Superweakly Interacting Massive Particles

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We investigate a new class of dark matter: superweakly-interacting massive particles (super-WIMPs). As with conventional WIMPs, superWIMPs appear in well-motivated particle theories with naturally the correct relic density. In contrast to WIMPs, however, superWIMPs are impossible to detect in all conventional dark matter searches. We consider the concrete examples of gravitino and graviton cold dark matter in models with supersymmetry and universal extra dimensions, respectively, and show that superWIMP dark matter satisfies stringent constraints from Big Bang nucleosynthesis and the cosmic microwave background.

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There is ample evidence that luminous matter makes up only a small fraction of all matter in the universe. Results from the Wilkinson Microwave Anisotropy Probe, combined with other data, constrain the non-baryonic dark matter density to $\Omega_{DM} = 0.23 \pm 0.04$, far in excess of the luminous matter density. We therefore live in interesting times: while the amount of dark matter is becoming precisely known, its identity remains a mystery.

WIMPs, weakly-interacting massive particles with weak-scale masses, are particularly attractive dark matter candidates. WIMPs have several virtues. First, their appearance in particle physics theories is independently motivated by the problem of electroweak symmetry breaking. Second, given standard cosmological assumptions, their thermal relic abundance is naturally that required for dark matter. Third, the requirement that WIMPs annihilate efficiently enough to give the desired relic density generically implies that WIMP-matter interactions are strong enough for dark matter to be discovered in current or near future experiments.

Here we consider a new class of non-baryonic cold dark matter: superweakly-interacting massive particles (superWIMPs or SWIMPs). As with WIMPs, superWIMPs appear in well-motivated theoretical frameworks, such as supersymmetry and extra dimensions, and their (non-thermal) relic density is also naturally in the desired range. In contrast to conventional WIMPs, however, they interact superweakly and so evade all direct and indirect dark matter detection experiments proposed to date.

For concreteness, we consider two specific super-WIMPs: gravitinos in supersymmetric theories, and Kaluza-Klein (KK) gravitons in theories with extra dimensions. Gravitino and graviton superWIMPs share many features, and we investigate them in parallel.

For gravitino superWIMPs, we consider supergravity, where the gravitino $\tilde{G}$ and all standard model (SM) superpartners have weak-scale masses. Assuming $R$-parity conservation, the lightest supersymmetric particle (LSP) is stable. In supergravity, the LSP is usually assumed to be a SM superpartner. Neutralino LSPs are excellent WIMP candidates, giving the desired thermal relic density for masses of 50 GeV to 2 TeV, depending on Higgsino content. In contrast, here we assume a $\tilde{G}$ LSP. The gravitinos considered here couple gravitationally and form cold dark matter, in contrast to the case in low-scale supersymmetry breaking models where light gravitinos couple more strongly and form warm dark matter.

We consider also the possibility of graviton dark matter in universal extra dimensions (UED), in which gravity and all SM fields propagate. We focus on $D = 5$ spacetime dimensions with coordinates $x^{\tilde{M}} = (x^5, y)$. The fifth dimension is compactified on the orbifold $S^1/Z_2$, where $S^1$ is a circle of radius $R$, and $Z_2$ corresponds to $y \to -y$. Unwanted massless fields are removed by requiring suitable transformations under $y \to -y$. For example, the 5D gauge field $V_M(x,y)$ transforms as $V_\mu(x,y) \to V_\mu(x,y)$ and $V_5(x,y) \to -V_5(x,y)$ under $y \to -y$, which preserves $V_\mu(x)$ and removes $V_5(x)$. Similar choices remove half of the fermionic degrees of freedom, producing chiral 4D fermions, and preserve the 4D graviton $h_{\mu\nu}^0(x)$ while removing $h_{5\mu\nu}^0(x)$ and $h_{55}^0(x)$. The gravi-scalar $h_{55}^0(x)$ remains; we assume that some other physics stabilizes this mode and generates a mass for it.

The orbifold compactification breaks KK number conservation, but preserves KK-parity. KK particles must therefore be produced in pairs, and current bounds require only $R^{-1} \gtrsim 200$ GeV. KK-parity conservation also makes the lightest KK particle (LKP) stable and a dark matter candidate. For $R^{-1} \sim$ TeV, weakly-interacting KK particles have thermal relic densities consistent with observations.

In particular, $B^1$, the first KK partner of the U(1) gauge boson, has been shown to be a viable WIMP dark matter candidate, with promising prospects for direct detection and also indirect detection in anti-matter searches, neutrino telescopes, and gamma ray detectors.

As in the case of supersymmetry, however, the lightest partner need not be a SM partner. In UED, the LKP could be $G^1$, the first KK partner of the graviton. $G^1$ is, in fact, perhaps the most natural LKP candidate — radiative contributions to KK masses, typically positive, are negligible for $G^1$. $G^1$ couplings are also gravitational,
and so highly suppressed.

Gravitinos and gravitons therefore naturally emerge as superWIMP candidates: stable massive particles with superweak interactions. Their weak gravitational interactions imply that they play no role in the thermal history of the early universe. (We assume inflation followed by reheating to a temperature low enough to avoid regenerating large numbers of superWIMPs.) Thus, if the next lightest supersymmetric particle (NLSP) or next lightest KK particle (NLKP) is weakly-interacting, it freezes out with a relic density of the desired magnitude. Much later, however, these WIMPs then decay to superWIMPs; as the WIMP and superWIMP masses are similar, the superWIMP then inherits the desired relic density.

Unlike WIMPs, however, superWIMPs are impossible to discover directly, and their annihilation rate is so suppressed that they also escape all indirect detection experiments. At the same time, unlike superheavy dark matter candidates with only gravitational interactions [11], superWIMPs inherit the desired relic density from a thermal abundance and arise from accessible electroweak physics. At colliders, WIMP decays to superWIMPs will occur long after the WIMP leaves the detector. If the NLSP or NLKP is neutral, its metastability will have no observable consequences. The discovery of a seemingly stable but charged NLSP or NLKP may, however, provide a strong hint for superWIMP dark matter.

We now investigate constraints on and alternative signals of superWIMP dark matter scenarios. The observable consequences of superWIMPs must rely on the decays of WIMPs to superWIMPs on cosmological time scales [12]. The NLSP or NLKP may be any SM partner. In supergravity, the lightest SM superpartner is often the Bino, the superpartner of the hypercharge gauge boson. In the minimal UED scenario [22], the lightest SM KK mode is often $B^1$. Motivated by these results, we now consider specific scenarios in which decays to superWIMPs are typically accompanied by photons, and we consider the impact of electromagnetic cascades.

In the supersymmetric photon superWIMP scenario, NLSP decay is governed by the coupling $-\sqrt{2} \tilde{G}_\mu \left[ \gamma^\mu, \gamma^0 \right] \gamma^\rho \tilde{F}_{\rho \nu} F_{\nu \rho}$, where $F$ is the U(1) field strength, and $M_* = (8\pi G_N)^{-1/2} \approx 2.4 \times 10^{18}$ GeV is the reduced Planck scale. The NLSP decay width is

$$\Gamma(B \to \tilde{G} \gamma) = \frac{\cos^2 \theta_W}{48\pi M_*^2} \frac{m_B^5}{m_{\tilde{G}}^3} \left[ 1 - \frac{m_{\tilde{G}}^2}{m_B^2} \right]^3 \left[ 1 + 3\frac{m_{\tilde{G}}^2}{m_B^2} \right]. \tag{1}$$

In models with low-scale supersymmetry breaking, $m_{\tilde{G}} \ll m_B$, and the gravitino couples dominantly through its $\pm \frac{1}{2}$ spin components. In the high-scale supersymmetry breaking scenarios considered here, however, the couplings of the $\pm \frac{1}{2}$ spin polarizations are of the same order and must be kept in deriving Eq. (1).

The properties of gravitinos in UED scenarios may be determined straightforwardly; details will be presented elsewhere [12]. Graviton superWIMPs couple to $B^1$ through $\sqrt{2} M_* G^\mu_\nu \left( -F^0_{\mu \rho} F^{1 \nu} F^0_{\rho \sigma} \right) + \frac{1}{2} \eta^{\mu \nu} F^0_{\rho \sigma} F^{1 \rho \sigma}$, where $F^0_{\mu \rho} \equiv \partial_\mu B^0_\rho - \partial_\rho B^0_\mu$. The $G^1 B^1 B^0$ vertex is identical to the $G^0 B^0 B^0$ vertex, and the longitudinal component of the massive $B^1$ plays no role. The NLKP decay width is

$$\Gamma(B^1 \to G^1 \gamma) = \frac{\cos^2 \theta_W}{72\pi M_*^2} \frac{m_B^3}{m_{\tilde{G}}^3} \left[ 1 - \frac{m_{\tilde{G}}^2}{m_B^2} \right]^3 \times \left[ 1 + \frac{3 m_{\tilde{G}}^2}{m_B^2} \right] \left[ 1 + \frac{m_{\tilde{G}}^4}{m_B^4} \right]. \tag{2}$$

The $\tilde{B}$ and $B^1$ lifetimes are given in Fig. 1. In both cases, in the limit $\Delta m \equiv m_{WIMP} - m_{SWIMP} \ll m_{SWIMP}$, the WIMP lifetime is proportional to $(\Delta m)^{-3}$ and is independent of the overall mass scale.

As we will see, the most relevant bounds constrain the total energy released in photons in WIMP decay, or more precisely, $\varepsilon \gamma Y_\gamma$, where $\varepsilon \gamma$ is the energy of the photons when created and $Y_\gamma = n_\gamma / n_{BG}$ is the number density of photons from WIMP decay normalized to the number density of background photons $n_{BG} = 2\varepsilon(3T^3/\pi^2)$, where $T$ is the temperature during WIMP decay. In the superWIMP scenario, WIMPs decay essentially at rest, and so $\varepsilon \gamma = (m_{WIMP}^2 - m_{SWIMP}^2)/(2m_{WIMP})$. Since a superWIMP is produced in association with each photon, $Y_\gamma = Y_{SWIMP}$, and the photon abundance is given by

$$Y_{SWIMP} \approx 3.0 \times 10^{-12} \left( \frac{\text{TeV}}{m_{SWIMP}} \right) \left( \frac{\Omega_{SWIMP}}{0.23} \right). \tag{3}$$

Predicted values for $\varepsilon \gamma Y_{SWIMP}$ are shown in Fig. 2.

As evident in Fig. 1, WIMP decays occur long after Big Bang nucleosynthesis (BBN), and so may, in principle, destroy the successful BBN predictions for light element abundances. The energy of photons produced in late decays is rapidly redistributed through scattering off background photons, $\gamma + \gamma \rightarrow e^+ e^-$, and inverse Compton scattering [13, 14]. As a result, the constraints of

![FIG. 1: Lifetimes for $\tilde{B} \to \tilde{G} \gamma$ (left) and $B^1 \to G^1 \gamma$ (right) for $\Delta m \equiv m_{WIMP} - m_{SWIMP}$ and $m_{SWIMP} = 0.1 \text{ TeV}$ (long dashed), 0.3 TeV (short dashed), and 1 TeV (solid).]
BBN are, to an excellent approximation, independent of the initial energy distribution of injected photons, and constrain only the total energy release.

Detailed analysis 12, demanding consistent predictions for deuterium, 3He, 4He, 6Li, and 7Li, excludes the region of parameter space shown in Fig. 2. The BBN constraint is weak for early decays: at early times, the universe is hot and the initial photon spectrum is rapidly softened, leaving few high energy photons above threshold to modify the light element abundances. We find that BBN excludes some of the relevant parameter region, but not all. In particular, for relatively short-lived WIMPs with \( \tau \lesssim 10^7 \) s and weak-scale masses, the requirement that superWIMPs form all of the dark matter is consistent with the constraints from BBN.

The cosmic microwave background (CMB) also imposes constraints 13,14. The injection of energy in the form of photons may distort the CMB from the observed black-body spectrum. Before redshifts of \( z \sim 10^7 \), elastic Compton scattering, bremsstrahlung \( eX \to eX \gamma \) (with \( X \) an ion), and double Compton scattering \( e^- \gamma \to e^- \gamma \gamma \) effectively thermalize injected energy. After \( z \sim 10^5 \), however, the photon number-changing interactions become ineffective, and the photon spectrum relaxes only to a Bose-Einstein distribution with chemical potential \( \mu \). After \( z \sim 10^5 \), even Compton scattering becomes ineffective, and deviations from the black-body spectrum may be parameterized by the Sunyaev-Zeldovich \( \delta \) parameter.

As with the BBN constraints, bounds from the CMB are largely independent of the injected energy spectrum, depending primarily on the total energy release. The bounds on \( \varepsilon \gamma Y_{\text{SWIMP}} \) scale linearly with the bounds on \( \mu \) and \( y \). We update the analysis of Ref. 15 to include the latest results \( |\mu| < 9 \times 10^{-3} \) and \( |y| < 1.2 \times 10^{-5} \) 17, with baryon density \( \Omega_B h^2 \approx 0.022 \) 1, where \( h \approx 0.71 \) is the normalized Hubble expansion rate. These bounds exclude energy releases above the CMB contours in Fig. 2. Remarkably, the CMB constraints are now so precise that they supersede BBN constraints for decay times \( \tau \sim 10^7 \) s and \( \tau \gtrsim 10^{10} \) s. Nevertheless, regions of superWIMP parameter space, including regions with weak-scale masses and mass splittings, remain viable.

Finally, for highly degenerate WIMP-SWIMP pairs, WIMPs decay very late to soft photons. The photon spectrum is not thermalized and may produce observable peaks in the diffuse photon spectrum. The present differential flux of photons from WIMP decay is

\[
\frac{d\Phi}{dE_\gamma} = \frac{c}{4\pi} \int_0^{t_0} \frac{dt}{\tau_{\text{WIMP}}} \frac{N_{\text{WIMP}}(t)}{V_0} \delta (E_\gamma - \varepsilon \gamma) ,
\]

where \( t_0 \approx 13.7 \) Gyr is the age of the universe 11, \( N_{\text{WIMP}}(t) = N_{\text{in}} \text{WIMP} e^{-t/\tau_{\text{WIMP}}} \), where \( N_{\text{in}} \text{WIMP} \) is the number of WIMPs at freeze-out, and \( V_0 \) is the present volume of the universe. The diffuse photon flux is a sensitive probe only when WIMPs decay in the matter-dominated era. We may then take \( 1 + z = (t_0/t)^{2/3} \), and

\[
\frac{d\Phi}{dE_\gamma} \approx \frac{3cN_{\text{in}}}{8\pi V_0 \varepsilon \gamma} \left( \frac{t_0}{\tau_{\text{WIMP}}} \right)^{\frac{2}{3}} F(a) \Theta (\varepsilon \gamma - E_\gamma) ,
\]

where \( F(a) = a^{1/2} e^{-a^{3/2}} \), \( a = (E_\gamma / \varepsilon \gamma)(t_0/\tau_{\text{WIMP}})^{2/3} \). \( F(a) \) is maximal at \( a = 3^{-2/3} \), and so, for \( \tau_{\text{WIMP}} < t_0 \), the differential photon flux reaches its maximal value at

\[
E_{\gamma}^{\text{max}} = \varepsilon \gamma \left( \frac{\tau_{\text{WIMP}}}{3t_0} \right)^{\frac{5}{6}} = 680 \text{ keV} \left( \frac{\text{GeV}}{\Delta m} \right)
\]

for gravitinos, and an energy 1.4 times smaller for gravitons. The \( \Delta m \) dependence follows from the redshifting of photons created with energy \( \Delta m \) by \( 1 + z \propto \tau_{\text{WIMP}}^{2/3} \propto (\Delta m)^2 \). For high degeneracies,

\[
\frac{N_{\text{in}}}{V_0} \text{WIMP} = 1.2 \times 10^{-9} \text{ cm}^{-3} \left( \frac{\text{TeV}}{m_{\text{WIMP}}} \right) \left( \frac{\Omega_{\text{WIMP}}}{0.23} \right) ,
\]

and so the maximal flux is

\[
\frac{d\Phi}{dE_\gamma} (E_{\gamma}^{\text{max}}) = 1.5 \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1} \text{ MeV}^{-1}
\]

\[
\times \left( \frac{\text{TeV}}{m_{\text{WIMP}}} \right) \left( \frac{\Delta m}{\text{GeV}} \right) \left( \frac{\Omega_{\text{WIMP}}}{0.23} \right)
\]

for gravitinos, and a factor of 1.4 larger for gravitons.

Representative photon energy spectra are shown in Fig. 3. Also shown are the measured diffuse fluxes from the observatories HEAO, OSSE, and COMPTEL 15, determined from observed fluxes by subtracting known point sources. The observed spectra fall rapidly with energy and so severely constrain the relatively hard photon spectra predicted by \( \Delta m \lesssim 10 \) GeV. Note that we have not included photon interactions which soften the photon
FIG. 3: Diffuse photon fluxes (solid) for $m_{\text{SWIMP}} = 1$ TeV, $\Omega_{\text{SWIMP}} = 0.23$, and $\Delta m = 1$ GeV (solid) and 10 GeV (long dashed), and upper bounds from observations (short dashed).

FIG. 4: Regions of the $(m_{\text{SWIMP}}, \Delta m)$ plane excluded by BBN, CMB, and diffuse photon constraints. The shaded regions and the regions below the CMB contours are excluded.

spectrum for $\Delta m \gtrsim$ few GeV \cite{19}; such effects can only enlarge the allowed parameter space discussed below.

In Fig. 3 we compile the constraints discussed above and show the allowed regions of the $(m_{\text{SWIMP}}, \Delta m)$ plane for the photon gravitino and graviton superWIMP dark matter scenarios. The BBN and CMB constraints are as discussed above. An additional region is excluded by the requirement that the diffuse photon flux never exceed the observed flux by $2\sigma$ for any energy. Although these data exclude some of the parameter space, the most well-motivated region with $m_{\text{SWIMP}}$, $\Delta m \sim 100$ GeV to 1 TeV remains an outstanding possibility.

In conclusion, we find that superWIMPs provide a qualitatively novel possibility for particle dark matter. Such particles appear in the form of gravitinos and gravitons in theories with supersymmetry and extra dimensions, and they naturally inherit the desired relic density from late-decaying weakly-interacting NLSPs or NLKPs.

SuperWIMPs satisfy existing constraints from BBN and CMB and evade all conventional dark matter experiments. On the surface, we have apparently committed Pauli’s “ultimate sin” by proposing a solution to the dark matter problem that has no observable consequences. However, improvements in the measurements discussed above may uncover anomalies. For example, a detailed study of the 15 keV to 10 MeV diffuse photon flux by INTEGRAL is currently underway \cite{20}. A pronounced bump in this spectrum could provide a striking signal of superWIMP dark matter. Finally, we note that neutrino, charged lepton, weak gauge boson, and Higgs boson NLSPs and NLKPs are all viable from the point of view of preserving the naturalness of the desired relic density. Some of these scenarios may be severely constrained by bounds on hadronic showers, but the remaining scenarios will have qualitatively different, and possibly very interesting, observational consequences.

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