Well Tempered Cosmology: Scales

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Well tempered cosmology provides a well defined path for obtaining cosmology with a low energy cosmic acceleration despite a high (Planck) energy cosmological constant $\Lambda$, through a scalar field dynamically canceling $\Lambda$. We explore relations between the mass scales entering the various Horndeski gravity terms, and focus on the cases of only one or only two mass scales, obtaining general solutions for the form of the action. The resulting cosmology can be natural and viable, and as one of the only paths to dealing with the cosmological constant problem it has a rationale to be a benchmark cosmology.

I. INTRODUCTION

The vast majority of theories proposed to explain current cosmic acceleration ignore the cosmological constant problem: that vacuum energy with energy density at the Planck scale or some other high energy phase transition should dominate the universe [1–3]. A low energy cosmological constant, or scalar fields, or even modified gravity do not generally address this critical issue.

However there are scalar-tensor modified gravity theories that can be devised to remove a high energy cosmological constant, through dynamical cancellation by the scalar field, and replace it by low energy cosmic acceleration. These include self tuning [4–13] and well tempered [14–19] cosmologies. Well tempered cosmology can in addition provide a full cosmic history, from radiation domination to matter domination to cosmic acceleration in a natural manner.

Well tempered cosmology uses particular relations between the terms of the Horndeski gravity Lagrangian in such a way that the scalar field and cosmic evolution equations become degenerate at a de Sitter attractor, “on shell”. Several studies have been made of these relations in special cases, i.e. assumptions about the functional form of one or another Lagrangian term [14–16] (even for a Minkowski background [17]). Perhaps the most general solution was obtained through a series expansion method, giving general relations between Lagrangian terms, as long as the functions were of polynomial form [18].

While polynomials (and in particular monomials) are an arbitrary choice, a natural constraint is shift symmetry, which brings some protection against quantum loop corrections. In this case the action terms only depend on functions of the quadratic kinetic variable $X \equiv \dot{\phi}^2/2$ of the scalar field $\phi$, and additive terms linear in $\phi$ (“tadpoles”). There then tend to be three mass scales in the theory (apart from the bare Planck mass, which is taken to be unity): one associated with the tadpole in the potential, one with the tadpole in the running Planck mass, and the scale of low energy cosmic acceleration, i.e. the asymptotic Hubble constant of the de Sitter state.

Here we examine cases where relations exist between these mass scales, aligning scales in the well tempered cosmology to leave only one or two free mass parameters. In addition, we seek to go beyond polynomial functional forms to derive general relations between the action terms that provide this cosmology.

In Section II we introduce the Friedmann and scalar field evolution equations in the well tempered Horndeski theory, and derive the relations between the free action functions under various conditions. We consider late time and early time behaviors in Section III, concluding in Section IV.

II. EVOLUTION EQUATIONS AND ACTION RELATIONS

Horndeski gravity is the most general scalar-tensor theory delivering second order equations of motion, and is composed of three action terms (we take the $G_5 = 0$ so that the speed of gravitational waves equals the speed of light), generalizing general relativity. Imposing shift symmetry, the three terms are

$$K(X, \phi) = K(X) - \lambda^3 \phi,$$

$$G_3(X, \phi) = G_3(X),$$

$$G_4 = \frac{1}{2} (M_{pl}^2 + M \phi).$$

We see there are two functions of $X$ (we write $K(X)$ to denote the $X$ dependent part of the generalized kinetic term; since we never again use the notation $K(X, \phi)$ this should not be confusing) and two mass scales — $M$ and $\lambda$ in the
tadpoles for the Planck mass coupling and scalar linear potential respectively.

For a flat Friedmann-Lemaître-Robertson-Walker cosmology, the evolution of the cosmic expansion rate \( H(t) = \dot{a}/a \) and scalar field are given by coupled equations of motion:

\[
3H^2(M_{pl}^2 + M\phi) = \rho_m + \Lambda + 2XK_X - K + 6H\dot{\phi}XG_{3X} - 3M\dot{\phi}
\]

\[
-2\dot{H}(M_{pl}^2 + M\phi) = \rho_m + P_m + \dot{\phi}(M - 2XG_{3X}) - H\dot{\phi}(M - 6XG_{3X}) + 2XK_X
\]

\[
0 = \dot{\phi} \left[ K_X + 2XK_{XX} + 6H\dot{\phi}(G_{3X} + XG_{3XX}) \right] + 3H\dot{\phi}K_X + \lambda^3 + 6XG_{3X}(H + 3H^2) - 3M(\dot{H} + 2H^2),
\]

where a subscript \( X \) denotes a derivative with respect to \( X \), and \( \rho_m \) and \( P_m \) are the matter energy density and pressure respectively.

Well tempering arises from degeneracy of the evolution equations “on shell”, in this case at a de Sitter asymptote \( H = H_{dS} \equiv h = \text{const} \). The de Sitter mass scale \( h \) is the third mass scale in the system. The degeneracy is satisfied by equating the coefficient of \( \dot{\phi} \) in one of Eq. (5) or Eq. (6) times the remaining terms in the other equation with the same product for the equation selection interchanged. That is, if \( A\ddot{\phi} + B = 0 = C\ddot{\phi} + D \) holds at the de Sitter asymptote then we solve both equations if \( AD = BC \), as seen by multiplying the equations by \( C \) and \( A \) respectively: \( AC\ddot{\phi} + BC = AC\ddot{\phi} + AD \) when \( AD = BC \). Thus the degeneracy condition is

\[
(M - 2g) \left\{ 3h\dot{\phi}K_X + 18h^2g - 6h^2M + \lambda^3 \right\}
\]

\[
= (K_X + 2XK_{XX} + 6h\dot{\phi}g_X) \left[ -h\dot{\phi}(M - 6g) + 2XK_X \right],
\]

where we simplify notation by writing \( g \equiv XG_{3X} \).

Various solutions to this equation for well tempering have been found in [16, 18] for polynomial forms of \( g \) or \( K_X \). The nonlinear nature of the equation has not allowed a general solution. However here we will adopt relations between mass scales that can allow for general solutions.

### A. Two Mass Scales \((\lambda^3 = 3h^2M)\)

An intriguing aspect of the degeneracy equation (7) is that when \( \lambda^3 = 3h^2M \) then the two terms \{ \} and [ ] are proportional. That is, when we relate the three mass scales in this particular way (one could also view this as \( h^2 = \lambda^3/(3M) \)), leaving only two mass scales in the system, then the equation greatly simplifies to

\[
(M - 2g)Z = \frac{\dot{\phi}}{3h} \left( K_X + 2XK_{XX} + 6h\dot{\phi}g_X \right) Z,
\]

where \( Z \equiv 3h\dot{\phi}K_X + 18h^2g - 3h^2M \), i.e. the \{ \} term (when \( \lambda^3 = 3h^2M \)). This then gives two cases of solutions, when \( Z = 0 \) and \( Z \neq 0 \).

The condition \( Z = 0 \) leads to

\[
g = \frac{M}{6} - \frac{\sqrt{2}X^{1/2}K_X}{6h}
\]

\[
K_X = \frac{hM}{\sqrt{2}}X^{-1/2} - \frac{6h}{\sqrt{2}}X^{-1/2}g;
\]

these are equivalent expressions, i.e. one can either specify \( K(X) \) and derive \( g \) or the other way around. Interestingly, on shell the \( H \) Friedmann equation (5) becomes \( 0 = \ddot{\phi}(M - 2g) \), implying that at the de Sitter asymptote the field coasts, \( \ddot{\phi} = 0 \). Note this is not the zero coefficient case called \((\check{\phi})\) in [16]; the equations of motion remain intact, and do not become trivial. (If we choose \( g = M/2 \), i.e. \( K_X = -hM\sqrt{2}X^{-1/2} \), then they do become trivial.) The choice \( \lambda^3 = 3h^2M \) is unique in the sense that with this relation, all \( \ddot{\phi} \) dependence drops out of the Friedmann equation (4) on-shell. We are left on-shell with an algebraic relation between \( X \) and the vacuum energy \( \Lambda \), which requires \( X \) constant, hence \( \ddot{\phi} = 0 \). This is a nice property, in that it has no runaway behavior.

For \( Z \neq 0 \), we can solve the degeneracy equation to give

\[
g = \frac{M}{2} + cX^{-1/2} + \frac{X^{-1/2}}{6h\sqrt{2}} [K(X) - 2XK_X]
\]

\[
K_X = bX^{-1/2} + \frac{3hM}{2\sqrt{2}}X^{-1/2}\ln X - 3h\sqrt{2}X^{-1/2}g - \frac{3h}{\sqrt{2}}X^{-1/2}\int \frac{dX}{X}g.
\]
Again, these two expressions are equivalent. Here \( b \) and \( c \) are arbitrary constants, and recall that \( K(X) \) refers to the \( X \) dependent part of \( K \), without the tadpole (since \( K(X) \) arises from integrating \( K_X \)). The solution Eq. (12) looks like the Branch B solution of [18], but here it is valid for all \( g(X) \), not simply a finite polynomial as was the case for that paper.

We emphasize that we can apply the above well tempering conditions to general functional forms, such as algebraic expressions.

From \( g \) and \( K_X \) we can readily obtain the action functions \( G_3 \) and \( K \). For \( Z = 0 \) these are

\[
G_3 = \frac{M}{6} \ln X - \frac{1}{3h\sqrt{2}} X^{-1/2} K(X) - \frac{1}{6h\sqrt{2}} \int dX X^{-3/2} K(X) \quad (13)
\]

\[
K = hM\sqrt{2}X^{1/2} - 3h\sqrt{2} \int dX X^{-1/2} g - 3h^2 M\phi . \quad (14)
\]

(\text{Note that a constant in } G_3 \text{ gives an ignorable total derivative in the action.}) For \( Z \neq 0 \) they are

\[
G_3 = \frac{M}{2} \ln X - 2cX^{-1/2} - \frac{1}{3h\sqrt{2}} X^{-1/2} K(X) \quad (15)
\]

\[
K = aX^{1/2} + \frac{3Mh\sqrt{2}}{2} X^{1/2} \ln X - 3h\sqrt{2} \int dX X^{-1/2} \left[ g + \frac{1}{2} \int \frac{dX'}{X'} g \right] - 3h^2 M\phi , \quad (16)
\]

where \( a \) is an arbitrary constant.

From the Friedmann expansion equation (4), on substituting back in the expressions for \( g \) we obtain

\[
3H^2 M^2_{pl} = \rho_m + \Lambda + 2XK_X \left( 1 - \frac{H}{h} \right) - K(X) - 2MH\sqrt{2}X^{1/2} + 3M\phi(h^2 - H^2) \quad [Z = 0] \quad (17)
\]

\[
3H^2 M^2_{pl} = \rho_m + \Lambda + 6c\sqrt{2} - [K(X) - 2XK_X] \left( 1 - \frac{H}{h} \right) + 3M\phi(h^2 - H^2) , \quad [Z \neq 0] \quad (18)
\]

for the \( Z = 0 \) and \( Z \neq 0 \) cases respectively. Recall \( K(X) \) has the tadpole \( 3h^2 M\phi \) removed (and moved to the last term). Thus the field \( \phi \) and its evolution \( X \) can dynamically cancel a bare cosmological constant everywhere (in the \( Z \neq 0 \) case, at the asymptotic limit \( H = h \) most terms vanish but the free \( c \) term can still cancel \( \Lambda \)).

We can also check the soundness of the theory through the no ghost and Laplace stability criteria. For the no ghost condition, it is convenient to work with the property functions [20], for which

\[
\alpha_M = \frac{\dot{\phi}}{H(M^2_{pl} + M\phi)} \quad (19)
\]

\[
\alpha_B = \frac{\dot{\phi}(2g - M)}{H(M^2_{pl} + M\phi)} \quad (20)
\]

\[
\alpha_K = \frac{2X(K_X + 2XK_{XX}) + 12\dot{\phi}Xg_X}{H^2(M^2_{pl} + M\phi)} \quad (21)
\]

for the Planck mass running, braiding, and kineticity respectively. When \( Z = 0 \),

\[
\alpha_K = \frac{2X(K_X + 2XK_{XX})(1 - H/h)}{H^2(M^2_{pl} + M\phi)} , \quad (22)
\]

i.e. the kineticity vanishes at the de Sitter asymptote (the same happens in quintessence and k-essence). When \( Z \neq 0 \),

\[
\alpha_K = \frac{-3h\sqrt{2}X^{1/2}(2g - M) - 12h\sqrt{2}X^{3/2}g_X (1 - H/h)}{H^2(M^2_{pl} + M\phi)} \quad (23)
\]

\[
= \frac{2X(K_X + 2XK_{XX})(1 - H/h) - 6Hc\sqrt{2} - (K - 2XK_X)H/h}{H^2(M^2_{pl} + M\phi)} . \quad (24)
\]

In the de Sitter limit, \( \alpha_K \rightarrow -3\alpha_B \).

The no ghost condition is

\[
\alpha_K + \frac{3}{2}\alpha_B \geq 0 . \quad (25)
\]

There is little we can say in general about this (or the Laplace stability condition, e.g. see Appendix A.2 of [16]), without adopting a specific \( K \) or \( G_3 \). However, in the de Sitter limit, we can see that for the \( Z = 0 \) case the theory is ghost free. In the \( Z \neq 0 \) case, in the de Sitter limit the no ghost condition becomes \(-3\alpha_B,ds/2)(2 - \alpha_B,ds) \geq 0 \), so we require \( \alpha_B,ds \leq 0 \) or \( \alpha_B,ds \geq 2 \).
B. Two Mass Scales ($M = 0$)

Let us consider instead two mass scales $\lambda$ and $h$, keeps the $G_4$ term at the minimal coupling of general relativity, and simplifies issues from nonminimal coupling such as discussed in Sec. 4 and Appendix C of [16]. The degeneracy equation becomes

$$4Xg_X + 2g \left( 1 + \frac{\lambda^3 \phi}{3h[2XK_X + 6h \phi \dot{g}]} \right) + \frac{\dot{\phi}}{3h} (K_X + 2XK_{XX}) = 0.$$  \hspace{1cm} (26)

We must have $Y \equiv [2XK_X + 6h \phi \dot{g}] \neq 0$ or the original degeneracy equation forces $\lambda = 0$. Being nonlinear, this equation cannot be solved in general, but one solution is, for $Y = c$ with $c$ a constant,

$$K_X = \frac{\lambda^3}{3h \sqrt{2X}} \left( 1 + \frac{\lambda^3 \sqrt{2X}}{3hc} \right)^{-1}, \quad K = c \ln \left( 1 + \frac{\lambda^3 \sqrt{2X}}{3hc} \right) - \lambda^3 \phi \hspace{1cm} (27)$$

$$g = \frac{c}{6h \sqrt{2X}} - \frac{\lambda^3 \sqrt{2X}}{18h^2} \left( 1 + \frac{\lambda^3 \sqrt{2X}}{3hc} \right)^{-1}. \hspace{1cm} (28)$$

Note these are algebraic functions, not covered by the polynomial solutions of [16]. This provides a well tempered solution with scalar field equation on shell

$$\ddot{\phi} = 3h \dot{\phi} + \frac{\lambda^3}{c} \dot{\phi}^2 \hspace{1cm} (29)$$

$$\dot{\phi} = -\frac{3hc}{\lambda^3} \frac{A e^{3ht}}{A e^{3ht} - 1}, \hspace{1cm} (30)$$

where $c < 0$. Thus again $\dot{\phi}$ approaches a constant, freezing $X$. Note that while $c$ is effectively another mass scale (it has the same dimensions as $X \sim m^2/t^2$), we could define it as some combination of $\lambda$ and $h$, e.g. $c \sim -(\lambda^3/h)^2$.

C. Single Mass Scale

We can further reduce the number of free mass scales by setting $M = \lambda = 0$. Since again the Planck mass does not run, this avoids the same problematic issues as mentioned in the previous subsection. This preserves all the results of Section II A, while simplifying some, e.g. the Laplace stability condition for the $Z = 0$ case in the de Sitter limit becomes

$$|g|_{\text{dS}} \leq \frac{hM_{\text{pl}}^2}{\dot{\phi}} \quad \text{equiv.} \quad |XK_X|_{\text{dS}} \leq 3h^2M_{\text{pl}}^2. \hspace{1cm} (31)$$

The single mass scale is $h$, which will be determined by the coefficients in $K$ or $G_3$, e.g. $g_n$ in $g = g_n X^n$.

III. ASYMPTOTIC BEHAVIORS

A. Late Time de Sitter Attractor

To check the attractor behavior to the de Sitter asymptote, we perturb $H = h + \delta h$, $\phi = \phi_0 + \delta \phi$. On shell, for the $Z = 0$ case we have $\phi_0 = 0$. In the $M = \lambda = 0$ single scale model, for simplicity, the solution to first order in the perturbations is

$$\delta h = \delta h_0 e^{-3K_X ht} \hspace{1cm} (32)$$

$$\delta \ddot{\phi} = 3 \delta \dot{\phi} \left( 1 + \frac{6M_{\text{pl}}^2 h^2}{\dot{\phi}_0^2} \right) \delta h \sim e^{-3K_X ht}. \hspace{1cm} (33)$$

As long as $K_X \geq 0$ in the de Sitter limit, the attractor is stable. Note that $\dot{\phi}$ approaches zero from below, i.e. the kinetic term $X$ slows to a constant, with the field evolving as

$$\dot{\phi}(t) = \dot{\phi}_0 \left[ 1 - \frac{\delta h_0}{hK_X} \left( 1 + \frac{6M_{\text{pl}}^2 h^2}{\dot{\phi}_0^2} \right) e^{-3K_X ht} \right]. \hspace{1cm} (34)$$
B. Early Time Matter Dominated Behavior

We want to insure a viable cosmology at early times, having a standard matter dominated epoch before late time acceleration. Consider Eq. (6) at early times when $H \gg h$. Keeping the leading order terms in $H$ gives

$$0 \approx 6 H g X \ddot{\phi} + 9 (1 - w_b) H^2 g - \frac{3}{2} (1 - 3 w_b) H^2 M$$

where we have used $\dot{H} = -(3/2)(1 + w_b) H^2$ for $w_b$ the early time background equation of state. This has the solution

$$g \sim c \frac{1 - w_b}{1 + w_b} t^{-(1 - w_b)/(1 + w_b)} + \frac{M}{6} \frac{1 - 3 w_b}{1 + w_b},$$

where $c$ is a constant. We see that $g$ gets large at early times. Through Eqs. (10) or (12) we see that this implies that $X^{1/2} K_X$ also is large at early times. From the Friedmann equation (4) we see that $\dot{\phi}$ grows as $t^{2/(1 + w_b)}$ so the field starts off slowly rolling. Thus, for a viable matter era we require $K_X \sim X < -1/2$. This also guarantees freedom from ghosts at these early times, and that the property functions all go to zero (restoring general relativity) at early times.

IV. CONCLUSIONS

As one of the few physical approaches that does not sweep the problem of a high energy cosmological constant under the rug, well tempered cosmology is an attractive avenue to deal with our universe, accelerating at a low energy scale. In well tempering, terms in the Horndeski action are related to each other, enabling a dynamical cancellation of the cosmological constant, while preserving matter and giving rise to low energy acceleration. Here we go further, and relate not just the forms of the action terms (while leaving them more general than previous solutions) but also their mass scales.

We present three such relations: removing one mass scale by a unique relation $\lambda^3 = 3 h^2 M$, by setting $M = 0$, or removing two mass scales by setting $M = 0 = \lambda$. Each has its own particular characteristics, benefits, and solutions. In the first case, one branch of the solution bounds the quadratic kinetic quantity $X = \dot{\phi}^2/2$, such that it goes to a constant in the de Sitter asymptotic rather than running away. This is related to $\phi$ (but not $X$) dependence dropping out of the Friedmann equation in that limit. The field behavior in the second branch depends on the particular form of $K(X)$ or $G_3(X)$. The second case can also give bounded solutions, and moreover possesses the minimal coupling of general relativity to the Ricci scalar. The third case has a blend of the above properties.

While with three mass scales, Horndeski terms in the Friedmann equation partially cancel due to the degeneracy condition, giving a constant piece to dynamically cancel a large cosmological constant, while $\dot{\phi}$ runs away to infinity at late times. With only two mass scales, the constant piece can be obtained by freezing $X$ in the de Sitter limit, so $\dot{\phi}$ goes to a constant.

In addition, we investigated the form of the property functions, more closely related to phenomenology, and the ghost free condition. Finally, we explored the late time de Sitter attractor nature, and the early time matter dominated behavior.

The new general functional solutions, not limited merely to the previous polynomial forms, and the discovery of cases with bounded scalar field evolution, show that well tempered cosmology is a rich field. Dynamical cancellation of a high energy cosmological constant and protection against quantum radiative corrections (due to shift symmetry) make such an approach a leading “benchmark” to use in testing our low energy accelerating universe.

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