Experimental signatures of a new dark matter WIMP

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Abstract –The WIMP proposed here yields the observed abundance of dark matter, and is consistent with the current limits from direct detection, indirect detection, and collider experiments, if its mass is $\sim 72 \text{ GeV} / c^2$. It is also consistent with analyses of the gamma rays observed by Fermi-LAT from the Galactic center (and other sources), and of the antiprotons observed by AMS-02, in which the excesses are attributed to dark matter annihilation. These successes are shared by the inert doublet model (IDM), but the phenomenology is very different: The dark matter candidate of the IDM has first-order gauge couplings to other new particles, whereas the present candidate does not. In addition to indirect detection through annihilation products, it appears that the present particle can be observed in the most sensitive direct-detection and collider experiments currently being planned.

As is well-known, there are a vast number of hypothetical dark matter candidates, most of which do not have well-defined masses or couplings, and many of which have already been ruled out by experiment. The most popular single candidate has been the lightest neutralino $\chi^0$ of supersymmetry (susy) [1]. But, as is also well-known, faith in this candidate has diminished during the past few years, because neither a susy dark matter WIMP nor other susy particles have been observed, despite strenuous efforts. In addition, the simplest susy models which have “natural” values for the parameters, and which are also compatible with limits from the LHC, are found to be in disagreement with both the abundance of dark matter and the limits from direct-detection experiments [2,4].

Here we propose an alternative candidate which in some respects resembles the neutralino, but which would open up a whole new sector with new particles and new physics which can be observed in the foreseeable future. We have named particles of this new kind “higgsons” [7–10], represented by $H$, to distinguish them from Higgs bosons $h$ and the higgsinos $\tilde{h}$ of susy. The lightest neutral particles in these three groups are $H^0$, $h^0$, and $\tilde{h}^0$. We will demonstrate below that these new particles (higgsons) have very favorable features – after the theory has been reformulated in the way described below.

Many of these features are shared by the inert doublet model (IDM) – introduced long ago [11] and later proposed as an explanation of dark matter [12] – for which the phenomenology has been very extensively explored in many papers. Only a representative sample can be cited here [13–27], but the basic idea is that the Standard Model Higgs doublet is supplemented by a second doublet of scalar fields which have odd parity under a postulated new $Z_2$ symmetry. To avoid confusion, these fields will be distinguished by a subscript $I$:

$$
\frac{1}{\sqrt{2}} \left( H^0_I + i A^0_I \right)
$$

Since each of these fields is odd with respect to the hypothetical $Z_2$, whereas all other fields are even, every term in the action must involve an even number of these fields. This makes the dark matter candidate $H^0$ stable, because it has the lowest mass. In addition, many of the successes described below for the present candidate, called $H^0$ here, are shared by $H^I$.

However, in the phenomenology of the IDM there are numerous important processes that involve first-order gauge couplings of $H^0$ to the other particles $A^I$ and $H^I$. See, e.g., Figs. 3 and 4 of [14]; Fig. 2 of [13]; Fig. 1 of [16]; Fig 4 of [18]; Figs. 9, 10, 11, 13, 14, 16, 17, 18, and 21 of [22]; Figs. 2 and 3 of [25]; Fig. 6 of [26]; Figs. 1 and 3 of [27] and the constraints on the IDM parameters in Section 4.1 of [21], which imply that the masses of $A^I$ and $H^I$ should be reasonably near that of $H^0$.

As will be seen below, the present dark matter candidate $H^0$ has only second-order gauge couplings – and these are
to only itself plus the gauge bosons. This results in a very different phenomenology, with a substantial reduction in the number of observable processes involved in annihilation, scattering, and creation. There are other differences: \(H^0\) is its own antiparticle, it does not require an extra \(\mathbb{Z}_2\) symmetry, and it is related to Higgs bosons in a very different way.

We begin with the primitive 2-component spin 1/2 bosonic fields \(\Phi_r\) introduced in our previous papers, which are joined in an \(n\)-component gauge multiplet \(\Phi_r\). (We name each multiplet \(\Phi_r\) after a typical member \(\Phi_r\) in order to simplify the notation.) Each \(\Phi_r\) transforms as a 2-component spinor under rotations, but is left unchanged by a boost.

It will be seen that the action for the physical fields \(\Phi_R\), as defined below, is invariant under both rotations and boosts. The final theory is in fact fully Lorentz invariant.

For simplicity, we begin with a single (e.g. grand-unified) gauge field having covariant derivative

\[
D_{\mu} = \partial_\mu - iA_\mu, \quad A_\mu = A^i_{\mu} v^i
\]

so that the coupling constant is temporarily absorbed into \(A_\mu\) and field strength

\[
F_{\mu \nu} = F^i_{\mu \nu} v^i, \quad F^i_{\mu \nu} = \partial_\mu A^i_\nu - \partial_\nu A^i_\mu + c^i_{ij} A^j_\mu A^k_\nu .
\]

Later we will specialize to the electroweak theory after symmetry breaking. Natural units \(\hbar = c = 1\) are used, with the \((-+++\) convention for the metric tensor. Summations are never implied over a repeated gauge index like \(r\) or \(R\), but are always implied over repeated coordinate indices like \(\mu, \nu = 0, 1, 2, 3\) and \(k = 1, 2, 3\), as well as the index \(i\) or \(j\) labeling gauge generators \(v^i\). The same name and symbol are used for a field and the particle which is an excitation of that field.

The initial action for the new fields \(\tilde{\Phi}_r\) is

\[
S_r = \int d^4x \tilde{\Phi}_r^\dagger(x) D^\mu D_\mu \tilde{\Phi}_r(x)
\]

\[
+ \int d^4x \tilde{\Phi}_r^\dagger(x) \vec{B}(x) \cdot \vec{\sigma} \tilde{\Phi}_r(x)
\]

where the 3-vectors \(\vec{\sigma}\) and \(\vec{B}\) respectively contain the Pauli matrices \(\sigma^i\) and the “magnetic” components of the field strength tensor \(F_{\mu \nu}\):

\[
F_{k k'} = -\varepsilon_{k k' l} B_{l k'} .
\]

This action is derived in Ref. [?], but here is taken to be a phenomenological postulate [?].) At this point we deviate from our earlier papers by requiring all fundamental physical fields \(\Phi_R\) to satisfy Lorentz invariance. This means that the anomalous second term in \([\text{?}]\) must disappear from the action. This can be achieved by requiring the physical fields to assume one of the two forms defined below, which are respectively called Higgs/amplitude fields and Higgs/field.

**Higgs/amplitude fields** are defined to be those for which

\[
\tilde{\Phi}_R^\dagger(x) \vec{\sigma} \tilde{\Phi}_R(x) = 0
\]

(where \(\vec{\sigma}\) always implicitly multiplies an appropriate identity matrix). These fields – analogous to the Higgs/amplitude modes observed in superconductors \([28]\) – are obtained by writing the 4n-component field \(\Phi_r\) in terms of 2n-component fields \(\Phi_r\) and \(\Phi_{r'}\) with opposite spins (and equal amplitudes):

\[
\tilde{\Phi}_r = \begin{pmatrix} \tilde{\Phi}_r \\ \tilde{\Phi}_{r'} \end{pmatrix}
\]

so that

\[
\tilde{\Phi}_r^\dagger(x) \vec{\sigma} \tilde{\Phi}_r(x) = \begin{pmatrix} \tilde{\Phi}_r^\dagger(x) \vec{\sigma} \tilde{\Phi}_{r'}(x) \\ \vec{\sigma} \tilde{\Phi}_r(x) \end{pmatrix} = 0 .
\]

In this case we write for each gauge component

\[
\Phi_R(x) = \phi_R(x) \xi_R \quad \text{with} \quad \xi_R^\dagger \xi_R = 1
\]

where \(\xi_R\) has 4 constant components and \(\phi_R(x)\) is a 1-component complex amplitude. Then \([\text{?}]\) and its counterpart for \(r'\) reduce to

\[
S_R = \int d^4x L_R , \quad L_R = \phi_R^\dagger(x) D^\mu D_\mu \phi_R(x)
\]

where \(\phi_R\) contains all the gauge components \(\phi_R\).

In the electroweak theory, we interpret \(\phi_R\) with components \(\phi_R^x\) and \(\phi_R^y\), as the usual Higgs doublet (in the Higgs basis), with \(\phi_R^0\) condensing and supporting Higgs boson excitations \(h_R^0\):

\[
\phi_R^0 = \psi_R^0 + h_R^0 , \quad \psi_R^0 = \langle \phi_R^0 \rangle .
\]

We have assumed a pair of bosonic doublets \(\Phi_r\) and \(\Phi_{r'}\) with the same gauge quantum numbers – i.e., (weak) isospin and hypercharge – in somewhat the same spirit as in standard susy. We choose the convention that \(\Phi_r\) has spin up and \(\Phi_{r'}\) spin down for the field \(\Phi_R\) of \([7]\) which condenses. There is then another independent field \(\Phi_R\) in which \(\Phi_r\) has spin down and \(\Phi_{r'}\) spin up. This field has no condensate. Only the kinetic term of \([\text{?}]\) has been treated at this point, without the potentially quite complicated terms that yield masses and further interactions.

**Higgs/field** are defined to be those for which

\[
\Phi_S^\dagger(x) \vec{B} \Phi_S(x) = 0
\]

(where \(\vec{B}\) implicitly multiplies a \(4 \times 4\) identity matrix, and the subscripts \(s\) and \(S\) will be used in this context to avoid confusion). This can be achieved by writing the
4-component $\Phi_S$ in terms of a 2-component field $\Phi_s$ and its charge conjugate $\Phi_s^c$:

$$\Phi_S = \frac{1}{\sqrt{2}} \left( \begin{array}{c} \Phi_s \\ \Phi_s^c \end{array} \right).$$  \hspace{1cm} (14)

From (3) and (5) we have

$$B_{k\ell'} = -\delta_{k\ell'} F_{kk'} , \quad F_{kk'} = F_{kk'}^{ij} t^j .$$  \hspace{1cm} (15)

Also, $\Phi_s$ and $\Phi_s^c$ have opposite gauge quantum numbers -- i.e. opposite expectation values for the generators $t^j$ (which are here treated as operators rather than matrices):

$$\Phi_s^{c\dagger} t^j \Phi_s^c = -\Phi_s^{\dagger} t^j \Phi_s .$$  \hspace{1cm} (16)

It follows that

$$\Phi_S^\dagger(x) \equiv \frac{\Phi_s}{\sqrt{2}} \Phi_s + \Phi_s^c \equiv \Phi_s \cdot \Phi_s + \Phi_s^c \cdot \Phi_s$$

$$= 0 .$$  \hspace{1cm} (17)

We have thus satisfied relativistic invariance by introducing the bosonic analog of Majorana fields.

At this point three comments are appropriate: (1) All the higgsion fields $\Phi_S$ are charge-neutral. (2) In the following we will consider only the $\Phi_S$ with no condensate. (3) The lowest-mass excitation of such a field, called $H^0$ below, is not the only stable higgson, but it will emerge from the early universe with the highest density: The more massive particles will fall out of equilibrium earlier and be rapidly thinned out by the subsequent expansion, and they also will have larger annihilation cross-sections.

Before the addition of mass and further interaction terms, the action for any higgsion field is

$$\int d^4x \Phi_S^\dagger(x) D^\mu D_\mu \Phi_S(x) .$$  \hspace{1cm} (20)

(This follows from the invariance of the first term in (3) under charge conjugation.) The mass eigenstates will then have an action

$$S_S = \int d^4x \Phi_S^\dagger(x) (D^\mu D_\mu - m^2) \Phi_S(x)$$

$$= \int d^4x \mathcal{L}_S$$

$$\mathcal{L}_S = -\left[ D^\mu \Phi_S(x) \right]^\dagger D_\mu \Phi_S(x) - \Phi_S^\dagger(x) m^2 \Phi_S(x)$$  \hspace{1cm} (21)

(22)

(23)

(after integration by parts with the assumption that the boundary terms vanish).

Since the effects of a spinor rotation cancel in (21), we can redefine $\Phi_S$ to have spin 0. (If $\Phi_s^{(0)}$ is the value of the $\Phi_S$ before a rotation, and the value afterward is $U_{rot} \Phi_s^{(0)}$, with $U_{rot}^\dagger U_{rot} = 1$, (21) is unchanged if we replace $\Phi_S$ by $\Phi_S^{(0)}$, and then rename the fields with $\Phi_S^{(0)} \to \Phi_S$.) $\Phi_S$ is then a scalar boson field (with a novel multicomponent structure), and the theory has the usual Lorentz invariance.

Quantization follows the same prescription as for ordinary scalar fields: For a general complex multicomponent bosonic free field with components $\Phi_i$, we have

$$\pi_i(x) = \frac{\partial \mathcal{L}}{\partial (\partial_0 \Phi_i^\dagger(x))} = -\partial_0 \Phi_i^\dagger(x) = \partial_0 \Phi_i(x) \quad \pi_i^\dagger(x) = \frac{\partial \mathcal{L}}{\partial (\partial_0 \Phi_i(x))} = -\partial_0 \Phi_i^\dagger(x)$$

$$\pi_i(x) = \partial_0 \Phi_i(x)$$  \hspace{1cm} (24)

$$\pi_i^\dagger(x) = \partial_0 \Phi_i^\dagger(x)$$  \hspace{1cm} (25)

(according to our $(-+++)$ convention for the metric tensor), so, with implied summation over $i$,

$$\mathcal{H} = \pi_i(x) \partial_0 \Phi_i(x) + \pi_i^\dagger(x) \partial_0 \Phi_i^\dagger(x) - \mathcal{L}$$

$$= \partial_0 \Phi_i(x) \partial_0 \Phi_i^\dagger(x) + \partial_0 \Phi_i^\dagger(x) \partial_0 \Phi_i(x)$$

$$+ \Phi_i(x) m^2 \Phi_i(x)$$  \hspace{1cm} (26)

$$+ \Phi_i^\dagger(x) m^2 \Phi_i^\dagger(x) .$$  \hspace{1cm} (27)

In the treatment above, $\Phi$ is a classical field. In the treatment below, to avoid complicating the notation, $\Phi$ is taken to be the corresponding quantum field. It is required to satisfy the usual equal-time commutation relations:

$$[\Phi_i(x), \pi_j(x')] = i \delta_{ij} \delta(x - x')$$  \hspace{1cm} (28)

with $\pi_j(x', x^0) = \partial_0 \Phi_j^\dagger(x', x^0)$, and with the other relation just being the Hermitian conjugate.

To achieve this, we follow the usual procedure, first writing $\Phi$ in terms of the usual destruction and creation operators $c_{pi}$ and $d_{pi}^\dagger$, where $s$ distinguishes states with the same 4-momentum $p$:

$$\Phi(x) = \int \frac{d^3 p}{(2\pi)^3} \frac{1}{\sqrt{2p_0}} \sum_s (c_{pi}^s U_p \phi \phi^p \phi + d_{pi}^\dagger \phi \phi^p \phi^p - \phi \phi^p \phi^p \phi)$$

$$\phi \phi^p \phi$$  \hspace{1cm} (29)

with $p_0$ on-shell. The particle and antiparticle states are respectively $\sqrt{2p_0} c_{pi}^s |0\rangle$ and $\sqrt{2p_0} d_{pi}^\dagger |0\rangle$.

We require as usual

$[c_{pi}^s, c_{pi'}^{s'}] = \delta^{ss'}(2\pi)^3 \delta (\vec{p} - \vec{p}')$  \hspace{1cm} (30)

$[d_{pi}^s, d_{pi'}^{s'}] = \delta^{ss'}(2\pi)^3 \delta (\vec{p} - \vec{p}')$  \hspace{1cm} (31)

where $\vec{p}$ is the 3-momentum, with the other commutators equaling zero. The $U_p$ and $V_p$ are orthonormal vectors satisfying

$$\sum_p U_p^\dagger U_p = 1 \hspace{1cm} \sum_p V_p^\dagger V_p = 1$$

(32)

where $1$ is the identity matrix. The above properties then imply that

$$\pi_i (x', x^0), \partial_0 \Phi_i^\dagger (x', x^0)$$

$$= \int \frac{d^3 p}{(2\pi)^3} \epsilon^{ij} (\vec{p} - \vec{p}') \frac{i p_0}{2 p_0} \sum_s (U_{pi}^s U_{pj}^s + V_{pi}^s V_{pj}^s)$$

$$= i \delta_{ij} \delta (\vec{p} - \vec{p}') .$$  \hspace{1cm} (33)
The Hamiltonian is obtained from

\[ H_\Phi = \int d^3 x \mathcal{H}(x) \]
\[ = \int \frac{d^3 p}{(2\pi)^3} \frac{p_0^2 + p_0^0}{2p_0} \sum_{s} (c^*_p c^*_p + d^*_p d^*_p) \]
\[ = \int \frac{d^3 p}{(2\pi)^3} p_0 \sum_{s} (\pi^*_p + \pi^*_p + 1) \]
\[ = \frac{\sqrt{2}}{4} \sum_{s} (\pi^*_p + \pi^*_p + 1) \]
\[ = \frac{\sqrt{2}}{4} \sum_{s} (\pi^*_p + \pi^*_p + 1) \]
\[ = \frac{\sqrt{2}}{4} \sum_{s} (\pi^*_p + \pi^*_p + 1) \]
\[ \Phi = H + iH' \]

For each such real field the above treatment can be modified to give the same basic results with

\[ d^*_p = c^*_p , \quad V^*_p = U^*_p \]

and a factor of 1/2 in for each of the two independent fields. Each higgson particle \( H \) or \( H' \) is thus its own antiparticle.

Let us now specialize to the electroweak theory. The full gauge covariant derivative after symmetry-breaking is

\[ D_\mu = \partial_\mu - i \frac{g}{\sqrt{2}} (W^+ T^+ + W^- T^-) \]
\[ - i \frac{g}{\cos \theta_w} Z_\mu (T^3 - \sin^2 \theta_w Q) - i e A_\mu Q \]

where \( A_\mu \) is now the electromagnetic field. First consider the coupling of \( \Phi_S \) to \( Z \), which results from the Lagrangian

\[ \Phi_S \left( \partial^\mu - ig Z^\mu T^3 \right) (\partial_\mu - ig Z^\mu T^3) \Phi_S \]
\[ = \Phi_S \left( \partial^\mu \partial_\mu - 2ig Z^\mu T^3 \partial_\mu - \frac{1}{4} g^2 Z^\mu Z_\mu \right) \Phi_S \]

(since \( (T^3)^2 = 1/4 \), where \( g_Z = g/\cos \theta_w \). There are no terms involving \( \partial^\mu \partial_\mu \) because

\[ \partial^\mu Z_\mu = 0 , \quad \partial^\mu W^\pm = 0 \]

follows from the equations of motion for the massive vector boson fields \( Z_\mu \) and \( W^\pm_\mu \). The simplest first-order interaction relevant to a neutral particle is then

\[ \mathcal{L}_Z^\Phi = -2i g_Z Z^\mu \Phi_S T^3 \partial_\mu \Phi_S . \]

For each plane-wave state \(-i \partial_\mu \Phi_S = p_\mu \Phi_S \), so \( \Phi_S T^3 \Phi_S = 0 \) implies \( \mathcal{L}_Z^\Phi = 0 \). The first-order couplings involving \( W^\mu \pm \) also vanish:

\[ \mathcal{L}_W^\Phi = -i \sqrt{2} g \Phi_S \left( W^\mu + T^+ + W^\mu - T^- \right) \partial_\mu \Phi_S = 0 . \]

We must then turn to the second-order interactions,

\[ \mathcal{L}_Z^\Phi = -\frac{g_Z^2}{4} \Phi_S Z^\mu Z_\mu \Phi_S , \quad \mathcal{L}_W^\Phi = -\frac{g_Z^2}{2} \Phi_S W^\mu W^-_\mu \Phi_S \]

Fig. 1: Annihilation into real \( W \) or \( Z \) boson pair, for \( m_A > 90 \) GeV. These processes would produce a relic abundance of dark matter that is more than an order of magnitude too low to agree with the observed dark matter density.

We have performed approximate calculations of the annihilation cross-sections for the lowest-order processes involving the second-order couplings of \( (17) \), shown in Figures 1 and 2 using standard methods [29,30]. The only new (and trivial) feature is the sums over higgson states using [32]. For the processes of Fig. 1 we make the approximation that the \( W, Z, \) and \( H \) masses are nearly equal (~80-100 GeV.). We find that the total annihilation cross-section is more than an order of magnitude too large for \( m_Z > m_H > m_W \), and about a factor of 2 larger still for \( m_H > m_Z \). (Without this approximation, the cross-sections would be even larger.) I.e., we obtain

\[ \langle \sigma v \rangle_S \gg \langle \sigma v \rangle_S \quad \text{for} \quad m_H > m_W \approx 80 \text{ GeV} \]

where \( \langle \sigma v \rangle_S = \approx 2.2 \times 10^{-26} \text{ cm}^3/\text{s} \) is the benchmark value obtained by Steigman et al. [31] for a WIMP with mass
above 10 GeV that is its own antiparticle, if the relic dark matter density is to agree with astronomical observations.

For the processes of Fig. 2, we make the approximation of neglecting the masses of the fermions (which are all much smaller than the masses of the particles involved). Right panel: Example of annihilation into one real and one virtual W boson. For the quarks with sufficiently low masses. The total annihilation cross-section from these processes will be consistent with the observed relic abundance for $m_H \sim 72$ GeV.

$$\langle \sigma v \rangle \sim \langle \sigma v \rangle_S \quad \text{with} \quad m_H \sim 72 \text{ GeV}.$$ (49)

This annihilation cross-section is consistent with the limits set by observation of gamma-ray emissions from dwarf spheroidal galaxies by Fermi-LAT [32, 54], if the annihilation is treated generally rather than simplistically assumed to proceed through a single channel. For the present particle, there are 29 annihilation processes, represented by Fig. 2.

This cross-section and mass are also consistent with analyses of the gamma-ray excess from dwarf spheroidal galaxies by Fermi-LAT [32, 54], and with analyses of the antiprotons observed by AMS [39, 43], which independently have been interpreted as potential evidence of dark matter annihilation – although there are, of course, competing interpretations based on backgrounds and alternative statistical approaches. The inferred values of the particle mass and annihilation cross-section are in fact remarkably similar to those obtained here; see e.g. the abstracts of Refs. [37] and [42], and Fig. 12 of Ref. [43].

The predictions of the present theory within this context are very similar to those for the IDM if the masses of $A^0_1$ and $H^0_2$ are well separated from that of $H^0_1$. A very detailed analysis of the Fermi-LAT gamma-ray data, and its comparison with IDM predictions, has been given in [21]. For annihilations of a dark matter WIMP with a mass of $\sim 72$ GeV the basic qualitative conclusions are the same as above.

Collider detection of dark matter particles often focuses on creation through the $Z$, Higgs, or some hypothetical new mediator [44]. For the present particle there is no first-order coupling to the $Z$ and there are no exotic new mediators, so for this kind of process only the Higgs portal remains a possibility. CMS and ATLAS have independently placed upper limits on the branching ratio for invisible Higgs decays to particles with a total mass of < 125 GeV [45, 46]. The present particle may have a small Higgs coupling, however, and the total mass of a pair should be $\sim 145$ GeV. It appears that the present particle is also consistent with other collider-detection limits, and that the best possibility for observation in a collider experiment is the process depicted in Fig. 3. The predicted signature for collider detection is then $\gtrsim 145$ GeV of missing transverse energy resulting from vector boson fusion (VBF). In the present context this is a weak process, but the results of Refs. [23] and [24] (for the IDM), and [47] and [48] (for double Higgs production) suggest that observation of this process may be barely possible with a sustained run of the high-luminosity LHC, if an integrated luminosity of up to 3000 fb$^{-1}$ can be achieved. (Definitive studies may have to await a 100 TeV collider.) If there is a contribution from Higgs coupling, of course, the signal for creation of these particles will be stronger.

The best hope for direct detection appears to be the one-loop processes of Fig. 4, which are the same as for the IDM in Fig. 1 of [19]. The results in Figs. 2 and 3 of that paper suggest that this mechanism – scattering with an exchange of two Z or W bosons – has a cross-section not far below 10$^{-11}$ pb for the present particle – perhaps barely within reach of sustained runs for the next generation of direct-detection experiments. Definitive studies would probably require even greater sensitivity, but with a cross-section still potentially attainable (and above the neutrino floor). Of course, there is again the possibility of some enhancement from Higgs exchange.

The present scenario is quite amenable to being tested by experiment and observation. For example, if the
the present theory this boson represents the lowest-energy
particle in 1905), so one might anticipate further richness
discovery of the electron in 1897 and the proposal of the
spin 1/2 fermions and spin 1 gauge bosons has turned out
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We conclude with a broad comment: The behavior of
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