Electroweak Baryogenesis in a Cold Universe

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ABSTRACT

We discuss the possibility of generating the baryon asymmetry of the Universe when the temperature of the Universe is much below the electroweak scale. In our model the evaporation of primordial black holes or the decay of massive particles re-heats the surrounding plasma to temperatures above the electroweak transition temperature leading to the restoration of electroweak symmetry locally. The symmetry is broken again spontaneously as the plasma cools and a baryon asymmetry is generated during the phase transition. This mechanism generates sufficient asymmetry for a second order electroweak phase transition. For a first order phase transition, sufficient asymmetry is generated if viscous effects slow down the heated plasma as it moves away from the black hole. In our scenario there is no wash-out of the asymmetry after the phase transition as the plasma rapidly cools to lower temperatures thereby shutting off the sphaleron processes.
I. INTRODUCTION

In recent years much effort has been devoted to formulating a mechanism for baryogenesis at the electroweak phase transition [1]. Electroweak baryogenesis is indeed a very exciting possibility as here one is working within the framework of a theory (or its extensions) which is reasonably well understood and the energy scale involved is accessible in laboratory experiments. However the requirement of a strong first order phase transition, large CP violation and a small Higgs mass to ensure that the created baryon asymmetry is not washed away after the phase transition places stringent constraints on these models and practically rules out electroweak baryogenesis in the context of the Standard Model. In this paper we discuss an alternative scenario for electroweak baryogenesis. Here the baryon asymmetry is created at temperatures much below the electroweak transition temperature during the evaporation of primordial black holes. When a black hole is evaporating it heats up the plasma around it to a temperature much higher than the ambient temperature for a short time. This can also happen due to the decay of massive particles. For appropriate black hole masses (or, particle masses) the temperature of the hot region rises above the electroweak transition temperature $T_{ew}$ and the electroweak symmetry is restored locally. Due to the transfer of energy out of this region the hot region will cool and the temperature will fall below $T_{ew}$. Thus in these hot regions the electroweak phase transition occurs again. Baryon asymmetry is then generated in these hot regions during this phase transition.

One motivation for our model is that a first order transition at a very early stage in the Universe, or density fluctuations in general, seem to inevitably lead to the formation of primordial black holes. Evaporation of these black holes will lead to local heating of the plasma and hence to the possibility of electroweak baryogenesis for appropriate black hole masses. Also, scalars found in string theory can decay after the electroweak transition and dilute any baryon asymmetry created earlier [2]. But in our scenario, we envision that the particles that dilute the baryon asymmetry during their decay can also recreate the baryon asymmetry after their decay.

The basic mechanism of our model can be realized either by particles emitted in the evaporation of a primordial black hole or by particles constituting the decay products of a heavy particle. In both cases, it is the energy of these emitted particles which thermalizes and leads to local heating of the surrounding plasma. For the black hole case the heated region will be spherically symmetric while for the particle decay case the region will be collimated. In this paper we will only discuss the black hole case. The particle decay case
is much more complicated. We will briefly comment on this case at the end of the paper but it is not clear to us at this stage if in that case one can get sufficient heating of the local regions. We hope to discuss the particle decay case in a future work.

A desirable feature of our scenario is that it works for a second order transition. In the standard cosmological scenario where the Universe cools due to Hubble expansion, a second order phase transition can not produce sufficient baryon asymmetry as the baryon asymmetry produced is proportional to $\dot{\phi}/\phi \sim H$, where $\phi$ is the Higgs field, and this quantity is too small. In our scenario the heated plasma cools much faster as it moves away from the black hole and $\dot{\phi}/\phi$ is large enough to produce the observed baryon asymmetry. For a first order phase transition, the asymmetry generated is suppressed by the Lorentz factor associated with the relativistic speeds at which the heated plasma moves away from the black hole. However if viscous effects in the plasma slow down the outward moving plasma this suppression will be decreased. In our scenario, since the heated plasma cools rapidly, the temperature after baryogenesis quickly becomes much smaller than the electroweak scale preventing any significant wash-out of the created asymmetry after the phase transition.

We divide the calculation of the baryon asymmetry in our model into two main steps. In Section 2, we estimate the volume of the plasma in which the electroweak symmetry is restored and which is relevant for baryogenesis in our scenario. In Section 3, we estimate the baryon asymmetry generated as these regions cool and undergo the electroweak transition. Then in Section 4 we point out why the upper bound on the Higgs mass does not exist in our model. We conclude in Section 5 with a brief discussion of how inhomogeneities in the baryon number distribution can be smoothed out before nucleosynthesis.

We briefly comment on the formation of primordial black holes in the early Universe. In pre-inflationary days it was argued that the initial spectrum of density inhomogeneities, invoked to explain the observed structure in the Universe, could also give rise to primordial black holes. It was also argued that white holes would be unstable and would convert to black holes. In first order phase transitions, particularly in first order inflationary scenarios, primordial black holes can be produced by collapsing regions of false vacuum trapped between bubbles of the true vacuum. Primordial black holes can

* However models of electroweak baryogenesis with topological defects are insensitive to the order of the electroweak phase transition and can create the observed asymmetry for a second order phase transition.
also be formed via the gravitational instability of inhomogeneities formed during bubble wall collisions in first order inflation \[9, 12\], by large amplitude density perturbations produced due to fluctuations in the inflaton field \[13\], by shrinking cosmic string loops \[14, 15\] and by expanding topological defects produced during inflation \[16\]. As our primary motivation is to demonstrate a qualitatively different scenario for implementing electroweak baryogenesis, we will not go into the specific mechanism which could give rise to the formation of black holes of required masses. However we point out that the formation of black holes which evaporate below the electroweak scale has been discussed in \[17\].

In ref. \[18\] a scenario which is similar to ours has been proposed. We comment later on this work. The presence of a hot plasma surrounding primordial black holes has also been discussed in refs. \[19–22\]. However, we point out that in these references the plasma consists of the black hole radiation and particles produced by the interaction of the black hole radiation with itself. In contrast, in our scenario we consider the interaction of the Hawking radiation with the ambient plasma surrounding the black hole, and the subsequent heating up of this plasma by the Hawking radiation. The possibility of obtaining the baryon asymmetry of the universe in the plasma around the black hole is mentioned in refs. \[13, 22\].

II. FORMATION OF HOT REGIONS

A black hole evaporates by emitting Hawking radiation with an associated temperature

\[
T_{bh} = \frac{M_{Pl}^2}{8\pi M_{bh}}
\]

where \(M_{bh}\) is the mass of black hole and \(M_{Pl} = 1.2 \times 10^{19} \text{GeV}\) is the Planck mass. We use natural units with \(\hbar = c = 1\). The rate of loss of mass by the evaporating black hole is given by

\[
\frac{dM_{bh}}{dt} = -\frac{\alpha M_{Pl}^4}{M_{bh}^2}
\]

Here, \(\alpha\) accounts for the scattering of emitted particles by the curvature and depends on \(T_{bh}\). For different values of \(T_{bh}\) values of \(\alpha\) have been tabulated in \[23\]. For \(T_{bh} = 1, 200\) and \(10^{15}\) MeV, the corresponding values of \(\alpha\) are \(3.6 \times 10^{-4}\), \(2.3 \times 10^{-3}\) and \(4.5 \times 10^{-3}\) respectively. As we will see later, the relevant value of \(T_{bh}\) for us will be higher than about \(10^6\) GeV. Therefore we shall set \(\alpha\) to be \(3 \times 10^{-3}\).

The lifetime \(\tau_{bh}\) of the black hole can be obtained by integrating eqn.(2). We get
\[ \tau_{bh} \simeq 10^2 M_0^{-4} M_p^3 \]  

(3)

where \( M_0 \) is the initial mass of the black hole. Eqn.(2) implies that very little energy is emitted until time of the order of \( \tau_{bh} \) which is when most of the energy of the black hole gets emitted. Thus for a black hole formed early in the Universe, it is reasonable to assume that the black hole essentially evaporates only when the age of the Universe is of order \( \tau_{bh} \).

Let us assume that black holes formed in the early Universe have masses so that they evaporate when the temperature of the Universe is \( T_U < T_{ew} \). The age of the Universe \( t_U \) when its temperature is \( T_U \) is

\[ t_U \simeq 0.3 g_*^{-1/2} M_p T_U^{-2} \]  

(4)

Here \( g_* \) is the number of degrees of freedom relevant at temperature \( T_U \). We are interested in black holes that decay after the electroweak phase transition but before the onset of nucleosynthesis at \( T_U = 1 \) MeV. For concreteness, we shall consider black holes decaying at \( T_U \) equal to 1 GeV and 10 GeV. For \( T_U \) in the range 1-100 GeV, \( g_* \) is equal to 100. From eqn.(3) we can then calculate \( M_0 \) so that \( \tau_{bh} = t_U \). We get (by using \( g_* = 100 \)),

\[ M_0 = 0.07 M_p^{5/3} T_U^{-2/3} \]  

(5)

The temperature of this black hole is

\[ T_{bh} = 0.6 M_p^{1/3} \tau_{bh}^{2/3} \]  

(6)

For \( T_U = 1 \) GeV we get \( M_0 = 4 \times 10^{11} M_p, T_{bh} = 1 \times 10^6 \) GeV and \( \tau_{bh} = 5 \times 10^{17} \) GeV\(^{-1} \) = 3 \times 10^{-7} \) s. The picture then is that these black holes emit particles with energies roughly equal to \( 10^6 \) GeV into the background plasma which is at a temperature \( T_U \) (\( \sim 1 \) GeV) to start with. These \( 10^6 \) GeV particles will scatter with the particles in the background plasma and will heat it up through their energy loss. For the black hole masses considered here only elementary particles will be emitted, such as quarks, gluons, photons, leptons, etc. (Emission of quarks and gluons by black holes is a non-trivial process as discussed in [25]. We will not worry about those details here.) For \( T_U = 10 \) GeV, \( M_0 = 8 \times 10^{10} M_p, T_{bh} = 6 \times 10^6 \) GeV and \( \tau_{bh} = 4 \times 10^{15} \) GeV\(^{-1} \) = 3 \times 10^{-9} \) s.

Obtaining the temperature profile outside a black hole radiating into an ambient plasma is non-trivial. We first show below how energy from the black hole is deposited
in the surrounding plasma thereby heating it up. It is necessary to include heat transfer to allow for the energy deposited close to the black hole to move out. We first use diffusion equations from stellar physics but then raise a concern that one does not have hydrostatic equilibrium in the plasma surrounding the black hole, unlike in stellar atmospheres. Therefore the plasma also moves out due to the pressure gradient. We model this outflow of the plasma using shells. However a more complicated numerical study may be necessary to understand the simultaneous heating of the plasma and energy transfer due to diffusion and bulk flow.

The energy loss of particles traversing a region of quark-gluon plasma has been discussed extensively in the literature. The energy loss per unit distance for an energetic quark with energy $E$ traversing a quark-gluon plasma at temperature $T$ is given by

$$-\frac{dE}{dx} \approx 0.1 T^2 \ln(\sqrt{E/T}),$$

(7)

where we have taken $\alpha_s \sim 0.1$. For the black hole of mass $4 \times 10^{11} M_{Pl}$, setting $T$ equal to the ambient temperature $T_U = 1 \text{ GeV}$ and $\Delta E = 10^6 \text{ GeV}$, we get $\Delta x \sim 10^6 \text{ GeV}^{-1}$, if we approximate $E(x)$ as $T_{bh}$. One could argue that $\Delta x$ is the distance in which the energy released by the black hole is thermalised as by $\Delta x$ an emitted particle has attained an energy comparable to the thermal energy of the particles in the surrounding plasma (loosely speaking, it has thermalised). Furthermore, it takes many collisions before the radiated particle comes to the same energy as the background, in which time the heat transferred along the way would also thermalise due to collisions. If one now equates $M_{bh} = \frac{\pi^2}{30} g_* T^{14/3} r_{st}^3$ where the stopping distance $r_{st} = \Delta x$ then $T = 400 \text{ GeV}$. However we are here grossly assuming that all the emitted energetic particles go through a plasma at 1 GeV and that all the black hole energy is thermalised to a uniform temperature. Instead let us consider the black hole evaporation in steps and equate $0.01 M_{bh}$ with the increase in energy of the ambient plasma within $r_{st}$. The temperature of the plasma rises to 100 GeV. The next 1% of the black hole mass released as Hawking radiation will now see an ambient plasma at 100 GeV and hence the new stopping distance $r'_{st}$ will be $200 \text{ GeV}^{-1}$. Thus a smaller region than before is heated to an even higher temperature. This continual process can give rise to a temperature gradient in the plasma surrounding the plasma. The above also shows that naively equating the energy emitted with $\rho \frac{4\pi}{3} r_{st}^3$, as we did above, will not

\footnote{This expression was derived for energies lower than $10^6 \text{ GeV}$ but since $\alpha_s$ is less for higher $E$ the perturbative derivation of the energy loss is even more accurate for our case.}
be correct. Though the plasma within $r_{st}$ rises to $T_{ew}$ even with the energy of 1% of the black hole we do not use this to estimate the volume in which the symmetry is restored as it is not clear what fraction of the black hole mass, emitted as radiation, sees a plasma at the initial ambient temperature of 1 GeV.

To obtain the temperature profile outside the black hole one must include heat transfer from the regions close to the black hole to outer regions. (In fact, if we ignore heat transfer the temperature of the plasma close to the black hole keeps rising until the plasma temperature becomes equal to the black hole temperature. After this the black hole can not evaporate any further until the energy deposited close to the black hole moves out.) To include the effects of diffusion we shall follow the calculation of the temperature profile of stellar interiors. In stars one has a hot stellar core in which nuclear reactions occur, and a stellar atmosphere surrounding the core. Energy moves through the stellar atmosphere via diffusion due to a thermal gradient. We have an analogous situation with a black hole radiating into a plasma around it and we shall assume that energy radiated by the black hole is quickly thermalised, so at any given time one has a hot black hole radiating into a plasma in which there is a thermal gradient. Energy diffuses from the hotter regions closer to the black hole to regions further away due to the thermal gradient. Under these assumptions we will use eqn. 5.11 of [27] (hereafter referred to as KW), namely,

$$\frac{\partial T}{\partial r} = -\frac{3}{16\pi ac} \frac{\kappa \rho l}{r^2 T^3}, \quad (8)$$

where $c = 1$ is the speed of light, $a = (\pi^2/30)g_\ast$. $\kappa \rho$ is the radiative cross-section per unit mass averaged over frequency times the energy density and equals $\sigma n = 1/\text{mean free path}$, where $\sigma$ is the cross section for scattering and $n$ is the number density of particles. The local luminosity $l(r)$ is defined as the net energy passing outward through a sphere of radius $r$ per second (see Section 4.2 of KW).

To integrate $\partial T/\partial r$, we need $\sigma n$ and $l(r)$. We use

$$\sigma n = (1/9)(2) \frac{4\pi \alpha_s^2}{3(4T^2)} \times (3/4)(1.2/\pi^2)g_\ast T^3. \quad (9)$$

The above expression for $\sigma$ is the cross section for quark-antiquark scattering (see eqn. 6.33 of [28]). The factors of $(1/9)(2)$ come from averaging over color and tracing over QCD generators, respectively. We have set the Mandelstam variable $s$ in the formula to equal $(2T)^2$, i.e., for incoming particles of equal (and opposite) momenta $T$. Then we find that $\sigma n = 2 \times 10^{-4}g_\ast T = 0.02T$, where we have used $\alpha_s = 0.1$.

In stars $l(r)$ can be a complicated function, depending on the distribution of sources and sinks of energy. Since we are interested in the steady state situation we shall take
the luminosity to be independent of \( r \) and equal to the energy radiated per unit time by the black hole. Therefore

\[
l(r) = -\frac{dM_{bh}}{dt} = \alpha M_{Pl}^4/M_{bh}^2.
\]  

(10)

Now one can integrate \( \partial T/\partial r \), with the boundary condition that \( T = T' \) at \( r = r' \). We shall initially take \( T' \) to be the ambient temperature \( T_U \). A priori one does not know \( r' \). Due to scattering of radiation in the plasma \( r' \) should be less than \( c\tau_{bh} \). We shall fix \( r' \) by requiring that the increase in the energy of the plasma for \( r \leq r' \) is equal to the radiated mass of the black hole. We get

\[
T^3 - T'^3 = \left( \frac{9}{16\pi ac} \right) \frac{\alpha M_{Pl}^4}{M_{bh}^2} \left( 0.02 \right) \left[ \frac{1}{r} - \frac{1}{r'} \right].
\]  

(11)

We shall apply this solution for \( r \geq 6R \), where \( R = 2M_{bh}/M_{Pl}^2 \) is the Schwarzschild radius of the black hole. \( 6R \) is the radius of the thermal atmosphere around the black hole \[29\] beyond which curvature effects of the black hole are negligible.

Our assumption of \( l(r) \) independent of \( r \) implies a steady state profile. However our choice of the boundary condition that the increase in energy of the plasma within \( r' \) equals the energy radiated by the black hole implies that all the black hole energy is used to create the temperature profile. More likely, the energy radiated from the black hole will first create a profile and then the remaining energy of the black hole will be transmitted through the plasma without further heating the plasma. This implies that the increase in energy of the plasma will be less than the radiated mass of the black hole. Therefore \( T(r) \) for any \( r \) will be less than what we obtain and our estimates of \( r_{100} \) in the diffusion picture will be an upper bound.

We start with a black hole of mass \( M_{bh} = 4 \times 10^{11} M_{Pl} \) and lifetime \( 5 \times 10^{17} \text{GeV}^{-1} \) which decays when the temperature of the universe is 1 GeV. Using the temperature profile in eqn.(11) we apply the energy constraint mentioned above to (numerically) obtain \( r' = 2 \times 10^9 \text{GeV}^{-1} \), and thus an expression for \( T(r) \). While applying the energy constraint we assume that only 90% of the black hole mass is radiated away. Using \( T(r) \) we then obtain \( r_{100} \), the radius below which the temperature is greater than 100 GeV, to be \( 4 \times 10^2 \text{GeV}^{-1} \).

Now we envision that the black hole has reduced in mass by 90% and so we have a black hole of mass \( 4 \times 10^{10} M_{Pl} \) and lifetime \( \tau = 5 \times 10^{14} \text{GeV}^{-1} \). We once again use the above equation and find that the energy constraint for a black hole that radiates 90% of its mass is consistent with \( r' = 1 \times 10^8 \text{GeV}^{-1} \) and \( T(r') = 2 \text{GeV} \). (We vary \( r' \) and \( T(r') \) so that we satisfy the energy condition for a value of \( r' \) such that \( T(r') \) is consistent with the
temperature profile for $M_{bh} = 4 \times 10^{11} M_{Pl}$. Also, the energy in the region $r \leq r'$ obtained from the temperature profile of the black hole at $M_{bh} = 4 \times 10^{11} M_{Pl}$ is much smaller than the energy deposited in this region by the black hole after its mass has reduced to $4 \times 10^{10} M_{Pl}$.) We now find that $r_{100} = 3 \times 10^4 \text{GeV}^{-1}$.

Now let the mass of the black hole be $4 \times 10^9 M_{Pl}$. Its lifetime is $5 \times 10^{11} \text{GeV}^{-1}$. We repeat the procedure above and obtain a consistent temperature profile with $T' = 30 \text{GeV}$ at $r' = 1 \times 10^6 \text{GeV}^{-1}$. We find that $r_{100} = 8 \times 10^5 \text{GeV}^{-1}$. Again the energy deposited greatly exceeds the energy already in the plasma for $r \leq r'$. We find that $r_{100} = 8 \times 10^5 \text{GeV}^{-1}$. Since $r_{100} \sim r'$ we can not get a larger $r_{100}$ as the black hole mass continues to reduce further. Thus for a black hole of mass $4 \times 10^{11} M_{Pl}$ radiating into a plasma (initially) at a temperature of 1 GeV we obtain $r_{100} = 8 \times 10^5 \text{GeV}^{-1}$.

Had we started with an initial ambient temperature of 10 GeV, and a black hole of mass $8 \times 10^{10} M_{Pl}$, temperature $6 \times 10^6 \text{GeV}$ and lifetime $4 \times 10^{15} \text{GeV}^{-1}$ or $3 \times 10^{-9} \text{s}$, we would first obtain $r' = 8 \times 10^7 \text{GeV}^{-1}$ for $T' = 10 \text{GeV}$ and $r_{100} = 8 \times 10^3 \text{GeV}^{-1}$. Allowing the black hole to further radiate with a mass of $8 \times 10^{9} \text{GeV}$ we find $r' = 5 \times 10^6 \text{GeV}$ with $T' = 10 \text{GeV}$. Then we obtain $r_{100}$ of $6 \times 10^5 \text{GeV}^{-1}$.

If one calculates the energy radiated per unit time into the black hole by the plasma at $r = 6R$, $P = \sigma(T(6R))^4 \pi (6R)^2$, where $\sigma = (1/4)(\pi^2/30)g_*$ is the Stefan-Boltzmann constant modified for more than the two photonic degrees of freedom, one finds that this is always a factor of $10^3$ less than $-dM_{bh}/dt|_{\text{vacuum}} = \alpha M_{Pl}^4/M_{bh}^2$. Therefore we can ignore any decrease in the net luminosity due to absorption by the black hole.

Though we have used the stellar diffusion equations to obtain the temperature profile there is one important difference between our situation and that of stellar atmospheres, namely, the existence of hydrostatic equilibrium. While gravity in stars provides hydrostatic equilibrium we show below that pressure gradients in the plasma will overcome the gravitational attraction of the black hole and the plasma will move out. This is to be expected as these black holes have much less mass and plasma temperatures are much larger as compared to stellar situations.

The gravitational force (per unit volume) on an element of the plasma of density $\rho$ and at a distance $r$ from the black hole is $(\rho + P)g$, where $g = G_N M/r^2$ is the gravitational acceleration and $P = \frac{1}{3} \rho$ is the pressure at a distance $r$ from the black hole. $M$ here includes the black hole mass and the energy plus pressure of the plasma within the sphere at radius $r$. One can show that the contribution of the plasma to $M$ is negligible. Therefore the gravitational force (per unit volume) is
The outward force (per unit volume) on the plasma element is given by $\nabla P = \partial P/\partial r$. Now $P = \frac{\pi^2}{90}g_*T^4$. Therefore

$$\partial P/\partial r = \frac{4\pi^2}{90}g_*T^3\partial T/\partial r.$$  \hspace{1cm} (13)

For $M_{bh} = 4 \times 10^{11}M_{Pl}$ we find that outward force due to the pressure gradient is much larger than the gravitational attractive force for regions where $T < 1 \times 10^5$ GeV. At $r_{100}$ $\partial P/\partial r$ is 9 orders of magnitude larger than the gravitational force. For $M_{bh} = 8 \times 10^{10}M_{Pl}$ gravity becomes sub-dominant at $T = 7 \times 10^5$ GeV and the disparity between the pressure gradient and the gravitational force at $r_{100}$ is even larger. Thus one should include effects of bulk flow in the modeling of the plasma outside the black hole.

We now calculate the distance at which heat transfer becomes dominated by bulk flow rather than diffusion. As we have shown, beyond some distance $r_a$ hydrostatic equilibrium is not maintained and the plasma acquires some outward bulk velocity. However bulk plasma motion will not be the dominant mode of heat transfer till the plasma velocity increases sufficiently. Let $r_c$ be the distance at which heat transfer due to bulk flow becomes dominant.

Starting from the Euler equation for a relativistic fluid \[30\] we obtain the following equation for steady-state situations (i.e. after setting time derivatives to zero)

$$v\gamma^2dv/dr = -1/(\rho + P)\partial P/\partial r,$$  \hspace{1cm} (14)

where $\gamma = 1/\sqrt{1-v^2}$ and $c = 1$. Using the expression for the pressure of a relativistic fluid, $P = (1/3)\rho$, we get

$$v/(1-v^2)dv/dr = -(1/T)(\partial T/\partial r).$$  \hspace{1cm} (15)

We would like to integrate the above equation with the range of integration for $v$ and $r$ being 0 to $v_c$ and $r_a$ to $r_c$ respectively. Since heat transfer in this region is still dominated by diffusion we use $\partial T/\partial r = -B/(v^2T^2)$ from eq. (8) above, where $B = (90/(1600\pi^3))0.06 \times 10^{-3}(M_{Pl}^4/M_{bh}^2)$. We obtain the expression for the temperature profile $T(r)$ as

$$T^3 = T_0^3 + 3B(1/r - 1/r_0),$$  \hspace{1cm} (16)

where $T_0 = \eta T_{bh}$ and $r_0 = 6R$. Since we are trying to ascertain the distance $r_c$ at which the diffusion approximation breaks down, we can not impose a boundary condition at large $r$,
i.e., at \( r' \) as we have done earlier. However, our ignorance of the plasma temperature at small \( r = 6R \) leads us to introduce the parameter \( \eta \). \( \eta \) is then a measure of how different the temperature at \( r = 6R \) is from the black hole temperature \( T_{bh} \). To know the precise value of \( \eta \) would require a complicated numerical simulation of the processes occurring close to the black hole. Therefore we instead solve for \( r_c \) for different values of \( \eta \).

For a black hole of mass \( 4 \times 10^{11} M_{Pl} \), if we substitute values of \( \eta \geq 0.12 \) in eq. (16) then even at \( r = \infty \) the temperature is greater than \( T_a = 1 \times 10^5 \text{ GeV} \). Since hydrostatic equilibrium breaks down for temperatures below this value this implies that hydrostatic equilibrium is maintained for all \( r \) and that the diffusion approximation is valid everywhere. However the large value of \( T \) at infinity for these values of \( \eta \) implies that the entire universe is heated up by a single black hole which violates energy conservation and shows that \( \eta \geq 0.12 \) should not be used. For \( \eta = 0.10 \), \( r_a \) is \( 6R = 4 \times 10^{-7} \text{ GeV}^{-1} \) and so for smaller values of \( \eta \) one also takes \( r_a \) equal to \( 6R \). For values of \( \eta \) between 0.12 and 0.13 \( r_a \) is greater than \( 6R \) and is obtained from eq. (16) to be \( 3.4 \times 10^{-11} (2.0 \times 10^{-3} - \eta^3)^{-1} \text{ GeV}^{-1} \).

\( v_c \) is then obtained from eq. (15) as

\[
v_c^2 = 1 - \exp\left(-\int_{r_a}^{r_c} \frac{2B}{r^2 T^3} dr\right).
\]  

(17)

\( r_c \) can be obtained by equating the flux due to bulk flow at \( r_c \), \( J(r_c) = \gamma^2 (\rho + P)(r_c) v_c \), with the total flux at \( r_c \), i.e.,

\[
\gamma^2 \frac{2 \pi^2}{45} g_* T_c^4 v_c = \frac{l}{(4 \pi r_c^2)}
\]

(18)

where \( T_c \) is the temperature at \( r_c \) and \( l \) is the luminosity at \( r_c \), which we take to be \(-dM_{bh}/dt\). This finally gives

\[
\gamma^2 \frac{2 \pi^2}{45} g_* T_c^4 r_c^2 v_c = \left(\alpha M_{Pl}^4\right)/(4 \pi M_{bh}^2)
\]

(19)

By substituting the expressions for \( v_c \) and \( T_c \) obtained in terms of \( r_c \), it is in principle possible to obtain \( r_c \). However the above equation cannot be solved analytically. Instead we numerically estimate \( r_c \) by plotting (LHS-RHS) of eqn.(19) as a function of \( r_c \) to ascertain at what point the graph crosses the \( r_c \)-axis. Substituting this value into eqs. (17) and (16) gives \( v_c \) and \( T_c \) respectively.
As mentioned earlier, we present results for certain values of $\eta$. For a black hole of mass $M_{bh} = 4 \times 10^{11} M_{Pl}$ evaporating when the ambient temperature of the Universe is $T_U = 1$ GeV we find $r_c = 9 \times 10^{-6}$ GeV$^{-1}$ and $3 \times 10^{-6}$ GeV$^{-1}$ for $\eta = 0.10$ and 0.08 respectively. The corresponding values of $v_c$ are 0.5 and 0.8 respectively. The corresponding values of $T_c$ are $1 \times 10^5$ GeV and $6 \times 10^4$ GeV respectively. For a black hole of mass $M_{bh} = 8 \times 10^{10} M_{Pl}$ evaporating when the ambient temperature of the Universe is $T_U = 10$ GeV we find $r_c = 2 \times 10^{-6}$ GeV$^{-1}$ and $5 \times 10^{-6}$ for $\eta = 0.10$ and 0.08 respectively. The corresponding values of $v_c$ are 0.5 and 0.8 respectively. The corresponding values of $T_c$ are $5 \times 10^5$ GeV and $3 \times 10^5$ GeV respectively. In all the above cases the temperature of the plasma at $r_c$ is much greater than 100 GeV indicating that $r_c$ is smaller than $r_{100}$. This shows that bulk flow dominates over diffusive transfer of energy in the plasma before the temperature of the plasma drops down to the electroweak transition temperature, thereby calling in question the validity of the diffusion picture.

To study the outward motion of the plasma for $r > r_c$ we shall consider the plasma as consisting of outgoing shells of infinitesimal thickness $dr$ (in the frame of the black hole). As these shells move out, they cool down due to the expansion of the plasma. Eventually these shells will reach a distance $r_{100}$ where the temperature (as defined in the frame of the plasma) becomes smaller than 100 GeV. At this time the electroweak transition will occur in these shells and, as we will discuss later, baryogenesis can take place there. Rewriting eqn.(15) and integrating we get

$$v^2 = 1 - \left(\frac{T}{T_c}\right)^2(1 - v_c^2),$$

which gives us a relation between the velocity of the outward moving plasma shell and the temperature of the plasma.

Since the temperature decreases with increasing $r$, the velocity of the plasma shell will keep increasing and will eventually become very close to the speed of light. The Lorentz factor is given by

$$\gamma = \frac{1}{\sqrt{1 - v_c^2}} \frac{T_c}{T}.$$  \hspace{1cm} (21)

For $T_U = 1$ GeV the Lorentz factor $\gamma_{100}$ at $r_{100}$ is 1000 for $\eta = 0.10$ and 0.08. For $T_U = 10$ GeV, $\gamma_{100}$ is 6000 and 5000 for $\eta = 0.10$ and 0.08 respectively. Below we shall take $v_{100}$, the velocity of the plasma shell at $r_{100}$, to be 1. Later we shall briefly mention the possible role of viscosity in slowing down the plasma and its consequences.

Assuming a steady state situation $l(r)$ is independent of $r$ and equal to $-dM_{bh}/dt$. Therefore for $r \geq r_c$
\[ r = \left[ \frac{l}{\gamma^2(8\pi^3 g_*/45)T^4 v} \right]^{\frac{1}{4}}. \] (22)

Above \( T \) is the temperature of the plasma shell at a distance \( r \). For a black hole of mass \( 4 \times 10^{11} M_{Pl} \), \( r_{100} \) is \( 7\gamma^{-1} \text{ GeV}^{-1} \) and it is \( 30\gamma^{-1} \text{ GeV}^{-1} \) for a black hole of mass \( 8 \times 10^{10} M_{Pl} \).

The volume relevant for baryogenesis must be calculated in the rest frame of the plasma since the sphaleron rates have been calculated in this frame. Consider a plasma shell of radius \( r_{100} \) in the black hole frame. There is no frame in which the entire shell is at rest. Instead one can divide the shell into plasma elements of area \( d\Omega r_{100}^2 \) and thickness \( dr \). In the rest frame of such a plasma element the volume element is \( d\Omega r_{100}^2 dr' \), where \( dr' = \gamma dr \) is the shell thickness in the frame of the plasma element. Therefore the volume relevant for baryogenesis contributed by one shell is

\[ dV = 4\pi r_{100}^2 dr' = 4\pi r_{100}^2 \gamma_{100} dr = 4\pi r_{100}^2 \gamma_{100} v_{100} dt. \] (23)

Plasma shells keep moving across \( r_{100} \) for a duration of the order of the lifetime \( \tau \) of the black hole. Thus the total volume relevant for baryogenesis is then given by

\[ V = 4\pi r_{100}^2 \gamma_{100} v_{100} \tau = 3\gamma^{-1}_{100} l \tau / (4\rho_{100}), \] (24)

where we have used eq. (22) for the second equality. Since \( l = -dM_{bh}/dt = \alpha M_{Pl}^4 / M_{bh}^2 \), we may also express the volume as

\[ V = 7 \times 10^{-11} \text{ GeV}^{-4} \gamma^{-1}_{100} M_{bh} = 5 \times 10^{-12} \text{ GeV}^{-4} \gamma^{-1}_{100} M_{Pl}^{5/3} / T_{Pl}^{2/3} \] (25)

where we have used eq. (2a) for the last equality. For a black hole of mass \( 4 \times 10^{11} M_{Pl} \) radiating into a plasma initially at 1 GeV, \( V \) is \( 3 \times 10^{20} \gamma^{-1}_{100} \text{ GeV}^{-3} \); for a black hole of mass \( 8 \times 10^{10} M_{Pl} \) and an initial ambient temperature of 10 GeV, \( V \) is \( 7 \times 10^{19} \gamma^{-1}_{100} \text{ GeV}^{-3} \).

Electroweak baryogenesis takes place via sphaleron processes with sphaleron size \( \sim (\alpha_W T_{ew})^{-1} \approx (3 \text{ GeV})^{-1} \). Since the thickness of the shell in the plasma frame for a temperature difference of 1 GeV, i.e., \( r'_{99} - r'_{100} \), is proportional to a Lorentz factor, which compensates for the \( \gamma \) factor in \( r_{100} \), one can easily fit the sphaleron in the plasma shell at \( r_{100} \).

### III. ESTIMATION OF THE BARYON ASYMMETRY

Let us assume that the number of excess baryons created around each black hole is \( N_B \). If the number density of black holes is \( n_{bh} \) then the number density \( n_B \) of baryon...
excess produced by these black holes is $n_B = N_B n_{bh}$. To estimate $n_{bh}$ we assume that the net energy density of black holes does not dominate the energy density of the Universe at the time when black holes decay. (With this assumption we are assured that black holes never dominate the dynamics of expansion of the Universe so the overall picture of the evolution of the Universe remains unchanged.) This gives the following expression for $n_{bh}$ (using eqn.(5)).

$$n_{bh} \lesssim 5 \times 10^{2} T_U^{14/3} M_{pl}^{-5/3}$$

(26)

For $T_U = 1$ GeV, $n_{bh} \lesssim 8 \times 10^{-30}$ GeV$^3$, while for $T_U = 10$ GeV, $n_{bh} \lesssim 4 \times 10^{-25}$ GeV$^3$. We shall use the maximum allowed value of $n_{bh}$ below.

The entropy density of the Universe (which will remain essentially unchanged even after all the black holes evaporate, since the black holes do not dominate the energy density of the Universe) is given by

$$s = \frac{2\pi^2}{45} g_* T_U^3 \simeq 40 T_U^3$$

(27)

We now proceed to estimate the baryon asymmetry produced in the hot regions by models of baryogenesis conventionally used at the electroweak scale. In the following we shall use $V$ to refer to the hot region where the temperature rises to at least $T_{ew}$ as well as to the volume of this region. Let us first assume that the electroweak transition is of first order. Before the evaporation of the black hole the Universe is in a state with the expectation value of $\phi$ at a value $v$. When the regions of volume $V$ heat up, the expectation value of $\phi$ in $V$ will evolve to the value 0. (Later we consider the possibility that $\phi$ may not roll to the minimum due to insufficient heating.) As the plasma in $V$ cools further, a local minimum at a non-zero value of $\phi$ appears in the effective potential. For $T < T_{ew}$, this minimum becomes the true minimum of the potential. Bubbles of true vacuum now appear in the hot region $V$ as the field tunnels from 0 to $v$. As these bubbles grow, they either produce a baryon asymmetry in the outer regions of a thick wall by “spontaneous baryogenesis” [31] or by the asymmetric reflection of particles in the symmetric phase from the surface of a thin bubble wall [32]. In the former case, the non-zero value of the time derivative of a field coupled to the baryon number acts as a chemical potential for baryon number, thereby biasing $B+L$ violating sphaleron processes in the wall to produce more baryons than anti-baryons. In the latter case, the asymmetry created in some quantum number not orthogonal to baryon number due to reflections off the bubble wall is converted into a baryon asymmetry by sphaleron processes. This baryon asymmetry is frozen after
it passes through the wall. A more sophisticated treatment of CP violation, particle transport and diffusion in the context of electroweak baryogenesis is discussed in refs. [33–37]. A baryon asymmetry can also be produced at the electroweak phase transition through the CP asymmetric interaction of the bubble wall with fluctuations in the winding number of gauge-Higgs fields configurations [38,39] and through the interference between electroweak sphaleron-induced baryon number violating processes and QCD sphaleron-induced CP violating processes [40]. In the standard electroweak baryogenesis scenario baryogenesis occurs in the entire Universe as the electroweak bubbles sweep through the Universe during the phase transition. In our case, our hot regions cover a volume fraction of $V_{n_{bh}} = 3 \times 10^{-9} \gamma^{-1} \text{GeV}^{-4} T_{U}^4$. However, the entropy density of the Universe at a temperature of $T_{U}$ is also lower by a factor of $10^6 \text{GeV}^3/T_{U}^3$ as compared to the entropy density at the electroweak phase transition. The ratio of the asymmetry that we obtain to that in the standard scenario is then volume fraction $\times$ entropy factor $= 3 \times 10^{-3} \gamma^{-1} T_{U} \text{GeV}^{-1}$.

The Standard Model does not have sufficient CP violation to create the observed baryon asymmetry of the Universe. Therefore we consider the results of ref. [36] for baryogenesis in supersymmetric models. With CP violating phases of the order of $10^{-4}$ they produce enough asymmetry for a first order electroweak phase transition. However, even with a larger value of the CP violating parameters the asymmetry that we obtain above for a first order electroweak phase transition will be insufficient because of the large value of $\gamma$.

However above we have completely ignored the effects of viscosity. Viscous effects can be of two types—bulk and shear [41]. The bulk viscosity coefficient for a (quark-gluon) relativistic plasma is approximately 0 [42]. However there is a contribution to the shear viscosity that is proportional to $\partial v/\partial r$, which will slow down the plasma as it moves out. The relativistic hydrodynamic equations with viscosity are very complicated and cannot be solved easily. If, however, the viscous effects decrease the plasma velocity at $r_{100}$ to even 0.9$c$, this will imply that $\gamma_{100}$ is 2. Then, with a value of CP violation $\sim 0.1$ or 0.01, one can obtain sufficient asymmetry in our scenario for black holes radiating into an ambient plasma at an initial temperature of 1 GeV or 10 GeV respectively.

Since our formalism above is not very rigorous we shall also estimate below the asymmetry in the diffusion picture ignoring any plasma motion due to pressure gradients. For a black hole of initial mass $4 \times 10^{11} M_{Pl}$ the volume fraction $V_{n_{bh}}$, where $V = \frac{4}{3} \pi r_{100}^3$ is $2 \times 10^{-11}$. Taking into account the entropy factor, the asymmetry is $2 \times 10^{-5}$ of the asymmetry in standard electroweak baryogenesis models. For a black hole of initial mass
$8 \times 10^{10} M_{Pl}$ the volume fraction is $4 \times 10^{-7}$. So the asymmetry is about $4 \times 10^{-4}$ of the asymmetry in standard electroweak baryogenesis models.

In the diffusion picture we have ignored the fact that once the black hole has evaporated completely energy from hotter regions will diffuse out and heat up cooler regions. This will effectively increase $r_{100}$. Furthermore, if we relax the constraint that the black holes do not dominate the Universe until they decay this may enhance the final asymmetry.

We point out an interesting phenomenon that can occur in the regions surrounding the black hole where the temperature rises but does not get much higher than $T_{ew}$. In these regions the effective potential changes from the zero temperature potential to a high temperature potential with a local minimum at $<\phi>=v$ as shown in fig. 1. In these regions the Universe is trapped in a state with $<\phi>=v$ and tunnels to the $<\phi>=0$ vacuum by bubble nucleation. Thus in these regions bubbles of $<\phi>=0$ are created. As long as the $<\phi>=v$ vacuum is higher than the $<\phi>=0$ vacuum these bubbles expand, as in fig. 2a. However as this region cools the two vacua become degenerate and then the state with $<\phi>=v$ becomes the true vacuum. If the $<\phi>=0$ bubbles have not already disappeared after collisions (leaving the Universe in the $<\phi>=0$ vacuum state) then the $<\phi>=0$ bubbles now begin to shrink, as in fig. 2b. Baryons are now created in the inner regions of the shrinking walls. Furthermore, $<\phi>=v$ bubbles may nucleate inside the $<\phi>=0$ bubbles (depending on the critical bubble size). As these expand baryons are created in their walls. The above phenomenon may occur in the diffusion picture for $r > r_{100}$ in the transition phase after the black hole evaporates away and energy from hotter regions moves out and heats up colder regions. In the case of heat transfer through bulk motion of the plasma this mechanism may occur at some distances away from the black hole in the early stages before the steady state is achieved.

We now consider a second order phase transition. In the hot regions, the symmetry is restored and the field settles at the minimum of the potential at $\phi=0$. As the region cools below $T_{ew}$ the minimum shifts from $\phi=0$. $<\dot{\phi}>$, in this case, is not due to Hubble expansion but is due to the cooling of the plasma as the plasma expands. Hence, unlike in the standard cosmology, $<\dot{\phi}>$ is large. The non-zero $<\dot{\phi}>$ acts as a chemical potential for baryon number during the phase transition and a net baryon number is created by $B+L$ violating sphaleron processes. The baryon asymmetry generated in a second order electroweak phase transition is given by [43]

$$n_B/s \approx 10^{-19} \epsilon [-\dot{T}/T]_{100}/H_{100}$$

(28)

where $H_{100}$ is the value of the Hubble constant when the temperature of the Universe
is 100 GeV, $\dot{T}/T$ is in the frame of the plasma when the temperature has fallen to 100 GeV and $\epsilon$ is a measure of CP violation.

To estimate $\dot{T}/T$ in the rest frame of the plasma when the temperature has fallen to 100 GeV we note that at $r_{100}$ in time $dt$ a shell of thickness $dr = v_{100} dt$ moves out at a temperature $T_{100}$. As the shell moves out $\dot{T} = \frac{dT}{dr} \frac{dr}{dt} \approx -T \gamma / r$, since eqs. (22) and (21) imply that $T$ is approximately proportional to $r^{-1}$. $dt'$ is the time differential in the plasma frame. Therefore $[-\dot{T}/T]_{100} = v_{100} \gamma_{100}/r_{100}$.

Using values of $r_{100}$ and $v_{100}$ obtained earlier we get $[-\dot{T}/T]_{100} = 1 \times 10^{-1} \gamma_{100}^2$ GeV and $3 \times 10^{-2} \gamma_{100}^2$ GeV for $T_U$ equal to 1 GeV and 10 GeV respectively. In the standard scenario $[-\dot{T}/T]_{100} = H_{100} = 10^{-14}$ GeV. Therefore the baryon asymmetry that we obtain in this scenario, taking into account the volume fraction and the entropy factor, is

$$n_B/s = 3 \times 10^{-9} \gamma_{100} \epsilon \text{ GeV}^{-1}$$ (29)

for $T_U = 1$ GeV and is

$$n_B/s = 9 \times 10^{-9} \gamma_{100} \epsilon \text{ GeV}^{-1}$$ (30)

for $T_U = 10$ GeV. Since $\gamma_{100} \sim 10^3$ for our choices of $\eta$ clearly one can obtain a sufficiently large baryon asymmetry in our scenario with a second order phase transition. It is interesting that while the asymmetry generated in a first order electroweak phase transition is suppressed by the Lorentz factor, the asymmetry is enhanced by the Lorentz factor for a second order phase transition. Also note that even if viscous effects reduce the value of $\gamma_{100}$ to 1 we can obtain sufficient asymmetry for $\epsilon \sim 10^{-1} - 10^{-2}$ (for $v_{100}$ not much smaller than 1). Further note that the possibility of creating a sufficient asymmetry is not very sensitive to the value of $\eta$ as long as $T_c$ is large enough ($\gtrsim 200$ GeV) to ensure that $v_{100}$ is not much smaller than 1.

Since the distance between black holes is much greater than $r_U$, the distance at which the outward moving plasma attains the ambient temperature, the heating of the plasma by one black hole does not directly affect the plasma around neighboring black holes. Once the outward moving plasma reaches $r_U$ it merges with the ambient plasma. As the energy density of the black holes does not dominate the energy density of the universe the heating of the ambient plasma between black holes, which could affect the temperature profile of neighboring black holes, is also negligible. Finally, baryons carried out by the outward moving plasma into the ambient plasma can subsequently enter the hot region around another black hole due to random Brownian motion across the intermediate region. However as the baryons cross $r_U$ of another black hole the plasma moving out from the
second black hole would push them out thereby preventing them from reaching $r_{100}$ of the second black hole, where sphaleron processes in equilibrium could destroy them.

We now briefly comment on the possibility of creating hot regions by the decay of massive particles. Massive particle decays will lead to the development of a collimated jet of particles, with each particle in the jet contributing to the heating of the plasma. From eqn.(7) for the energy loss of a particle created by the decay of this massive particle it seems very difficult to get sufficient energy to heat the region to electroweak scale. However, note that the expression for the energy loss as given in eqn.(7) has been derived by neglecting very hard collisions. Inclusion of hard collisions may lead to a rapid branching of the initial decay product which may significantly increase the energy loss. Also note that baryogenesis will only take place in regions that are at least as wide as the sphaleron size.

**IV. CONSTRAINTS ON THE HIGGS MASS**

In the electroweak baryogenesis scenario in the context of the Standard Model if the high temperature vev in the broken phase is small then the sphaleron processes are still active. These sphaleron processes will destroy the baryon asymmetry created during the phase transition. A large enough high temperature Higgs vev in the broken phase to avoid the wash-out of the asymmetry implies a small Higgs mass which is tightly constrained by the experimental lower bound on the Higgs mass. In our scenario, on the other hand, the requirement of a small Higgs mass does not apply, as we show below. Of course, this is relevant only if viscous forces are sufficient to decrease the suppression of the asymmetry due to the Lorentz factor for a first order electroweak phase transition. A large Higgs mass also implies that the electroweak phase transition is likely to be of second order in which case the asymmetry generated at the electroweak phase transition in standard electroweak baryogenesis is very small [43]. As we have already indicated, we generate sufficient asymmetry with a second order phase transition in our scenario.

The baryon number density after a first order electroweak phase transition is given by

$$n_B = n_{B_i} \exp(-\int \Gamma dt) = n_{B_i} \exp(-\int_{100}^{T'} \frac{\Gamma}{\dot{T}}dT)$$  \hspace{1cm} (31)$$

where $n_{B_i}$ is the baryon number density created during the phase transition, $\Gamma(T)$ is the rate per unit time of B violation and we choose the upper limit $T'$ to be $T_U$, the ambient temperature of the universe. More precisely, one could put $T'$ as the temperature at which the B violation rate becomes less than the Hubble expansion rate but the above choice is
sufficient. For $\dot{T}$ we use $-(\gamma v/r)T$ and $r = r_{100}T_{100}/T$. $\gamma$ is given by eqn. (21) and we let $v = 1$.

The perturbative rate of baryon number violation in the broken phase is given by

$$\Gamma = yf\alpha_W^4 T \left[ \frac{M_W^7}{\alpha_W^4 T^7} \right] \exp[-E_{\text{sph}}/T]$$

(32)

where $E_{\text{sph}} = 2M_W(T)/\alpha_W$, with $M_W(T) = M_W(0)\sqrt{1 - T^2/T_{ew}^2}$ and $T_{ew} \sim 100\text{ GeV}$ is the critical temperature. The above is valid for $2M_W(T) \ll T \ll 2M_W(T)/\alpha_W$, i.e., it breaks down close to $T_{ew}$. $y$ is obtained by comparing the B violation rate to the sphaleron rate per unit volume in eqns. (2.17) and (2.14) of ref. [44]. Choosing B-L=0 and three families we get $y=20$. $f$ includes other factors found in eq. (2.19) of ref. [44].

The non-perturbative lattice calculation of the sphaleron rate per unit volume in the broken phase close to the critical temperature $T_{ew}$ is given by

$$\Gamma_{\text{sph}}/V = x\alpha_W^4 T^4,$$

(33)

where $x$ varies from $10^{-8}$ to $10^{-5}$ for different values of the couplings in the theory. Therefore $\Gamma(T)$ is given by

$$\Gamma = xy\alpha_W^4 T.$$  

(34)

We can model the rate of baryon number violation for the range $T = 100\text{ GeV} - T_U$ as

$$\Gamma(T) = xy\alpha_W^4 T \exp[-E_{\text{sph}}/T]T_{ew}^7/T^7$$

(35)

where we ignore the temperature dependence in $M_W$. This should hence give us an upper bound on the actual rate.

We then find that the reduction in the baryon number density due to sphaleron processes is negligible. (The fraction of baryons destroyed before the temperature of the expanding plasma becomes $T_U$ is $10^{-10}$ and $10^{-11}$ for $T_U$ equal to 1 GeV and 10 GeV, respectively.) This is true even for much larger values of $x$, including the larger perturbative estimates of $x$ shown in fig. 4 of ref. [45].

The expression for the asymmetry in the second order phase transition case, adapted from the results of ref. [43], takes into account any destruction of the asymmetry when the value of the asymmetry at any time becomes greater than the equilibrium value.

The above arguments, coupled with the observation that our model works for a second order electroweak phase transition, implies that the constraint of a lower bound on the Higgs mass does not apply to our scenario. However, as we have pointed out earlier, insufficient CP violation in the Standard Model requires one to consider extensions of the Standard Model.
V. CONCLUSION

The baryon asymmetry that is created in our scenario is inhomogeneously distributed in the universe. However an estimate of the distance traversed by the baryons beyond $r_U$ (in the plasma motion case) due to random Brownian motion indicates that the baryons produced in our scenario homogenize themselves on time scales of the order of the lifetime of the black holes. Particles traveling through a plasma suffer collisions and the actual distance traveled by a particle in the time $\tau$ is $\sqrt{Nl_p}$, where $N$ is the number of collisions and $l_p$ is the mean free path [46]. Now $N = \tau/(l_p/c)$. Therefore $\sqrt{Nl_p} = (\tau c l_p)^{1/2}$. The mean free path=$1/\left(\sigma n\right)$ and, as stated earlier, $\sigma n \sim 0.02T$. If we substitute the black hole lifetime for $\tau$ then the distance traversed by the baryons beyond $r_U$ due to random Brownian motion is $5 \times 10^9 \text{GeV}^{-1}$ and $1 \times 10^8 \text{GeV}^{-1}$ for ambient temperatures of 1 GeV and 10 GeV respectively. Since the distance between black holes is $5 \times 10^9 \text{GeV}^{-1}$ and $1 \times 10^8 \text{GeV}^{-1}$ for ambient temperatures of 1 GeV and 10 GeV respectively and the inverse Hubble parameter is of the order of the black hole lifetime this indicates that the baryon density inhomogeneities can be wiped out on time scales of the order of few times the black hole lifetime.

Neutrino inflation and baryon diffusion can also homogenize any baryon number density fluctuation prior to the onset of primordial nucleosynthesis. Neutrino inflation is the process of heating up of overdense baryon density lumps in the early Universe by neutrinos which causes the lumps to expand, or inflate, thereby reducing their overdensity. Besides this mechanism, neutrons and protons also diffuse out of overdense baryon density lumps. Mechanisms of diluting overdense baryon density lumps are discussed in ref. [47]. The above arguments also apply to the diffusion scenario where the size of the baryon density inhomogeneity is given by $r_{100}$.

As we mentioned earlier, the scenario in ref. [18] has similar features to our work. In this paper one obtains the temperature profile outside the black hole in a manner similar to that used for stellar interiors. As we have pointed out, the assumption that heat transfer is primarily due to diffusion, as in stellar interiors, is not necessarily valid for the plasma surrounding a primordial black hole.

There are many open questions which remain to be explored. The issue of black hole formation with the required number density and mass spectrum for our scenario requires to be discussed in the context of some realistic model such as a first order phase transition occurring earlier in the Universe. Also one must keep in mind observational constraints on massive black holes [48–50] which may lead to strong restrictions on the spectrum of
masses of primordial black holes. A more complete numerical calculation can give us an estimate of the temperature of the plasma close to the black hole and fix \( \eta \). Finally viscous effects can be included in the relativistic Euler equation to more accurately ascertain the velocity of the plasma.

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**FIGURE CAPTIONS**

Fig.1: Effective potential when the region superheats with \(< \phi > = v\) becoming metastable.

Fig.2: (a) Nucleation of \(< \phi > = 0\) bubble in the metastable phase with \(< \phi > = v\). (b) Nucleation of \(< \phi > = v\) bubble inside a shrinking \(< \phi > = 0\) bubble.
FIG. 2.