1. Introduction

The typical theoretical prediction for an extragalactic neutrino flux at $E_\nu \sim 10^{16} - 10^{18}$ eV is the generic upper bound on the diffuse neutrino spectrum from AGN jets obtained by Mannheim et al.\cite{Mannheim} In their model AGN source is not completely thin for a CR flux, and the normalization of the resulting CR spectrum is such that the CR intensity does not exceed the proton spectrum obtained from observations. At $E_\nu \sim 10^{16}$ eV this bound is

$$j(E)E^2 < 3 \times 10^{-7} \text{ GeV} \cdot \text{cm}^{-2} \cdot \text{sr}^{-1} \cdot \text{s}^{-1}$$

($j(E)$ is the differential diffuse neutrino spectrum). More conservative bound\cite{Mannheim} corresponds to the case when the sources are completely transparent for CRs: $j(E)E^2 < 2 \times 10^{-8}$ (independently on the neutrino energy if $j(E) \propto E^{-2}$).

At higher energies, $E_\nu \sim 10^{20}$ eV, extragalactic neutrino flux can have several components: the guaranteed component is the “cosmogenic” neutrino flux resulting from interactions of CRs with relic photons of CMB (cosmic microwave background). According to the most optimistic predictions\cite{Mannheim}, the value $j(E)E^2$ for the cosmogenic flux may be as large as

$$3 \times 10^{-7} \text{ GeV} \cdot \text{cm}^{-2} \cdot \text{sr}^{-1} \cdot \text{s}^{-1}$$

at $E_\nu \sim 10^{20}$ eV. Other components of neutrino flux at ultrahigh energies are hypothetical: e.g., neutrino from decays of topological defects, neutrino from interactions of new (exotic) hadrons with mass $\sim (2 - 5)$ GeV in CR flux\cite{Mannheim}, with relic photons etc.

Experimental bounds on the extragalactic diffuse neutrino flux are on the level

$$j(E)E^2 \approx 10^{-6} \text{ GeV} \cdot \text{cm}^{-2} \cdot \text{sr}^{-1} \cdot \text{s}^{-1}$$

obtained, at $E \approx 10^{15}$ eV, in experiments on neutrino telescopes, and, at higher energies, on large air shower arrays (see, e.g., review\cite{Mannheim}).
2. Neutrino-nucleon interactions

2.1. $\nu N$ DIS in Standard Model

Generically, the differential cross section for the process $\nu_l N \rightarrow lX$ at large neutrino energies strongly depends on the behavior of parton distribution functions (pdfs) at small values of Bjorken variable $x$ and large values of $Q^2$. Besides, the gauge boson propagator effects are sufficient (leading to a nonlinear rise of $\sigma_{\nu N}(E_{\nu})$) when $Q$ becomes comparable to the electroweak scale, $Q^2 \gg M_W^2$. In the energy region which we are interested in, $E_{\nu} > 10^7$ GeV, one has

$$Q^2_{\text{char}} \approx M_W^2, \quad x_{\text{char}} < 10^{-4} - 10^{-5}.$$ 

In QCD-improved parton model perturbative QCD corrections to the nonperturbative pdfs (taken from experiment) are calculated using two approaches: DGLAP scheme (resummations of the perturbative expansion retaining leading terms in $\ln(Q^2/\Lambda^2_{\text{QCD}})$ and BFKL scheme (resummations retaining leading terms in $\ln(1/x)$). Both these logarithms become large in the relevant kinematic region. Since, very roughly, $x$-dependence of pdfs of sea quarks has a power-law behavior, $\propto x^{-0.3}$, for $x \ll 1$, it follows that

$$\sigma_{\nu N}(E_{\nu}) \propto E_{\nu}^{0.3}.$$ 

The power-law rise of the cross section implies a violation of unitarity at very high energies.

It is rather evident that there must be non-linear (“higher twist”) QCD corrections to the standard calculation, those which are beyond the linear evolution DGLAP and BFKL schemes. Physically, these non-linear corrections arise due to the growth of parton (gluon) density in the nucleon target at small $x$. Taking into account the recombination gluon effects leads to a taming of the fast rise of the cross section, in agreement with unitarity requirements.

There are different approaches accounting these gluon screening effects. To feel the order of magnitude of corresponding corrections to $\sigma_{\nu N}(E_{\nu})$ it is enough to compare the results of two models: unified BFKL/DGLAP + screening model\cite{6} and phenomenological colour-dipole model\cite{7}. In the latter model, the DIS is considered in the laboratory system: virtual weak gauge boson at some distance from the target fluctuates into $qq$-pair interacting with the target via multiple gluon exchanges. This approach, in its original form, does not take into account QCD evolution at all. It appears that at extremely high neutrino energies, $E_{\nu} \approx 10^{12}$ GeV, the difference in the predictions of these models for the total cross section does not exceed factor $2 - 3$ (see Ref.\cite{6}).

The $\sigma_{\nu N}(E_{\nu})$ scales, approximately, as $E_{\nu}^{0.36}\text{eV}$ for $E_{\nu} > 10^{16}$ eV. Using this one can easily obtain the estimate for the rate of contained events in neutrino telescopes\cite{8}. According to this estimate, the value $j(E)E^2$ should be larger than

$$10^{-7} \text{GeV} \cdot \text{cm}^{-2} \cdot \text{sr}^{-1} \cdot \text{s}^{-1}$$

to be measurable at $E_{\nu} > 10^{19}$ eV at 1 km$^3$ detector, if the SM prediction for the total cross section is correct.
2.2. TeV scale gravity models for $\sigma_{\nu N}$

It is easy to show that the gravitational scattering of any two particles (via graviton exchange) becomes strong if $\sqrt{s} \gg M_{Pl} \approx 10^{19}$ GeV, i.e., at inaccessibly high energies of colliding particles.

In TeV-scale quantum gravity models proposed by Arkani-Hamed et al.\textsuperscript{9} there are two most important independent parameters: $M_D$ (fundamental gravitational, or Planck, scale) and $n$ – the number of extra (compact) dimensions. The value $M_D$ is connected with the 4-dimensional Newton constant through the compactification radius $R$. If the compactification volume, $V = 2\pi R$, is large, $M_D$, which is equal to $\left(\frac{M_{Pl}^2}{V^n}\right)^{\frac{1}{n+2}}$, is much smaller than $M_{Pl}$. In TeV-scale gravity models one has $M_D \approx 1$ TeV.

In these models, because of the compactification, the extra $n$ components of the graviton momentum are quantized. To an observer in the usual 4-dimensional space-time, the graviton would appear to be a massive Kaluza-Klein (KK) state. The result of the summation over KK states in the expression for the amplitude of gravitational scattering is

$$A \approx \frac{s^2}{M_D^4} \left(\frac{\Lambda}{M_D}\right)^{n-2},$$

where $\Lambda (\approx M_D$) is the cut-off constant (the Born amplitude is divergent but after eikonalization the amplitude becomes finite). We see from this expression that the amplitude is large if $\sqrt{s} \geq M_D \approx 1$ TeV, rather than $\sqrt{s} > M_{Pl}$, as in the 4-dimensional theory. Due to $s^2$-dependence, the gravitational scattering amplitude becomes, at some energy, larger than any other amplitude.

The gravitational amplitude and cross section can be calculated reliably if the process is semiclassical, i.e., when the Schwarzschild radius ($R_S$) of the colliding system (which is the classic value) is much larger than quantum lengths, $\lambda_p$ (it is higher dimensional Planck length) and $\lambda_B$ (de Broglie wave length of the colliding particles). It is easy to see, that for the semiclassical picture of the gravitational scattering the inequality $\sqrt{s} \gg M_D$ must take place (the “trans-Planckian energy regime”). In this regime gravitational interaction dominates over other gauge interactions and is a semiclassical process.

The important characteristic of the gravitational scattering is the impact parameter $b$. If $b \gg R_S$ one has $-\frac{1}{2} \ll 1$ ( $t$ is the square of the momentum transfer) and it corresponds to a small scattering angle limit (and elastic scattering). If, in opposite case, $b < R_S$, the collision leads to a production of a black hole with Schwarzschild radius $R_S$ and mass $M_{BH}$ equal, roughly, to cms energy $\sqrt{s}$.

The condition $\sqrt{s} \gg M_D$ corresponds to the existence of the minimum value of BH mass which can be reliably produced in the collision:

$$M_{BH}^{min} = \sqrt{s_{min}} = \alpha M_D, \quad \alpha \gg 1,$$

$\alpha$ is the model parameter. The cross section of the BH production in $\nu N$ collision is equal to

$$\sigma_{\nu N} = \sum_i \int_0^1 dx f_i(x, \mu) \sigma_{\nu i}^{BH}(xs),$$
where \( f_i(x, \mu) \) is pdf for the parton of type \( i \), \( \mu \approx R^{-1} \), and \( \sigma_{\nu i}^{BH} \) is the cross section of BH production in \( \nu \)-parton collision at cms energy \( \sqrt{xs} \),

\[
\sigma_{\nu i}^{BH}(xs) \approx \pi R_s^2(M_{BH} = \sqrt{xs})\Theta(\sqrt{xs} - \alpha M_D).
\]

The differential elastic gravitational scattering of neutrino on nucleon, \( \frac{d\sigma}{dy} \), is calculated (using multiple graviton exchange) as a function of \( y = \frac{E_\nu - E_\nu'}{E_\nu} \). This cross section grows as \( y \) decreases. The small \( y \) region corresponds to long distance processes where neutrino interacts with a parton and transfers only a small portion of its energy, surviving after interaction.

### 2.3. Neutrino interactions in string theory

If cms energies of collisions are close to the fundamental scale of gravity, i.e., if \( \sqrt{s} \approx M_D \) ("Planckian region"), the classical description cannot be trusted. String theory provides the best hope for understanding the regime of strong quantum gravity, and for computing amplitudes at energies close to \( M_D \). So, the models with extra dimensions must be embedded in realistic string models, in which the unification of gravity with the SM takes place.

In string theory the massless graviton is the zero mode of a closed string, whereas the gauge bosons of SM are the lightest modes of an open string. The scattering amplitudes in string theory are amplitudes of string exchanges and are described by the formulas of Veneziano’s (exchange of an open string) or Virasoro’s (exchange of a closed string) type. Correspondingly, one expects the presence of string Regge (SR) excitations of the graviton and gauge bosons (analogously to resonances in Veneziano model). By duality arguments, if there are amplitudes with SR excitations exchanges in the \( t \)-channel, having corresponding \( t \)-channel poles, these amplitudes will also have \( s \)-channel poles.

In lepton-quark scattering the \( s \)-channel resonances are leptoquarks. From the resonance amplitude

\[
\nu + q \rightarrow X_n^J \rightarrow \nu + q
\]

(\( J \) is a spin of the resonance) one can obtain the cross section \( \Pi \sigma_n^J(\nu q) \equiv \sigma (\nu q \rightarrow X_n^J) \), and in the narrow width approximation one has

\[
\sigma_n^J(\nu q) \sim g^2 \delta (s - nM_S^2),
\]

where \( g \) is the gauge coupling constant and \( M_S \) is the string scale. One can assume that \( M_S \sim M_D \sim 1 \) TeV, and, in this case, the SR resonances are "TeV strings". So, we have, in fact, the GUT unification at TeV scale.

### 2.4. Search for a new physics

Neutrino experiments with neutrino of natural origin can be effectively used for a search of a physics beyond the standard model. There are different possibilities.
High Energy Neutrinos as a Probe for New Physics and Astrophysics

1. **Study of contained events in large neutrino telescopes.** At large enough energies of neutrino, say, $E_\nu > 500 - 1000$ TeV, most part of such events may be due to gravitational elastic scattering of neutrino on the nucleon target or BH production by neutrinos. One can study the angular dependence, up-down ratio as well as the energetic distribution.\(^1\)

2. **Study of through-going muons with $E > E_0$, produced in BH production processes in large neutrino telescopes.**\(^2\)

3. **Study of quasi-horizontal air showers which occur at rates exceeding the predictions of SM and have distinct characteristics** (see, e.g., Ref. \(^3\)). Although greatly reduced by hypothetical BH production process (the cross section of which can be, in TeV gravity models, about $10^4$ nb at $10^{20}$ eV), neutrino interaction lengths, $L = 1.7 \times 10^7$ km w.e. ($\sigma/\sigma$), are still far larger than the Earth’s atmospheric depth, which is $L = 0.36$ km w.e., when traversed horizontally. Neutrinos therefore produce black holes uniformly at all atmospheric depths. As a result, the most promising signal of BH production by CRs is quasi-horizontal showers initiated by neutrinos deep in atmosphere. Distinct characteristics of the showers are following: anomalous electromagnetic component;\(^4\) large multiplicity; large $\mu/\tau$ ratio; they look nucleus-like (for not a small $X_{\text{max}}$); they have a curved front, with particles well spread in time.

4. **Earth-skimming idea.** The large detectors of air showers can register the fluxes of UHE neutrinos and, simultaneously, will be able to measure $\sigma_{\nu N}$ at energies as high as $10^{20}$ eV. Several proposed experiments plan to detect UHE neutrinos by observation of nearly horizontal air showers (HAS) resulting in atmosphere from $\nu$-air interactions. At $E \approx 10^{20}$ eV the $\sigma_{\nu N}$ in SM is about $10^{-31}$ cm$^2$. The air shower probability of HAS production is proportional to $\sigma_{\nu N}$. In addition to HAS, the experiments can also observe up-going air showers (UAS) initiated by muon and tau leptons produced by neutrinos interacting just below the surface of the Earth. The expected rate of UAS depend non-linearly on the cross section (if $\sigma_{\nu N}$ larger than $10^{-32}$ cm$^2$, UAS probability is inversely proportional to $\sigma_{\nu N}$). Thus, by comparing the HAS and UAS rates, the cross section can be determined. Taken together, these rates ensure a total event rate, weakly depending of the value of $\sigma_{\nu N}$. If there is a new physics (e.g., BH production), this idea is also useful. Really, a large rate of quasi-horizontal showers may be attributed to either an enhanced neutrino flux or an enhanced BH cross section. While an enhanced flux increases these rates, a large BH cross section will suppress them, since the hadronic decay products of BH evaporation will not escape the Earth’s crust.\(^4\)

3. Neutrinos and UHECR puzzle

3.1. **Z-burst model**

It had been proposed\(^1\) that the primaries which propagate across distances above the GZK zone ($\approx 50$ Mpc) are neutrinos, which then annihilate with relic neutrinos within this zone to create a flux of nucleons and photons with energies above $E_{\text{GZK}}$. The annihilation cross section is relatively large ($\approx 10^{-5}$ mb) near $Z$-resonance and the resonance neutrino energy depends on the neutrino mass.
There are three main difficulties of the Z-burst model.

1. Primary protons have to be accelerated to extremely high energies, \( E > 10^{23} \text{ eV} \), in order to produce in astrophysical sources (via \( pp \) and \( p\gamma \) reactions) UHE neutrinos. The photons produced in the same reactions have to be absorbed inside the source otherwise the diffuse background of MeV–GeV photons will be too large. Hidden sources are unable to produce neutrinos of such high energies because the decay length of pions becomes equal to its scattering length. The luminosities of sources required in this model are too high, \( \approx (10^{45} - 10^{47}) \text{ erg/s} \).

2. Even if one assumes that neutrinos originate from decays of superheavy dark matter particles through the only channel, \( X - \nu \bar{\nu} \), the problems with diffuse photon background are not avoided because higher order corrections to this process give rise to electroweak cascades transferring large fraction of energy to photons and electrons.

3. The required UHE neutrino fluxes in the Z-burst model are very large: they are almost excluded by the new experimental limits from FORTE and GLUE experiments (as well as by the new limit from EGRET).

It was shown in the work of Gelmini et al., using the AGASA data, that the nonobservation of CR events at \( E^{\text{CR}} > 2 \times 10^{20} \text{ eV} \) implies a lower bound \( \approx 0.3 \text{ eV} \) on the neutrino mass. Since this value exceeds \( \sqrt{\Delta m_{ij}^2} \) from neutrino oscillation experiments, the bound applies to all three neutrino masses. If there is neutrino mass hierarchy then it follows from SK data that the mass of the heavier neutrino is \( \approx 0.04 \text{ eV} \). It is argued by Gelmini et al., that AGASA data are incompatible with such a low value because in this case the resonance energy is too high and this predicts too many CRs beyond the AGASA end point, \( 2 \times 10^{20} \text{ eV} \). So, the bound leaves only a small interval for neutrino mass around \( \approx 0.3 \text{ eV} \) (if Z-burst model is correct).

The flux requirement obtained in the work of Gelmini et al., is about \( 7 \times 10^{-36} \text{ (eV \cdot m^2 \cdot sr \cdot s)^{-1}} \), and such a high value is marginally excluded by GLUE and FORTE experiments. The relic neutrino density could be enhanced by either gravitational clustering or by a large lepton asymmetry. Calculations argue against of such a clustering of relic neutrinos if neutrino mass is so small (\( m_\nu \approx 0.3 \text{ eV} \)): the overdensity \( \delta \) of neutrinos in our Local Group of galaxies is \( < 10 \), on a length scales \( \geq 1 \text{Mpc} \) (see Ref. [21]). Also, BBN physics gives the bound on a lepton asymmetry, \( |\xi_\nu| < 0.1 \), from which it follows that the extra contribution to \( \rho_\nu \) from degeneracy is very small.

3.2. Neutrinos as CR primaries

It was noted above that within the SM the neutrino-nucleon interaction cross section at \( 10^{20} \text{ eV} \) is about \( 10^{-4} \text{ mb} \), i.e., on five orders of magnitude smaller than necessary to produce air showers starting high in the atmosphere. It appears also that models with extra dimensions predict cross sections which on a factor \( \approx 10^2 \) larger than that in SM. It is clearly not enough for explanation of the GZK puzzle.

It was argued recently that the cross section at \( E > 10^{18} \text{ eV} \) can be greatly (on a factor \( 10^5 - 10^6 \)) enhanced by nonperturbative electroweak instanton contributions. The pro-
cesses induced by electroweak instantons (which represent tunneling transitions between topologically inequivalent vacua) violate baryon + lepton (B+L) number and are characterized by the large multiplicity of final state particles (quarks, gauge and Higgs bosons). The transition rate of the tunneling process is exponentially suppressed at low energies when
\[ E_{CMS} \ll 4E_{sph} \approx 4\pi \frac{M_W}{\alpha_W} \approx 30 \text{ TeV}. \]

Here, \( E_{sph} \approx 8 \text{ TeV} \) is the sphaleron energy. At high temperatures and high energies the transition rate can be unsuppressed, leading to observable effects in CR experiments. According to Ref. [23] at \( E_{\text{lab}} \approx 10^{20} \text{ eV} \) the cross section is about 3 mb.

According to the calculation of Bezrukov et al. [24] based on a generalized semi-classical approach, a severe exponential suppression of the cross section is extended up to energy \( 30E_{sph} \approx 250 \text{ TeV} \). Bezrukov et al. did not calculate the cross section, but estimated only the upper limit of the exponential function. Therefore one cannot obtain the reliable cross section from their paper although it is tempting to do this. [25]

4. Neutrino from dark matter

The existence of dark matter (DM) is the experimental evidence for new physics beyond the SM. The most recent WMAP data [26] give the amount of cold dark matter (CDM) as
\[ 0.095 < \Omega_{CDM} h^2 < 0.129, \quad (2\sigma \text{ C.L.}) \]
and the total matter density
\[ 0.126 < \Omega_m h^2 < 0.143. \]

There are many particle dark matter candidates for a stable WIMPs (weakly interacting massive particles) which could have the relevant relic density.

Indirect DM searches have been proposed [27] to observe the products of DM annihilation including neutrinos. WIMPs, which scatter elastically in the Sun or Earth, may become gravitationally bound and, over the age of the solar system, they may accumulate in these objects, greatly enhancing their annihilation rate. Neutrinos can escape the Sun or Earth and can be registered in a detector.

4.1. SUSY DM

All superpartners are charged under a discrete symmetry called \( R \)-parity (+1 for SM particles and −1 for their superpartners). \( R \)-parity conservation guarantees that the lightest superpartner, being odd under \( R \)-parity, is absolutely stable and becomes a DM candidate.

The desired WIMP is the lightest neutralino \( \tilde{\chi}_1^0 \) which is a mixture of the superpartners of the hypercharge gauge boson (\( \tilde{b}^0 \)), the neutral \( SU(2)_W \) gauge boson (\( \tilde{\omega}^0 \)) and the two neutral Higgs bosons,
\[ \tilde{\chi}_1^0 = a_1 \tilde{b}^0 + a_2 \tilde{\omega}^0 + a_3 \tilde{h}_u^0 + a_4 \tilde{h}_d^0. \]
The recent progress in particle physics and cosmology favors the region of the SUSY theory parameters where the neutralino is a Bino-Higgsino mixture rather than Bino-like.

The rate of neutrino production in WIMP annihilations is highly model dependent. Neutralinos are Majorana fermions, and the Pauli exclusion principle suppresses the direct production in the annihilations of pairs of light fermions, so only indirect channels of neutrino production are available, through the decay of particles produced in the processes \( \chi \chi \rightarrow \tau^+ \tau^-, b\bar{b}, W^+ W^-, ZZ, HH, t\bar{t}. \)

If neutralino is Bino-like, the production of \( SU(2)_W \) gauge bosons is suppressed. If neutralino exists and is heavy enough and if it is really a mixture Bino-Higgsino, the neutrino signal from the Sun direction can be detected by large neutrino telescopes (see, e.g., Ref.\(^{28}\)). The limits on neutralino mass from experiments and from calculations of relic number and mass density are very weak: \( 20 \text{ GeV} \lesssim m_\chi \leq 600 \text{ GeV} \).

4.2. Kaluza-Klein DM

There is one class of models with extra dimensions, in which all of the fields of the SM propagate in these dimensions, not only gravitons. Each SM field has an infinite tower of KK partners with identical spins and couplings and masses of order \( n/R \). Naturally, in such models the conservation of the momentum along the extra dimensions and, as a consequence, the conservation of KK number, takes place. Radiative corrections break KK number down to a discrete conserved quantity, KK-parity: \( (-1)^n \), where \( n \) is the number of the KK level. KK-parity ensures that the lightest KK-partner (LKP) at level one, being odd under KK-parity, is stable, and can be a DM candidate.

In the concrete model\(^{30}\) the LKP is a neutral WIMP which is a linear combination of the first KK mode \( B_1 \) of the hypercharge gauge boson and the first KK mode \( W^0_1 \) of the neutral \( SU(2)_W \) gauge boson.

The calculation of the relic density of KK DM based on WMAP’s results for \( \Omega_{\text{CDM}} \) predicts the value of LKP mass: \( m_{\text{LKP}} \approx (600 - 1200) \text{ GeV} \).

In contrast with LSP, LKP is a vector particle, and the helicity suppression is absent, so, LKP can directly annihilate into neutrinos. Besides, the annihilation into pairs of top-quarks is sufficient due to a large value of LKP mass. Top-quarks almost always decay in the channel \( t \rightarrow Wb \). The \( W \), in turn, decays with equal branching ratios into \( \nu_e \nu_e, \nu_\mu \nu_\mu, \nu_\tau \nu_\tau \). It is important that both decays, \( t \)'s and \( W \)'s (and also \( \tau \)'s) take place without any energy loss. So, in this model, the Sun is effective source of high energy (\( E_\nu \approx 100 \text{ GeV} - 1 \text{ TeV} \)) tau-neutrinos and (due to oscillations) muon neutrinos.

4.3. Superheavy DM

It is well known that the assumption that DM is a thermal relic is too restrictive. Unitarity bound on the annihilation cross section and the assumption of thermal equilibrium in the
early universe lead to the bound,
\[ \Omega h^2 \geq 0.1 \left( \frac{M}{10^5 \text{ GeV}} \right)^2, \]
where \( \Omega \) is the ratio of the DM mass density to \( \rho_c \) today. However, if WIMPs are not thermal relics, their masses are not thermodynamically constrained, i.e., they may be even superheavy (“wimpzillas”). Moreover, their present day abundance does not depend on whether they have strong (“simpzillas”), weak, electromagnetic, or only gravitational interactions.

Such particles can be, e.g., gravitationally produced near the end of inflation, due to non-adiabatic change of the scale factor during the transition from the de-Sitter to the radiation dominated phase.

High energy neutrinos are produced by the simpzilla annihilations, which produce a quark or gluon pair which then fragment into hadronic jets. Neutrino fluxes from the Sun were calculated by Albuquerque et al.\textsuperscript{34} It was shown that \( \nu_\tau \) and \( \nu_\mu \) fluxes can be rather large and measurable by neutrino telescopes, with \( E_\nu \) in the interval \((50 \text{ GeV} – 1 \text{ TeV})\), the spectrum is very similar with the case of KK DM (the channel \( t \to Wb \) also works). One should note that the typical \( E_\nu \) value in the spectrum is much lower than the simpzilla mass \((M \approx 10^8 – 10^{16} \text{ GeV})\), due to the large numbers of neutrinos produced per annihilation.

Unfortunately, direct dark matter search experiments exclude\textsuperscript{35} the most natural values of simpzilla mass (those which are comparable with the inflaton mass in chaotic inflation models \((\approx 10^{12} \text{ GeV})\)).

References
1. K. Mannheim, R. J. Protheroe and J. P. Rachen, \textit{Phys. Rev.} \textbf{D63}, 023003 (2001).
2. E. Waxman and J. N. Bahcall \textit{Phys. Rev.} \textbf{D59}, 023002 (1999).
3. D. V. Semikoz and G. Sigl, \textit{JCAP} \textbf{04}, 003 (2004) [\texttt{hep-ph/0309128}].
4. M. Kachelrieß, D. V. Semikoz and M. A. Tortola, \textit{Phys. Rev.} \textbf{D68}, 043005 (2003).
5. D. F. Torres and L. A. Anchordoqui, \textit{Rept. Prog. Phys.} \textbf{67}, 1663 (2004) [\texttt{astro-ph/0402371}].
6. K. Katak and J. Kwieciński, \textit{Eur. Phys. J.} \textbf{C29}, 521 (2003).
7. K. Golek-Biernat and M. Wusthoff, \textit{Phys. Rev.} \textbf{D59}, 014017 (1999).
8. G. Sigl, \texttt{hep-ph/0408165}.
9. N. Arkani-Hamed, S. Dimopoulos and G. Dvali, \textit{Phys. Lett.} \textbf{B429}, 263 (1998).
10. R. Emparan \textit{et al.}, \textit{Phys. Rev.} \textbf{D65}, 064023 (2002).
11. F. Cornet \textit{et al.}, \textit{Phys. Rev. Lett.} \textbf{86}, 4235 (2001).
12. J. I. Illana \textit{et al.}, \textit{Phys. Rev. Lett.} \textbf{93}, 151102 (2004) [\texttt{hep-ph/0402279}]; P. Jain \textit{et al.}, \textit{Phys. Rev. Lett.} \textbf{D66}, 065018 (2002) [\texttt{hep-ph/0205052}].
13. H. Tu, \textit{Surveys High Energ. Phys.} \textbf{17}, 149 (2002) [\texttt{hep-ph/0205024}].
14. J. Feng and A. Shapere, \textit{Phys. Rev. Lett.} \textbf{88}, 021303 (2002).
15. L. Anchordoqui and H. Goldberg, \textit{Phys. Rev.} \textbf{D65}, 047502 (2002).
16. J. L. Feng \textit{et al.}, \textit{Phys. Rev. Lett.} \textbf{88}, 161102 (2002) [\texttt{hep-ph/0105067}]; A. Kusenko and T. J. Weiler, \textit{Phys. Rev. Lett.} \textbf{88}, 161101 (2002) [\texttt{hep-ph/0106071}].
17. D. Fargion \textit{et al.}, \textit{Astrophys. J.} \textbf{517} 725 (1999); T. J. Weiler \textit{Astropart. Phys.} \textbf{11}, 303 (1999).
18. V. Berezinsky, M. Kachelrieß and S. Ostapchenko, \textit{Phys. Rev. Lett.} \textbf{89}, 171802 (2002).
19. A. Strong \textit{et al.}, [\texttt{astro-ph/0306345}].
20. G. Gelmini \textit{et al.}, \textit{Phys. Rev.} \textbf{D70}, 113005 (2004) [\texttt{hep-ph/0404272}].
21. S. Singh and C. Ma, \textit{Phys. Rev.} \textbf{D67}, 023506 (2003).
22. A. D. Dolgov et al., *Nucl. Phys.* **B632**, 363 (2002).
23. A. Ringwald, *JHEP* **10**, 008 (2003) [hep-ph/0307034]; Z. Fodor et al., *Phys. Lett.* **B561**, 191 (2003) [hep-ph/0303080].
24. F. Bezrukov et al., *Phys. Rev.* **D68**, 036005 (2003) [hep-ph/0304180]; *Phys. Lett.* **B574**, 75 (2003) [hep-ph/0305300].
25. T. Han and D. Hooper, *Phys. Lett.* **B582**, 21 (2004).
26. C. Benneth et al., *Astrophys. J. Supp.* **148**, 1 (2003).
27. J. Silk and M. Srednicki, *Phys. Rev. Lett.* **53**, 624 (1984).
28. D. Hooper and J. Silk, *New J. Phys.* **6**, 023 (2004) [hep-ph/0311367].
29. J. Ellis, *Phys. Scripta* **T85**, 221 (2000).
30. H. Cheng et al., *Phys. Rev.* **D66**, 056006 (2002).
31. G. Servant and T. Tait, *Nucl. Phys.* **B650**, 391 (2003).
32. K. Griest and M. Kamionkowski, *Phys. Rev. Lett.* **64**, 615 (1990).
33. D. J. H. Chung et al., *Phys. Rev.* **D59**, 023501 (1999) [hep-ph/9802238]; V. Kuzmin and I. Tkachev, *JETP Lett.* **68**, 271 (1998) [hep-ph/9802238].
34. I. Albuquerque et al., *Phys. Rev.* **D64**, 083504 (2001).
35. I. Albuquerque and L. Baudis, *Phys. Rev. Lett.*, **90**, 221301 (2003).