PHASE TRANSITIONS IN THE CORE OF GLOBAL EMBEDDED DEFECTS

Minos Axenides
Institute of Nuclear Physics, N.C.R.P.S. Demokritos
153 10, Athens, Greece

Leandros Perivolaropoulos
Department of Physics
University of Crete
71003 Heraklion, Greece and

Mark Trodden
Particle Astrophysics Theory Group
Department of Physics
Case Western Reserve University
Cleveland, OH 44106-7079

Abstract

We demonstrate the existence of global monopole and vortex configurations whose core exhibits a phase structure. We determine the critical values of parameters for which the transition from the symmetric to the non-symmetric phase occurs and discuss the novel dynamics implied by the non-symmetric cores for defect interactions. We model phase transitions in the core of global embedded topological defects by identifying the relevant parameters with the vacuum expectation value of a dynamical scalar field. Finally, we argue that superheavy defects that undergo a core phase transition in the very early universe provide a novel realization for topological inflation.

\[ \text{E-mail address: axenides@gr3801.nrcps.ariadne-t.gr} \]
\[ \text{E-mail address: leandros@physics.uch.gr} \]
\[ \text{E-mail address: trodden@theory1.phys.cwru.edu} \]
1 Introduction

The existence of topological defects [1, 2] with a non-trivial core phase structure has recently been demonstrated for embedded global domain walls and vortices [3] (see also [4] for earlier studies). The analogous phenomenon has also been found to occur in experiments performed in superfluid $^3$He (see e.g. Ref. [5]). In that case, a phase transition in the core of superfluid $^3$He – $^3$He $^B$ vortices was observed to correspond to different spin-orbital states of their core. In the case of the field theory, the partial breaking of global symmetries [3] deforms the vacuum manifold in such a way as to reduce the dimensionality of the non-contractible configurations (points, circles, spheres etc.) that are otherwise admitted, giving rise to ”embedded” topological defects (domain walls-vortices). In these defects a core with a nontrivial symmetry phase structure generically appears as a consequence of the embedding. This structure is associated with either the zero value of a scalar field in the core (symmetric phase) or a nonzero value (broken phase) for different regions of the enlarged parameter space that the vacuum manifold deformation affords. Such embedded defects may be viewed [3] as interpolating between “texture” type defects [6, 7, 8], where the field boundary conditions are uniform, and symmetric core defects [4, 10], where the topological charge is due to the field variation at the boundaries. Superconducting strings [11] are a special case of this type of defect, in which the core acquires a nonzero vacuum expectation value (VEV) for a particular range of parameters. The existence of a non-symmetric core implies a large variety of new properties and phenomena for such topological defects. Some of these effects were in fact discussed in Ref. [3], where it was shown, for example, that in the range of parameters for which a non-symmetric core is favored, an initially spherical domain wall bubble is unstable towards planar collapse. That study focused on the stability and properties of domain walls and vortices with a non-symmetric core.

In the present work we complete the list of such defects by presenting a model for an “embedded” global monopole [12]. Moreover we carefully study the implications of non-symmetric cores for defect interactions such as those between vortices and those between monopoles.

In addition, we investigate the stability of new types of core structures with non-trivial winding in the vortex core describing vortices within vortices. However, as we show, these structures are unstable. In the next
section, we construct an $O(4)$ model that admits global monopoles with a core exhibiting both symmetric and nonsymmetric false vacuum phases for appropriate values of the free parameters, and identify their range of values for which this core structure is stable. We also analyze the vortex-vortex and monopole-monopole interactions for defects with non-symmetric cores. We show that the non-trivial core structure induces an additional interaction potential which can be attractive or repulsive depending on the relative orientation of the scalar field value in the defect cores.

An interesting implication of our results is that the transition from a symmetric to the non-symmetric false vacuum in the defect core can be viewed as a genuine phase transition if the parameters of our models are rendered new scalar external degrees of freedom with their own dynamics. Symmetric or non-symmetric cores correspond to the zero or nonzero value of a scalar field condensate developing in the defect core. In a topological inflation scenario, an inflating string core enters a re-heating period if a non-zero VEV develops in its core, rapidly decreasing the vacuum energy there. This possibility is studied in section 4.

## 2 Vortex Core Dynamics

Before discussing the monopole case, we briefly review and extend the case of the vortex with non-trivial core structure first discussed in Ref. [3]. Consider the following Lagrangian density describing the symmetry breaking $SU(2) \rightarrow U(1) \rightarrow I$, where the first breaking is explicit and the second is spontaneous:

$$\mathcal{L} = \frac{1}{2} \partial_\mu \Phi^\dagger \partial^\mu \Phi + \frac{M^2}{2} \Phi^\dagger \Phi + \frac{m^2}{2} \Phi^\dagger \tau_3 \Phi - \frac{h}{4} (\Phi^\dagger \Phi)^2.$$  \hspace{1cm} (1)

Here, $\Phi = (\Phi_1, \Phi_2)$ is a complex scalar doublet and $\tau_3$ is the $2 \times 2$ Pauli matrix. Using the rescaling

$$\Phi \rightarrow \frac{M}{\sqrt{h}} \Phi \hspace{1cm} (2)$$
$$x \rightarrow \frac{1}{M} x \hspace{1cm} (3)$$
$$m \rightarrow \alpha M \hspace{1cm} (4)$$
the potential reduces to

\[ V(\Phi) = -\frac{M^4}{2\hbar}((\Phi^\dagger\Phi) + \alpha^2(|\Phi_1|^2 - |\Phi_2|^2) - \frac{1}{2}(\Phi^\dagger\Phi)^2) . \]  

(5)

The ansatz

\[ \Phi = \begin{pmatrix} \Phi_1 \\ \Phi_2 \end{pmatrix} = \begin{pmatrix} f(\rho)e^{im\theta} \\ g(\rho) \end{pmatrix} , \]  

(6)

with boundary conditions

\[ \lim_{\rho\to0} f(\rho) = 0, \quad \lim_{\rho\to0} g'(\rho) = 0 \]  

(7)

\[ \lim_{\rho\to\infty} f(\rho) = (\alpha^2 + 1)^{1/2}, \quad \lim_{\rho\to\infty} g(\rho) = 0 , \]  

(8)

leads to the field equations

\[ f'' + \frac{f'}{\rho} - \frac{m^2}{\rho^2} f + (1 + \alpha^2)f - (f^2 + g^2)f = 0 \]  

(9)

\[ g'' + \frac{g'}{\rho} + (1 - \alpha^2)g - (f^2 + g^2)g = 0 . \]  

(10)

It is straightforward \[3\] to use a relaxation algorithm to numerically solve the system (9,10) with the boundary conditions (7,8). For \( \alpha > \alpha_{cr} \approx 0.37 \) the solution relaxes to a form that is symmetric at the defect core \( (g(0) = 0) \), while for \( \alpha < \alpha_{cr} \) a non-zero VEV at the core is energetically favored. As a check of the numerical results, we have also verified the stability of the symmetric \( (g(\rho) = 0) \) solution by solving the linearized eigenvalue problem of equation (10) around \( g(\rho) = 0 \) and showing that there are no negative modes for \( \alpha > \alpha_{cr} \).

A more interesting core structure is obtained through the ansatz

\[ \Phi = \begin{pmatrix} \Phi_1 \\ \Phi_2 \end{pmatrix} = \begin{pmatrix} f(\rho)e^{im\theta} \\ g(\rho)e^{in\phi} \end{pmatrix} . \]  

(11)

The resulting configuration corresponds to the formation of an additional \( \Phi_2 \) vortex localized within the core of the \( \Phi_1 \) vortex. This ‘vortex within vortex’ configuration can be metastable in the parameter region where a non-zero VEV of the \( \Phi_2 \) component is favored in the core of the \( \Phi_1 \) vortex. Clearly,
since the true vacuum outside the vortex has $\Phi_2 = 0$, the winding can only persist within the core of the $\Phi_1$ vortex. The ansatz (11) leads to the same field equation (9) for $f(\rho)$, while the equation for $g(\rho)$ becomes
\begin{equation}
 g'' + \frac{g'}{\rho} - \frac{n^2}{\rho^2}g^2 + (1 - \alpha^2)g - (f^2 + g^2)g = 0 , \tag{12}
\end{equation}
with boundary conditions
\begin{align}
 \lim_{\rho \to 0} f(\rho) &= 0, & \lim_{\rho \to 0} g(\rho) &= 0 \tag{13} \\
 \lim_{\rho \to \infty} f(\rho) &= (\alpha^2 + 1)^{1/2}, & \lim_{\rho \to \infty} g(\rho) &= 0 . \tag{14}
\end{align}
Using a relaxation method\cite{14} again, we find that for all $\alpha > 0$, the system (9), (12) relaxes to a solution with $g(\rho) = 0$. Thus, a stable non-trivial winding does not develop within the vortex core. We have verified this result by examining the field equation solution
\begin{equation}
 \Phi = \begin{pmatrix} \Phi_1 \\ \Phi_2 \end{pmatrix} = \begin{pmatrix} f(\rho)e^{im\theta} \\ 0 \end{pmatrix} \tag{15}
\end{equation}
for unstable modes towards the configuration (11). This amounts to solving the linearized eigenvalue problem
\begin{equation}
 -g'' - \frac{g'}{\rho} + \frac{n^2}{\rho^2}g^2 - (1 - \alpha^2)g + f^2g = \omega^2g , \tag{16}
\end{equation}
where $f$ satisfies equation (9) with $g = 0$, and looking for negative eigenvalues $\omega^2$. Using a shooting method \cite{14}, with initial conditions $g(0) = 0$ and $g'(0) = 1$ (required for non-trivial winding), we showed that, for nonzero winding, $n$, $g$ has no instability for any $\alpha$. The stability of such ‘vortices within vortices’ however, is expected to improve by introducing a $U(1)$ gauge field. This effectively screens the repulsive centrifugal barrier $\frac{n^2}{\rho^2}g^2$ which prohibits the development of a negative mode and drives the system towards a ‘vortex within vortex’ configuration. A detailed general study of defects within defects is currently in progress\cite{15}.

The existence of non-trivial structures within the core of defects has interesting implications for their interactions. For example, consider a vortex produced by an explicit breaking of $O(3)$ symmetry to $O(2)$, with subsequent spontaneous breaking to $I$. The corresponding vortex ansatz is of the form
\[ \Phi = \begin{pmatrix} \Phi_1 \\ \Phi_2 \\ \Phi_3 \end{pmatrix} = \begin{pmatrix} f(\rho) \cos \phi \\ \pm f(\rho) \sin \phi \\ g(\rho) \end{pmatrix}, \]  \hspace{1cm} (17)

with boundary conditions

\[ \lim_{\rho \to 0} f(\rho) = 0, \quad \lim_{\rho \to 0} g(\rho) = \pm c \]  \hspace{1cm} (18)
\[ \lim_{\rho \to \infty} f(\rho) = \text{const}, \quad \lim_{\rho \to \infty} g(\rho) = 0. \]  \hspace{1cm} (19)

Here, the + (−) (in 17) corresponds to a vortex (antivortex), and c is determined dynamically and depends on the parameters of the Lagrangian.

Consider now a field configuration corresponding to a pair of vortices at a large distance \( l \) from each other, and let \( \rho_1, \rho_2 \) be the distances of any point from the centers of the two vortices respectively. The corresponding field configuration is of the form

\[ \Phi = \begin{pmatrix} \Phi_1 \\ \Phi_2 \\ \Phi_3 \end{pmatrix} = \begin{pmatrix} f(\rho_1)f(\rho_2) \cos(\phi_1 \pm \phi_2) \\ f(\rho_1)f(\rho_2) \sin(\phi_1 \pm \phi_2) \\ g(\rho_1, \rho_2) \end{pmatrix}, \]  \hspace{1cm} (20)

with \(|g(0, l)| = |g(l, 0)| = c\) and the +, (−) corresponds to a vortex-vortex (vortex-antivortex) pair. Consider first the case of a vortex-antivortex pair with \( g_2(0, l) = -g_2(l, 0) = \pm c \). Assuming a linear interpolation of \( g_2 \) between \( g_2(0, l) \) and \( g_2(l, 0) \), and substituting the total energy of the sum of two isolated vortices for the total energy of the interacting configuration (see also Ref. \[16\]) we find, for large separations \( l \), the interaction energy

\[ E^{V-V}_{\text{int}} \simeq 2\pi \eta (\log(\eta l) - \log(\eta L)) + \frac{A_2 c^2}{l}. \]  \hspace{1cm} (21)

Here, \( \eta \) is the scale of symmetry breaking, \( A \) is a positive constant with dimensions of length squared and \( L \) is a cutoff length related to the size of the system. We have also assumed that the cores remain undistorted during the interaction. The form of \( E^{V-V}_{\text{int}} \) implies that the attractive vortex-antivortex force could be balanced by the repulsive force originating from the gradient energy of the \( g \) component at a critical distance

\[ l_{cr} = \frac{A c^2}{\eta \pi}. \]  \hspace{1cm} (22)
This however cannot lead to a vortex-anti-vortex static bound state, because, according to Derrick’s theorem, no static, finite energy scalar field configuration exists in two dimensions with non-zero potential energy. This means that the above expression for the interaction energy breaks down at distances larger than $l_{cr}$.

A similar interaction term can be obtained by a vortex-vortex pair with $g_2(l,0) = g_2(0,l) = \pm c$. In this case it is straightforward to show that $E_{int-v} \simeq -E_{int-v}$ and there is an attractive gradient-reducing term balancing the standard repulsive vortex-vortex interaction at the same critical distance given by (22). We have shown however [15], that despite the fact that this configuration has infinite energy and Derrick’s theorem cannot be applied, defect bound states do not exist in this case either. The detailed investigation of these issues, with and without gauge fields is currently in progress [15] using numerical simulations of defect evolution.

3 Monopole Core Dynamics

The above study of the core structure and interactions of vortices can easily be extended to monopoles. As a concrete example we consider a model with $O(4)$ symmetry explicitly broken to $O(3)$, which is subsequently broken spontaneously to $O(2)$. This is described by the Lagrangian

$$\mathcal{L} = \frac{1}{2} \partial_{\mu} \Phi^* \partial^{\mu} \Phi + \frac{M^2}{2} \Phi^* \Phi + \frac{m^2}{2} (|\Phi_1|^2 + |\Phi_2|^2 + |\Phi_3|^2) - \frac{m^2}{2} |\Phi_4|^2 - \frac{h}{4} (\Phi^* \Phi)^2,$$

(23)

with the monopole ansatz

$$\Phi = \begin{pmatrix} \Phi_1 \\ \Phi_2 \\ \Phi_3 \\ \Phi_4 \end{pmatrix} = \begin{pmatrix} f(r) \cos \theta \sin \phi \\ f(r) \cos \theta \cos \phi \\ f(r) \sin \theta \\ g(r) \end{pmatrix}.$$  

(24)

After a rescaling similar to the one of equations (2,3,4), we obtain the field equations for $f(r)$ and $g(r)$ as

$$f'' + \frac{2f'}{r} - \frac{2f}{r^2} + (1 + \alpha^2) - (f^2 + g^2)f = 0$$

(25)
Figure 1: Field configuration for a symmetric-core global monopole with $\alpha = 0.3$.

$$g'' + \frac{2g'}{r} + (1 - \alpha^2)g - (f^2 + g^2)g = 0.$$  \hspace{1cm} (26)

Using the same methods as for vortices, we find that the development of a non-zero VEV for $g(r)$ at the monopole core is energetically favored for

$$\alpha < \alpha_{cr} \simeq 0.27$$ \hspace{1cm} (27)

As in the case of vortices this result was verified by two methods:

- Using a relaxation method to solve the system (25,26) numerically with boundary conditions

$$\lim_{r\to0} f(r) = 0, \quad \lim_{\rho\to0} g'(r) = 0 \quad \lim_{r\to\infty} f(r) = (\alpha^2 + 1)^{1/2}, \quad \lim_{r\to\infty} g(r) = 0.$$ \hspace{1cm} (28)

The relaxed solution had $g(0) = 0$ for $\alpha > \alpha_{cr}$ (Fig. 1) and $g(0) \neq 0$ for $\alpha < \alpha_{cr}$ (Fig. 2).

- Solving equation (25) numerically for $g = 0$ and substituting in the solution to (26), linearized in $g$. The resulting Schroedinger type equation was then solved numerically with initial conditions $g'(0) = 0$ and $g(0) = 1$. For $\alpha > \alpha_{cr} \simeq 0.27$ it was shown to have no negative eigenvalues. We found negative eigenvalues for $\alpha < \alpha_{cr}$ thus showing the instability of the $g(r) = 0$ solution towards the solution obtained by the relaxation method ($g(0) = c \neq 0$).
The monopole-antimonopole interactions can also be studied in a similar way to vortex interactions. Using similar notation we find that the interacting configuration ansatz for a monopole pair along the z axis is of the form

$$\Phi = \begin{pmatrix} \Phi_1 \\ \Phi_2 \\ \Phi_3 \\ \Phi_4 \end{pmatrix} = \begin{pmatrix} f(r_1)f(r_2)\cos(\theta_1 \pm \theta_2)\sin \phi \\ f(r_1)f(r_2)\cos(\theta_1 \pm \theta_2)\cos \phi \\ f(r_1)f(r_2)\sin(\theta_1 \pm \theta_2) \\ g(r_1,r_2) \end{pmatrix}, \quad (30)$$

where the $+$ ($-$) sign corresponds to a configuration that would be purely repelling (attracting) if it were not for the non-zero VEV of $g(r_1,r_2)$ in the monopole cores. Notice that in this case both configurations correspond to a monopole-antimonopole pair as the defects have opposite topological charges for both signs [16]. An attracting configuration with $g(0,l) = -g(l,0) = \pm c$ has interaction energy (for a derivation in the case of symmetric core see Ref. [16]), at large separations $l$,

$$E_{\text{int}}^- \simeq \eta^2((-4\pi^2 l + \frac{4\pi L}{3}) + A\frac{2c^2}{l}, \quad (31)$$

with $A$ a constant of dimension length squared. Similarly for a repelling configuration with $g(l,0) = g(0,l) = \pm c$ the interaction energy is [16]

$$E_{\text{int}}^+ \simeq \eta^2((8\pi^2 l + 8\pi L) - A\frac{2c^2}{l} \quad (32)$$
The precise form of these results requires numerical simulations of defect evolution which are currently in progress.

4 Dynamical Core Symmetry Breaking and Topological Inflation

It has recently been realized that superheavy topological defects (those for which the symmetry breaking scale is approximately the Planck scale) can act as seeds for inflation. This behavior is known as topological inflation and occurs because at Planck scales, the core of a defect consists of trapped vacuum energy density over a horizon-sized region and thus the slow-roll conditions are satisfied. In these scenarios, it is natural to ask how inflation ends, since for traditional defects, if the slow-roll conditions are satisfied initially, then they will be satisfied for all time as long as the evolution is classical.

However, for the types of defects we consider here, there is an alternative method for topological inflation to end. This occurs if the defect is initially formed with a symmetric core (and hence has trapped energy density there), but at some later time the fields evolve in such a way that a non-symmetric core becomes favored. In such a situation topological inflation should begin and then terminate as a field evolves through some critical value and there ceases to be vacuum energy density trapped in the core.

As a simple example, consider a modest extension of the vortex model mentioned earlier. The conditions necessary for topological inflation to take place are that the core be in the false vacuum and that the symmetry breaking scale be of order the Planck scale. We assume the latter condition is true for the remainder of this section. Introduce a new singlet scalar field, $\chi$, with potential $U(\chi)$ and a $\chi^2 \Phi^\dagger \tau_3 \Phi$ coupling to the $\Phi_i$. Note that we could introduce further couplings between the fields, consistent with the symmetries, but choose not to for the sake of simplicity. The Lagrangian density is
\[ \mathcal{L} = \frac{1}{2} \partial_{\mu} \Phi^\dagger \partial^{\mu} \Phi + \frac{1}{2} (\partial_{\mu} \chi) \partial^{\mu} \chi - \left( \frac{\lambda}{2} \chi^2 + \frac{m^2}{2} \right) \Phi^\dagger \tau_3 \Phi + \frac{M^2}{2} \Phi^\dagger \Phi - \frac{h}{4} (\Phi^\dagger \Phi)^2 - U(\chi) , \] (33)

where, here, \( \alpha \equiv (m/M) > \alpha_{cr} \simeq 0.37 \). We write the potential for \( \chi \) as

\[ U(\chi) = \frac{a}{4} \chi^4 - \frac{b}{3} \chi^3 + \frac{c}{2} \chi^2 , \] (34)

with \( a(T), b(T), c(T) > 0 \), with \( T \) temperature, and perform the rescaling in (\ref{eq:rescaling_1}) along with

\[ \chi \rightarrow \frac{M}{\sqrt{h}} \chi , \quad \lambda \rightarrow \lambda h , \quad a \rightarrow ah , \]
\[ b \rightarrow bM \sqrt{h} , \quad c \rightarrow cM^2 , \quad \mathcal{L} \rightarrow \frac{M^4}{h} \mathcal{L} . \] (35)

This yields the re-scaled Lagrangian density as

\[ \mathcal{L} = \frac{1}{2} \partial_{\mu} \Phi^\dagger \partial^{\mu} \Phi + \frac{1}{2} (\partial_{\mu} \chi) \partial^{\mu} \chi + \frac{1}{2} \left( \alpha^2 - \lambda \chi^2 \right) \left( |\Phi_1|^2 - |\Phi_2|^2 \right) + \frac{1}{2} \Phi^\dagger \Phi - \frac{1}{4} (\Phi^\dagger \Phi)^2 - \frac{a}{4} \chi^4 + \frac{b}{3} \chi^3 - \frac{c}{2} \chi^2 . \] (36)

the equations of motion for the three fields \( \Phi_1, \Phi_2 \) and \( \chi \) are

\[ \partial^{\mu} \partial_{\mu} \Phi_1 - (\alpha^2 - \lambda \chi^2 + 1) \Phi_1 + (\Phi^\dagger \Phi) \Phi_1 = 0 \] (37)
\[ \partial^{\mu} \partial_{\mu} \Phi_2 - (1 - \alpha^2 + \lambda \chi^2) \Phi_2 + (\Phi^\dagger \Phi) \Phi_2 = 0 \] (38)
\[ \partial^{\mu} \partial_{\mu} \chi + [\lambda (\Phi^\dagger \tau_3 \Phi) + c] \chi + a\chi^3 - b\chi^2 = 0 . \] (39)

Now, let us focus on the dynamics in the core of the defect. This is the important region for topological inflation. At high temperatures, we expect the potential for \( \chi \) to have a global minimum at the origin. However, at low temperatures, we can arrange the potential to be a simple quartic with local minimum at \( \langle \chi \rangle = 0 \) and, for appropriate parameter choices, a global minimum at
\( \chi^* \equiv \langle \chi \rangle_{\text{min}} \equiv \frac{b}{2a} \left[ 1 + \left( 1 - \frac{4ac}{b^2} \right)^{1/2} \right] . \) \hfill (40)

Let us choose \( a, b \) and \( c \) at low temperatures so that

\[
\alpha^2 - \lambda \chi^2 < \alpha_{\text{cr}} \quad \text{and} \quad |c| > \lambda v^2 , \hfill (41, 42)
\]

where \( v \) is the values of \( \Phi_2 \) in the non-symmetric core (note that \( \langle \Phi_1 \rangle = 0 \) in the core always).

The evolution of this system in the early universe should be as follows. At a phase transition at around the Planck temperature, superheavy topological defects are formed, with a symmetric core and the auxiliary field, \( \chi \), trapped in its local minimum at the origin. The core of such a defect satisfies the conditions for topological inflation and the space-time inside the core begins to expand exponentially. As the temperature drops, the potential for \( \chi \) evolves and at some time, \( \chi = 0 \) becomes a quasi-stable state and at a later time tunnels and rolls to its global minimum at \( \chi = \chi^* \). At that point the core of the defect undergoes a first order phase transition from symmetric to non-symmetric phase. When this happens, the conditions for topological inflation cease to be satisfied and inflation ends rapidly. Notice that this scenario predicts the existence of non-zero cosmological constant \( \Lambda \) because the defect core never reaches the global minimum of the \( \Phi \) potential.

This is a new scenario for ending topological inflation and takes advantage of the new core structures we have introduced. Similarly we can model a second order phase transition in the core of our defect by choosing \( b = 0 \). At high temperatures the core is in the symmetric phase \( \langle \chi \rangle = 0 \) which is a global minimum as before. At low enough temperatures \( T < T_{\text{crit}} \) with the parameter \( c \) relaxing continuously to negative values \( c(T) < 0 \) a nonzero scalar condensate \( \langle \chi \rangle_{\text{min}} = -\frac{c}{a} \) develops continuously in the core. Its trapped vacuum energy relaxes in contrast to the previous scenario continuously to the nonzero value determined by the \( \Phi \) field.
5 Conclusions

We have demonstrated the existence of global vortex and monopole configurations which arise from vacuum manifolds $S^2$ and $S^3$ which are deformed by partial breaking of their global symmetries to $S^1$ and $S^2$ respectively. Furthermore, these embedded topological defects possess either symmetric or non-symmetric cores for appropriately specified values of their parameters. We have also shown that, without gauge fields, it is not possible to support specific metastable topological defects confined in the false vacuum of other defects. The existence of such confined defects in the presence of gauge fields is an open issue under investigation. Further, we have used arguments based on energetics to show that non-trivial structures in defect cores can lead to additional types of interaction potentials. Finally, we have extended our models in such a way that the transition from symmetric to non-symmetric phase in the core is dynamical, with a scalar field condensate being the effective order parameter. In that case, if topological inflation occurs in the symmetric defect core, a phase transition in the core may cause the vacuum energy there to be reduced. This dynamical ingredient for topological inflation provides a novel exit mechanism for the defect core from its otherwise eternally inflating state and thus is a new mechanism for terminating inflation.

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