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INFLATION AT THE ELECTROWEAK SCALE

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ABSTRACT: We present a simple model for slow-rollover inflation where the vacuum energy that drives inflation is of the order of G_F^{-2}; unlike most models, the conversion of vacuum energy to radiation ("reheating") is moderately efficient. The scalar field responsible for inflation is a standard-model singlet, develops a vacuum expectation value of the order of 4×10^6 GeV, has a mass of order 1 GeV, and can play a role in electroweak phenomena. We also discuss models where the energy scale of inflation is somewhat larger than the electroweak scale, but still well below the unification scale.

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Over the past decade cosmologists have come to realize that elementary-particle physics plays a very important role in cosmology: Microphysical events that took place during the earliest moments of the Universe ($t \ll 10^{-5}$ sec) and involved very high energies ($E \gg$ GeV) likely hold the key to understanding some of the most puzzling features of the Universe today. For example, baryon-number, C, and CP violating interactions occurring early on can explain the net baryon number of the Universe (baryogenesis [1]); the ubiquitous dark matter may be comprised of relic elementary particles (particle dark matter [2]); an early period of rapid expansion may account for the smoothness and spatial flatness of the Universe (inflation [3]); and a variety of early Universe scenarios have been proposed to explain the origin of the density inhomogeneities necessary to seed the formation of structure in the Universe (inflation, cosmic strings [4], and textures [5]).

Until recently it appeared that the "input microphysics" for these intriguing speculations involved energies of the order of 10^{14} GeV or larger (grand-unification scale), well beyond the "reach" of terrestrial experiments. However, scenarios for baryogenesis based upon physics at the electroweak scale have been put forth [6], and here we propose a simple model for inflation at the electroweak scale. The appeal of early Universe scenarios based upon physics at the electroweak scale, of course, is the possibility that the underlying physics can be tested in the near future, e.g., at LEP (CERN), at the Tevatron (Fermilab), or at the SSC.

Historically, inflation [7] developed from an attempt to solve the monopole problem associated with grand-unified theories (extreme overproduction of magnetic monopoles during the GUT phase transition [8]), and thus involved unification-scale energies. Further, because the baryon asymmetry of the

Universe must be produced *after* inflation and most scenarios for baryogenesis involve superheavy particles and unification-scale physics, it seemed necessary that inflation involve a very-high energy scale. Indeed, in essentially all models of inflation the vacuum energy that drives inflation is of the order of $(10^{14}\,\mathrm{GeV})^4$ [9]. Moreover, in some models of inflation—chaotic inflation [10], inflation based upon a simple supergravity model [11], and extended inflation [12]—the energy scale of inflation is set by requiring density perturbations of an appropriate size, and in these models that energy scale *must* be of the order of the unification scale.

In this *Letter* we discuss a simple model of inflation where the vacuum energy that drives inflation can be as small as the electroweak scale ($\approx 1\,\mathrm{TeV}$). We begin the description of our model by reviewing the requirements that a "successful" model of inflation must satisfy [3, 13]:

(1) Sufficient inflation to solve the horizon and flatness problems. This corresponds to $N \sim 30 + \ln(T_{\mathrm{RH}}/1\,\mathrm{TeV})$ e-foldings of the cosmic-scale factor during inflation, where T_{RH} is the temperature at the beginning of the post-inflation, radiation-dominated epoch. When the energy scale of inflation is smaller, the required amount of inflation is less.

(2) Density perturbations of appropriate size: $\delta\rho/\rho \approx 10^{-5}$ (most difficult requirement to satisfy). Density perturbations must be large enough to initiate structure formation and small enough to be consistent with the smoothness of the cosmic background radiation (CBR). Moreover, the recent detection of temperature anisotropies in the CBR on angular scales larger than about 10° by the COBE DMR [14] allows us to be more precise about the amplitude of the density perturbations.

(3) Sufficiently-high reheat temperature. The Universe must be radiation

dominated by the epoch of primordial nucleosynthesis ($T_{\rm RH} \gg 1$ MeV) so that nucleosynthesis proceeds in the usual way, and, hot enough after inflation for baryogenesis to take place, as any pre-inflation baryon asymmetry is diluted exponentially by the enormous entropy release associated with reheating. While it was thought that baryogenesis required temperatures in excess of 10^{10} GeV or so, interesting models now exist where baryogenesis occurs at the electroweak scale [6] and temperatures as low as 1 GeV [15, 16].

(4) The abundance of unwanted, massive relics such as monopoles, gravitinos, and oscillating scalar fields produced after inflation (e.g., during reheating) must be very small. In order that such nonrelativistic relics not contribute too much mass density today, their energy density after inflation must be less than ($10^{-8}\,{\rm GeV}/T_{\rm RH}$) times that of radiation. This is easier to satisfy when the energy scale of inflation is lower: Not only is the constraint less stringent, but many of the dangerous relics are too heavy to be produced at such a low energy scale. Monopoles provide a good example: In unification-scale models of inflation there is the concern that GUT symmetry breaking occurs after inflation, so that the monopole problem is not solved.

(5) An integral part of a sensible particle-physics model—or better yet, a testable part!

We denote the scalar field responsible for inflation by ϕ; as is well appreciated, in slow-rollover inflation ϕ must be very weakly coupled in order to satisfy the density-perturbation constraint [3]. At the energy scale of interest, ϕ must be a gauge singlet of the effective Lagrangian [17]. For simplicity, we take its scalar potential to be of the Coleman-Weinberg type [18], where the symmetry-breaking minimum is generated by radiative corrections,

$$V(\phi) = \frac{B\sigma^4}{2} + B\phi^4 \left[\ln(\phi^2/\sigma^2) - \frac{1}{2}\right]; \qquad (1)$$

other simple polynomial potentials can also be used (e.g., $V = V_0 - \alpha\phi^4 + \beta\phi^5$ [19]). Here σ is the global minimum of the potential and B is a dimensionless coupling whose value must be about $10^{-15} - 10^{-14}$ to achieve density perturbations of the appropriate size. (*All* models of inflation have such a small coupling constant whose fundamental understanding is still lacking.) Coleman-Weinberg potentials are very flat near $\phi = 0$, $V \simeq \mathcal{M}^4 - b\phi^4$ where $\mathcal{M}^4 = B\sigma^4/2$, $b = |\ln(\phi^2/\sigma^2)|\, B$, and for this reason have been used often in models of inflation [20].

If, for the moment, we ignore the coupling of ϕ to other fields, the equation of motion for ϕ in the expanding Universe is,

$$\ddot{\phi} + 3H\dot{\phi} + V' = 0; \tag{2}$$

where we have also assumed that the ϕ field is homogeneous (or at least constant over a region of space of order the Hubble radius). During inflation ϕ "rolls" very slowly—but inevitably—toward $\phi = \sigma$, and as it does its potential energy dominates the energy density of the Universe driving a nearly constant expansion rate,

$$H^2 = \frac{8\pi V(\phi)}{3{m_{\rm Pl}}^2} \simeq \frac{4\pi B\sigma^4}{3{m_{\rm Pl}}^2}; \tag{3}$$

where $m_{\rm Pl} \equiv G^{-1/2} = 1.22 \times 10^{19}\,{\rm GeV}$ is the Planck mass. During the slow-roll phase, when ϕ is near the origin ($\phi \lesssim \phi_e$ [21]), the $\ddot{\phi}$ term can be neglected so that $\dot{\phi} \simeq -V'/3H$. Using this approximation, it follows that during the time it takes the scalar field to evolve from ϕ to the minimum of its potential the cosmic-scale factor grows by $N(\phi)$ e-foldings:

$$N(\phi) \simeq \frac{8\pi}{{m_{\rm Pl}}^2} \int_{\phi}^{\phi_e} \frac{V(\phi)d\phi}{-V'} \simeq \frac{\pi}{2|\ln(\phi^2/\sigma^2)|}\, \frac{\sigma^4}{{m_{\rm Pl}}^2\phi^2}; \tag{4}$$

where $|\ln(\phi^2/\sigma^2)| \approx 60$ is approximately constant during the slow roll. In order to achieve the 30 or so e-foldings of inflation required the initial value of the scalar field must be less than $\sigma^2/30m_{\rm Pl} \simeq 10^{-14}\sigma(\mathcal{M}/\,\mathrm{TeV})$; this is the least attractive feature of our model.

During the slow-roll phase density fluctuations arise due to quantum fluctuations in the scalar field ϕ. The amplitude of the perturbation on a given scale λ, when that scale crosses inside the horizon, is roughly [3]

$$\left(\frac{\delta\rho}{\rho}\right)_{\mathrm{HOR},\lambda} \sim \left(\frac{H^3}{V'}\right)_{N_\lambda} \sim \sqrt{B}N_\lambda^{3/2}; \tag{5}$$

where subscript N_λ indicates that the quantity is to be evaluated when the scale of interest crossed outside the horizon during inflation, which occurs $N_\lambda \simeq 21 + \ln(T_{\rm RH}/1\,\mathrm{TeV}) + \ln(\lambda/\,\mathrm{Mpc})$ e-foldings before the end of inflation. To achieve $\delta\rho/\rho \approx 10^{-5}$, B must be of order 10^{-15}.

The quadrupole anisotropy in the CBR temperature detected by the COBE DMR [14] allows us to be more precise about the value of B. Expanding the CBR temperature on the sky in spherical harmonics, the quadrupole temperature anisotropy is related to $a_2^2 \equiv \sum_m \langle a_{2m}^2\rangle$ (the average, over all observations positions in the Universe, of the sum of the $l = 2$ spherical-harmonic amplitudes squared) and the inflationary potential:

$$\left(\frac{\Delta T}{T}\right)_Q^2 = \frac{a_2^2}{4\pi} = \frac{32\pi}{45}\frac{V^3}{V'^2 m_{\rm Pl}{}^6} \approx \frac{2|\ln(\phi^2/\sigma^2)|B}{45\pi^2}N_\lambda^3. \tag{6}$$

Setting $N_\lambda \sim 30$, the scale of relevance for the quadrupole anisotropy, and taking $(\Delta T/T)_Q \simeq 6\times 10^{-6}$, we find that $B = 6\times 10^{-15}$. This result is relatively insensitive to the scale of inflation—for $\sigma = 10^{16}\,\mathrm{GeV}$, $|\ln(\phi^2/\sigma^2)| \sim 15$ and $N_\lambda \sim 50$, which leads to $B \simeq 3\times 10^{-15}$—but very sensitive to the value of $(\Delta T/T)_Q$, which is probably uncertain by a factor of two.

One last remark about density perturbations; from Eq. (5) we see that the perturbations are not quite scale invariant, $(\delta\rho/\rho)_{\mathrm{HOR},\lambda} \propto N_\lambda^{3/2}$. Expanding $(\delta\rho/\rho)_{\mathrm{HOR},\lambda}$ about the mean of the galaxy scale (1 Mpc) and the present horizon scale (10^4 Mpc) we find that $(\delta\rho/\rho)_{\mathrm{HOR},\lambda} \propto \lambda^{0.06}$ (corresponding to a power spectrum $|\delta_k|^2 \propto k^n$ with $n = 0.88$). This has the effect of depressing perturbations on small scales relative to large scales by about a factor of two, and may be important, as some numerical simulations indicate that an exactly scale-invariant spectrum of density perturbations normalized to the COBE DMR quadrupole has too much power on small scales [22].

Quantum fluctuations during inflation also give rise to a spectrum of gravitational waves [23]; these gravitational waves cross the horizon after inflation with an amplitude of the order of $H/m_{\mathrm{Pl}} \sim 2 \times 10^{-32} (\mathcal{M}/\,\mathrm{TeV})^2$, orders of magnitude smaller than in models where the scale of inflation is of the order of the unification scale—and far too small to be detected.

Finally, consider reheating, the conversion of the vacuum energy to thermal radiation. After its slow roll, the ϕ field begins to oscillate about the minimum of its potential, and the vacuum energy that drives inflation is converted into coherent scalar-field oscillations (corresponding to a condensate of nonrelativistic ϕ particles). Reheating takes place when the ϕ particles decay into light fields, which, through their decays and interactions, eventually produce a thermal bath of radiation. During the epoch of coherent ϕ oscillations the Universe is matter dominated and the energy density trapped in the ϕ field decreases as the cube of the scale factor. The reheat temperature is determined by the decay time of the scalar field oscillations, which is given by the inverse of the decay width Γ of the ϕ [3]. If $\Gamma \lesssim H$, the coherent oscillation phase is relatively long and the reheat temperature $T_{\mathrm{RH}} \simeq \sqrt{m_{\mathrm{Pl}}\Gamma} < \mathcal{M}$,

corresponding to less than 100% conversion of vacuum energy to radiation. Inefficient reheating is the rule for slow-rollover inflation. On the other hand, if $\Gamma \gtrsim H$, ϕ oscillations decay rapidly, and $T_{\rm RH} \simeq \mathcal{M}$, corresponding to 100% conversion of vacuum energy to radiation. Next we discuss why reheating is typically very inefficient in slow-rollover inflation, and how it becomes more efficient as the scale of inflation is decreased.

Suppose the ϕ field couples to a light, Majorana fermion with Yukawa coupling g; its decay width $\Gamma = g^2 m_\phi/4\pi$, where $m_\phi^2 = V''(\sigma) = 8\sqrt{2B}\mathcal{M}^2 \simeq \mathrm{GeV}^2\,(\mathcal{M}/\,\mathrm{TeV})^2$. The condition for efficient reheating is

$$\frac{\Gamma}{H} = \sqrt{\frac{3g^4}{8\pi^3}}\,\frac{m_{\rm Pl}}{\sigma} \simeq \left(\frac{g}{2\times 10^{-6}}\right)^2 \frac{\mathrm{TeV}}{\mathcal{M}} \gtrsim 1. \tag{7}$$

The condition for efficient reheating depends upon the scale of inflation: The larger the scale of inflation, the larger the value of g required for efficient reheating; for $\mathcal{M} = 10^{14}\,\mathrm{GeV}$, good reheating requires $g \gtrsim 0.5$.

Next, consider the other constraints to the Yukawa coupling g. In order not to spoil the flatness of $V(\phi)$, the radiative corrections due to the fermion that couples to ϕ must be small: This requires $g^4 \ll B$ or $g \ll 3\times 10^{-4}$. Further, the coupling to the ϕ field will give it a mass of order $m_f \sim g\sigma$, which must be less than half the mass of the ϕ. This provides the stricter constraint, $g \lesssim \sqrt{2B}$, and illustrates how reheating and density perturbations work at cross purposes: Reheating is better for a larger value of B, but density perturbations require a very small value for B.

By saturating the bound $g \lesssim \sqrt{2B} \sim 10^{-7}$, we can express the maximum achievable reheat temperature as a function of the scale of inflation:

$$\frac{T_{\rm RH}(\max)}{\mathcal{M}} \sim \frac{\sqrt{\Gamma m_{\rm Pl}}}{\mathcal{M}} \simeq B^{5/8}\sqrt{\frac{m_{\rm Pl}}{\mathcal{M}}} \approx 0.1\sqrt{\frac{1\,\mathrm{TeV}}{\mathcal{M}}}. \tag{8}$$

For $\mathcal{M} \sim$ TeV, $T_{\rm RH}(\max)$ is of order 100 GeV, and $T_{\rm RH}(\max)$ grows only as $\sqrt{\mathcal{M}}$: for the canonical scale of inflation, $\mathcal{M} \sim 10^{14}$ GeV, $T_{\rm RH}(\max)$ is only 3×10^7 GeV. Other modes of reheating are possible; e.g., $\phi \to 2\chi$ (χ is another scalar field), through interaction terms of the form, $\mathcal{L}_{\rm int} = \beta\phi\chi^2$ or $\lambda\phi^2\chi^2$. Up to factors of order unity, the same result obtains for the maximum achievable reheat temperature, i.e., Eq. (8) [24, 25].

Let us squarely address the least attractive feature of our model, the small initial value of ϕ required for sufficient inflation, $\phi_i \lesssim \sigma^2/30m_{\rm Pl} \simeq 10^{-14}\sigma\,(\mathcal{M}/\,{\rm TeV})$. Many models of slow-rollover inflation require a small initial value for ϕ; the very small value required here traces to the very-low energy scale of inflation: For comparison, taking $\mathcal{M} \sim 10^{14}$ GeV, $\phi_i \lesssim 10^{-3}\sigma$. This problem can be mitigated by degrees by increasing the scale of inflation.

In order to achieve inflation in our model there must be regions of the Universe where the value of the ϕ field is very small; such regions will undergo inflation. In regions where the value of the ϕ field is not small, there will be no inflation. After inflation, the regions where ϕ was sufficiently small have grown exponentially in size—and with plausible assumptions about the distribution of the initial value of ϕ the inflated regions should occupy most of the physical volume of the Universe.

Such a small initial value for ϕ is not spoiled by the quantum fluctuations in ϕ, which are of the order of $H/2\pi \sim 2 \times 10^{-7}\sigma^2/m_{\rm Pl} \sim 10^{-19}\sigma$. Thermal fluctuations will spoil such localization: $\langle\phi^2\rangle_T^{1/2} \sim T \simeq$ TeV $\sim 10^{-4}\sigma$. However, it can be argued that ϕ is so weakly coupled it is not in thermal contact with the Universe; indeed, this argument has been used for other models of inflation [11].

Another way of insuring that the small initial value of the ϕ field is not

spoiled by thermal fluctuations is to arrange that inflation begin "cold," $T \ll 1\,\mathrm{TeV}$. There are plausible ways that this might occur. If the Universe, or a small portion of it, were so negatively curved that it became curvature dominated early on, say at a temperature T_{CD}, then the temperature when inflation begins is $T_{\mathrm{inflate}} \sim \mathrm{TeV}^2/T_{\mathrm{CD}}$, which can easily be small enough to render the thermal fluctuations impotent. (Within the spirit of "generic" initial conditions, one would expect the curvature radius at the Planck epoch to be of the order of the Hubble radius, in which case $T_{\mathrm{CD}} \sim m_{\mathrm{Pl}}$ and $T_{\mathrm{inflate}} \sim 10^{-13}\,\mathrm{GeV}$.) Or, the Universe can become matter dominated long before inflation, e.g., by monopoles produced at the GUT phase transition, or other massive relics produced copiously in the early Universe. And of course, it is not necessary that the Universe have any radiation in it prior to inflation: It could have begun cold.

Finally, we comment briefly on the phenomenology of our model. Because ϕ is very weakly coupled, it must be an $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$ singlet; however, it can *indirectly* influence electroweak physics. The vacuum expectation value (VEV) of ϕ, $\sigma \sim 4 \times 10^6\,\mathrm{GeV}$, can induce a negative mass-squared for the Higgs field (call it ψ) that does lead to electroweak symmetry breaking, through a coupling $\lambda\psi^2\phi^2$ [26]. A negative mass-squared of order $1\,\mathrm{TeV}$ requires $\lambda \sim 10^{-7}$. Since the radiative corrections to the $V(\phi)$ due to ψ are of order $\lambda^2/4\pi^2 \sim 10^{-15} \sim B$, they are about the right size to account for the ϕ field's symmetry-breaking potential. The VEV of ϕ can give rise to particle masses, e.g., righthanded neutrinos; in this case reheating can take place by ϕ decays into righthanded neutrinos and their subsequent decays into light leptons. If the scale of inflation is raised slightly, $\mathcal{M} \sim 200\,\mathrm{TeV}$, the mass of the ϕ particle is of order several hundred GeV. In this case,

reheating can take place through ϕ decays into electroweak Higgs and their subsequent decays into the particles of the standard model.

In sum, we have presented a simple model of slow-rollover inflation where the vacuum energy that drives inflation can be as small as the electroweak scale—orders of magnitude smaller than in previous models. Inflation at a low-energy scale has a number of attractive features: reheating is more efficient; the monopole problem is more easily solved; and last, but not least, such a model is potentially testable.

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[26] In this regard we have followed the models of Shafi and Vilenkin [9] and Pi [9] where the VEV of the ϕ was used to induce GUT symmetry breaking.