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Ernest Ma and Utpal Sarkar

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Scalar Neutrino as Dark Matter: Radiative Neutrino Mass and Leptogenesis

Ernest Ma^a and Utpal Sarkar^{b,c}

^a *Department of Physics and Astronomy, University of California,
Riverside, California 92521, USA*

^b *Physical Reserach Laboratory, Ahmedabad 380009, India*

^c *Physics Department and McDonnell Center for the Space Sciences,
Washington University, St. Louis, Missouri 63130, USA*

Abstract

In the Minimal Supersymmetric Standard Model (MSSM), the scalar neutrino $\tilde{\nu}_L$ has odd R parity, yet it has long been eliminated as a dark-matter candidate because it scatters elastically off nuclei through the Z boson, yielding a cross section many orders of magnitude above the experimental limit. We show how it can be reinstated as a dark-matter candidate by splitting the masses of its real and imaginary parts in an extension of the MSSM with scalar triplets. As a result, radiative neutrino masses are generated. This severely constrains the $U(1)$ and $SU(2)$ gaugino masses as a function of the $\tilde{\nu}_L$ mass, which is bounded by the 2011 XENON100 data to be above 125 GeV. Whereas a $\tilde{\nu} - \tilde{\nu}^*$ asymmetry is created from the decay of a heavy scalar triplet together with a lepton asymmetry, it gets washed out by $\tilde{\nu} - \tilde{\nu}^*$ oscillations after the electroweak phase transition. However, $\tilde{\nu}_2$ then decays into $\tilde{\nu}_1 \nu \bar{\nu}$ fast enough so that only $\tilde{\nu}_1$ survives as dark matter today.

The imposition of R parity, i.e. $R \equiv (-1)^{3B+L+2j}$, in the Minimal Supersymmetric Standard Model (MSSM) of particle interactions serves at least two purposes. One is to avoid proton decay in its renormalizable interactions; the other is to establish a dark-matter candidate which is neutral and stable, i.e. odd under R . This candidate particle may be a boson or a fermion. If it is a boson, then it should be the lightest of three scalar neutrinos $\tilde{\nu}_L$. If it is a fermion, then it should be the lightest of four neutralinos, i.e. the $U(1)$ and neutral $SU(2)$ gauginos and the two neutral higgsinos. However, a scalar neutrino scatters elastically off nuclei with an amplitude mediated by the Z boson, yielding a cross section many orders of magnitude above the present experimental limit, so it was eliminated as a dark-matter candidate many years ago. As for the lightest neutralino, which is a Majorana linear combination of gauginos and higgsinos, it is still considered as the canonical candidate for dark matter.

To reinstate the scalar neutrino as a dark-matter candidate, its elastic scattering with nuclei must be suppressed and this is easily achieved by splitting the mass of its real and imaginary components. The reason is that the coupling of the vector Z boson to $\tilde{\nu}_L = (\tilde{\nu}_1 + i\tilde{\nu}_2)/\sqrt{2}$ is of the form $Z\tilde{\nu}_1\tilde{\nu}_2$, so if the mass gap is greater than about 100 keV, this process is forbidden by kinematics in the nuclear elastic recoil experiments.

There are now two issues to be considered. (1) How is this splitting achieved? A mass splitting term $\tilde{\nu}_L\tilde{\nu}_L$ cannot be put in by hand, because it is not invariant under the $SU(2)_L \times U(1)_Y$ gauge symmetry of the MSSM. If it is simply assumed to be an effective term without specifying its underlying origin, then it cannot be guaranteed that whatever conclusion is drawn from its existence will not be affected by the actual dynamics which generated it in the first place. Here we assume that it comes from the gauge-invariant term $\Delta_1^0\tilde{\nu}_L\tilde{\nu}_L - \sqrt{2}\Delta_1^+\tilde{\nu}_L\tilde{e}_L - \Delta_1^{++}\tilde{e}_L\tilde{e}_L$, where $\Delta_1 = (\Delta_1^{++}, \Delta_1^+, \Delta_1^0)$ is a scalar triplet, with a vacuum expectation value $\langle\Delta_1^0\rangle = u_1$. (2) Once the specific origin of this splitting is identified,

what are its physical consequences? The first is of course neutrino mass. Since $\tilde{\nu}_L$ carries lepton number L , the induced mass splitting term $\tilde{\nu}_L\tilde{\nu}_L$ breaks L to $(-1)^L$. The observed neutrinos must then have Majorana masses and a radiative contribution must exist through the exchange of $\tilde{\nu}_L$ and neutralinos in one loop. More importantly, the scalar triplet Δ_1 should also couple to the neutrinos directly which then obtain masses through u_1 in the well-known manner of the Type II seesaw. This latter would imply a very small u_1 , much less than 100 keV, thus invalidating the interpretation of $\tilde{\nu}_1$ as dark matter.

In the following we overcome the above objection by forbidding the dimension-four term $\Delta_1^0\nu_L\nu_L$. We do this by assigning $L = 0$ to $\Delta_{1,2}$ where $\Delta_2 = (\Delta_2^0, \Delta_2^-, \Delta_2^{--})$ and insisting that L be conserved by all dimension-four terms of the supersymmetric Lagrangian of this model. We then break the supersymmetry by soft terms which are allowed to break L to $(-1)^L$ as well, i.e. the dimension-three term $\Delta_1^0\tilde{\nu}_L\tilde{\nu}_L$. We then show that neutrinos do acquire radiative Majorana masses [1] in this case, but they are only compatible with $\tilde{\nu}_1$ as dark matter if the $U(1)$ and $SU(2)$ gaugino masses have opposite signs, a phenomenological possibility that has been largely overlooked. We also show how the decays of $\Delta_{1,2}$ result [2] in both a lepton asymmetry and an asymmetry in $\tilde{\nu}_L$, with its relic density determined by the subsequent annihilation of $\tilde{\nu}_L\tilde{\nu}_L$ into $\nu_L\nu_L$. Note that the mass splitting of the scalar neutrino is not induced by heavy singlet (right-handed) neutrino superfields through mixing. If it were [3, 4], then there would also be a tree-level neutrino mass from the Type I seesaw. If the inverse seesaw mechanism were used instead [5], then again there would be both a tree-level mass and a loop-induced mass. In our case, only the latter occurs and as we show later in Eq. (3), this is a crucial condition for $\tilde{\nu}_1$ to be a viable dark-matter candidate.

The superpotential of this model is given by

$$\begin{aligned}
W = & \mu\hat{\Phi}_1\hat{\Phi}_2 + f_{ij}^e\hat{\Phi}_1\hat{L}_i\hat{e}_j^c + f_{ij}^d\hat{\Phi}_1\hat{Q}_i\hat{d}_j^c + f_{ij}^u\hat{\Phi}_2\hat{Q}_i\hat{u}_j^c \\
& + M\hat{\Delta}_1\hat{\Delta}_2 + f_1\hat{\Delta}_1\hat{\Phi}_1\hat{\Phi}_1 + f_2\hat{\Delta}_2\hat{\Phi}_2\hat{\Phi}_2,
\end{aligned} \tag{1}$$

where $\hat{\Phi}_1 \sim (1, 2, -1/2)$, $\hat{\Phi}_2 \sim (1, 2, 1/2)$, $\hat{L} \sim (1, 2, -1/2)$, $\hat{e}^c \sim (1, 1, 1)$, $\hat{Q} \sim (3, 2, 1/6)$, $\hat{d}^c \sim (3^*, 1, 1/3)$, $\hat{u}^c \sim (3^*, 1, -2/3)$, as in the MSSM. The Higgs triplet superfields are $\hat{\Delta}_1 \sim (1, 3, 1)$ and $\hat{\Delta}_2 \sim (1, 3, -1)$, which have been assigned lepton number $L = 0$, so that the terms $\hat{\Delta}_1 \hat{L}_i \hat{L}_j$ are forbidden.

We allow L to be broken by soft terms which also break the supersymmetry of the model, but only by two units, i.e. $\Delta L = \pm 2$. This would forbid the bilinear $\tilde{L}_i \Phi_2$ and trilinear $\tilde{L}_i \tilde{L}_j \tilde{e}_k^c$, $\tilde{L}_i \tilde{Q}_j \tilde{d}^c$ terms, but allow the trilinear $\Delta_{1,2} \tilde{L}_i \tilde{L}_j$ terms. This pattern is stable because it is maintained by the residual Z_2 symmetry $(-1)^L$. Since $\Delta_{1,2}$ mix through the soft $BM\Delta_1\Delta_2$ term, the two resulting mass eigenstates both decay into states of $L = 2$ as well as $L = 0$, i.e. $\Phi_{1,2}$. Leptogenesis [2] is then possible. Details of this supersymmetric scenario has been worked out previously [6, 7, 8]. For the present model we consider the resonance condition [7] $B = \Gamma_{\pm}$, where Γ_{\pm} are the decay widths of the physical triplet Higgs scalars, to obtain the required lepton asymmetry before the electroweak phase transition is over.

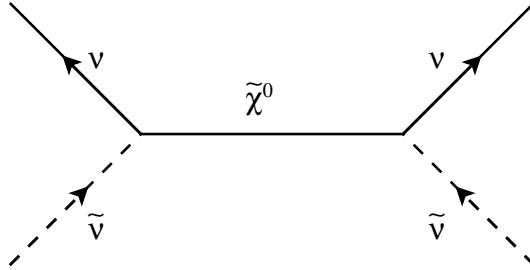


Figure 1: Annihilation of $\tilde{\nu}\tilde{\nu} \rightarrow \nu\nu$ via neutralino exchange.

As a lepton asymmetry is established, there is also an asymmetry of the scalar neutrinos $\tilde{\nu}$. This connection between visible and dark matter has been explored previously [9, 10, 11]. It is also possible in the context of the radiative seesaw model of neutrino mass [1] with the addition of heavy scalar triplets, as proposed recently [12]. As the Universe cools below $m_{\tilde{\nu}}$, the relic abundance of $\tilde{\nu}$ is determined by its annihilation cross sections with $\tilde{\nu}^*$ and

with itself. The former is very large, which means that there is nothing left but for the fact there is a $\tilde{\nu} - \tilde{\nu}^*$ asymmetry. The latter is rather small, but it still diminishes the $\tilde{\nu}$ number density, until the process goes out of thermal equilibrium due to the expansion of the Universe. In that case, it may be a possible candidate for the dark matter of the Universe. The $\tilde{\nu}\tilde{\nu}$ annihilation proceeds through neutralino exchange, as shown in Fig. 1. In the 4×4 neutralino mass matrix, if the higgsino mass parameter μ is large, then the 2×2 gaugino mass matrix does not mix significantly with the 2×2 higgsino mass matrix, resulting in approximate mass eigenvalues $m_{1,2}$ for the $U(1)$ and $SU(2)$ gauginos. In that case, this cross section \times relative velocity is given by

$$\langle\sigma v\rangle = \frac{g^4}{128\pi c^4 m_{\tilde{\nu}}^2} \left(\frac{s^2 y_1}{y_1^2 + 1} + \frac{c^2 y_2}{y_2^2 + 1} \right)^2, \quad (2)$$

where $s = \sin \theta_W$, $c = \cos \theta_W$, and $y_{1,2} = m_{1,2}/m_{\tilde{\nu}}$.

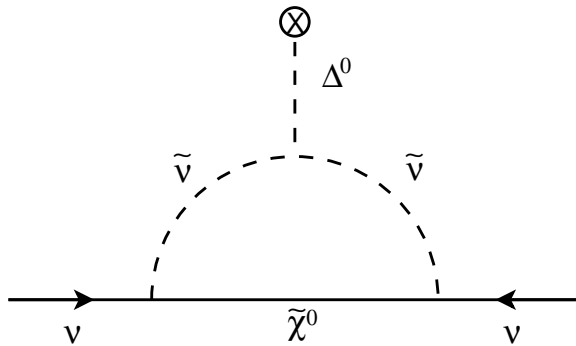


Figure 2: One-loop radiative Majorana neutrino mass via neutralino exchange. Lepton number L becomes $(-1)^L$ through the soft $\Delta^0 \tilde{\nu}\tilde{\nu}$ term.

For each $\tilde{\nu}$, a radiative neutrino mass is also generated in one loop by neutralino exchange, as shown in Fig. 2, resulting in [1, 3]

$$\frac{m_\nu}{\Delta m_{\tilde{\nu}}} = \frac{g^2}{32\pi^2 c^2} \left[\frac{s^2 y_1}{y_1^2 - 1} \left(1 - \frac{y_1^2}{y_1^2 - 1} \ln y_1^2 \right) + \frac{c^2 y_2}{y_2^2 - 1} \left(1 - \frac{y_2^2}{y_2^2 - 1} \ln y_2^2 \right) \right]. \quad (3)$$

Since $\Delta m_{\tilde{\nu}} > 100$ keV is needed to suppress the interaction of $\tilde{\nu}_1$ with nuclei in underground direct-search experiments, and $m_\nu < 1$ eV for neutrino mass, the above ratio should be

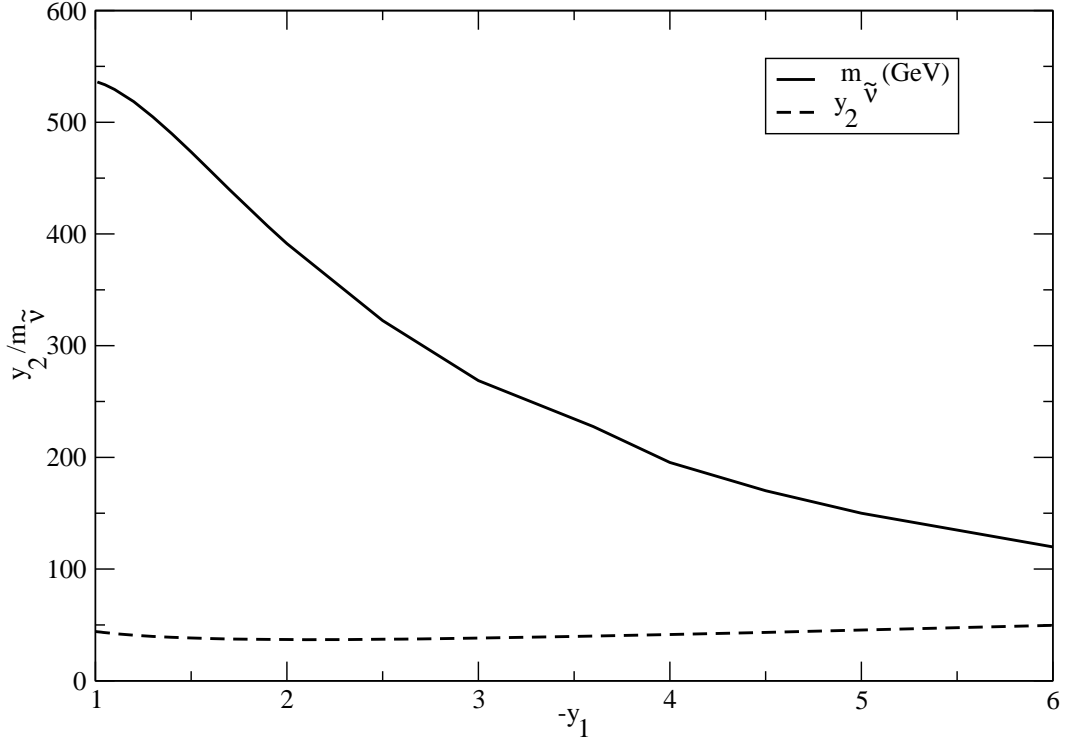


Figure 3: y_2 and $m_{\tilde{\nu}}$ are plotted against $-y_1$, which correspond to $m_\nu/\Delta m_{\tilde{\nu}} = 0$ and $\langle\sigma v\rangle = 0.01$ pb.

less than 10^{-5} . We assume that $|y_2| > |y_1| > 1$ and plot y_2 as a function of $-y_1$ so that $m_\nu/\Delta m_{\tilde{\nu}} = 0$ in Fig. 3. This is a severe constraint and a specific testable prediction of our model at the Large Hadron Collider (LHC). We also use Eq. (2) to plot $m_{\tilde{\nu}}$ as a function of $-y_1$ for the particular value of $\langle\sigma v\rangle = 0.01$ pb. To understand this choice, we note that for dark matter, since an asymmetry in the scalar neutrinos is generated along with the lepton asymmetry, the dark matter abundance is determined by the CP amplitude. As shown in Ref. [7], the CP asymmetry could be of order 0.01 or larger when the resonance condition is satisfied, so the dark matter abundance estimate is suppressed by 0.01. But this is inversely proportional to the annihilation cross section, so we take 0.01 pb instead of the usual 1 pb for our estimate. This is an order-of-magnitude estimate of the equivalent value for $\tilde{\nu}$ because its number density before freeze-out is diminished by its asymmetry. It should only be taken

as a guide to the resulting range of allowed $\tilde{\nu}$ mass values. We note that $m_{\tilde{\nu}} < m_Z/2$ is ruled out experimentally, because $Z \rightarrow \tilde{\nu}_1 \tilde{\nu}_2$ would then contribute to its invisible width, which already agrees very well with what is expected from the three known neutrinos of the Standard Model.

In underground dark-matter direct-search experiments, the spin-independent elastic cross section for $\tilde{\nu}_1$ scattering off a nucleus of Z protons and $A - Z$ neutrons normalized to one nucleon is given by

$$\sigma_0 = \frac{1}{\pi} \left(\frac{m_N}{m_{\tilde{\nu}} + Am_N} \right)^2 \left| \frac{Zf_p + (A - Z)f_n}{A} \right|^2, \quad (4)$$

where m_N is the mass of a nucleon, and $f_{p,n}$ come from Higgs exchange [13]:

$$\frac{f_p}{m_p} = \left(-\frac{0.075}{4} - \frac{0.925(3.51)}{54} \right) \frac{g^2}{c^2 m_\phi^2}, \quad (5)$$

$$\frac{f_n}{m_n} = \left(-\frac{0.078}{4} - \frac{0.922(3.51)}{54} \right) \frac{g^2}{c^2 m_\phi^2}. \quad (6)$$

Assuming an effective $m_\phi = 125$ GeV and using $Z = 54$ and $A - Z = 77$ for ^{131}Xe , we plot σ_0 as a function of $m_{\tilde{\nu}}$ in Fig. 4. Note for $m_{\tilde{\nu}} = 130$ GeV, σ_0 is about 1.2×10^{-8} pb, which is not far below the upper limit of the 2011 XENON100 exclusion [14]. The allowed range for $m_{\tilde{\nu}}$ is above 125 GeV.

Before the electroweak phase transition, the dark-matter asymmetry of $\tilde{\nu}$ is maintained after freeze-out, but after the electroweak phase transition, $\tilde{\nu} - \tilde{\nu}^*$ oscillations become possible and this asymmetry is washed out. In other words, the number densities of $\tilde{\nu}$ and $\tilde{\nu}^*$ become equal. However, they are now organized into the mass eigenstates $\tilde{\nu}_1 = \sqrt{2}\text{Re}(\tilde{\nu})$ and $\tilde{\nu}_2 = \sqrt{2}\text{Im}(\tilde{\nu})$, and if $\Delta m_{\tilde{\nu}} < 2m_e$, the decay of $\tilde{\nu}_2$ is only to $\tilde{\nu}_1 \nu \bar{\nu}$ through Z exchange. This decay width is given by

$$\Gamma(\tilde{\nu}_2 \rightarrow \tilde{\nu}_1 \nu \bar{\nu}) = \frac{G_F^2 (\Delta m_{\tilde{\nu}})^5}{60\pi^3}. \quad (7)$$

For $100 \text{ keV} < \Delta m_{\tilde{\nu}} < 1 \text{ MeV}$, the corresponding lifetime is of order 10^9 to 10^4 seconds.

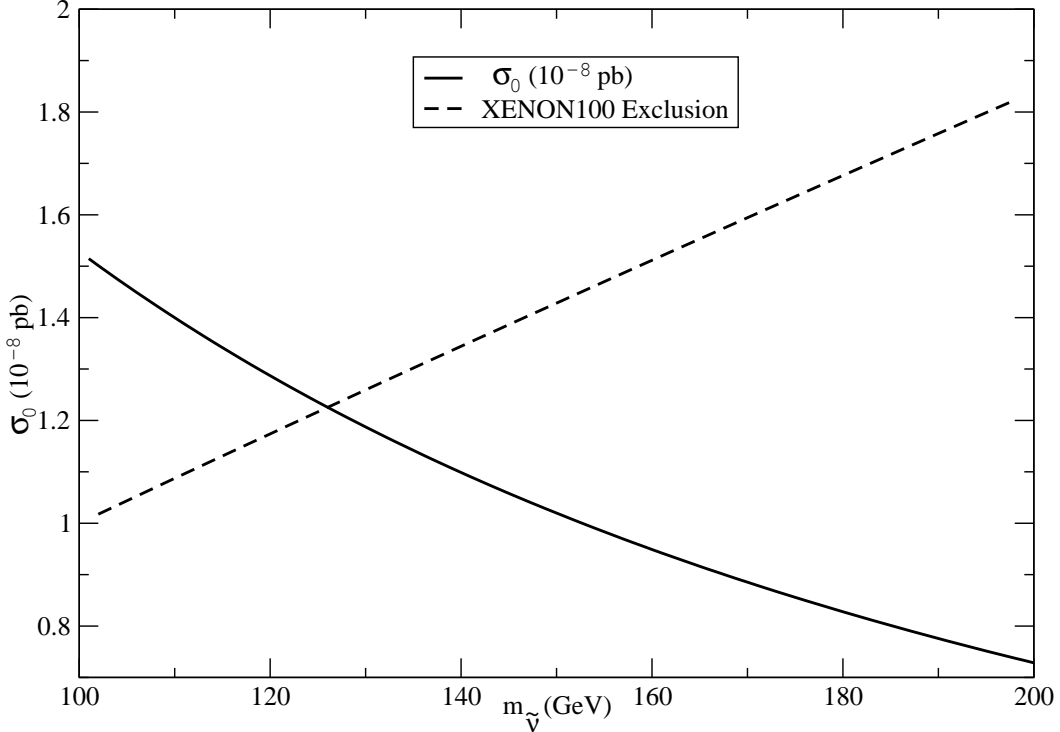


Figure 4: Spin-independent elastic scattering cross section of $\tilde{\nu}_1$ with ^{131}Xe through Higgs exchange is plotted together with the present experimental bound from the direct-search experiment XENON100, as a function of $m_{\tilde{\nu}}$.

This means that at present only $\tilde{\nu}_1$ survives as dark matter. If a concentration of $\tilde{\nu}_1$ has accumulated inside the sun or the earth, $\tilde{\nu}_1\tilde{\nu}_1$ annihilation to two monoenergetic neutrinos [15] would be a spectacular indication of this scenario.

If $\tilde{\nu}_1$ is dark matter, its production at the Large Hadron Collider (LHC) must always be accompanied by a lepton. If it comes from the decay of the $U(1)$ gaugino, i.e. $\tilde{B} \rightarrow \bar{\nu}\tilde{\nu}$, then it is completely invisible. If it comes from a chargino, i.e. $\tilde{\chi}^+ \rightarrow e^+\tilde{\nu}$, then it may be discovered through its missing energy and large mass. More detailed study of this scenario is required.

The Higgs triplet scalars $\Delta_{1,2}$ are the common origins of both the observed baryon asymmetry and the $\tilde{\nu}_L$ dark matter of the Universe. They are likely to be very heavy, say of order

10^8 GeV, in which case they are not accessible at the LHC. On the other hand, resonant leptogenesis may occur naturally in this scenario [8], which would allow them to be at the TeV scale. In that case, the direct decay $\Delta_1^{++} \rightarrow \tilde{e}_i^+ \tilde{e}_j^+$ would serve to map out the neutrino mass matrix, in analogy to the previously proposed simple scenario [16, 17], where $\Delta^{++} \rightarrow e_i^+ e_j^+$.

In conclusion, we have proposed that the dark matter of the Universe is the real component of the lightest scalar neutrino $\tilde{\nu}_1$ in the Minimal Supersymmetric Standard Model. To implement this unconventional scenario, we add Higgs triplet superfields $\hat{\Delta}_{1,2}$, so that the observed neutrinos acquire small radiative Majorana masses from the mass splitting terms $\Delta_1 \tilde{\nu}_L \tilde{\nu}_L$. The fact that $m_\nu / \Delta m_{\tilde{\nu}} < 10^{-5}$ forces the $U(1)$ and $SU(2)$ gaugino masses to have opposite signs, and fixed values of their masses as a function of $\tilde{\nu}$ mass. Using the latest XENON100 data, the allowed mass range of $\tilde{\nu}_1$ is above 125 GeV, assuming a Higgs-boson mass of 125 GeV.

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