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Neutralino dark matter annihilation to monoenergetic gamma rays as a signal of low mass superstrings

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Abstract

We consider extensions of the standard model based on open strings ending on D-branes, in which gauge bosons and their associated gauginos exist as strings attached to stacks of D-branes, and chiral matter exists as strings stretching between intersecting D-branes. Under the assumptions that the fundamental string scale is in the TeV range and the theory is weakly coupled, we study models of supersymmetry for which signals of annihilating neutralino dark matter are observable. In particular, we construct a model with a supersymmetric R-symmetry violating (but R-parity conserving) effective Lagrangian that allows for the s-wave annihilation of neutralinos, once gauginos acquire mass through an unspecified mechanism. The model yields bino-like neutralinos (with the measured relic abundance) that annihilate to a $\gamma\gamma$ final state with a substantial branching fraction ($\sim 10\%$) that is orders of magnitude larger than in the minimal supersymmetric standard model. A very bright gamma-ray spectral line could be observed by gamma-ray telescopes.

Superstring theory is a promising candidate to explain the underlying symmetries of nature, e.g., the probable existence and breaking of supersymmetry (SUSY). In particular, TeV-scale superstring theory provides a brane-world description of the standard model, which is localized on hyperplanes extending in $p+3$ spatial dimensions, the so-called D-branes. Gauge interactions emerge as excitations of open strings with endpoints attached on the D-branes, whereas gravitational interactions are described by closed strings that can propagate in all nine spatial dimensions of string theory (these comprise parallel dimensions extended along the $(p+3)$-branes and transverse dimensions). The apparent weakness of gravity at energies below a few TeV can then be understood as a consequence of the gravitational force "leaking" into the transverse compact dimensions of spacetime. This is possible only if the intrinsic scale of string excitations is also of order a few TeV. Should nature be so cooperative, one would expect to see a few string states produced at the LHC, most distinctly manifest in the dijet [1] and γ+jet [2] spectra resulting from their decay.

An attractive feature of broken SUSY is that with R-parity conservation the lightest supersymmetric particle (LSP) is a possible candidate for cold dark matter [3]. Requiring the relic abundance to conform with cosmological dark matter measurements serves to constrain the underlying theory. A consequence is that it may be possible to detect the annihilation products of such particles, such as gamma rays, charged leptons, and neutrinos.

In this Letter, we propose new processes, based in brane-world string theory, for the efficient annihilation of neutralino LSP's (χ^0's) into monochromatic gamma rays, Z-bosons, charged Ws and pairs of gluons (via $\chi^0\chi^0 \to \gamma\gamma$, γZ, ZZ, W^+W^-, and gg). By requiring that the total annihilation rate generate the measured dark matter abundance [4], we constrain the parameters of the model: the string scale, the neutralino mass, the string coupling constant, and two unknown dimensionless parameters that depend on the details of compactification. We then calculate the gamma-ray spectrum from neutralino annihilation in the central region of the Milky Way, and explore the prospects for discovery with present and future gamma-ray telescopes.

The basic unit of gauge invariance for D-brane constructions is a $U(1)$ field, so that a stack of N identical D-branes generate a $U(N)$ theory with the associated $U(N)$ gauge group. Gauge bosons and associated gauginos (in a supersymmetric theory) arise from strings terminating on *one* stack. For simplicity we consider a model with 3 stacks corresponding to gauge groups $U(3) \times U(2) \times U(1)$, labeled stacks a, b, and c, respectively.

We consider the introduction of new operators, based on superstring theory, that avoid p-wave suppression by permitting s-wave annihilation into gauge bosons at an adequate rate. To create an s-wave, both gauginos must be in the same helicity state, either left- or right-handed. Such gaugino pair annihilation violates R-symmetry by $\Delta r = \pm 2$, and is therefore forbidden in supersymmetric Yang-Mills theory, at least at the perturbative level. However, it can appear in conjunction with a SUSY-breaking gaugino mass generation mechanism.

In superstring theory, there are no conserved charges associated with continuous global symmetries; consequently *even with unbroken SUSY*, R-symmetry can be violated by higher-dimensional operators, although only at certain orders of perturbation theory. The R-charge deficit Δr is related to the Euler characteristic of the string worldsheet, $\chi = 2-2g-h$, where g is the genus and h is the number of boundaries: $|\Delta r| \leq -2\chi$, with $\chi \leq 0$. Of course, only SUSY-preserving interactions are allowed in string perturbation theory. Note that if SUSY is broken at the string level, these additional restrictions are lifted. As an example, consider the disk worldsheet with $g=0,\ h=1$, hence $\chi=1$, which incorporates the effects of all tree-level interactions, including the exchanges of virtual Regge excitations. Recall that gaugino

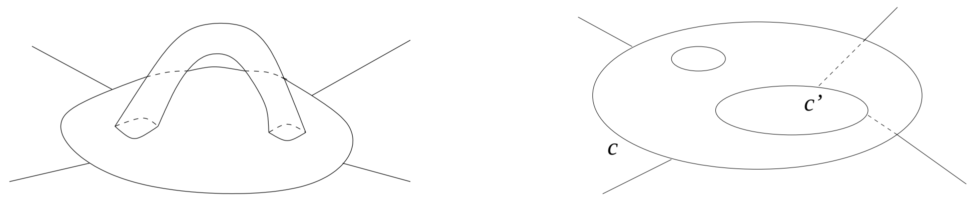

FIG. 1: Left: "Genus 3/2" worldsheet with one handle, with four vertices inserted at the boundary. Right: Two-loop open string worldsheet with two vertices inserted at the boundary c and two at c' while the third one is "empty".

and gauge boson vertex operators are inserted at the disk boundary "attached" to the associated stack of D-branes. In this case, $\Delta r = 0$, so that two like-helicity gauginos cannot annihilate into gauge bosons and the amplitude for $\lambda^{\pm}\lambda^{\pm}$ annihilation into an arbitrary number of gauge bosons vanishes at the disk level, to all orders in $\alpha' = 1/M_s^2$. This can be confirmed by using SUSY Ward identities along the lines of Ref. [5]. For a gaugino pair to annihilate into gauge bosons one needs a worldsheet with Euler characteristic $\chi = -1$. It can be realized in two ways: a "genus 3/2" worldsheet with $g = 1,\ h = 1$ [6, 7], which is essentially a disk with a closed string handle depicted in the left-hand side of Fig. 1, and $g = 0,\ h = 3$, which is a two-loop open string worldsheet [8] depicted in the right-hand side of Fig. 1. The case of a worldsheet with three boundaries is particularly interesting [8]. If one inserts two vertex operators creating gauge bosons or gauginos associated with one stack of D-branes, say c, at a single worldsheet boundary, two vertex operators associated with stack c' at another boundary, while keeping "empty" the third boundary as in the right-hand side of Fig. 1, one obtains a non-vanishing contribution to the four-point scattering amplitude. The corresponding effective interaction is described by the supersymmetric F-term,

$$\mathcal{L}_{\rm int} = 3\ g_s^3\ N\ M_s^{-3}\, \widetilde{F}^{(0,3)}\, ({\rm Tr}\, W^c_\alpha \epsilon^{\alpha\beta} W^c_\beta)({\rm Tr}\, W^{c'}_\gamma \epsilon^{\gamma\delta} W^{c'}_\delta)\big|_{\theta^2}\ +\ c.c., \tag{1}$$

where W are the usual chiral superfields with field strengths associated with appropriate gauge groups and the traces are taken in the fundamental representations. Here, g_s is the string coupling, and N is the number of D-branes attached to the empty boundary. (A total of six possibilities in the three-stack model under consideration.) The factor of 3 is the number of choices of the empty boundary. If the stack $c' \neq c$, the above contribution yields the full amplitude. The factor $F^{(0,3)} = 3N\widetilde{F}^{(0,3)}$ is the genus zero topological partition function on a worldsheet with $h = 3$ boundaries. It depends on the moduli of the compact space and takes into account various string configurations in six internal dimensions. The corresponding amplitudes are called "topological" because they are determined by the topology of the compact dimensions and, unlike standard amplitudes, they do *not* contain any kinematical singularities associated with Regge excitations. On the other hand, if $c' = c$, the four-point amplitude also receives non-topological contributions from all four vertices inserted at the same boundary [9].[1] In this Letter, we are mainly interested in the former case, with all bino-like gauginos associated with the $U(1)$ stack and gauge bosons associated with all three

[1] The amplitudes induced by the "genus 3/2" worldsheets, which involve closed strings propagating in the handle, can be analyzed in a similar way. These amplitudes are related to (and, in some sense, they are "square roots" of) the genus 2 topological amplitudes in type II string theory [9]. Although they are not

stacks. For our purposes, it is sufficient to focus on the effective interaction term of Eq. (1). In the case where the annihilation occurs through emission of a pair of $U(1)$ gauge bosons, the interaction will include unknown non-topological and "genus 3/2" contributions. Note that such a term can be induced not only by string physics, but also by any R-symmetry violating extension of the standard model at a sufficiently low energy scale.

The interaction term contained in Eq. (1), relevant to s-wave gaugino annihilation, is

$$\mathcal{L}_{\lambda\lambda} = \frac{\mathcal{T}}{8}\, M_s^{-3}\, (\delta_{c_1 c_2} \lambda^{c_1}_{\alpha_1} \epsilon^{\alpha_1 \alpha_2} \lambda^{c_2}_{\alpha_2})(\delta_{c'_3 c'_4} F^{c'_3 \beta_3}_{\alpha_3} \epsilon^{\alpha_3 \alpha_4} \epsilon_{\beta_3 \beta_4} F^{c'_4 \beta_4}_{\alpha_4}) \; + \; c.c., \tag{2}$$

with lower case Latin and Greek letters labeling gauge and spinor indices, respectively, and the dimensionless coupling constant

$$\mathcal{T} = 3N\, g_s^3\; \widetilde{F}^{(0,3)}\,. \tag{3}$$

In Eq. (2), F_α^β denote self-dual gauge field strengths in the spinorial representation,

$$F_\alpha^\beta = \frac{1}{2} F_{\mu\nu} \sigma^\mu_{\alpha\dot{\alpha}} \bar{\sigma}^{\nu\dot{\alpha}\beta} = F_{\mu\nu} \sigma^{\mu\nu\beta}_{\alpha} : \qquad F_\alpha^\beta F_\beta^\alpha = -F^{\mu\nu} F_{\mu\nu} + \frac{i}{2}\, \epsilon_{\mu\nu\rho\lambda} F^{\mu\nu} F^{\rho\lambda}. \tag{4}$$

In order to compute the amplitude for the gaugino pair annihilation into two gauge bosons, we need the wavefunctions of all particles in addition to the interaction term (2). We consider the case of massive gauginos, with Majorana mass m, without addressing details of the SUSY breaking mechanism. In the center of mass frame, two gauginos moving along the z-axis with three-momenta $\mathbf{k}$ and $-\mathbf{k}$, respectively, are described by the Majorana spinors:

$$\begin{pmatrix} u \\ \bar{u} \end{pmatrix} \quad \text{with} \quad u^+(\pm\mathbf{k}) = \sqrt{m} \begin{pmatrix} e^{\mp\eta/2} \\ 0 \end{pmatrix} \;\; , \quad u^-(\pm\mathbf{k}) = \sqrt{m} \begin{pmatrix} 0 \\ e^{\pm\eta/2} \end{pmatrix}, \tag{5}$$

where $u^\pm$ refer to spin up and down, respectively, and the rapidity

$$\eta = \sinh^{-1} \left(\frac{|\mathbf{k}|}{m} \right). \tag{6}$$

On the other hand, the polarization vectors for the gauge bosons are

$$\epsilon_\mu^\pm (k, q) = \pm \frac{\langle q^\mp | \gamma_\mu | k^\mp \rangle}{\sqrt{2} \langle q^\mp | k^\pm \rangle}\,, \tag{7}$$

where $\epsilon^\pm$ refer to helicities, k is the momentum and q is an arbitrary reference vector. In Eq. (7) and below, we use the notation of Ref. [10]. Since,

$$F_\alpha^\beta |_{F_{\mu\nu} = \epsilon^+_\mu k_\nu - \epsilon^+_\nu k_\mu} = 0\,, \tag{8}$$

the interaction term written explicitly in Eq. (2) couples only to $(--)$ gauge boson helicity configurations while its complex conjugate couples only to $(++)$ configurations. In all, there are only four non-vanishing helicity amplitudes:

$$\mathcal{M}(\lambda_-^{c_1} \lambda_-^{c_2} \;\rightarrow\; g_+^{c'_3} g_+^{c'_4}) \;\equiv\; \mathcal{M}_{--} \quad , \quad \mathcal{M}(\lambda_+^{c_1} \lambda_+^{c_2} \;\rightarrow\; g_+^{c'_3} g_+^{c'_4}) \;\equiv\; \mathcal{M}_{+-}\;, \tag{9}$$

strictly topological, one does not expect kinematical singularities to appear in the intermediate channels. The only difference from the two-loop open string worldsheets is that in the genus 3/2 case gauginos and gauge bosons must belong to the same stack of D-branes.

$$\mathcal{M}(\lambda_+^{c_1}\lambda_+^{c_2} \;\to\; g_-^{c'_3}g_-^{c'_4}) \;\equiv\; \mathcal{M}_{++} \quad , \quad \mathcal{M}(\lambda_-^{c_1}\lambda_-^{c_2} \;\to\; g_-^{c'_3}g_-^{c'_4}) \;\equiv\; \mathcal{M}_{-+} \;, \tag{10}$$

where the notation corresponds to all particles incoming.

The effective interaction F-term (1) yields

$$\mathcal{M}_{--} = 2\,\mathcal{T} M_s^{-3}\,\delta^{c_1c_2}\delta^{c'_3c'_4}\epsilon^{\alpha_1\alpha_2}u^-_{\alpha_1}(\mathbf{k}_1=\mathbf{k})\,u^+_{\alpha_2}(\mathbf{k}_2=-\mathbf{k})\,\epsilon^-_{3\mu}k_{3\nu}\epsilon^-_{4\rho}k_{4\lambda}\mathrm{Tr}(\sigma^{\mu\nu}\sigma^{\rho\lambda})\;. \tag{11}$$

Using Eqs. (5) and (7), we obtain

$$\mathcal{M}_{--} \;= \mathcal{T} M_s^{-3} m\, e^{\eta}\langle 34\rangle^2\,, \tag{12}$$

where we omitted the trivial $\delta^{c_1c_2}\delta^{c'_3c'_4}$ group factor enforcing identical gauge charges of the annihilating gauginos as well as those of the created gauge bosons. Similarly, for the process with fermion helicities reversed,

$$\mathcal{M}_{+-} \;= -\mathcal{T} M_s^{-3} m\, e^{-\eta}\,\langle 34\rangle^2\,. \tag{13}$$

The two remaining amplitudes, $\mathcal{M}_{++}$ and $\mathcal{M}_{-+}$ are obtained by complex conjugating $\mathcal{M}_{--}$ and $\mathcal{M}_{+-}$, respectively.

At this point we focus on one specific assignment of stacks to boundaries. With a choice of binos (hypercharge gauge bosons) as our LSP, and with the assumption of relatively small mixing with the other $U(1)$ subgroups in stacks a and b, the bino is largely associated with the $U(1)$ stack c. Under the preceding assumption of small mixing, we note that each photon (Z) vertex introduces a factor of approximately $\sin\theta_W$ ($\cos\theta_W$) if inserted at the boundary associated with the $U(2)$ stack b, and $\cos\theta_W$ ($\sin\theta_W$) if inserted at the boundary associated with stack c. To retain the purely topological structure of the amplitude we attach the second boundary to stacks $a,\ b,\ c$ and leave the third boundary empty.

In order to compute the annihilation rate, we need the sum of squared amplitudes, averaged over the helicities and gauge indices of initial gauginos and summed over the helicities and gauge indices of final gauge bosons:

$$\begin{aligned} |\mathcal{M}(\chi^0\chi^0\to WW)|^2 \;&=\; \frac{N_{c'}^2-1}{4N_c^2}\big(\,|\mathcal{M}_{--}|^2+|\mathcal{M}_{+-}|^2+|\mathcal{M}_{++}|^2+|\mathcal{M}_{-+}|^2\big) \\ &=\; \frac{3}{2}\,\mathcal{T}^2\,\frac{s^2(s-2m_{\chi^0}^2)}{M_s^6}\;, \end{aligned} \tag{14}$$

where the Mandelstam variable, $s=(k_1+k_2)^2$, and WW denotes final states including $W^+W^-,\ ZZ,\ \gamma Z$, or $\gamma\gamma$. Near threshold ($s\simeq 4m_{\chi^0}^2$), the total annihilation rate into the three $SU(2)$ gauge vector bosons is

$$\sigma v|_{WW} = \frac{3c}{4\pi}\,\mathcal{T}^2\left(\frac{\hbar}{M_s\;c}\right)^2\rho^4\,, \tag{15}$$

where $\rho\equiv m_{\chi^0}/M_s$. In a similar manner,

$$\sigma v|_{gg} = \frac{8c}{4\pi}\,\mathcal{T}^2\left(\frac{\hbar}{M_s\;c}\right)^2\rho^4\,, \tag{16}$$

and

$$\sigma v|_{BB} = \zeta^2\frac{c}{4\pi}\,\mathcal{T}^2\left(\frac{\hbar}{M_s\;c}\right)^2\rho^4\,. \tag{17}$$

The factor ζ (which in principle can take any real value) parameterizes the uncertainty in the $\chi\chi \to BB$ amplitude because of the aforementioned non-topological components in the matrix element, where all four vertices are attached to the same boundary, or due to scattering in the "genus 3/2" configuration. Dominance of the topological component corresponds to $\zeta \simeq +1$.

We now constrain a combination of the free parameters of the model by requiring that neutralinos have the measured dark matter abundance [4]. The density of neutralinos that survives after freezing out of thermal equlibrium in the early universe is given by

$$\Omega_{\chi^0} h^2 \simeq 0.1 \left(\frac{x_{\rm FO}}{20}\right) \left(\frac{g_\star}{80}\right)^{-1/2} \left(\frac{\langle \sigma v \rangle_{\rm eff}}{3 \times 10^{-26}\,{\rm cm}^3/{\rm s}}\right)^{-1}, \tag{18}$$

where $x_{\rm FO}$ is the neutralino mass divided by the freeze-out temperature, $g_\star$ is the number of external degrees of freedom available at the freeze-out temperature, and $\langle \sigma v \rangle_{\rm eff}$ is the effective neutralino annihilation cross section evaluated at the freeze-out temperature. The desired effective annihilation rate, $\sigma v|_{WW} + \sigma v|_{gg} + \sigma v|_{BB} = \langle \sigma v \rangle_{\rm eff} \simeq 3 \times 10^{-26}$ cm^3/s,[2] required to generate the measured relic density, $\Omega_{\rm CDM} h^2 = 0.113 \pm 0.003$, is obtained if

$$\left(1 + 0.083(\zeta^2 - 1)\right) \left(\frac{\widetilde{F}^{(0,3)}}{2.8}\right)^2 \left(\frac{g_s}{0.2}\right)^6 \left(\frac{\rho}{0.5}\right)^4 \left(\frac{2\ {\rm TeV}}{M_s}\right)^2 \simeq 1\,. \tag{19}$$

A sizable value for $\widetilde{F}^{(0,3)}$ is not implausible. As an example, consider the magnetized brane model whose partition function is given by Eqs. (5.28)-(5.30) of Ref. [8]. Crudely replacing the sum over discrete lattice momenta with integrals, one finds that $\widetilde{F}^{(0,3)}$ is proportional a product of three wrapping numbers of a D9 brane around three 2-tori. This number can in principle be large, thus widening the available m_{χ^0}–M_s parameter space.[3] For example, if $\widetilde{F}^{(0,3)} = 6$, then for $m_{\chi^0} = 2$ TeV, M_s can be probed to 4 TeV.

These results have important implications for ongoing and future gamma ray searches for dark matter. In particular, neutralinos annihilating in the Milky Way halo to final states containing a photon (such as $\gamma\gamma$ or γZ) lead to a very distinctive gamma-ray line which if sufficiently bright could provide a "smoking gun" signature of annihilating dark matter.

If we assume little mixing with $U(1)$'s from stacks a and b, the projection onto any photon (Z) in the W^3W^3 final state entails a mixing angle $\sin\theta_W$ ($\cos\theta_W$).[4] For annihilation into the various channels we find,

$$\sigma v|_{\gamma\gamma} = \tfrac{1}{3}\, \sigma v|_{WW} \left(\sin^2\theta_W \ + \ \zeta \cos^2\theta_W\right)^2 , \tag{20}$$

$$\sigma v|_{ZZ} = \tfrac{1}{3}\, \sigma v|_{WW} \left(\cos^2\theta_W \ + \ \zeta \sin^2\theta_W\right)^2 , \tag{21}$$

$$\sigma v|_{\gamma Z} = \tfrac{1}{3}\, \sigma v|_{WW} \ 2\cos^2\theta_W \sin^2\theta_W \ (1-\zeta)^2 , \tag{22}$$

$$\sigma v|_{W^+W^-} = \tfrac{2}{3}\, \sigma v|_{WW} \,. \tag{23}$$

[2] In addition to neutralino self-annihilation, $\langle \sigma v \rangle_{\rm eff}$ can potentially include the effects of coannihilation between neutralinos and other superparticles of similar mass. We neglect contributions from these processes.

[3] This includes the range of string scales consistent with the correct weak mixing angle found in the $U(3) \times U(2) \times U(1)$ quiver model [11].

[4] We note in passing that in the minimal extension of the standard model this mixing angle is fixed and introduces a multiplicative factor of 0.96 into the right-hand-side of Eqs. (14)–(17) [12].

($\sigma v|\, gg$ is given in Eq. 16 above.) The vanishing of $\sigma v|_{\gamma Z}$ for $\zeta = +1$ is a reflection of the symmetry in Eq. (1) in the topological case, where the coupling is independent of the choice of assignment of the stacks on the boundaries. For $\zeta = +1$, these cross sections yield an 8.3% branching fraction to $\gamma\gamma$. The $\gamma\gamma$ fraction is much larger than is predicted by the existing one-loop broken SUSY calculations [13, 14]. For all parameter space satisfying the measured dark matter abundance [4], the standard annihilation rates to $\gamma\gamma$ or γZ are typically smaller than about 10^{-28} cm^3/s. In contrast, our model predicts $\sigma v|_{\gamma\gamma} \sim 3 \times 10^{-27}$ cm^3/s, which is more than an order of magnitude larger than the standard SUSY result.

For neutralinos with masses above a few hundred GeV, the H.E.S.S. observations of the Galactic Center (GC) [15] can be used to probe the dark matter annihilation cross section. The flux of gamma rays from dark matter annihilation in the GC is given by

$$\Phi_\gamma(E_\gamma, \psi) \simeq \frac{\sigma v|_i}{8\pi} \left. \frac{dN_\gamma}{dE_\gamma}\right|_i \int_{\text{l.o.s.}} n^2_{\chi^0}(r)\; dl(\psi)\; d\psi\,, \tag{24}$$

where i denotes the final state, ψ is the angle observed away from the GC, dN_γ/dE_γ is the spectrum of gamma rays produced per annihilation, and $n_{\chi^0}(r)$ is the number density of dark matter particles as a function of the distance from the GC. The integral is performed over the observed line-of-sight assuming a dark matter distribution which follows the Navarro-Frenk-White (NFW) halo profile [16]. In Fig. 2, the dotted curve is the gamma-ray spectrum corresponding to a 1 TeV neutralino with a total annihilation rate $\sigma v|_{\text{tot}} = 3 \times 10^{-26}$ cm^3/s, that annihilates to $\gamma\gamma$ and γZ with branching fractions of 0.1%, which is typical for a TeV neutralino in the the minimal supersymmetric standard model (MSSM). For significantly larger branching fractions to $\gamma\gamma$ or γZ, the prospects for detection are greatly improved. The solid curve in Fig. 2, corresponding to $\zeta = +1$, is the gamma-ray spectrum for a neutralino that annihilates to $\gamma\gamma$ with an 8.3% branching fraction, and does not annihilate to γZ. Unlike the case of a typical neutralino, a very bright and potentially observable gamma-ray feature is predicted. For example, the suggestive structure at 2.5 TeV in H.E.S.S. data from 2004 [15] can be easily accommodated within this model. If an experiment were to detect a strong gamma-ray line without a corresponding continuum signal from the cascades of other annihilation products, it could indicate the presence of a low string scale.

In summary, within the context of D-brane TeV-scale string compactifications, we constructed a model that generates a supersymmetric R-symmetry violating effective Lagrangian which allows for the s-wave annihilation of neutralinos, once gauginos acquire mass through an unspecified mechanism. The model allows for a neutralino relic abundance consistent with the measured dark matter density. The branching fractions to monochromatic gamma rays is orders of magnitude larger than in the MSSM. A very bright and distinctive gamma-ray line that may lie within the reach of current or next-generation gamma-ray telescopes is predicted. A flux near the limit presently imposed by the H.E.S.S. data would strongly support a near purely topological origin for the R-symmetry violating effective Lagrangian.

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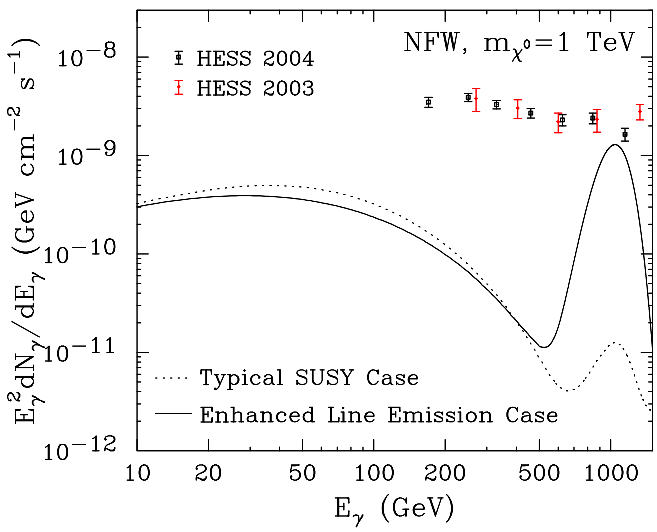

FIG. 2: The gamma-ray spectrum from neutralino dark matter annihilating in the Galactic Center (within a solid angle of 10^{-3} sr), assuming the NFW halo profile. The spectrum has been convolved with a gaussian of width $\Delta E_\gamma/E_\gamma$ =15%, the typical energy resolution of H.E.S.S. and other ground based gamma-ray telescopes. The solid curve corresponds to dark matter annihilation with $\zeta = +1$, for which the $\gamma\gamma$ final state has a branching fraction of 8.3%. The dotted curve corresponds to 0.1% branching fractions to $\gamma\gamma$ and γZ, typical of neutralino annihilation in the MSSM. In both cases, we considered a 1 TeV mass and a total annihilation cross section of 3×10^{-26} cm^3/s. The continuum portion of the spectrum arises from the decay products of the W and Z bosons, and gluons as calculated using Pythia. Also shown for comparison are the H.E.S.S. data [15] which are generally interpreted to be of astrophysical origin [17].

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