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Aspects of the Electroweak Phase Transition*

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ABSTRACT

The electroweak phase transition is reviewed in light of some recent developments. Emphasis is on the issue whether the transition is first or second order and its possible role in the generation of the baryon asymmetry of the universe.

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1. Introduction

Key features of the observable universe are often thought to be the results of transient phenomena which occurred in the course of its evolution. The electroweak phase transition (EWPT) is such a phenomenon. However, until recently, it was thought to bear no consequences for today's universe. In a classic paper, Kuzmin, Rubakov and Shaposhnikov showed otherwise.[1] They uncovered the possibility that the matter antimatter asymmetry of today's world could have been produced during the EWPT through non-perturbative physics. In their pioneering work, they established the necessity of a first order phase transition. Various groups have since proposed a wealth of explicit mechanisms.[2-14] It has become imperative to broaden our understanding of the EWPT in order to transform a scenario of baryogenesis into an actual prediction of the baryon asymmetry of the universe (BAU) that can be confronted with observations. In the following, I will describe recent progress made in this direction using the minimal standard model as a prototype.

2. First Order vs. Second Order

2.1. The Effective Potential

The standard method of determining the order of the EWPT is to compute the effective potential $V(\phi,T)$ for the higgs vev ϕ, taking into account the coupling of the vacuum to a thermal bath of particles at a temperature of about 100 GeV. The calculations are usually done in the imaginary time formalism. Here is a real time picture. Let us construct a fictitious ϕ-wall interpolating between the two phases and held steady in the plasma by adequate sources. The goal is to compute $V(\phi,T)$ across the wall, that is, minus the total pressure in the plasma. The unbroken phase is filled with a gas of relativistic particles whose pressure is known to be $g^* \frac{\pi^2}{90} T^4$ with $g^* \sim 100$. The pressure in a given region of space is the latter supplemented with a pure higgs contribution $-V(\phi)$ and the total momentum exchanged between the wall and the plasma integrated up to this point. The latter is easily computed for a single particle via conservation of energy, $k^2 + m(\phi)^2 =$ constant.[9] This yields the following result

$$V(\phi,T) = V(\phi) - g^* \frac{\pi^2}{90} T^4 - \sum_{species} \int_0^{m(\phi)^2} dm^2 \int \frac{d^3k}{(2\pi)^3} \frac{n(E)}{2E} + \cdots \qquad (1)$$

$$V(\phi,T) = D(T^2 - T_o^2)\phi^2 - ET\phi^3 + \frac{\lambda_T}{4}\phi^4 + \cdots \qquad (2)$$

with D, E, λ_T and T_o functions of the couplings of the theory. The cubic term originates from the bosonic sector: $E \sim m_W^3, m_Z^3$. To see that, note that the Bose-Einstein distribution $n(E)$ behaves as $\frac{T}{m}$, as $k \to 0$. Inputting this information in Eq. (1) readily generates a term $\sim m^3$. The dots indicate that Eq. (1) corresponds to a one loop calculation. It is known that there are infrared divergences at higher order in the bosonic sector which contribute effectively as $(\frac{gT}{m})^n$. One can worry whether, once properly taken into account, they wouldn't wash away the one loop infrared effect. An answer, using imaginary time techniques, was recently proposed.[6,7] Its real time counterpart goes as follows. Gauge interactions in the plasma affect the low momentum behavior of the gauge bosons, giving them an additional mass Π which is non vanishing in the limit of small ϕ. Π is the polarization tensor of the zero modes of the gauge bosons computed in *imaginary time*. This amounts to reducing the population of the long wavelength excitations of the gauge fields from T/m to the smaller value $T/\sqrt{m^2+\Pi}$ by screening them with a cloud of $SU(2)$ charged particles, a dramatic effect in the limit $\phi \to 0$. An infrared improved effective potential $V^I(\phi, T)$ can then be computed by the above method with the replacement $n(E) \to n^{\Pi}(E) = (\exp\beta\sqrt{k^2+m^2+\Pi} - 1)^{-1}$. One obtains Eq. (2) with the substitution

$$-ET\phi^3 \to -\frac{T}{12\pi} \sum_{W_\parallel, Z_\parallel, W_\perp, Z_\perp} (m^2(\phi) + \Pi)^{3/2} - (\Pi)^{3/2} \tag{3}$$

$\Pi_{W_\parallel, Z_\parallel}$ can be computed perturbatively. In leading order, it is $\sim g^2T^2$ and dominates the mass term in the interpolating wall between the two phases; in such a case, its contribution to Eq. (3) is polynomial in ϕ^2 which only slightly corrects the parameters D and λ_T. $\Pi_{W_\perp, Z_\perp}$ is not computable perturbatively but is believed[a] to be at most $\sim g^4T^2$; in such a case, it can be ignored[b] and the RHS of Eq. (3) turns into a term cubic in ϕ. As a result, one obtains

$$V^I(\phi, T) = D'(T^2 - T_o^2)\phi^2 - \frac{2}{3}E'T\phi^3 + \frac{\lambda_T'}{4}\phi^4 + \dots \tag{4}$$

where now the dots refer to perturbative corrections. The presence of the cubic term in this improved effective potential implies a first order phase transition. This conclusion is subject to caveats. (1) The phase transition would be of 2^{nd} order if the magnetic mass, $\Pi_\perp$, were found to be significantly larger. (2) Gleiser and Kolb[10] argued, on the basis of an ϵ expansion applied to a simpler system with similar small ϕ behavior, that Eq. (4) doesn't describe properly long range fluctuations in the scalar sector.[b] They expect a more weakly first order phase transition.

2.2 Completion of the Phase Transition

The universe supercools until thermal fluctuations are able to destabilize the system at a significant rate. This corresponds to the nucleation of "critical bubbles". These bubbles evolve to a macroscopic size $\sim 10^{-5}\frac{M_{planck}}{T^2}$ before they collide and fill the universe. During this short period ($\frac{\Delta t}{t} \sim 10^{-2}$), baryogenesis takes place in the propagating bubble walls. Alternatives to the scenario above have been proposed. For instance, Kolb and Gleiser[9] have argued that long range fluctuations are so rapid at the phase transition that the universe is more adequately described by an emulsion of subcritical domains of both vacua which smoothly interpolates between the two phases as the universe cools down. This scenario has been strongly criticized[11] and shown to be possibly relevant for a range of parameters orthogonal

[a] Except in a region of small ϕ which doesn't affect the qualitative behavior of $V^I(\phi, T)$.[8]

[b] These fluctuations are ignored in Eq. (4). An assumption believed to be reasonable if the higgs mass is below 150 GeV.

to the ones which allow baryogenesis. This comes about by requiring the freezing-out of the baryon violating processes in the broken phase in order to prevent the washing out of the BAU; in the standard minimal model[7] and with the use of Eq. (4), this requires a higgs mass no larger than 40 GeV, far below the experimental limit.

3. Bubble Wall Dynamics

3.1 Bubble Wall Velocity

All the scenarios of baryogenesis make convenient assumptions on the shape and velocity of the wall. In the scenarios making use of the quantum mechanical reflection of top quarks,[4] the wall thickness is assumed to be of the order of the Compton wavelength of the reflected particles in order to prevent an excessive suppression. This condition requires particular conditions, the wall thickness being typically one or two orders of magnitude too large. In the scenarios of baryogenesis inside the wall,[3] the velocity is assumed to be large enough to prevent the BAU to diffuse and be washed out in the unbroken phase but it is assumed to be slow enough to maximize the production rate.[12,3] The physics of the damping of the wall was only recently understood.[13,7] The moving wall sets the plasma out of equilibrium $n(E) \to n(E) + \delta n(E)$ by an amount proportional to the velocity. This excess $\delta n(E)$ generates, in turn, according to Eq. (1) an additional component to the pressure which grows until it balances the difference in pressure across the wall. This condition fixes the velocity v_w to be in the range of 0.05 to 1, depending on the parameters. A range favorable for baryogenesis.

3.2 Bubble Wall Stability

A velocity smaller than the speed of sound ~ 0.6 is characteristic of a deflagration process. It is common belief that a deflagration front is unstable under perturbations whose size is large enough to overcome the surface tension $\sim \frac{\sigma}{T^3}$. These instabilities have been contemplated in the context of both the EWPT[14] and the QCD phase transition.[15] However, Landau's original stability analysis[17] was designed for violent macroscopic phenomena. A recent linear stability analysis[16] was tailored for more general phenomena and in particular, for the EWPT where it was shown that no perturbation can destabilize the moving wall in the allowed range of velocity. The reason goes as follows. A perturbation of the front triggers fluctuations in the temperature and velocity of the plasma in both phases, all of which are entangled by conservation of energy momentum. Additional information on the microscopic dynamics has to be input to determine completely the subsequent evolution of the perturbation. Landau assumed that the relative velocity wall plasma is unaffected by the fluctuations. However, we learned above that, in the EW case, the wall plasma velocity is proportional to $V^I(\phi, T)$, a sensitive function of the temperature, and, consequently, varying significantly as the temperature fluctuates. This effect tends to oppose to the growth of the perturbation. This sensitivity is measured by a dimensionless parameter which turns out to be so large $\sim \frac{1}{v_w}$ that it prevents the perturbation from growing at all.

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