 \log\left(\frac{1 + \beta}{1 - \beta}\right),$$

where $\beta = \sqrt{1 - (m_W/m_\chi)^2}$.

In the limit of large $\chi$ masses (above, say, 500 GeV) this becomes

$$M_{im} \sim \log\left(\frac{4m^2_\chi}{m^2_W}\right).$$

We have checked this result using still another method of calculation that relies on the generalized optical theorem, writing (the integral is over the $WW$ phase space)

$$M_{im} = \frac{1}{2} \int M_{\chi\chi \rightarrow WW} M_{\gamma\gamma \rightarrow WW}^*,$$

(14)
where this time we employ the chargino propagator rather than the pointlike approximation for the $\chi\chi WW$ vertex. This gives the same result and thus in addition proves the assertion that for the calculation of the imaginary part, the approximation of a pointlike effective vertex is exact whenever $m_{\chi^\pm} > m^0_{\chi}$.

Once the imaginary (absorptive) part has been found, the real (dispersive) part can be calculated using dispersion relation techniques. However, for the pointlike vertex a subtraction will be needed, corresponding to a constant term resulting from the chargino propagator cutoff in the diagrams that generated the effective vertex. For very large $\chi$ masses the leading logarithmic terms will dominate and can be calculated:

$$M^{LL}_{im} = \log\left(\frac{s}{m_W^2}\right), \quad (15)$$

$$M^{LL}_{re} = \frac{1}{2\pi} \log^2\left(\frac{s}{m_W^2}\right), \quad (16)$$

where $s = 4m_{\chi}^2$.

In the following we will use the imaginary part to get a strict lower bound for the rate estimates. To get an idea of the full result we will use the leading log estimate (16) for the real part.

For the application we are considering, annihilation of heavy non-relativistic dark matter particles in the galactic halo, the $\gamma$ ray line coming from $\chi\chi \rightarrow \gamma\gamma$ cannot realistically be discriminated from the corresponding line from $\chi\chi \rightarrow Z^0\gamma$. For example, the $\gamma$ energy for $m_{\chi} = 1$ TeV is 1 TeV for the $\gamma\gamma$ mode and $(4m_{\chi}^2 - m_Z^2)/(4m_{\chi}) = 998$ GeV for the $Z^0\gamma$ mode, i.e. within the $10^{-3} - 10^{-2}$ spread of velocities of the annihilating particles (and of course well within the $5 - 10\%$ energy resolution that at most can be achieved with present-day Air Cherenkov Telescopes). We should therefore add the $\gamma$ ray luminosities from the two sources. In fact, as we now shall see, the line strength from $Z^0\gamma$ is expected to be a few times larger than that from the $\gamma\gamma$ final state.

In the gauge we have chosen and to leading logarithmic accuracy, the difference between the $\gamma\gamma$ and $Z^0\gamma$ calculations is simply the replacement $e^2 \rightarrow e^2 \cot \theta_w$ in Eq. (11). In the rate, this amounts to an enhancement by the factor $\cot^2 \theta_w \sim 3.4$. When calculating the intensity of the line signal, the fact that each $\gamma\gamma$ event gives two photons is compensated by the division of the symmetry factor ($= 2$) due to the presence of two identical particles in the final state.

In Fig. 2 we show our lower bound for the average number of $\gamma$ line photons of energy $\sim m_{\chi}$ per $W^+W^-$ annihilation:

$$F_{\gamma} \equiv \frac{2\sigma v(\chi\chi \rightarrow \gamma\gamma) + \sigma v(\chi\chi \rightarrow Z^0\gamma)}{\sigma v(\chi\chi \rightarrow W^+W^-)} = \alpha_{e.m.}^2 \left(1 + \cot^2 \theta_w\right) \log^2\left(\frac{4m_{\chi}^2}{m_W^2}\right), \quad (17)$$

as well as the value for this rate resulting from the adoption of (16) for the dispersive part of the amplitude. When converting this into an effective line flux
per $\chi \chi$ annihilation one has to correct by a factor giving the branching ratio of $\chi \chi \rightarrow W^+W^-$ normalized to the total annihilation cross section. This typically means a reduction of a factor of 2 for a Higgsino (due mainly to the $ZZ$ channel), but no reduction for a species (e.g. a triplet neutrino [28]) that does not couple to $Z$ bosons at tree level. The remarkable feature of our result [17] is its very large size compared to the case of, e.g. a pure Bino [18, 19]. As can be seen in Fig. 2, the effective branching ratio to line photons can easily be as large as $10^{-2}$. There will of course also be a continuous, diffuse flux of lower energy $\gamma$s coming from the fragmentation of the $W^+W^-$ and $ZZ$ final states (see [18]). This distribution however lacks conspicuous features and has its main contribution at very low energies where the background $\gamma$ flux is much more severe and, what is worse, its shape is largely unknown. It may therefore prove very difficult to detect dark matter candidates through this "soft" photon flux. If detectors (e.g. ACTs) customized to achieve high energy resolution are built, a line signal would be an obvious feature to search for in the data. If such a line is found, there is certainly no known astrophysical background that could account for it, and the energy would simply correspond to the mass of the dark matter particle.

Although we have focused on the Higgsino, one of the presently favoured dark matter candidates, our results are in fact of much greater generality for the following reasons. Any massive (TeV scale) neutral cold dark matter fermion may be expected to have a substantial annihilation strength into $W^+W^-$ pairs (although exceptions can be found, such as pure Binos). The fermion can be of Majorana or Dirac nature. For Majoranas, our analysis can be directly applied. The difference with Dirac fermions is that the CP properties do not forbid the slow fermions in the halo to annihilate in the triplet $S$ (vector) state. However, due to gauge invariance, a massive spin 1 state can not decay into two photons, so the only contribution comes from the singlet $S$ state which is the one we have already treated. Thus, we expect the logarithmic enhancement to be a quite general feature. The actual branching ratio line photons may, however, be lower by a factor of a few due to the new triplet contributions to the total annihilation rate.

In the case of a heavy ($m_N > m_W$) weak triplet Majorana neutrino (which was considered in [28] for $m_N < m_W$), our results need no modification, since $WW$ is the only important tree level annihilation mode. That model is an interesting example of a viable particle dark matter model where direct detection is essentially impossible (since the neutrino does not have tree level $Z$ couplings), but where $\gamma$ ray lines are predicted at a substantial level.

We now turn to some estimates of the rates for the gamma line processes. Since the annihilation rate depends on the square of the neutralino number density, the rate and angular distribution of the $\gamma$ line flux will depend on the model for the Galactic halo. In particular, there are some arguments that there may be a core region of the galaxy where the dark matter density is substantially higher than average [11]. In the model by Berezinsky et al [12] the density profile falls
as $1/r^{1.8}$ all the way from 5 Mpc down to the innermost part of the galaxy, where a black hole is thought to be residing. This model gives an enormous enhancement (several orders of magnitude) of the $\gamma$ ray flux from $\chi$ annihilations in the central region of the galaxy. Assuming this halo density profile, already existing measurements of high energy cosmic gamma ray fluxes constrain very strongly the supersymmetric models [29] (for a recent criticism of this model for the halo, see [30]).

On the more conventional side, it has recently been pointed out [14] that the Large Magellanic Cloud may have its own halo (albeit with only partially known parameters), which possibly could give an enhanced flux from that direction. We will consider these cases, noting as in [18] that for a smooth halo the rates we find are too small to be detectable with present technology. On the other hand, one of the the main advantages for turning the attention to the $\gamma$ ray signatures is just that this is a promising way of mapping the density distribution of the halo in the case it is not completely homogeneous. For instance, there is a possibility that string or texture generated perturbations in the cold dark matter density could have gone non-linear early and survived tidal disruptions to provide a very clumpy halo [13].

For masses up to a few TeV, we can neglect galactic absorption of the $\gamma$ rays. According to [31] the optical depth at the 10 kpc length scale is around $10^{-3}$. (Remember that for the high mass range we are considering, even for the $Z\gamma$ channel, $E_\gamma = m_\chi$ to an accuracy of within a percent.)

For a localized source of dark matter annihilations at the center of the galaxy, the gamma line flux is given by

$$F_\gamma = \frac{(\sigma v)_\gamma}{r_{\odot}^2} \int_0^{r_0} \left( \frac{\rho(r)}{m_\chi} \right)^2 r^2 dr,$$

where $r_\odot \approx 8.5$ kpc is the distance of the solar system from the center of the galaxy, and $\rho(r)$ is the density profile within the core radius $r_0$. For a generic dark matter particle that gives a substantial contribution to $\Omega$ of the universe, $(\sigma v)_{tot}$ is of the order of $10^{-26}$ cm$^3$s$^{-1}$. Since the annihilation rate of neutralinos into $W^+W^-$ does not vanish nor become small when $< v > \to 0$, $(\sigma v)_{tot}$ is expected to be of the same order of magnitude in the galactic halo. As an example, this happens for a Higgsino of the minimal supersymmetric extension of the Standard Model of a mass around 1 TeV. Since we see from Fig. 2 that the number of line photons per annihilation is roughly of the order of a percent, we can use the value $(\sigma v)_{line} \sim 10^{-28}$ cm$^3$s$^{-1}$ in our estimates.

In fact, for the case of a pure Higgsino, we can calculate the absolute value of the lower bound of the line cross section. In that case $f = 1$, $\epsilon_0 = \pm 1$ in Eq. (1), and we find (assuming near Higgsino-chargino degeneracy, which is generally an excellent approximation for a nearly pure, massive Higgsino in minimal supersymmetric models)
\[ 2\sigma_{\gamma\gamma}(\chi\chi \to \gamma\gamma) + \sigma_{Z\gamma}(\chi\chi \to Z^0 \gamma) > \frac{\pi \alpha_m^4}{8m_\chi^2 \sin^6 \theta_w} \log^2 \left( \frac{4m_\chi^2}{m_W^2} \right), \quad (19) \]

which corresponds to \(0.6 \cdot 10^{-28} \text{ cm}^3\text{s}^{-1}\) for \(m_\chi = 1 \text{ TeV}\). Assuming, lacking a full calculation, a contribution from the dispersive part of the amplitudes of the same order of magnitude, we see that in this case the annihilation rate into line photons is indeed around \(10^{-28} \text{ cm}^3\text{s}^{-1}\).

Using the Ipser-Sikivie model for the core of the halo \[11\], a flux of around 200 photons per year is predicted in an ACT of area 20 000 m\(^2\) during a typical observation ”year” of \(2 \cdot 10^6 \text{ s}\), with an estimated (but poorly known) background about an order of magnitude lower (including misidentified protons and electrons; see \[18\] for further discussions). This rate for Higgsino-like particles is two orders of magnitude larger than that expected for pure Binos \[18, 19\]. In the model by Berezinsky et al \[12\], the flux can be still at least two orders of magnitude larger. However, if one assumes a smooth halo distribution \(\rho(r) = \rho_\odot (r_\odot^2 + a^2)/(r^2 + a^2)\) with \(a \sim r_\odot\), the rate falls below one event per year in our generic ACT detector in the direction of the galactic center. In the LMC model of Gondolo \[14\] the rate could be enhanced by an order of magnitude making it just about reaching the detectability limit. It should be noted, though, that the high threshold (200-300 GeV) of present day ACTs make these processes rate limited. Going to larger areas and lower thresholds could greatly increase the discovery potential of this type of detector.

To conclude, we have shown that the \(\gamma\gamma\) and \(Z\gamma\) annihilation rates are quite high for any dark matter candidate that couples with full strength to \(W^\pm\). A typical case is provided by a pure Higgsino, for which we have shown that the \(\gamma\) line strength is very much larger than that of a pure Bino. It should be noticed, though, that our calculations indicate this large \(\gamma\) line rate to be generic for any dark matter candidate that couples to \(W^\pm\) with ordinary electroweak strength. If the galactic halo contains regions with higher dark matter density than the local (solar neighbourhood) value, the detection of high energy monoenergetic photons could at the same time determine the mass of the dark matter particle and map its galactic density distribution.

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References

[1] See, e.g. M.S. Turner, Physica Scripta T36 (1991) 167.

[2] C. Alcock et al, Nature 359 (1993) 393; E. Aubourg et al., Nature 365 (1993) 623.

[3] E. Gates and M.S. Turner, FERMILAB-PUB-93-357-A (1993).

[4] H.E. Haber and G.L. Kane, Phys. Rep. 117 (1985) 75.

[5] U. Amaldi, W. de Boer and H. Furstenau, Phys. Lett. B260 (1991) 447.

[6] See, e.g., P.K.F. Greider, Europhysics News 23 (1992) 167.

[7] See, e.g., D.M. Lowder et al., Nature 353 (1991) 331.

[8] J. Silk, K. Olive and M. Srednicki, Phys. Rev. Lett. 55 (1985) 257; G.B. Gelmini, P. Gondolo and E. Roulet, Nucl. Phys. B351 (1991) 623; F. Halzen, M. Kamionkowski and T. Stelzer, Phys. Rev. D45 (1992) 4439.

[9] M. Mori et al., Phys. Rev. D48 (1993) 5505.

[10] J. Silk and M. Srednicki, Phys. Rev. Lett. 53 (1984) 624; F. Stecker, S. Rudaz and T. Walsh, Phys. Rev. Lett. 55 (1985) 2622; J. Ellis et al, Phys. Lett. B214 (1989) 403; M. Turner and F. Wilczek, Phys. Rev. D42 (1990) 1001.

[11] J.R. Ipser and P. Sikivie, Phys. Rev. D35 (1987) 3695.

[12] V.S. Berezinsky, A.V. Gurevich and K.P. Zybin, Phys. Lett. B294 (1992) 221.

[13] J. Silk and A. Stebbins, Astrophysical Journal 411 (1993) 439.

[14] P. Gondolo, contributed paper to TAUP 93, Gran Sasso, Italy; to appear in the Proceedings.

[15] L. Bergström and H. Snellman, Phys. Rev. D37 (1988) 3737; S. Rudaz, Phys. Rev. D39 (1989) 3549; A. Bouquet, P. Salati and J. Silk, Phys. Rev. D40 (1989) 3168.

[16] R. Barbieri and G. F. Giudice, Nucl. Phys. B306 (1988) 63; L. Bergström, Phys. Lett. 225B (1989) 372.

[17] V.S. Berezinsky, A. Bottino and V. de Alfaro, Phys. Lett. B274 (1992) 122.

[18] M. Urban, A. Bouquet, B. Degrange, P. Fleury, J. Kaplan, A.L. Melchior and E. Paré, Phys. Lett. B293 (1992) 149.
[19] Note that there is an error in Eq. (20) of [18], entirely due to one of the present authors (JK). The prefactor should read $8 \times 10^{-31}$ cm$^3$s$^{-1}$ instead of $4.2 \times 10^{-29}$ cm$^3$s$^{-1}$.

[20] K. Griest, M. Kamionkowski and M. Turner, Phys. Rev. D41 (1990) 3565.

[21] J.H. Kühn, J. Kaplan and O. Safiani, Nucl. Phys. B157 (1979) 125.

[22] L. Bergström, Phys. Lett. 225B (1989) 372.

[23] K. Fujikawa, Phys. Rev. D7 (1973) 393.

[24] M.B. Gavela, G. Girardi, C. Malleville and P. Sorba, Nucl. Phys. B193 (1981) 257.

[25] L. Begström and G. Hulth, Nucl. Phys. B259 (1985) 137.

[26] N.M. Monyonko and J.H. Reid, Prog. Theor. Phys. 73 (1984) 734.

[27] H.E. Haber and D. Wyler, Nucl. Phys. B323 (1989) 267.

[28] P. Chardonnet and P. Salati, Phys. Lett. B262 (1991) 307.

[29] M. Urban, A. Bouquet, B. Degrange, M. Chantell, P. Fleury, J. Kaplan, A.L. Melchior, E. Paré, X. Sarazin and T. Weekes, Proceedings of the International Conference on Cosmic Ray Physics (Calgary, July 1993), preprint PAR-LPTHE 93-19; V.S. Berezinsky, A. Bottino and G. Mignola, Gran Sasso preprint LNGS 94/90 (1994).

[30] R.A. Flores and J. Primack, UC Santa Cruz preprint SCIPP 93/01 (1994).

[31] V.S. Berezinsky et al., *Astrophysics of Cosmic Rays*, North-Holland, 1990.
Figure Captions

1. Some diagrams relevant to the annihilation $\chi\chi \rightarrow \gamma\gamma$. For slow Majorana particle annihilation, only the diagram in (a) contributes (the diagrams obtained by crossing the photon lines are not shown in the figure).

2. The average number $F_\gamma$ (Eq. (17)) of $\gamma$ line photons from $\chi\chi \rightarrow \gamma\gamma$ and $Z^0\gamma$ normalized to the $W^+W^-$ annihilation rate, as a function of the $\chi$ mass. The solid line is the lower bound obtained by using only the imaginary part of the amplitude which is given by unitarity. The dashed line is an estimate of the full rate obtained by taking the leading logarithmic result of Eq. (16) for the real (dispersive) part of the amplitude.
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