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Weak-scale hidden sector and electroweak Q-balls

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Abstract

By extending the scalar sector of the Standard Model (SM) by a $U(1)$ singlet, we show that the electroweak symmetry breaking enables the formation of a stable, electrically neutral, colorless Q-ball which couples to the SM particle spectrum solely through the Higgs boson. This Q-ball has mainly weak and gravitational interactions, and behaves as a collection of weakly interacting massive particles. Therefore, it can be a candidate for the dark matter in the universe.

Classification of the total particle spectrum as those residing in the 'hidden sector' and those of the 'observable sector' has played a crucial role in high-scale supersymmetric theories for the implementation of the supersymmetry breaking at low energies. Indeed, the soft supersymmetry breaking terms in the low-energy globally supersymmetric Lagrangian are mediated by either gravity [1] (supergravity models) or messenger fields [2] (gauge-mediated supersymmetry breaking models) from a hidden sector containing heavy singlets. Similar ideas, with essential differences from the supersymmetric case, also have been proposed for the scalar sector of the Standard Model (SM) [3]. In the latter case one extends the scalar sector by a gauge singlet which interacts only with the SM Higgs doublet and, of course, gravity. Unlike the supersymmetric theories, in such models singlet field and the SM particle spectrum are not necessarily at diversely different mass scales; one can take both observable and hidden sectors around the same energy scale, that is, weak scale. The other important difference lies in the fact that the hidden sector in the SM Lagrangian is designed to account for the non-observation of the Higgs particle at the colliders as detailed in [4]. To utilize the large N expansion technique, and to increase the available invisible decay channels for the Higgs boson, in [3,4] use has been made of an $O(N)$ symmetric SM singlet. However, as long as one is not interested in the calculation of specific scattering processes involving the Higgs boson and the singlet the requirement of an $O(N)$ singlet can be relaxed as we shall do below.

Inspired from the idea of a hidden Higgs sector [3,4], below we extend the SM Lagrangian by a complex SM singlet scalar which interacts only with the usual Higgs doublet. Namely we extend the SM gauge group by an extra global Abelian group factor, $U(1)_s$, under which none of the SM particle spectrum is charged. Below we shall show that this is the simplest extension of the SM Higgs sector which accommodates the non-topological solitons [5,6], or in particular, Coleman's Q-balls [7]. Non-topological solitons are extended objects with finite mass and spatial extension, and arise in scalar field theories when there is an exact continuous symmetry and some kind of attractive interaction, as classified by Coleman [7]. In what follows we show that electroweak symmetry breaking produces the necessary interactions between the physical Higgs field and the singlet so that electrically neutral, colorless non-topological solitons naturally appear in the true electroweak vacuum.

After analyzing the properties of this extended scalar sector with subsequent discussion of the Q-balls it will be seen that: (1) In the electroweak vacuum the physical Higgs field and the singlet combine to form an absolutely stable Q-ball whose interactions with the SM particle spectrum is provided solely by the Higgs field. (2) Due to the extension of the scalar sector, Lagrangian necessarily obtains extra unknown parameters which, however, remain embedded in the expressions for the extensive parameters of the Q-ball, and do not have any direct effect on the interaction between the SM particle spectrum and the Q-ball. (3) Q-ball forms the state of minimal energy in the scalar sector of the SM, and thus, any scalar produced in the true vacuum via some collision process immediately escapes to the Q-ball implying that the collider search for the Higgs particle may not obtain a significant signal. (4) Photon, gluon and light fermions, due to their loop suppressed couplings to the Higgs particle, interact very weakly with the Q-ball compared to the massive electroweak bosons and heavy fermions. In this sense the Q-ball behaves as a stable collection of weakly interacting massive particles, and thus forms a candidate for the dark matter in the universe.

Below we work out the extended SM Higgs sector and show that the above-mentioned

points do naturally follow. The model Lagrangian is invariant under the SM gauge group $G_{SM} = SU(3)_c \times SU(2) \times U(1)_Y$ and an extra global Abelian group $U(1)_s$. Suppressing the contributions of the gauge bosons and fermions, the Lagrangian reads

$$\mathcal{L} = \frac{1}{2} \partial_\mu S^* \partial^\mu S + (\partial_\mu \Phi)^\dagger (\partial^\mu \Phi) - V(|S|, |\Phi|). \quad (1)$$

Here the potential $V(|S|, |\Phi|)$ is defined by

$$V(|S|, |\Phi|) = \frac{1}{2} m_s^2 |S|^2 + \frac{1}{4} \lambda_s |S|^4 + m^2 |\Phi|^2 + \lambda |\Phi|^2 - \kappa |S|^2 |\Phi|^2. \quad (2)$$

where Φ is the SM Higgs doublet and S is the SM- singlet. As the form of the potential suggests enlargement of the SM spectrum by a singlet brings about three new parameters m_s^2 , λ_s and κ . For the potential to be bounded from below in S and Φ directions it is necessary to have $\lambda_s > 0$, $\lambda > 0$, and $\sqrt{\lambda \lambda_s} - \kappa > 0$. In addition to these, for $U(1)_s$ symmetry to remain unbroken one needs $m_s^2 > 0$. The sign of κ is irrelevant for calculating the Higgs decay width in the true vacuum of the theory, however, for the formation of Q-balls it should be positive. In passing we note that for $\kappa \sim \lambda \sim \lambda_s$ contribution of the quartic terms resembles that of the D-terms in supersymmetric theories.

As usual, when m^2 is negative, Φ develops a non-zero vacuum expectation value $\langle \Phi \rangle = (v/\sqrt{2}, 0)$ where $v = \sqrt{-m^2/\lambda} = 246$ GeV. Around this vacuum expectation value Φ can be expanded as $\Phi = (1/\sqrt{2})(v + h + i\pi_Z, \pi_W^1 + i\pi_W^2)$ where π_Z ($\pi_W^{1,2}$) are the Golstone bosons swallowed by Z (W^\pm) to acquire a mass. Therefore, in this minimum of the potential the SM gauge symmetry is broken and the Lagrangian (1) takes the form

$$\mathcal{L} = \frac{1}{2} \partial_\mu S^* \partial^\mu S + \frac{1}{2} \partial_\mu h \partial^\mu h - V(|S|, h) \quad (3)$$

with

$$V(|S|, h) = \frac{1}{2} (m_s^2 - \kappa v^2) |S|^2 + \frac{1}{4} \lambda_s |S|^4 + \frac{1}{2} m_h^2 h^2 + \lambda v h^3 + \frac{\lambda}{4} h^4 - \kappa v h |S|^2 - \frac{1}{2} \kappa h^2 |S|^2, \quad (4)$$

where h is the physical Higgs boson with mass $m_h = \sqrt{-2m^2}$ as usual, and we subtracted the vacuum energy to make $V(0, 0) = 0$. As the form of $V(|S|, h)$ suggests, electroweak symmetry breaking modifies the mass of the singlet as $m_s^2 \rightarrow m_s^2 - \kappa v^2$, which must be positive to leave $U(1)_s$ unbroken. This can always be satisfied by choosing m_s^2 large enough.

There are some important points deserving discussion about the Lagrangian (3). Firstly, passage from (1) to (3) involves the electroweak symmetry breaking $G_{SM} \rightarrow SU(3)_c \times U(1)_{EM}$ so that experimentally well-established SM particle spectrum arises. During this transition neither the Higgs field, h , nor the other particles are affected by the presence of the singlet field, that is, electroweak phase transition proceeds independently of the interactions with the singlet. Secondly, the last two terms in (4) dictate the decay and scattering of the Higgs particle to invisible matter. Indeed these operators realize the processes $h \rightarrow SS$ and $hh \rightarrow SS$ whose signatures are, of course, outside the experimental detection. It is the $h \rightarrow SS$ decay [4] that is stressed in the literature to account for the non-observation of the Higgs particle at the colliders. Finally, both Lagrangians (1) and (3) are invariant

under $U(1)_s$ so that $U(1)_s$ is a global continuous symmetry of the theory before and after the electroweak phase transition.

There is a large amount of literature on non-topological solitons (see [5] and references therein). Although there are various types of such solitonic solutions [5,6] here we are interested in non-topological solitons of Q-ball type [7]. In recent years interest in the physical implications of Q-balls has accelerated after observing that they generically exist in supersymmetric theories [8,9]. In general, Q-balls exist in scalar field theories having a continuous symmetry and some kind of attractive interaction among the scalars. They are absolutely stable as long as the symmetry group is exact. Q-balls can arise from the self interactions of a single scalar field [7] as well as from the interactions among various scalar fields with different flavour [8]. In fact, demonstration of the existence of Q-ball type solutions for the Lagrangian (3) is nothing but a special case of the corresponding analyses in the supersymmetric models [8,9] which contain Higgs fields and scalar fermions. Other closely related works are [10,6] in which Q-ball formation in a model containing one real and one complex field was investigated. Although the Lagrangians employed in these works are similar to (3) the real scalar field there has nothing to do with the Higgs boson of the SM. Guided by the existing literature, below we present a short description of the Q-ball formation in the electroweak Lagrangian (3) referring the reader to [7,6,8,10] for details. $U(1)_s$, being an exact symmetry of the Lagrangians (1) and (3), has the conserved charge

$$Q = \frac{1}{2i} \int d^3\vec{x} S^*(\vec{x}, t) \overleftrightarrow{\partial}_t S(\vec{x}, t) . \quad (5)$$

Since Q vanishes identically for the trivial solution $S(\vec{x}, t) \equiv 0$, the field configuration that minimizes the energy

$$E = \int d^3\vec{x} \left\{ \frac{1}{2} |\partial_t S|^2 + \frac{1}{2} |\nabla S|^2 + \frac{1}{2} (\partial_t h)^2 + \frac{1}{2} (\nabla h)^2 + V(|S|, h) \right\} \quad (6)$$

for $Q \neq 0$, must have non-vanishing values in a finite domain. This constrained minimization can be accomplished by introducing a Lagrange multiplier ω and minimizing

$$\mathcal{E} = E + \omega(Q - \frac{1}{2i} \int d^3\vec{x} S^*(\vec{x}, t) \overleftrightarrow{\partial}_t S(\vec{x}, t)) \quad (7)$$

with respect to fields and ω , independently. A close inspection of \mathcal{E} shows that those terms having an explicit time dependence can be eliminated by requiring fields to rotate in the internal $U(1)_s$ space with angular velocities proportional to their $U(1)_s$ charges [7], that is, $S(\vec{x}, t) = e^{i\omega t} \bar{S}(\vec{x})$ and $h(\vec{x}, t) = \bar{h}(\vec{x})$. Here \bar{S} and \bar{h} are real and time-independent fields. With these redefinitions \mathcal{E} becomes

$$\mathcal{E} = \int d^3\vec{x} \left\{ \frac{1}{2} |\nabla \bar{S}|^2 + \frac{1}{2} (\nabla \bar{h})^2 + V_\omega(\bar{S}, \bar{h}) \right\} + \omega Q \quad (8)$$

where the effective potential V_ω is given by

$$V_\omega(\bar{S}, \bar{h}) = V(\bar{S}, \bar{h}) - \frac{1}{2} \omega^2 \bar{S}^2 . \quad (9)$$

Consequently, the requirement of a finite $U(1)_s$ charge leads one to a new potential V_ω which, unlike the original potential V which has its global minimum at $(S = 0, h = 0)$, can develop

a global minimum at some field configuration away from the origin because of the $1/2\omega^2\bar{S}^2$ term. In the true electroweak vacuum there is a perturbative particle spectrum consisting of, in addition to the usual SM spectrum, scalar bosons S with unit $U(1)_s$ charge and mass $\mu(0, 0) = (m_s^2 - \kappa v^2)^{1/2}$, where

$$\mu^2(\bar{S}, \bar{h}) \equiv \frac{2V(\bar{S}, \bar{h})}{\bar{S}^2}. \quad (10)$$

It is in this perturbative scheme that invisible decay rate of the Higgs particle has been computed [3,4]. However, the minimization of the energy for a finite $U(1)_s$ charge modifies the original potential as in V_ω and gives rise to appearance of new particles in the spectrum [7]. Indeed, if μ^2 is minimized for some field configuration $(\bar{S} = S_0 \neq 0, \bar{h} = h_0 \neq 0)$, that is,

$$\omega_0^2 \equiv \frac{2V(S_0, h_0)}{S_0^2} = \text{Min}[\mu^2(\bar{S}, \bar{h})] < \mu^2(0, 0) \quad (11)$$

then there exists non-dispersive solutions of the field equations ($\delta\mathcal{E}/\delta\bar{S} = 0$, $\delta\mathcal{E}/\delta\bar{h} = 0$) which are the absolute minima of the energy for fixed Q [7]. For $\omega = \omega_0$ the effective potential V_ω has two degenerate minima one at the origin $(\bar{S} = 0, \bar{h} = 0)$, and the other at $(\bar{S} = S_0 \neq 0, \bar{h} = h_0 \neq 0)$. To determine S_0 and h_0 it is convenient to introduce a polar coordinate system in two-dimensional field space by defining $\tan\theta \equiv \bar{h}/\bar{S}$ and $H \equiv \sqrt{\bar{S}^2 + \bar{h}^2}$ [8]. Expressing μ^2 (10) in terms of H and θ one obtains

$$\mu^2(H) = \frac{1}{\cos^2\theta}\{m_H^2 + A_H H + \frac{1}{2}\lambda_H H^2\} \quad (12)$$

where

$$\begin{aligned} m_H^2 &= (m_s^2 - \kappa v^2)\cos^2\theta + m_h^2\sin^2\theta \\ A_H &= (\lambda\sin^2\theta - \kappa\cos^2\theta)v\sin\theta \\ \lambda_H &= \lambda_s\cos^4\theta + \lambda\sin^4\theta - 2\kappa\sin^2\theta\cos^2\theta. \end{aligned} \quad (13)$$

It is easy to see that μ^2 is minimized for $H = H_0 \equiv -2A_H/\lambda_H$ which remains non-vanishing as long as $A_H \neq 0$. Moreover it is always positive if κ is sufficiently large, $\kappa \gtrsim \lambda\tan^2\theta$. This proves that there is a field configuration (S_0, h_0) away from the origin and minimizing the quantity μ^2 (12). Correspondingly, one has

$$\omega_0 = \frac{1}{\cos\theta}(m_H^2 - \frac{2A_H^2}{\lambda_H})^{1/2} \quad (14)$$

which is always real as long as $U(1)_s$ is an exact symmetry of the true electroweak vacuum, that is, $A_H^2 < 4\lambda_H m_H^2/9$. For $\omega = \omega_0$, the effective potential has two minima at $H = 0$ and $H = H_0$ between which it is maximized at $H = H_0/2$ with a value $A_H^4/4\lambda_H^3$. As ω takes values larger than ω_0 , the minimum at $H = H_0$ becomes the negative global minimum of the potential. The transition of the system from the original minimum $H = 0$ to the new global minimum proceeds through the quantum tunneling. The corresponding equations of motion can be identified with those of a bounce [11]. As long as $\omega_0 \leq \omega \leq \sqrt{V''(H = 0)}$ there is

always a bounce solution [13] having spherical symmetry $H(\vec{x}) = H(r)$, $r = \sqrt{\vec{x}^2}$ [12] and a localized nature $H(r) \rightarrow 0$ and $r \rightarrow \infty$. The resulting object is a lump of H matter with finite mass $M(Q) = \hat{\mu}(Q)Q$ where $\hat{\mu}(Q) < m_H$ always. Moreover, $\hat{\mu} \rightarrow \mu(S_0, h_0)$ for $Q \rightarrow \infty$ [7], and has a Q -dependent expression for smaller values of Q [13]. Irrespective of the detailed expressions for its extensive parameters, the Q-ball in question is a spherically symmetric object with a finite spatial extension characterized by its radius $R(Q)$, and a finite mass $M(Q)$ [11,7,13,8].

The effective potential V_{ω_0} depends only on λ_H and A_H , in particular, it is independent of m_H^2 . This m_H^2 independence of V_{ω_0} leaves m_s^2 largely free as it does not affect the possible transitions between the two degenerate minima. For the calculation of the Higgs invisible decay rate $\Gamma(h \rightarrow SS)$ in the true electroweak vacuum it is necessary to have $m_h^2 \geq 4(m_s^2 - \kappa v^2)$. However, formation of the electroweak Q-ball does not have such a strong requirement, it requires only the right hand side of this inequality be positive in order to leave $U(1)_s$ unbroken. Unlike the effective potential V_{ω_0} , however, ω_0 (14) has an explicit dependence on m_s^2 . The physically relevant interval for ω is $\omega_0 \leq \omega \leq m_H$ [13], and mass per unit charge $\hat{\mu}(Q)$ interpolates between the two extremes of ω . For $m_s^2 \gg v^2$, the allowed interval for ω shrinks to a narrow interval just below m_H . In this large m_s^2 limit the total charge Q necessarily becomes vanishingly small in which case the semiclassical analysis presented above can be invalidated. To avoid such an extreme (and also useless) case we assume that m_s^2 and v^2 are not at diversely different mass scales.

It is the non-vanishing of the trilinear coupling A_H that guarantees the existence of Q-balls in the electroweak vacuum. As is seen from (13) A_H is proportional to $\sin \theta$ and v . Vanishing of $\sin \theta$ means the neglect of the Higgs field in constructing the Q-ball solution. Similarly, vanishing of v means the absence of the electroweak phase transition. In both cases the electroweak Lagrangian (3) cannot accomodate a Q-ball solution. Hence, the existence of the Q-ball type solitons in the Lagrangian (3) is triggered essentially by the electroweak phase transition so we call them electroweak Q-balls from now on.

The analysis above proves that the electroweak Q-balls exist. To extract the necessary information about their relevance to the real world one should investigate its interactions with the environment. First we discuss the stability of the electroweak Q-ball. By construction, Q-ball is the state of minimal energy in the scalar sector of the theory. Hence, by energy conservation, it cannot decay to bosons, in particular, its constituents. More importantly, any scalar produced in the true vacuum by some scattering process rapidly escapes to the Q-ball as a statement of the Bose statistics. Importance of this statement for the Higgs phenomenology is that, even if the SM Higgs particle is produced by some future collider it can immediately escape to the Q-ball on its way to the detector. This observation comes by no surprise because such models were proposed to account for the non-observation of Higgs at collider searches by its large invisible decay rate [3,4]. After making these observations about the bosonic sector of the SM, it remains to discuss fermionic instability channels. First of all, one notes that there is no fermion (and also boson) in the SM particle spectrum which has nonzero $U(1)_s$ charge. This proves that there is no decay mode which can cause an erosion of the charge contained in the Q-matter. Since the charge remains unchanged always, one concludes that Q-ball is absolutely stable [7]. The electroweak Q-ball would evaporate if $U(1)_s$ symmetry were identified with $U(1)_{B,L}$ as in the supersymmetric theories

[8] or other models [15]. This very stability of the electroweak Q-ball implies that it can survive without dispersion for rather long time intervals at the cosmic scales.

To extract further information about the physical properties of the electroweak Q-ball it may be convenient to study the scattering of the observable particles from its bulk. One notes that the Q-matter in the electroweak Q-ball has two components; the SM Higgs \bar{h} and the singlet \bar{S} both having nonzero values over the spatial extension of the soliton. While the singlet provides the absolute stability of the electroweak Q-ball, the Higgs component is responsible for communication with the observable sector of the SM, that is, fermions and gauge bosons. The remarkable thing about the electroweak Q-ball is that any observable particle incident on it gets reflected through its interactions only with the Higgs boson, without feeling the presence of the singlet. The coupling between the SM particles and Higgs is known for every species [14]: all fermions generically couple as $gm_f/2M_W$ and the massive vector boson W (Z) as gM_W ($gM_Z/\cos\theta_W$). Obviously photon and gluon do not have tree-level couplings due to the electric neutrality and colorlessness of the Q-ball. Therefore, massless fermions, photon and gluon can have couplings only at the loop level through effective $h\bar{f}f$, $h\gamma\gamma$ and hgg vertices [14]. Apparently, during all these scattering events momentum conservation is provided by the emission of sound waves from the electroweak Q-ball [7]. Due to the asymptotic freedom, quarks and gluon cannot have isolated free-particle states and their coupling strengths to Higgs boson are relevant only for studying the scatterings of hadrons from the Q-ball. The gross observation about the interaction between electroweak Q-ball and the SM particle spectrum is that a typical scattering process

$$\text{particle} + \text{Q-ball} \rightarrow \text{particle} + \text{Q-ball+sound waves} \quad (15)$$

proceeds essentially with the weak interactions because electromagnetic and strong interactions can arise only at the loop level. Hence, the electroweak Q-ball has essentially weak and gravitational interactions, and from the view point of the SM particle spectrum, it is a weakly interacting stable lump of non-baryonic matter. These properties of the electroweak Q-ball reminds at once one of the missing mass in the universe, that is, the dark matter.

There is strong evidence from a variety of sources for a large amount of dark matter in the universe [16]. There is also extensive evidence that a substantial amount of the dark matter is non-baryonic. Models of galaxy formation classify the non-baryonic dark matter as hot and cold depending on if the constituents have relativistic or non-relativistic velocities, respectively. This classification can be reduced to the language of masses of the dark matter particles with a dividing line $m_{DM} \sim 1 \text{ keV}$. If the dark matter particles are their own anti-particles, and they are in thermal equilibrium with the radiation then their relic abundance is determined mainly by their annihilation cross section. The value of the annihilation cross section needed to make the relic abundance of the dark matter close to unity is remarkably close to one would expect for a weakly interacting massive particle (WIMP) with a mass $m_{DM} \sim M_Z$. The two best known and most studied cold dark matter candidates are neutralino [17] and axion [18] both qualifying to be WIMP's. If R-parity is conserved the neutralino with a mass in hundreds of GeV is a WIMP candidate. Similarly, axion arises in extensions of the SM to solve strong CP problem, and it has a rather small mass of the order of 10^{-5} eV . Recently, L-balls occurring in the scalar sectors of the supersymmetric theories are identified as dark matter candidates [19,20]. If the scalar potential is flat the L-ball can

be large enough to survive until the present time in spite of the evaporation to light leptons. Recalling the properties of the electroweak Q-ball one observes that it behaves as a collection of some Q WIMP's each with mass $M(Q)/Q = \hat{\mu}(Q)$. For consistency one needs $\hat{\mu}(Q) \sim M_Z$, which establishes another requirement for having m_s^2 at the weak scale. Unlike the L-balls of the supersymmetric models, the electroweak Q-ball does not suffer from evaporation to light fermions so it is a stable dark matter candidate. Guided by the analysis of [10] one would say that the small electroweak Q-balls can be produced copiously in the early universe during the electroweak phase transition, and they subsequently merge to form big Q-balls [19]. Due to their absolute stability they can survive until the present time and contribute to the total mass density of the universe in the form of dark matter [19,21]. It is the future astronomical observations about the dark matter that will specify its luminous, baryonic and non-baryonic proportions. After having such quantitative information about the dark matter that it will be possible to test the predictions of the electroweak Q-ball. However, at the present precision of the observations, dark matter is essentially non-luminous and non-baryonic so that the electroweak Q-ball can be a candidate to explain its existence in the universe.

As usual proton is absolutely stable because there is no $U(1)_{L,B}$ violating interactions in the SM Lagrangian. However, in the presence of the SM singlet one can introduce a non-renormalizable interaction of the form

$$\Delta\mathcal{L} = \kappa' \frac{|S|^2}{M_X} \bar{Q}L + h.c. \quad (16)$$

where $M_X \gg m_s$ is a large mass scale, κ' is a Yukawa coupling, and Q and L are light quark and lepton with $m_Q > m_L$. $\Delta\mathcal{L}$ breaks the B- and L-symmetries of the SM Lagrangian but preserves the $U(1)_s$ symmetry. In the true electroweak vacuum $Q \rightarrow L$ transition occurs only at the two loop level and its rate is small [22]. However, inside the electroweak Q-ball S has a non-vanishing VEV and $Q \rightarrow L$ type transitions can occur at the tree level as illustrated by (15). More interestingly, nucleon scattering from the Q-ball can realize $Q \rightarrow L$ transition with a probability $\sim \kappa'^2 \frac{\bar{S}^4}{m_Q^2 M_X^2}$ [22]. The ultimate existence of such L- and B-violating interactions will be tested with future experiments on the dark matter.

With the ending of LEP2 period without a signal, search for the Higgs boson will continue at the LHC [23]. In near future, the LHC will be searching for the Higgs signal in the mass range $M_Z \lesssim m_h \lesssim 2M_Z$ expecting to observe the Higgs resonance in gluon-gluon fusion to photon pairs. In this search strategy observation of the Higgs resonance is essential to extract the Higgs signal from the large irreducible background. In future, if the LHC fails to find a Higgs signal, possibility of a weak-scale hidden sector will be strengthened [3,4,24]. Besides the Higgs discovery potential of the LHC, there are strong theoretical arguments stating that a linear collider working at the center-of-mass energy $\sqrt{s} = 500$ GeV with an integrated annual luminosity of 500 fb^{-1} , such as the TESLA collider [26], will definitely find a signal for the Higgs boson [25] independent of the complexity of the Higgs sector and Higgs boson decay modes. Hence, the no-lose theorem of [25], armed with the recently proposed TESLA collider [26], forms a testing ground for the existence of a perturbative Higgs sector. If the results of [25] cannot be confirmed at a future linear collider then the present status of the symmetry breaking sector of the SM will be questionable, and a possibility of having a hidden electroweak sector will increase.

In this work we investigated the implications of a weak-scale hidden sector for dark matter searches, proton-to-lepton transitions, and Higgs phenomenology at the future colliders. The main object of the discussion has been the electroweak Q-ball, a non-topological, extended, absolutely stable object having mainly weak and gravitational interactions. The formation of the electroweak Q-ball is triggered by the electroweak phase transition. The indispensable component of the analysis, that is, the SM singlet has a mass at the electroweak scale, as required by both the Q-ball formation and dark matter phenomenology. If the future collider search fails to find an observable signal for the Higgs boson, possibility of a weak-scale hidden sector, like the one discussed above, will increase. As discussed in the text, supersymmetric theories generically and naturally accomodate Q-balls corresponding to the exact global symmetries, $U(1)_{B,L}$, of the low-energy theory. However, it is known that these Q-balls necessarily evaporate by emitting lepton or baryon number from their surface. Thus, despite the existence of Higgs-sfermion trilinear couplings in the Lagrangian, only in the case of flat potentials that one can construct big enough Q-balls that can survive until the present time after being formed in the early universe. Unlike the $U(1)_{B,L}$ symmetries leading naturally to Q-balls in supersymmetric theories, the electroweak Q-ball is constructed by postulating an extra $U(1)$ whose origin is unknown. However, it is with this $U(1)_s$ group that the electroweak Q-ball does not suffer from evaporation, and is an absolutely stable object [$U(1)_s$ has no relation to $U(1)_{B,L}$]. Despite the unknown nature of the singlet and its couplings, the electroweak Q-ball interacts with the visible matter only through the Higgs field, and couplings of the SM particle spectrum to Higgs are already known. Therefore, existence of the electroweak Q-balls can be directly tested with the future astronomical observations on the dark matter. Indeed, the experimental data concerning the luminous, baryonic and non-baryonic proportions of the dark matter should be derivable from the coupling strengths of the observable matter to the Q-ball. For Q-ball formation, the SM Higgs sector is not the only possiblity, in fact, a similar analysis can be carried out for two-doublet models in which there are two CP-even and a CP-odd scalar Higgs bosons which can contribute to the associated Q-matter.

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REFERENCES

- [1] S. Soni and A. Weldon, Phys. Lett. **B126**, 215 (1983); G. F. Giudice and A. Masiero, Phys. Lett. **B206**, 4480 (1988).
- [2] M. Dine, W. Fischler, and M. Srednicki, Nucl. Phys. **B189**, 575 (1981); L. Alvarez-Gaume, M. Claudson, and M. B. Wise, Nucl. Phys. **B207**, 96 (1982).
- [3] J. D. Bjorken, Int. J. Mod. Phys. **A7**, 4189 (1992); R. S. Chivukula, M. Golden, D. Kominis, and M. V. Ramana, Phys. Lett. **B293**, 400 (1992); R. S. Chivukula and M. Golden, Phys. Lett. **B267**, 233 (1991).
- [4] T. Binoth and J. J. van der Bij, Z. Phys. **C75**, 17 (1997).
- [5] T. D. Lee and Y. Pang, Phys. Rep. **221**, 251 (1992).
- [6] R. Freidberg, T. D. Lee, A. Serlin, Phys. Rev. **D13**, 2739 (1976).
- [7] S. Coleman, Nucl. Phys. **B262**, 263 (1985).
- [8] A. Kusenko, Phys. Lett. **B405**, 108 (1997).
- [9] K. Enqvist and J. McDonald, Phys. Lett. **B425** (1998)309.
- [10] K. Griest, E. W. Kolb, and A. Massarotti, Phys. Rev. **D40**, 3529 (1989).
- [11] S. Coleman, Phys. Rev. **D15**, 2929 (1977); C. G. Callan, and S. Coleman, Phys. Rev. **D16**, 1762 (1977).
- [12] S. Coleman, V. Glaser, and A. Martin, Comm. Math. Phys. **58**, 211 (1978).
- [13] A. Kusenko, Phys. Lett. **B404**, 285 (1997).
- [14] J. F. Gunion, H. E. Haber, G. Kane, and S. Dawson, *The Higgs Hunter's Guide*, Addison-Wesley, New York, (1990).
- [15] A. Cohen, S. Coleman, H. Georgi, and A. Manohar, Nucl. Phys. **B272**, 301 (1986).
- [16] *Particle Physics and Cosmology*, M. Srednicki, ed., North-Holland, Amsterdam, (1990).
- [17] G. Jungman, M. Kamionkowski, and K. Greist, Phys. Rep. **267**, 196 (1995).
- [18] M. S. Turner, Phys. Rep. **197**, 67 (1990).
- [19] A. Kusenko and M. Shaposhnikov, Phys. Lett. **B418**, 46 (1998).
- [20] K. Enqvist and J. McDonald, hep-ph/9807269.
- [21] K. Greist and E. W. Kolb, Phys. Rev. **D40**, 3231 (1989); J. A. Freiman, A. V. Olinto, M. Gleiser, and C. Alcock, Phys. Rev. **D40**, 3241 (1989).
- [22] G. Dvali, A. Kusenko, and M. Shaposhnikov, Phys. Lett. **B417**, 99 (1998).
- [23] N. V. Krasnikov and V. A. Matveev, Phys. Part. Nucl. **28**, 441 (1997).
- [24] A. Ghinculov, T. Binoth, and J. J. van der Bij, Phys. Lett. **B427**, 343 (1998).
- [25] J. R. Espinosa and J. F. Gunion, hep-ph/9807275.
- [26] R. Brinkman, TESLA Collaboration, DESY-M-97-04, (1997); Turk. J. Phys. **22**, 661 (1998).