Collider Signatures of Goldstone Bosons

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Abstract

Recently Weinberg suggested that Goldstone bosons arising from the spontaneous breakdown of some global hidden symmetries can interact weakly in the early Universe and account for a fraction of the effective number of neutrino species $N_{\text{eff}}$, which has been reported persistently $1\sigma$ away from its expected value of three. In this work, we study in some details a number of experimental constraints on this interesting idea based on the simplest possibility of a global $U(1)$, as studied by Weinberg. We work out the decay branching ratios of the associated light scalar field $\sigma$ and suggest a possible collider signature at the Large Hadron Collider (LHC). In some corners of the parameter space, the scalar field $\sigma$ can decay into a pair of pions with a branching ratio of order $O(1)\%$ while the rest is mostly a pair of Goldstone bosons. The collider signature would be gluon fusion into the standard model Higgs boson $gg \to H$ or associated production with a $W$ gauge boson $qq' \to HW$, followed by $H \to \sigma\sigma \to (\pi\pi)(\alpha\alpha)$ where $\alpha$ is the Goldstone boson.
I. INTRODUCTION

Existence of Goldstone boson(s) is a manifestation of the spontaneous breakdown of some exact or nearly-exact global continuous symmetry in Nature [1]. Such Goldstone or pseudo-Goldstone bosons would be exactly or nearly massless. The well known example in the standard model (SM) is the pion which in the modern view can be interpreted as the Goldstone boson of spontaneous breakdown of the chiral $SU(2) \times SU(2)$ symmetry. Another logical possibility is the presence of a global hidden symmetry that the usual SM particles do not experience. The simplest choice is a global hidden $U(1)$ symmetry associated with a new quantum number $W$ of which all the hidden particles carry non-vanishing $W$ charges while all SM particles are neutral.

Weinberg [2] showed that such a simple extension to the SM could bring the Goldstone boson into weak interactions with the SM particles via a Higgs portal, $g(S^\dagger S)(\Phi^\dagger \Phi)$, where $S(x)$ is a complex singlet scalar field neutral under the SM symmetries with a nonzero $W$ quantum number, and $\Phi$ is the SM Higgs doublet with $W = 0$. Thus, the Goldstone bosons could remain in thermal equilibrium in the early Universe until they went out of equilibrium at a temperature above but not much above the muon mass. In this way, the Goldstone boson could contribute a fraction of 0.39 to the effective number $N_{\text{eff}}$ of neutrino species present in the era before recombination [2]. The requirement for this to happen is that the interactions of the Goldstone boson with the SM particles should be strong enough to bring it into thermal equilibrium and also weak enough such that it decouples close to the neutrino-decoupling temperature. The nature of derivative couplings of Goldstone bosons can easily satisfy this requirement.

There are a number of constraints on the model, namely on the Goldstone boson and the massive scalar $\sigma$ field associated with the Goldstone boson. Since they are weakly coupled to the Higgs boson, they would contribute to the invisible decay width of the Higgs boson [2, 3]. There are a number of other constraints in existing data [3], e.g., search for invisible particles in hadron decays, quarkonium decays, etc. In particular, here we point out that the invisible Higgs search at LEP-II will give the most stringent constraint on the mixing angle. The detail will be given in the next section.

In this work, besides working out the constraints on the model, we point out it may be possible to detect the $\sigma$ field and the Goldstone boson of the model at the LHC, via the
visible decay mode of the $\sigma$ field, namely $\sigma \to \pi\pi$, especially when the modulus of the field $S$ takes on a large vacuum expectation value (VEV). This is the main result of this work. We also estimate the event rates at the LHC-8 and LHC-14.

Studies of Goldstone bosons at the LHC in other context can be found in Ref. [4]. The related dark matter phenomenology has also been studied in Ref. [5].

II. THE MODEL

The model [2] is based on adding a complex singlet field $S$ to the SM Higgs doublet, through which the singlet field interacts with the SM particles. The renormalizable Lagrangian density is given by

$$
L = (\partial_\mu S^\dagger)(\partial^\mu S) + \mu^2 S^\dagger S - \lambda(S^\dagger S)^2 - g(S^\dagger S)(\Phi^\dagger \Phi) + L_{\text{sm}}
$$

where the Higgs sector in the $L_{\text{sm}}$ is

$$
L_{\text{sm}} \supset (D_\mu \Phi)^\dagger (D^\mu \Phi) + \mu_{\text{sm}}^2 \Phi^\dagger \Phi - \lambda_{\text{sm}}(\Phi^\dagger \Phi)^2.
$$

(1)

To respect the low energy theorem, we follow Ref. [2] to write $S$ in term of a radial field $r(x)$ and a Goldstone field $\alpha(x)$ as

$$
S(x) = \frac{1}{\sqrt{2}} (\langle r \rangle + r(x)) e^{i2\alpha(x)}
$$

(3)

in which the radial field develops a VEV $\langle r \rangle$ where the field $S$ is expanding around. Note that one can always set $\langle \alpha(x) \rangle = 0$ by field redefinition. The SM Higgs doublet field $\Phi$ is expanded about the VEV as

$$
\Phi(x) = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ \langle \phi \rangle + \phi(x) \end{pmatrix}
$$

(4)

in the unitary gauge, and $\langle \phi \rangle \approx 246$ GeV. Expanding the Lagrangian in Eq. (1) around the VEVs and replacing $\alpha(x) \to \alpha(x)/(2\langle r \rangle)$ in order to achieve a canonical kinetic term of the $\alpha(x)$ field describing a scalar of dimension 1, we obtain

$$
L \supset \frac{1}{2} (\partial_\mu r)(\partial^\mu r) + \frac{1}{2} \frac{(\langle r \rangle + r)^2}{\langle r \rangle^2} (\partial_\mu \alpha)(\partial^\mu \alpha) + \frac{\mu^2}{2} (\langle r \rangle + r)^2 - \frac{\lambda}{4} (\langle r \rangle + r)^4
$$

$$
+ \frac{1}{2} (\partial_\mu \phi)(\partial^\mu \phi) + \frac{\mu_{\text{sm}}^2}{2} (\langle \phi \rangle + \phi)^2 - \frac{\lambda_{\text{sm}}}{4} (\langle \phi \rangle + \phi)^4
$$

$$
- \frac{g}{4} (\langle r \rangle + r)^2 (\langle \phi \rangle + \phi)^2.
$$

(5)

1 We have normalized the kinetic energy term of a complex scalar field in the canonical form with the coefficient equals to 1.
Two tadpole conditions can be written down using \( \partial V/\partial r = 0 \) and \( \partial V/\partial \phi = 0 \) where \( V \) is the scalar potential part of Eq.(5):

\[
\langle \phi \rangle^2 = \frac{4\lambda_{\text{sm}}^2 - 2g\mu^2}{4\lambda_{\text{sm}} - g^2},
\]

\[
\langle r \rangle^2 = \frac{4\lambda_{\text{sm}}\mu^2 - 2g^2\mu_{\text{sm}}}{4\lambda_{\text{sm}} - g^2}.
\]

Taking the decoupling limit \( g \to 0 \) from the above equations, we recover the SM condition of \( \langle \phi \rangle^2 = \mu_{\text{sm}}^2/\lambda_{\text{sm}} \) as well as \( \langle r \rangle^2 = \mu^2/\lambda \).

The interaction fields \( r(x) \) and \( \phi(x) \) are no longer mass eigenstates because of the mixing term proportional to \( g \). The mass term is

\[
\mathcal{L}_m = -\frac{1}{2}(\phi(x) \ r(x))^\dagger \begin{pmatrix} 2\lambda_{\text{sm}}\langle \phi \rangle^2 & g\langle r \rangle \langle \phi \rangle \\ g\langle r \rangle \langle \phi \rangle & 2\lambda\langle r \rangle^2 \end{pmatrix} \begin{pmatrix} \phi(x) \\ r(x) \end{pmatrix}.
\]

We rotate \( (\phi(x) \ r(x))^T \) by an angle \( \theta \) into physical fields:

\[
\begin{pmatrix} H(x) \\ \sigma(x) \end{pmatrix} = \begin{pmatrix} \cos \theta & \sin \theta \\ -\sin \theta & \cos \theta \end{pmatrix} \begin{pmatrix} \phi(x) \\ r(x) \end{pmatrix}.
\]

The physical masses of the \( H(x) \) and \( \sigma(x) \), and the mixing angle are given by

\[
m_H^2 = 2\lambda_{\text{sm}}\langle \phi \rangle^2 \cos^2 \theta + 2\lambda\langle r \rangle^2 \sin^2 \theta + g\langle r \rangle \langle \phi \rangle \sin 2\theta,
\]

\[
m_\sigma^2 = 2\lambda\langle r \rangle^2 \cos^2 \theta + 2\lambda_{\text{sm}}\langle \phi \rangle^2 \sin^2 \theta - g\langle r \rangle \langle \phi \rangle \sin 2\theta,
\]

\[
\tan 2\theta = \frac{g\langle r \rangle \langle \phi \rangle}{\lambda_{\text{sm}}\langle \phi \rangle^2 - \lambda\langle r \rangle^2}.
\]

In the small \( \theta \) limit \( (\theta \lesssim 0.01 \text{ as will be shown later}) \),

\[
m_H^2 \approx 2\lambda_{\text{sm}}\langle \phi \rangle^2,
\]

\[
m_\sigma^2 \approx 2\lambda\langle r \rangle^2,
\]

\[
\theta \approx \frac{g\langle r \rangle \langle \phi \rangle}{m_H^2 - m_\sigma^2}.
\]

We can now write down the interactions terms in the limits of \( \theta \ll 1 \) and \( m_\sigma \ll m_H^2 \): \(^2\)

\[
\mathcal{L}_{H\alpha\alpha} = \frac{\theta}{\langle r \rangle} H (\partial_\mu \alpha)(\partial^\mu \alpha),
\]

\[
\mathcal{L}_{\sigma\alpha\alpha} = \frac{1}{\langle r \rangle} \sigma (\partial_\mu \alpha)(\partial^\mu \alpha),
\]

\[
\mathcal{L}_{H\sigma\sigma} = -\frac{g}{2} \langle \phi \rangle H \sigma^2.
\]

\(^2\) In the coupling of \( H\sigma\sigma \), the next-to-leading term in \( \theta \) is \( g\theta \langle r \rangle \), which is suppressed by a factor of \( (\langle r \rangle/\langle \phi \rangle)\theta \) relative to the leading term. However, when the ratio \( \langle r \rangle/\langle \phi \rangle \) is large, this next-to-leading term could be a sizable correction to the leading term.
III. CURRENT CONSTRAINTS ON THE $\sigma$ FIELD AND GOLDSTONE BOSON

Constraints from LEP searches for an invisibly decaying Higgs boson. The $\sigma$ field has a suggested mass of about 500 MeV and it can be produced via the mixing with the Higgs field [2]. Such a light scalar boson can be readily produced in hadron decays, quarkonium decays, as well as in the $Z$ boson decays, and at $e^+e^-$ collisions. The OPAL collaboration [6] searched for an invisibly decaying Higgs boson for the whole mass range from 1 GeV to 108 GeV at LEPII where the limit of the following ratio

$$\frac{\sigma(Zh)B(h \rightarrow \chi^0\chi^0)}{\sigma(ZH_{sm})}$$

was obtained. We can extrapolate the mass of the $h$ boson to be below 1 GeV, and the ratio being excluded is down to almost $10^{-4}$ (See Fig. 5 of [6]).

In the present context, the production cross section of $Z\sigma$ is

$$\sigma(Z\sigma) = \theta^2 \times \sigma(ZH_{sm})$$

Assuming the $\sigma$ field decays entirely into Goldstone bosons and disappears, we can constrain the mixing angle $\theta$

$$\theta \lesssim 10^{-2} \quad . \quad (13)$$

A less stringent constraint $\theta < 0.27$ has also been obtained recently in [3] from $B(\Upsilon \rightarrow \gamma E) < 2 \times 10^{-6}$ [7] via Wilczek mechanism [8] with the one-loop QCD correction [9].

Invisible width of the Higgs boson. The invisible decay of the Higgs boson goes through two processes:

$$H \rightarrow \alpha\alpha, \quad H \rightarrow \sigma\sigma \rightarrow 4\alpha$$

The decay partial widths are given by

$$\Gamma(H \rightarrow \alpha\alpha) = \frac{1}{32\pi} \frac{m_H^3}{\langle \phi \rangle^2 \langle r \rangle^2} \theta^2, \quad (14)$$

$$\Gamma(H \rightarrow \sigma\sigma) \approx \frac{1}{32\pi} \frac{m_H^3}{\langle \phi \rangle^2 \langle r \rangle^2} \theta^2, \quad (15)$$

in which we have assumed $m_{\sigma} \ll m_H$. Since the $\sigma$ field decays mostly into the Goldstone bosons (see the next section), we add both channels to obtain the invisible width of the Higgs boson [10]

$$\Gamma_{inv}(H) = \frac{2}{32\pi} \frac{m_H^3}{\langle \phi \rangle^2 \langle r \rangle^2} \theta^2. \quad (16)$$
One of the global fits to all the SM Higgs boson signal strength has constrained the non-standard decay width of the Higgs boson to be less than 1.2 MeV (branching ratio about 22%) at 95% CL [11]. The constraint on $\langle \phi \rangle / \langle r \rangle$ from the invisible width is similar to the constraint on $g$ obtained in Ref. [2]. Note that we also account for the decay of $H \rightarrow \sigma \sigma$ as a part of the invisible width of the Higgs boson. Numerically, from Eq.(16), we have

$$\theta \frac{\langle \phi \rangle}{\langle r \rangle} \leq 0.043$$

We use this constraint to rule out the parameter space in the plane of $(\theta, \langle \phi \rangle / \langle r \rangle)$, shown by the shaded region in Fig. 1.

![Parameter Space](image)

**FIG. 1.** Parameter space of $(\theta, \langle \phi \rangle / \langle r \rangle)$. The shaded region is ruled out by the invisible width of the Higgs boson to be less than 1.2 MeV. The condition for muon decoupling and the ratio $f$ are given in Eqs. (17) and (21), respectively, and are shown here for $m_\sigma = 500$ MeV. The upper limit of $\theta$ is taken to be 0.01 constrained by the search for invisibly decaying Higgs boson at LEP-II.

**Muon decoupling.** In the early Universe, the Goldstone bosons are at thermal equilibrium with the SM particles. As the Universe cooled down from its Hubble expansion, the Goldstone bosons would go out of the equilibrium since its weak interaction with the SM particles could no longer keep up with the Hubble expansion. It was argued in [2] that the
best scenario for the Goldstone bosons to go out of equilibrium is at a temperature still above the muon and electron masses but below all other masses of the SM. After decoupling, the Goldstone bosons were free and its temperature $T$ would then just fall off like the inverse of the Friedmann-Roberston-Walker scale factor $a$. Since the total cosmic entropy is conserved during the adiabatic expansion, after the muon annihilation, the constancy of $T a$ for the Goldstone bosons implies they behave like neutrino impostors contributing to the measured $\Delta N_{\text{eff}} = (4/7)(43/57)^{4/3} = 0.39 \ [2]$, which is consistent with the recent Planck result [12]. For this scenario to work, the annihilation rate of $\alpha \alpha \leftrightarrow \mu^+ \mu^-$ must be of the same order of the Hubble expansion rate at the temperature $k_B T \approx m_\mu$, i.e. [2]

$$
g^2 m_\mu^7 m_{\text{PL}}^4 \approx 1, \tag{17}
$$

where $m_{\text{PL}}$ is the Planck mass. From Eq. (11), one can express $g^2$ in terms of $\theta$, $\langle \phi \rangle / \langle r \rangle$ and $m_\sigma$. However, its dependence on $m_\sigma$ is rather weak for $m_\phi \gg m_\sigma$. This muon decoupling condition of Eq. (17) is shown in Fig. 1 for $m_\sigma = 500$ MeV. Note that this muon decoupling is not a constraint, but rather an interesting condition at the early universe for the Goldstone boson to explain $\Delta N_{\text{eff}}$.

IV. DECAY OF THE $\sigma$ FIELD

Because of the constraint from the Higgs invisible width and condition for muon decoupling in Eq. (17), the mass range of $\sigma$ cannot be much larger than $O(1)$ GeV [2]. We therefore show the mass range from 1 MeV to 1000 MeV for $\sigma$ from now on, and use 500 MeV when we need a typical value. The decay modes of such a light $\sigma$ field are very similar to those of a very light Higgs boson ($\lesssim 1$ GeV) [13]. The $\sigma$ can decay into a pair of electrons, muons, photons, pions and Goldstone bosons.

The formulas for the decays into $e^+ e^-$, $\mu^+ \mu^-$ and $\gamma \gamma$ are the same as the Higgs boson, up to a mixing angle. Thus, for the $f \bar{f}$ final state, we have

$$
\Gamma(\sigma \to f \bar{f}) = \theta^2 \frac{m_f^2 m_\sigma}{8\pi \langle \phi \rangle^2} \left[ 1 - \frac{4 m_f^2}{m_\sigma^2} \right]^{3/2}, \tag{18}
$$

in the small $\theta$ limit. For $m_\sigma < 1$ GeV the only possibility for fermionic decays are $f = e, \mu$. The decay width for $\sigma \to \gamma \gamma$ is the same as the one for the SM Higgs boson, up to $\theta^2$. We do not repeat the formula here except noting that the loop formulas for the light quarks are not trustworthy due to non-perturbative effects. However the $\sigma \to \gamma \gamma$ mode is not important.
Since the $\sigma$ field is very light and close to the hadronic scale $\Lambda_{QCD}$, its decay into two gluons is not applicable because of the non-perturbative hadronic effects, in contrast to the SM Higgs boson. The only hadrons that the $\sigma$ field can decay into is a pair of pions $\pi^+ \pi^-$ and $\pi^0 \pi^0$. The decay width of $\sigma \to \pi \pi$ summing over the isospin channels is given by [13]

$$\Gamma(\sigma \to \pi \pi) = \theta^2 \frac{1}{216 \pi} \frac{m_\sigma^3}{\langle \phi \rangle^2} \left( 1 - \frac{4 \sigma^2}{m_\sigma^2} \right)^{1/2} \left( 1 + \frac{11 \sigma^2}{2 m_\sigma^2} \right)^2. \quad (19)$$

The major difference is the decay mode $\sigma \to \alpha \alpha$ of the Goldstone boson. The partial width is

$$\Gamma(\sigma \to \alpha \alpha) = \frac{m_\sigma^3}{32 \pi \langle r \rangle^2}. \quad (20)$$

It is easy to see that the visible partial widths are all proportional to $\theta^2$ because the decay into visible particles is only possible via the mixing with the SM Higgs boson. On the other hand, the decay into a pair of Goldstone bosons is not suppressed by the mixing angle, but inversely proportional to the square of $\langle r \rangle$. We show the branching ratios of the $\sigma$ field for $\langle r \rangle = 1$ and 7 TeV in Fig. 2.

![Fig. 2](image-url)

**FIG. 2.** Decay branching ratios for the $\sigma$ field for (a) $\langle r \rangle = 1$ TeV and (b) $\langle r \rangle = 7$ TeV. The mode $\pi \pi$ includes $\pi^+ \pi^-$ and $\pi^0 \pi^0$. The mixing parameter $\theta$ is set at 0.01.

Mostly, the $\sigma$ field decays invisibly into the Goldstone bosons. It essentially adds to the invisible width of the SM Higgs boson, as the Higgs boson can decay via $H \to \sigma \sigma \to 4\alpha$ and $H \to \alpha \alpha$. However, when $\langle r \rangle$ goes to a very large value, say 7 TeV, as remarked already in [2], the decay of $\sigma \to \pi \pi$ can be as large as 2%.

We show the ratio

$$f \equiv \frac{\Gamma(\sigma \to \pi \pi)}{\Gamma(\sigma \to \alpha \alpha)} = \theta^2 \frac{4 \langle r \rangle^2}{27 \langle \phi \rangle^2} \left( 1 - \frac{4 \sigma^2}{m_\sigma^2} \right)^{1/2} \left( 1 + \frac{11 \sigma^2}{2 m_\sigma^2} \right)^2. \quad (21)$$
in Fig. 1 for $m_{\sigma} = 500$ MeV. In most part of the allowed region, the ratio $f$ is well below $10^{-4}$, thus mostly the $\sigma$ field decays into Goldstone boson. Nevertheless, if one goes to the corner where $\langle \sigma \rangle / \langle \pi \rangle$ is very small, we can achieve $f \approx 10^{-2}$. Such a value of $f$ would imply very interesting signatures for the $\sigma$ field and the Goldstone boson.

V. COLLIDER SIGNATURES

When the branching ratio $B(\sigma \to \pi \pi) \approx 2\%$, the collider signature would be very interesting. The dominant production of the $\sigma$ field is via the decay of the Higgs boson, followed by the decays of the two $\sigma$ fields. We can look for one $\sigma$ decaying invisibly into a pair of Goldstone bosons while the other one decays visibly into a pair of pions. Therefore, we expect

$$gg \to H \to \sigma\sigma \to (\pi\pi)(\alpha\alpha), \quad (22)$$

where the invariant mass of the pion pair is located right at $m_{\sigma}$. The signature would be a distinguished pion pair with $m_{\pi\pi} \approx m_{\sigma}$ plus a large missing energy carried away by the Goldstone bosons $\alpha$.

We perform a rough estimate of event rate here. The production cross section of the SM Higgs boson the LHC-8 is about 19 pb [14], and the non-standard decay branching ratio of the Higgs boson is limited to be less than about 20\% [11]. Therefore, using the analysis above we choose a currently allowed branching ratio of the Higgs boson:

$$B(H \to \sigma\sigma) \lesssim 10\%. \quad (23)$$

The cross section at the LHC-8 with $\langle r \rangle = 7$ TeV would be

$$\sigma(gg \to H) \times B(H \to \sigma\sigma) \times B(\sigma \to \pi\pi) \times B(\sigma \to \alpha\alpha) \times 2$$

$$\approx 19 \text{ pb} \times 0.1 \times 0.02 \times 0.97 \times 2 \approx 73 \text{ fb} . \quad (24)$$

For LHC-14, one should multiply the above number by a factor of 2.8.

Since the intermediate $\sigma$ boson is only $O(1)$ GeV, its decay products would be very collimated. The two $\alpha$’s become missing energies, while the two pions are very collimated, which appear to be a “microjet”, and experimentally it looks like a $\tau$ jet. The final state then consists of a microjet jet, which is made up of two pions, and a large missing energy. We first discuss the case when the two pions are charged pions. Ideally, we would like
to separate the two charged pions with an angular separation between them of order $\sim 2m_\sigma/p_{T_\sigma} = 1\text{ GeV}/60\text{ GeV} \approx 0.015$ which is rather small. Only the pixel detector inside the LHC experiments has some chances of separating them. The pixel tracker of the CMS detector [15] consists of three barrel layers with radii 4.4, 7.3 and 10.2 cm, and two endcap disks on each side of the barrel section. The spatial resolution ranges from 20 $\mu$m to about 100 $\mu$m, depending on the direction. Taking conservatively 100 $\mu$m as the spatial resolution and divide it by the average radius of the pixel detector, say 5 cm, we obtain an angular resolution of $2 \times 10^{-3}$ [3]. This is smaller than the average angular separation between the two charged pions estimated above by almost an order of magnitude. Thus it seems quite plausible to separate the two collimated charged pions. However, there is no guarantee that the pattern recognition algorithms would be able to reconstruct two distinct tracks, especially in the presence of large number of pile-up events. In the next phase of the CMS, another layer will be added to the pixel detectors at a radius of 16 cm [16]. The angular resolution will be further improved and the likelihood of separating the tracks of the two charged pions will be increased.

If the experiment cannot resolve the two charged pions, then the final state will look like a single jet consisting of some hadrons, plus missing energy. It is similar to signatures of many new models beyond the SM. In this case, one can make use of the associated production of the Higgs boson with a $W$ (or $Z$) boson, followed by the leptonic decay of the $W$ and the same decay mode of the Higgs boson:

$$pp \to WH \to (\ell\nu)(\sigma\sigma) \longrightarrow (\ell\nu)(\pi\pi + \alpha\alpha) .$$

The final state then consists of a charged lepton, a single jet of two unresolved charged pions, plus missing energy. The charged lepton is an efficient trigger of the events. The major SM background is the production of $W + 1$ jet, which could be orders of magnitude larger [17]. It presents an extreme challenge for experimentalists, although we may make use of the missing energy spectrum, because the signal also receives missing energy from $\sigma \to \alpha\alpha$ decay in addition to the neutrino from the $W$ decay. One may also use the feature of a microjet (similar to a $\tau$ jet) that is somewhat “thin” compared to the usual hadronic jet to separate the signal from backgrounds. In the case that one of the $\sigma$’s decays into

\[3\] With both outer trackers and pixel detectors, the resolution could be $2 - 10$ times better than $2 \times 10^{-3}$ for pions with $p_T > 10$ GeV.
two neutral pions, the process can give rise to 4 photons collimated as one “fat” photon. The final state would be a charged lepton, a “fat” photon plus missing energy, challenged by the major SM background of $W \gamma$ production which has a much larger cross section [18]. However, one can make use of the fact that the photon in the signal is “fat” to distinguish it from the background one.

Therefore, using both the gluon fusion and associated production with a $W$ we have provided more options to explore this model. However, in all situations that we studied above, they present great challenges to the experimentalists. Detailed detection simulations are needed in order to settle down if the proposed search is feasible or not. In the following, we will be contented by performing rough estimates on the signal cross section at the LHC-8 and LHC-14, so that experimentalists can have some ideas how large the signal cross section that one can obtain.

At $\langle r \rangle = 7$ TeV, the branching ratio of $\sigma$ into $\pi\pi$ is as large as 2%. For smaller $\langle r \rangle = 3 - 7$ TeV, the branching ratio into $\pi\pi$ ranges from about 0.4% to 2%, for which we may have enough cross sections for detection. We perform parton-level Monte Carlo simulations to estimate the event cross sections at the LHC-8 and LHC-14 for $\langle r \rangle = 3 - 7$ TeV. We normalize the uncut gluon fusion cross sections and the associated production cross sections to those given in the LHC Physics Web site [14]. For both the pions and charged lepton, we impose the same $p_T$ and rapidity cuts as $p_T > 30$ GeV and $|\eta| < 2.5$ respectively. We show the cross sections after the cuts in Table I. We have multiplied the cross sections by the branching ratios $B(H \rightarrow \sigma\sigma) \times B(\sigma \rightarrow \pi\pi) \times B(\sigma \rightarrow \alpha\alpha) \times 2$ to the Higgs boson decay, and $B(W \rightarrow \ell\nu) = 2/9$ to the $W$ boson decay. At the LHC-8 with about 20 fb$^{-1}$, the gluon fusion can produce a handful of events against the background if the two pions can be resolved. Nevertheless, if the pions cannot be resolved the associated production only has a cross section of order $O(0.05)$ fb, which may not be enough for detection. At the LHC-14 with a projected luminosity of $O(100)$ fb$^{-1}$, both the gluon fusion and associated production give sizable event rates whether or not the two pions can be resolved. Here, as mentioned above, the most important experimental issue is resolving the two pions. Although our rough estimate of angular separation by the pixel and tracking detectors indicates one may be able to resolve the pions, difficulties coming from the pile-up, pattern recognition, and track reconstruction post real challenges for our experimentalists. A proper detector simulation is called for before any realistic conclusion can be drawn.
TABLE I. Cross sections in fb for the gluon fusion process $pp \to H \to \sigma \sigma \to (\pi \pi)(\alpha \alpha)$ and the associated process $pp \to WH \to (\ell \nu)(\sigma \sigma) \to (\ell \nu)(\pi \pi + \alpha \alpha)$ at the LHC-8 and LHC-14 with the selection cuts described in the text. We choose $m_\sigma = 500$ MeV.

| $(r)$ (TeV) | $B(\sigma \to \pi \pi)$ | Cross Section (fb) LHC-8 | Cross Section (fb) LHC-14 |
|-----------|----------------|----------------|----------------|
|           | gluon fusion | $WH$ | gluon fusion | $WH$ |
| 3         | $3.72 \times 10^{-3}$ | 0.16 | 0.013 | 0.39 | 0.024 |
| 4         | $6.58 \times 10^{-3}$ | 0.27 | 0.022 | 0.68 | 0.043 |
| 5         | $1.02 \times 10^{-2}$ | 0.42 | 0.034 | 1.05 | 0.067 |
| 6         | $1.46 \times 10^{-2}$ | 0.60 | 0.049 | 1.50 | 0.095 |
| 7         | $1.97 \times 10^{-2}$ | 0.80 | 0.065 | 2.00 | 0.13 |

To summarize, the logical possibility of the existence of a hidden sector of Goldstone bosons masquerading as fractional cosmic neutrinos and communicate to our visible world through the Higgs portal as suggested recently by Weinberg [2] is explored further phenomenologically here. We have studied the constraints from the invisible Higgs search at LEP-II, the invisible Higgs width derived from global fittings using all the LHC signal strength data, and the condition of muon decoupling from evolution of our Universe. We also studied Higgs decays into a pair of $\sigma$ and its various decay modes. This interesting idea of Goldstone bosons as cosmic neutrino impostors can be tested by searching for the process of $gg \to H \to \sigma \sigma \to (\pi \pi)(\alpha \alpha)$ and the associated production $WH \to (\ell \nu)(\sigma \sigma) \to (\ell \nu)(\pi \pi + \alpha \alpha)$ at the LHC-8 and LHC-14.
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it will contribute to the invisible Higgs decay width as $\Gamma_{H \rightarrow \pi^-\pi^-} = \frac{1}{16\pi}(fg\langle \phi \rangle \langle r \rangle / (m_H^2 - m_\pi^2))^2(m_H^2 - 4m_\pi^2)^{3/2}$. The exponent $3/2$ reflects the P-wave feature of the decay amplitude.

We do not include this in our analysis since more free parameters are necessary to bring in.

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