Cosmological Implication of Electroweak Monopole

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Abstract. We estimate the remnant electroweak monopole density of the standard model in the present universe. We show that, although the electroweak phase transition is of the first order, the monopole production comes from the thermal fluctuations of the Higgs field after the phase transition, not the vacuum bubble collisions during the phase transition. Moreover, most of the monopoles produced initially are annihilated as soon as created, and this annihilation continues very long time, longer than the muon pair annihilation time. As the result the remnant monopole density at present universe becomes very small, of 10^{-11} of the critical density, too small to be the dark matter. We discuss the physical implications of our results on the ongoing monopole detection experiments.

1 Introduction

Ever since Dirac has proposed the Dirac monopole generalizing the Maxwell's theory, the monopole has become an obsession in physics [1]. After the Dirac monopole we have had the Wu-Yang monopole [2], the 't Hooft-Polyakov monopole [3], the grand unification (Dokos-Tomaras) monopole [4], and the electroweak (Cho-Maison) monopole [5, 6]. But these monopoles, except the electroweak one are unrealistic.

Indeed the Dirac monopole in electrodynamics should transform to the electroweak monopole after the unification of the electromagnetic and weak interactions, and the Wu-Yang monopole in QCD is supposed to make the monopole condensation to confine the color. Moreover, the 'tHooft-Polyakov monopole exists only in an hypothetical theory, and the grand unification monopole which could have been amply produced at the grand unification scale in the early universe probably has become completely irrelevant at present universe after the inflation.

This makes the electroweak monopole the only realistic monopole we could ever hope to detect. This has made the experimental confirmation of the electroweak monopole one of the most urgent issues in the standard model after the discovery of the Higgs particle, and the newest MoEDAL ("the magnificent seventh") detector at LHC is actively searching for the monopole [7, 8]. On the other hand, the 14 TeV LHC may have no chance to produce the monopole if the monopole mass becomes

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larger than 7 TeV. In this case we must try to detect the remnant monopole at present universe produced in the early universe. so it is fortunate that IceCube, Antares, Auger, and similar experiments are searching for them [9–11]

To detect the remnant monopoles, however, we need to know how many of them are left over in the present universe. There have been discussions on the cosmological production of monopoles, but most of them have been on the grand unification monopoles [12–15]. The general consensus is that the grand unification monopoles would have overclosed the universe without the inflation, but the inflation might have completely diluted them in such a way that they could have no visible impact on the present universe [16].

For the electroweak monopole, however, the situation was not so clear [17–20]. Zeldovich studied the cosmological impact of the electroweak monopole of mass around 10 TeV, and suggested that the electroweak monopoles could also overclose the universe [18]. But certainly we need a more accurate discussion on this issue. *The purpose of this letter is to discuss the cosmological production and successive evolution of the electroweak monopole to predict the remnant monopole density of the standard model at present universe.*

Our result in this paper strongly implies that, the electroweak monopoles might have been amply produced during the electroweak phase transition, but unlike the grand unification monopole they do not alter the standard cosmology in any significant way. Moreover, most of them could have been quickly annihilated. Nevertheless enough of them, roughly 1.31×10^{-9} of the baryon density, might have survived at the present universe. This implies that, although the remnant density of the electroweak monopole is too small to be the dark matter, there are enough of them left over which could be detected.

Before we estimate the electroweak monopole density at present universe, it is important for us to clarify t he existing confusions and misunderstandings on the electroweak monopole. The most serious misunderstanding is that the standard model has no monopole [21]. The basis for this “no-go theorem" is that the monopole topology in spontaneously broken gauge theory is given by $\pi_2(G/H)$, where G is the full symmetry and H is the unbroken symmetry. So, with $G = SU(2) \times U(1)_Y$ and $H = U(1)_{em}$, the second homotopy becomes trivial. According to this logic the standard model has no monopole topology.

The problem with this logic is that $\pi_2(G/H)$ is not the only possible monopole topology. Strictly speaking this logic applies only when G is simple. But in the standard model G is made of direct product of two groups $SU(2)$ and $U(1)_Y$, both of which admit the monopole topology. Moreover, the standard model can be viewed as a gauged CP^1 model which has the monopole topology. This tells that the standard model has a different type of monopole topology, the combination of the non-Abelian $SU(2)$ monopole topology and the Abelian $U(1)_Y$ monopole topology [5, 22]. This makes the electroweak monopole a hybrid between Dirac and ’tHooft-Polyakov.

Another way to see the existence of the monopole topology is that the standard model (after the symmetry breaking) contains the electromagnetic part which is essentially Maxwell’s theory, except the $U(1)_{em}$ becomes non-trivial. And the non-trivial $U(1)_{em}$ is well known to have the Abelian monopole topology. This assures that the standard model does have the monopole.

The above argument tells the followings. First, the electroweak monopole is the electroweak generalization of Dirac monopole. But unlike the Dirac monopole which is optional, the electroweak monopole must exist because the electromagnetic U(1) in the standard model becomes non-trivial

[5, 6]. This makes the discovery of the monopole, not the Higgs particle, the final (and topological) test of the standard model.

Second, there is a fundamental difference between the Dirac monopole and the electroweak monopole. It carries the magnetic charge twice bigger than the Dirac monopole. This is because, in the course of the electroweak unification, the period of electromagnetic U(1) becomes 4π, not 2π [5, 6]. So we could tell if it is the Dirac monopole or the electroweak monopole, when we discover it.

Third, unlike the Dirac monopole, we can estimate the mass of the electroweak monopole to be of the order of several TeV, roughly $1/\alpha$ times the W-boson mass [22–24]. This is because the monopole mass essentially comes from the same Higgs mechanism which generates the W-boson mass, except that the monopole potential couples to the Higgs multiplet magnetically. In fact, thos was exactly the reason for Zeldovich to predict the mass of the electroweak monopole to be several TeV [18]. In the following the estimate of mass of the electroweak monopole plays a crucial role for us to estimate the remnant monopole density.

2 Electroweak Phase Transition

Now we discuss the phase transition at the electroweak scale. To do this, we have to find the temperature-dependent effective potential of the Higgs field. Fortunately this has been known [17, 19],

$$
\begin{aligned}
V_T(\rho) &= V_0(\rho) - \frac{C_1}{12\pi}\rho^3\, T + \frac{C_2}{2}\rho^2\, T^2 - \frac{\pi^2}{90} N T^4 + \delta V_T, \\
V_0(\rho) &= \frac{\lambda}{8}(\rho^2 - \rho_0^2)^2, \\
C_1 &= \frac{6M_W^3 + 3M_Z^3}{\rho_0^3} \simeq 0.36, \\
C_2 &= \frac{4M_W^2 + 2M_Z^2 + M_H^2 + 4m_t^2}{8\rho_0^2} \simeq 0.36,
\end{aligned} \tag{1}
$$

where V_0 is the zero-temperature potential, N is the total number of distinct helicity states of the particles with mass smaller than T (counting fermions with the factor 7/8), C_1 and C_2 are the contributions from the ρ dependent part of the masses of gauge bosons, Higgs field, and fermions, M_W, M_Z, M_H, and m_t are the W-boson, Z-boson, Higgs boson, and the top quark masses, and δV_T is the slow-varying logarithmic corrections and the lighter quark contributions which we will neglect from now on.

The effective potential (with $\delta V_T = 0$) is schematically drawn in Fig. **??**. It has three local extrema at

$$
\begin{aligned}
\rho_s &= 0, \\
\rho_\pm(T) &= \frac{C_1}{4\pi\lambda}\, T \pm \sqrt{\Big(\frac{C_1}{4\pi\lambda}\Big)^2\, T^2 + \rho_0^2 - \frac{2C_2}{\lambda}\, T^2}.
\end{aligned} \tag{2}
$$

The first extremum $\rho_s = 0$ represents the Higgs vacuum of the symmetric (unbroken) phase, and the second extremum $\rho_-(T)$ represents the local maximum, and the third extremum $\rho_+(T)$ represent the local minimum Higgs vacuum of the broken phase. But notice that these two extrema $\rho_\pm$ appear only

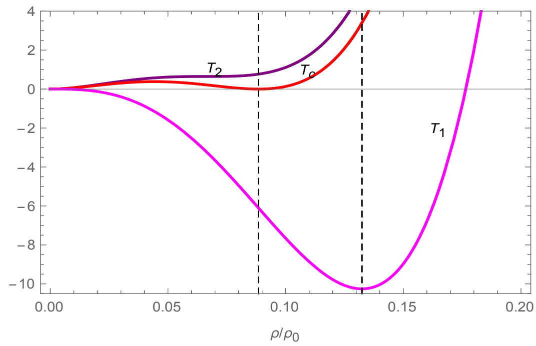

Figure 1. The amplified effective thermal potential for $T = T_2, T_c$, and T_1. Here the vertical axis represents $V_{eff} - V_0$ in the unit $10^{-5} \times V_0$, and the vertical dashed lines indicate the positions of the local minima of the potentials.

when T becomes smaller than T_2

$$T_2 = \frac{T_1}{\sqrt{1 - C_1^2/16\pi^2\lambda C_2}} \simeq 146.7 \text{ GeV},$$
$$T_1 = \sqrt{\frac{\lambda}{2C_2}}\,\rho_0 \simeq 146.4 \text{ GeV}. \tag{3}$$

So above this temperature only $\rho_s = 0$ becomes the true vacuum of the effective potential, and the electroweak symmetry remains unbroken.

At $T = T_2$ we have

$$\rho_- = \rho_+ = (C_1/4\pi\lambda)\, T_2 \simeq 16.3 \text{ GeV}, \tag{4}$$

but as temperature cools down below T_2 we have two local minima at ρ_s and ρ_+ with $V_T(0) < V_T(\rho_+)$, until T reaches the critical temperature T_c where $V_T(0)$ becomes equal to $V_T(\rho_+)$,

$$T_c = \frac{T_1}{\sqrt{1 - C_1^2/18\pi^2\lambda C_2}} \simeq 146.6 \text{ GeV},$$
$$\rho_+(T_c) = \frac{C_1}{3\pi\lambda}\, T_c \simeq 21.8 \text{ GeV}. \tag{5}$$

So $\rho_s = 0$ remains the minimum of the effective potential for $T > T_c$. Notice that $T_c/\rho_0 \simeq 0.6$ but $\rho_+/\rho_0 \simeq 0.09$.

Below this critical temperature ρ_+ becomes the true minimum of the effective potential, but $\rho_s = 0$ remains a local (unstable) minimum till the temperature reaches T_1. But at $T = T_c$ the new vacuum bubbles start to nucleate at $\rho = \rho_+$, which takes over the unstable vacuum $\rho_s = 0$ completely at $T = T_1$ where $\rho_+(T_1)$ becomes around 32.6 GeV. From this point ρ_+ becomes the only (true) minimum, which transforms to the well-known Higgs vacuum ρ_0 at zero temperature. The effective potential is shown in Fig. 1.

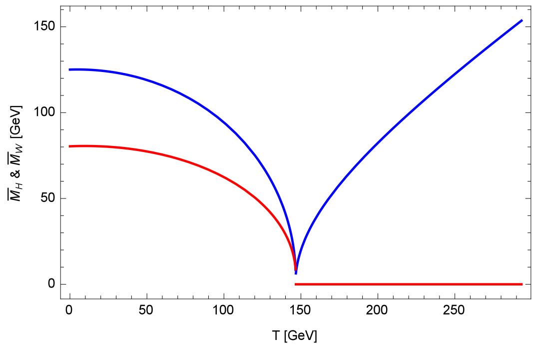

Figure 2. The temperature-dependent Higgs and W-boson masses. The blue and red curves represent the Higgs and W-boson masses.

This tells that two local minima ρ_s and ρ_+ are separated by an energy barrier, which means that the electroweak phase transition is of the first order. On the other hand, the energy barrier is very small. Moreover, the barrier lasts only for short period since the temperature difference from T_c to T_1 is very small, $\delta = (T_c - T_1)/T_c \simeq 0.001$. This strongly implies that the phase transition is very mildly the first order, almost the second order.

From (1) we can clculate the temperature-dependent Higgs mass $\bar{M}_H$,

$$\bar{M}_H^2 = \frac{d^2 V_{eff}}{d\rho^2}\Big|_{\rho_{min}}$$
$$= \begin{cases} [(T/T_1)^2 - 1]M_H^2/2, & T \geq T_c, \\ [(\rho_+/\rho_0)^2 + 1 - (T/T_1)^2]M_H^2/2, & T < T_c. \end{cases} \tag{6}$$

So $\bar{M}_H$ acquires its minimum value 5.53 GeV at $T = T_c$ and becomes 11.7 GeV at $T = T_1$, and approaches to the zero temperature value 125.7 GeV as the universe cools down.

Similarly, the W-boson which was massless at high temperature becomes massive at $T = T_c$ when the Higgs field acquires the non-vanishing vacuum expectation value. So, as the Higgs vacuum $\rho = 0$ starts to tunnel to $\rho = \rho_+$ at T_c, the W-boson starts to become massive toward the value $g\rho_+(T_c)/2 \simeq$ 7.0 GeV. And it acquires the mass 10.5 GeV at $T = T_1$, which approaches the well-known zero-temperature value at $T = 0$. The temperature-dependent Higgs and W-boson masses are shown in Fig. 2. They play an important role because they determine the correlation length $\xi = \bar{M}_H^{-1}$ and the monopole mass $M_m \simeq (1/\alpha) \times \bar{M}_W$.

3 Cosmological Production of Electroweak Monopole

The cosmological production of monopole in the second order phase transition was estimated by Kibble and later improved by Zurek [12, 15]. Kibble observed that the fundamental scale which determines the initial monopole density in the second order phase transition is the horizon distance at the time of the transition, and set a upper bound for the initial monopole density. Zurek improved this bound taking into account the fact that the phase transition does not take place instantaneously but continuously.

But this picture does not apply to the electroweak monopole production, because the electroweak phase transition is the first order. In this case, the correlation length becomes finite because the Higgs mass becomes non-zero at the critical temperature. Moreover, what is important in the monopole production is the time of monopole formation. To produce the monopoles we need the change of topology induced by the zero points of the Higgs field, which comes from the thermal fluctuations. And obviously these zero points, the seeds of the monopoles, do not appear instantaneously at the critical temperature, but come to exist some time after the phase transition.

It is well known that the temperature which governs this thermal fluctuation is the Ginzburg temperature, not the critical temperature. This means that even in the second order phase transition, the relevant scale in the monopole production is not the horizon distance at the critical temperature but the correlation length around the Ginzburg temperature.

Of course, in the strongly first order phase transition, the symmetric vacuum becomes meta-stable below the critical temperature, and the vacuum bubbles nucleate during the phase transition. So the vacuum bubble collisions during the phase transition becomes the monopole production mechanism [14]. However, the electroweak phase transition is very mildly first order, so that the bubble formation becomes unimportant. So the monopole formation takes place by the thermal fluctuation of the Higgs field after the phase transition, not by the bubble collisions during the phase transition.

Just below the critical temperature the Higgs field is still subject to large fluctuations which bring $\langle\rho\rangle$ back to zero. This is possible so long as

$$\xi^3 \Delta F \leq T, \tag{7}$$

where $\xi(T)$ is the correlation length and $\Delta F(T) = V(\rho_s) - V(\rho_+)$ is the difference in free energy density between two phases. This large fluctuation disappears when the equality holds, at the Ginzburg temperature. So we have to find the Ginzburg temperature first.

Unfortunately the Ginzburg temperature T_G of the potential (1) can not be expressed in a simple form, but we can find T_G numerically directly from (1) and (6),

$$T_G \simeq 57.6\ \text{GeV}, \quad \xi_G \simeq 1.7 \times 10^{-16}\ \text{cm}. \tag{8}$$

With this we may assume that the period of the monopole formation is between T_1 and T_G, and identify the temperature and correlation length of the initial monopole production to be the mean value of T_1 and T_G,

$$\begin{aligned} T_i &= \frac{T_1 + T_G}{2} \simeq 102.0\ \text{GeV}, \\ \xi_i &= \frac{\xi(T_1) + \xi(T_G)}{2} \simeq 9.3 \times 10^{-16}\ \text{cm}. \end{aligned} \tag{9}$$

From this we can estimate the initial density of the monopoles

$$\left(\frac{n_m}{T^3}\right)_i \simeq \frac{g_P}{\xi_i^3 T_i^3} \simeq 9.1 \times 10^{-4}, \tag{10}$$

where g_P is the probability that one monopole is produced in one correlation volume, which we assume to be around 10 %.

To see how much energy we need to produce the electroweak monopoles, we calculate the energy density of the monopoles from (10)

$$\varrho_{mo}(T_i) = M_m(n_m)_i \simeq 9.0\times 10^{-3}\ T_i^4\ \Big(\frac{M_m}{1\ \text{TeV}}\Big),$$
$$\frac{\varrho_{mo}(T_i)}{\varrho(T_i)} \simeq 2.5\times 10^{-4}\Big(\frac{M_m}{1\ \text{TeV}}\Big). \tag{11}$$

So, assuming the initial monopole mass to be around 1 TeV, we may conclude that the universe need to consume only a tiny fraction (about 0.025 %) of the total energy to produce the monopoles. This assures that the electroweak monopole production does not alter the standard cosmology.

In fact, we can say that even (10) is an overestimation. The reason is that, as we will see soon, the monopole-antimonopole capture radius becomes much bigger the correlation length ξ_i, so that most of the monopoles annihilate with the anti-monopoles as soon as they are produced.

4 Cosmic Evolution of Electroweak Monopole

Although the monopoles are topological and absolutely stable, there are two factors which change the initial monopole density, the Hubble expansion and the annihilation of monopole-antimonopole pairs. So the change of monopole density n_m is determined by the Boltzmann equation [12, 13]

$$\frac{dn_m}{dt} + 3Hn_m = -\sigma n_m^2. \tag{12}$$

where H and σ are the Hubble parameter and the monopole annihilation cross section.

There are two things which affect the annihilation of the monopoles, the thermal Brownian motion (random walk) of the monopole in hot plasma of charged particles and the attraction between monopoles and anti-monopoles. After the creation the monopoles diffuse in hot plasma by the Brownian motion with the mean free path $l_{\rm free}$ given by

$$l_{\rm free} = v_t t_{\rm free} \simeq \frac{1}{BT}\sqrt{\frac{M_m}{T}},\quad B = \frac{1}{T}\sum_i n_i\sigma_i, \tag{13}$$

where $v_t \simeq \sqrt{T/M_m}$ and $t_{\rm free}$ are the thermal velocity and the mean free time of the monopoles, n_i and σ_i are the number density and the cross section of the i-th relativistic charged particles and the sum is the sum over all spin states [13].

With

$$n_i \simeq \frac{3\zeta(3)}{4\pi^2}T^3,\quad \sigma_i \simeq \Big(\frac{g_m q_i}{4\pi}\Big)^2\frac{1}{T^2} = \Big(\frac{q_i}{e}\Big)^2\frac{1}{T^2}, \tag{14}$$

we have

$$B \simeq \frac{3\zeta(3)}{4\pi^2}\sum_i\Big(\frac{q_i}{e}\Big)^2 \simeq 0.09\times\sum_i\Big(\frac{q_i}{e}\Big)^2, \tag{15}$$

where q_i is the electric charge of the i-th particle and $\zeta(3) = 1.202...$ is the Riemann zeta function. Since the charged particles in the plasma are the leptons and quarks, we may put $B \simeq 3$. From this we have $v_t \simeq 0.31\times(1\ \text{TeV}/M_m)^{1/2}\ c$ and $l_{\rm free} \simeq 2.0\times(M_m/1\ \text{TeV})^{1/2}\times 10^{-16}$ cm around T_i.

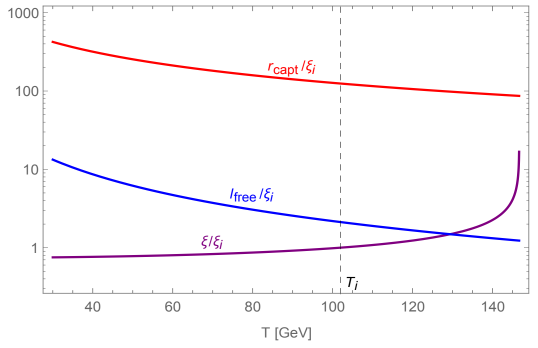

Figure 3. The relevant scales, ξ in purple, l_{free} in blue, and r_{capt} in red, against T. They are normalized by the correlation length ξ_i at T_i. Here we set $M_m = 5$ TeV.

Now, against the thermal random walk of the monopoles, we have the attractive Coulomb force between monopoles and anti-monopoles which makes them drift towards each other. The drift velocity v_d of the monopole at a distance r from the anti-monopole is given by [25]

$$v_d \simeq \frac{1}{\alpha} \times \frac{1}{BT^2 r^2} = \frac{\alpha_m}{BT^2 r^2}, \tag{16}$$

where α_m is the monopole fine structure constant defined by $\alpha_m = g_m^2/4\pi = 1/\alpha$. Notice that $v_d(r = \xi_i) \simeq 2\ c$. This, of course, is unrealistic but tells that $v_t \ll v_d$. This shows that the Coulomb attraction between the monopoles and anti-monopoles is much stronger than the thermal diffusion.

Moreover, the monopole-antimonopole capture radius given by $r_{\text{capt}} \simeq \alpha_m/T$ becomes about 2.7×10^{-14} cm at T_i, which is much bigger than l_{free} at this temperature. This strongly implies that the monopole-antimonopole annihilation is much more important around T_i, so that as soon as they are created, they annihilate each other. In Fig. 3 we plot the relevant scales ξ, l_{free}, and r_{capt} against T for comparison. This clearly shows that the capture radius is much bigger than the mean free length and correlation length in a wide range of T. This assures that (10) is indeed an overestimation.

Now, if we let the mean distance between the monopole and anti-monopole r be $r \simeq n_m^{-1/3}$, the capture time is given by

$$t_{\text{capt}} \simeq \frac{r}{v_d} \simeq \alpha \times \frac{BT^2}{n_m}. \tag{17}$$

From this we have the monopole-antimonopole annihilation cross section

$$\sigma \simeq \frac{1}{t_{\text{capt}} n_m} \simeq \frac{\alpha_m}{BT^2}. \tag{18}$$

With this we can solve the Boltzmann equation (12),

$$\begin{aligned} \frac{n_m}{T^3} &= \frac{1}{A(M_m/T - M_m/T_i) + B}, \\ A &= \alpha_m \times \frac{Cm_p}{BM_m}, \quad B = \left(\frac{n_m}{T^3}\right)_i^{-1}. \end{aligned} \tag{19}$$

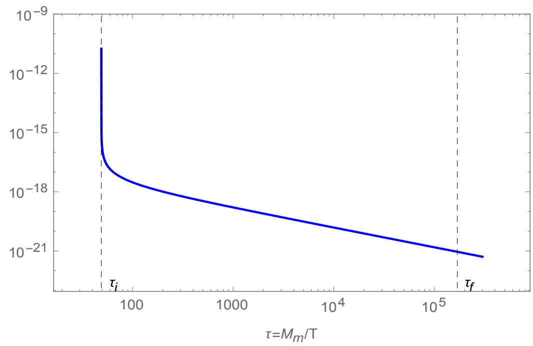

Figure 4. The evolution of the monopole density n_m/T^3 against $\tau = M_m/T$. The final value of the monopole density is independent of the initial value.

Here M_m is treated as a constant, but in reality it depends on time. So it should be understood as a mean value.

The evolution of the monopole density n_m/T^3 against M_m/T is shown in Fig. 4, where we have put $M_m = 5$ TeV. Notice that, regardless of the initial condition the monopole density approaches to a constant as the universe is cooled down.

The diffusive capture process is effective only when $\ell_{\text{free}} < r_{\text{capt}}$, which determines the temperature T_f below which the monopole-antimonopole annihilation ceases,

$$T_f \simeq \alpha^2 \times \frac{M_m}{B^2} \simeq 5.9 \times \Big(\frac{M_m}{1 \text{ TeV}}\Big) \text{ MeV}. \tag{20}$$

With $M_m = 5$ TeV, we have $T_f \simeq 29.5$ MeV, which is below the muon decoupling temperature. Actually, around this temperature B becomes very small so that the capture radius becomes smaller. This is because the only charged particles remaining in the plasma are electrons and positrons. But the above analysis clearly shows that the annihilation continues very long time.

The monopole density after annihilation becomes

$$\Big(\frac{n_m}{T^3}\Big)_f \simeq \alpha^3 \times \frac{M_m}{BCm_p} \simeq 1.8 \times 10^{-22}\Big(\frac{M_m}{1 \text{ TeV}}\Big). \tag{21}$$

Obviously this value is much lower than the initial density given by (10). The number of monopole within the comoving volume is conserved thereafter. But they still interact with the electron pairs in the hot plasma before decouple around $T_d \simeq 0.5$ MeV, when the electron pairs disappear and the interaction rate becomes less than the Hubble expansion rate.

Assuming that the expansion is adiabatic, the current number density and the energy density of the monopole is given by

$$\begin{aligned} \Big(\frac{n_m}{T^3}\Big)_0 &= \Big(\frac{g_{s,0}}{g_{s,f}}\Big)\Big(\frac{n_m}{T^3}\Big)_f, \\ \rho_{mo,0} = M_m n_{m,0} &= M_m\Big(\frac{g_{s,0}}{g_{s,f}}\Big)\Big(\frac{n_m}{T^3}\Big)_f\, T_0^3, \end{aligned} \tag{22}$$

where T_f is the temperature of universe when the monopole annihilation process effectively ceases and $T_0 = 2.73$ K $= 2.35 \times 10^{-13}$ GeV is the temperature of the universe today. Here g_s is the effective number of degrees of freedom in entropy.

So the current density parameter of monopole can be written

$$\Omega_{mo}\, h^2 = \frac{\rho_{mo,0}\, h^2}{\rho_{c,0}} \simeq 1.2 \times 10^{-12} \times \Big(\frac{M_m}{1\ \text{TeV}}\Big)^2, \tag{23}$$

where $\rho_{c,0} = 3H_0^2/8\pi G$ is the critical density of present universe and $h \simeq 0.678$ is the scaled Hubble constant in the unit $H_0/(100$ km s^{-1} Mpc$^{-1})$. With $h \simeq 0.678$ and $M_m \simeq 5$ TeV, we have $\Omega_{mo} \simeq 6.53 \times 10^{-11}$. This is about 1.31×10^{-9} of the baryon density, which assures that the electroweak monopole cannot be a dark matter candidate.

In terms of the number density, this translates to about 6.1×10^{-20} *cm*$^{-3}$*, or about* $2.3 \times 10^{-13}\ n_b$*, where* $n_b \simeq 2.5 \times 10^{-7}$ *cm*$^{-3}$ *is the number density of the baryons. Intuitively, this means that there are roughly* 6.6×10^7 *monopoles per unit volume of the earth.* This is a small but significant number which suggests that there are enough electroweak monopoles left over in the universe that we could detect.

Notice that, since the decoupling temperature of the electroweak monopole is around 0.5 MeV, the free streaming monopoles startas non-relativistic. But eventually they are accelerated by the inter-galactic magnetic field, and they become extremely relativistic.

5 Discussion

In this talk we have studied the cosmic production and the successive evolution of the electroweak monopole, and estimated the remnant monopole density in the present universe. Our result tells that, unlike the grand unification monopole, the electroweak monopole can not alter the standard cosmology in any significant way. Nevertherless there are enough remnant monopoles in the present universe that we can detect.

Actually the real remnant monopole density could be much less than the above estimate (23) for the following reasons. First, as the only heavy stable particle with mass about 10^4 times heavier than the proton, they can easily generate the density perturbation and become an excellent candidate for the seed of the large scale structures in the early universe. So many of them might have been buried at galactic centers.

Second, the monopoles have a very short penetration length in the matter because they have strong magnetic interaction. In fact a relativistic electroweak monopole can travel only a few meters (less than 10 m) in Aluminum before they are trapped [26]. This means that most of the monopoles left over which did not become the seed of large scale structure in the early universe could have been trapped and filtered out by the stellar objects, when they collide with them. This strongly implies that the actual remnant monopole density could be much less that (23).

Our result provides a useful tip for the remnant monopole detection experiments. Certainly there are enough remnant electroweak monopoles for these experiments. But the problem is that when the monopoles pass through the earth atmospheric sphere, they loose most of the kinetic energy and slow down quickly when they reach earth surface. So most of them are trapped near the earth surface. This again is because they interact strongly with matter.

This implies two things. First, it would be extremely difficult for the monopoles to arrive at detectors buried 2 km under the iceburg at the south pole or under the medeterranean sea. Second, a best way to detect the remnant monopole is to locate the detector at high altitude. The remnant monopole detection experiments should keep these points in mind.

As importantly our analysis provides an important information for the monopole production at LHC. A central issue at LHC is what is the monopole production mechanism at LHC, and whether LHC could produce the monopole. It has generally been believed that LHC could produce the monopole pair by Drell-Yan process ($p\bar{p} \rightarrow \gamma \rightarrow M\bar{M}$) and/or photon fusion process ($p\bar{p} \rightarrow p\bar{p} + \gamma\gamma \rightarrow p\bar{p} + M\bar{M}$) [27]. This would be the case if we treat the monopole as a point particle. But since the monopole is a topological particle, it is not clear if this is correct.

Our analysis suggests that the monopole production mechanism at LHC should be the topological production during the electroweak phase transition. This should be contrasted with the above view popular in high energy community.

But for LHC to produce the monopole it should satisfy three criteria. First, LHC should have enough energy to produce the monopole pair. This means that if the monopole mass should be less than than 6.5 TeV for the present LHC to produce them. On the other hand our analysis shows that in the electroweak phase transition, we need only about 1 TeV to create the monopole seed (the baby monopole). This implies that the energy constraint on the 13 TeV LHC to produce the electroweak monopole pair may not be so strong obstacle to create the monopole pair.

Second, the size of the collision core should be large enough, larger than the correlation length, for the topological singularity to form. Third, the collision time sholud be longer than the thermal fluctuation time of the Higgs field. Fortunately, LHC easily satisfies these criteria [28]. So there should be no problem for LHC to produce the monopole, if LHC satisfies the energy constraint.

Obviously the electroweak monopole has important implications in physics, in particular in cosmology. First, it could have been the source of the large scale structure in the universe, as we have remarked. Second, as the heaviest relativistic magnetically charged particles in the universe, they become a natural source of ultra-high energy cosmic rays. Third, they could might generate the magnetic field in the universe and become the source of the intergalactic magnetic field. Moreover, they could play an important role in the electroweak baryogenesis.

But the most important point of the electroweak monopole is that it must exist. we hope our result could be useful to detect them. The details of our discussions will be presented in a separate publication [28].

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