The Abdus Salam International Centre for Theoretical Physics

# DSU2023 (The Dark Side of the Universe) | (SMR 3863)

10 Jul 2023 - 14 Jul 2023
Outside, Kigali, Rwanda

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T01

# Primordial Black Holes generated from SUSY Breaking in Non-oscillatory Inflation Models

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Supersymmetric models of quintessential inflation provide an interesting framework that accounts for both the cosmological inflation at the early universe as well as the dark energy at the present time. We study the effects of SUSY breaking in a class of sgoldstinoless models of non-oscillatory inflation, on the dynamics of both inflation and late time expansion. In these models we show that SUSY breaking effects vanish on the dark energy tail, without any fine tuning. On the other hand integrating out the sgoldstino leads to a sizeable back-reaction on the inflation dynamics and observables. We show that the peculiar form of the potential in the large supersymmetry breaking scale limit can generate peaks in the scalar power spectrum produced from inflation. This may lead to the formation of primordial black holes with various masses in the early Universe, where the perturbation modes may re-enter the horizon during or after the kination phase. This can account for the observed dark matter abundance.

[1] L. Heurtier, A. Moursy and L. Wacquez, [arXiv:2207.11502 [hep-th]].

T02

# Light singlino DM of the natural NMSSM

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Inspired by the fact that relatively small values of the effective higgsino mass parameter of the $Z_3$-symmetric Next-to-Minimal Supersymmetric Standard Model (NMSSM) could render the scenario 'natural', we explore the plausibility of having relatively light neutralinos and charginos (the electroweakinos or the ewinos) in such a scenario with a rather light singlino-like Lightest Supersymmetric Particle (LSP), which is a Dark Matter (DM) candidate, and singlet-dominated scalar excitations. By first confirming the indications in the existing literature that finding simultaneous compliance with results from the Large Hadron Collider (LHC) and those from various DM experiments with such light states is, in general, a difficult ask, we proceed to demonstrate, with the help of a few representative benchmark points, how exactly and to what extent could such a highly motivated 'natural' setup with a singlino-like DM candidate still remains plausible.

T03

# Abstract for...17TH INTERNATIONAL WORKSHOP ON THE DARK SIDE OF THE UNIVERSE (DSU 2023)...

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The behavior of scalar fields in a thermal plasma plays an important role in the study of the early universe, and it is relevant for addressing problems in astrophysics and cosmology. In this talk, we present a calculation of the dissipation rate in a simple scalar model at a finite temperature. We discuss the impact of thermal masses on the quasi-particle kinematics and show the different allowed regimes for different processes, such as decay, inverse decay, and Landau damping. We point out and correct an error in an earlier computation in [1]. For some parameter choices, our correction can significantly change the evolution of the system. These results could be implemented in models of inflation, and have implications for reheating, baryogenesis, and dark matter, among others.

[1] M. Drewes, J. U. Kang, Nucl. Phys. B 875 (2013) 315–350, arXiv.1305.0267 [hep-ph]

# Cosmological constraints of diffusive dark-fluid models

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In this work, we consider an interacting dark-fluid cosmological model in which energy exchange between dark matter and dark energy occurs through diffusion. After solving the background expansion history, we attempt to constrain the cosmological parameters by comparing simulated values of the model against cosmological data. We consider four different cases and compare them against the $\Lambda$CDM model as the "true model". We will show that the diffusive model in which dark energy flows to dark matter is the most likely alternative to $\Lambda$CDM model. We will also demonstrate how this interacting-fluid model can give a potential explanation to the so-called Hubble tension.

[1]Presenting author.

T05

# Dark Sectors in String Theory.

T06

# Confronting the Chaplygin gas with data: background and perturbed cosmic dynamics

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**Abstract**

In this paper, we undertake a unified study of background dynamics and cosmological perturbations in the presence of the Chaplygin gas. This is done by first constraining the background cosmological parameters of different Chaplygin gas models with SNIa data, and then feeding these observationally constrained parameters in the analysis of cosmological perturbations. Based on the statistical criteria we followed, none of the models has a substantial observational support but we show that the so-called 'original' and 'extended/generalized' Chaplygin gas models have *some observational support* and *less observational support*, respectively, whereas the 'modified' and 'modified generalized' Chaplygin gas models miss out on the *less observational support* category but cannot be ruled out. The so-called 'generalized cosmic Chaplygin gas' model, on the other hand, falls under the *no observational support* category of the statistical criterion and can be ruled out. We follow the $1 + 3$ covariant formalism of perturbation theory and derive the evolution equations of the fluctuations in the matter density contrast of the matter-Chaplygin gas system for the models with some or less statistical support. The solutions to these coupled systems of equations are then computed in both short-wavelength and long-wavelength modes.

T07



# Nonstationary energy in general relativity

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(Received 21 November 2019; published 15 January 2020)

Using the time evolution equations of (cosmological) general relativity in the first order Fischer-Marsden form, we construct an integral that measures the amount of nonstationary energy on a given spacelike hypersurface in $D$ dimensions. The integral vanishes for stationary spacetimes; and with a further assumption, reduces to Dain's invariant on the boundary of the hypersurface which is defined with the Einstein constraints and a fourth order equation defining approximate Killing symmetries.

DOI: 10.1103/PhysRevD.101.024035

## I. INTRODUCTION

Dain [1] constructed a geometric invariant that measures the *nonstationary* energy for an asymptotically flat hypersurface in $3+1$ dimensions for the case of time-symmetric initial data which, for vacuum, is an invariant that quantifies the total energy of the gravitational radiation. So this invariant is a component of the total Arnowitt-Deser-Misner (ADM) energy [2] assigned to an asymptotically flat hypersurface. That construction was extended to the time-nonsymmetric case recently in [3]. To give an example of how useful such a geometric invariant can be when constructing initial data for the gravitational field, let us recall the first observation of the merger of two black holes [4]. According to this observation, two initial black holes with masses (approximately) $36M_{\odot}$ and $29M_{\odot}$ merged to produce a single stationary black hole of mass $62M_{\odot}$ plus gravitational radiation of total energy equivalent to $3M_{\odot}$. Assuming this system to be isolated in an asymptotically flat spacetime, the total initial ADM energy of $65M_{\odot}$ is certainly conserved. But this total ADM energy of the initial data needs a refinement as it clearly has a nonstationary part equal to $3M_{\odot}$. The important question is to identify this nonstationary energy in the initial data.

Dain's construction and its extension to the nontime symmetric case by Kroon and Williams [3] are based on several earlier crucial works one of which is the Killing initial data (KID) concept of Moncrief [5] and Beig-Chruściel [6]; and a fourth order operator defined by Bartnik [7]. Of course all of the discussion is related to the Cauchy problem in general relativity and the related issue of constructing initial data for the time evolution equations. Here by using the time-evolution equations, in the form given by Fischer and Marsden [8], we construct a new representation of the nonstationary energy in generic $D$ dimensional spacetimes with or without a cosmological constant.

The outline of the paper is as follows: in Sec. II we briefly summarize Dain's construction using the constraints and present a new approach using the evolution equations. In Sec. III we give the details of the relevant computations in $D$ dimensions. The Appendix is devoted to the ADM decomposition.

## II. DAIN'S INVARIANT IN BRIEF AND A NEW FORMULATION

Leaving the details of the construction to the next section, let us first briefly summarize the ingredients needed to define Dain's invariant on a spacelike hypersurface $\Sigma$ of the spacetime $\mathcal{M} = \mathbb{R} \times \Sigma$. Then we shall discuss our new formulation via the evolution equations.

The initial data on the hypersurface is defined by the Riemannian metric $\gamma_{ij}$ and the extrinsic curvature $K_{ij}$ in local coordinates. Denoting $D_i$ to be the covariant derivative compatible with $\gamma_{ij}$ and assuming the usual ADM decomposition of the spacetime metric $g_{\mu\nu}$, the line element reads

$$ds^2 = (N_i N^i - N^2)dt^2 + 2N_i dt dx^i + \gamma_{ij} dx^i dx^j, \quad (1)$$

while the extrinsic curvature becomes[1]

[1]Our definition of the extrinsic curvature is as follows: given $(X, Y)$ two vectors on the tangent space $T_p\Sigma$ and $n$ be the unit normal to $\Sigma$, then $K(X, Y) \coloneqq g(\nabla_X n, Y)$ with $\nabla$ being the covariant derivative compatible with the spacetime metric $g$.

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$$K_{ij} = \frac{1}{2N}(\dot{\gamma}_{ij} - D_i N_j - D_j N_i), \tag{2}$$

with the lapse function $N = N(t, x^i)$ and the shift vector $N^i = N^i(t, x^i)$. The spatial indices can be raised and lowered with the $D-1$ dimensional spatial metric $\gamma$; over dot denotes the derivative with respect to $t$, and the Latin letters are used for the spatial dimensions, $i, j, k, \ldots = 1, 2, 3, \ldots D-1$, whereas the Greek letters are used to denote the spacetime dimensions, $\mu, \nu, \rho, \ldots = 0, 1, 2, 3, \ldots D-1$. All the relevant details of the ADM decomposition are given in the Appendix.

Under the above decomposition of spacetime, the $D$-dimensional Einstein equations

$$R_{\mu\nu} - \frac{1}{2} R g_{\mu\nu} + \Lambda g_{\mu\nu} = \kappa T_{\mu\nu} \tag{3}$$

yield the Hamiltonian and momentum constraints on the hypersurface $\Sigma$ as

$$\begin{aligned} \Phi_0(\gamma, K) &:= -{}^{\Sigma}R - K^2 + K_{ij}^2 + 2\Lambda - 2\kappa T_{nn} = 0, \\ \Phi_i(\gamma, K) &:= -2D_k K_i^k + 2D_i K - 2\kappa T_{ni} = 0, \end{aligned} \tag{4}$$

where $K := \gamma^{ij} K_{ij}$ and $K_{ij}^2 := K^{ij} K_{ij}$. From now on we shall work in vacuum, hence $T_{\mu\nu} = 0$. Denoting $\Phi(\gamma, K)$ to be the constraint covector with components $(\Phi_0, \Phi_i)$ and $D\Phi(\gamma, K)$ to be its linearization about a given solution $(\gamma, K)$ to the constraints and $D\Phi^*(\gamma, K)$ to be the formal adjoint map, then following Bartnik [7], one defines another operator $\mathcal{P}$:

$$\mathcal{P} := D\Phi(\gamma, K) \circ \begin{pmatrix} 1 & 0 \\ 0 & -D^m \end{pmatrix}. \tag{5}$$

The reason why we need this operator will be clear below. Using the formal adjoint $\mathcal{P}^*$ of Bartnik's operator, Dain [1] defines the following integral over the hypersurface

$$\mathcal{I}(N, N^i) := \int_{\Sigma} dV \mathcal{P}^* \begin{pmatrix} N \\ N^k \end{pmatrix} \cdot \mathcal{P}^* \begin{pmatrix} N \\ N^k \end{pmatrix}, \tag{6}$$

where the multiplication is defined as

$$\begin{pmatrix} N \\ N^i \end{pmatrix} \cdot \begin{pmatrix} A \\ B_i \end{pmatrix} := NA + N^i B_i. \tag{7}$$

The integral (6) is to be evaluated for specific vectors $\xi := (N, N^i)$ that satisfy the fourth-order equation

$$\mathcal{P} \circ \mathcal{P}^*(\xi) = 0, \tag{8}$$

which Dain called the *approximate Killing initial data* (KID) equation. It is clear that if $\xi$ satisfies the lower derivative equation $\mathcal{P}^*(\xi) = 0$, then it also satisfies (8). Moreover, these particular solutions, together with an assumption on their decay at infinity, also solve the KID equations which are simply $D\Phi^*(\gamma, K)(\xi) = 0$. In fact this point is crucial but well-established: Moncrief [5] proved that $\xi$ is a spacetime Killing vector satisfying $\nabla_\mu \xi_\nu + \nabla_\nu \xi_\mu = 0$ if and only if it satisfies the KID equations. Namely one has

$$\nabla_\mu \xi_\nu + \nabla_\nu \xi_\mu = 0 \Leftrightarrow D\Phi^*(\gamma, K)(\xi) = 0, \tag{9}$$

with $(N, N^i)$ being the projections off and onto the hypersurface of the Killing vector field $\xi$. The physical picture is clear: initial data on the hypersurface clearly encode the spacetime symmetries. There have been rigorous works on the KIDs in [6,9,10] which we shall employ in what follows.

Observe that for any Killing vector field $\mathcal{I}(N, N^i)$ vanishes identically. So by design, Dain's invariant identically vanishes for initial data with exact symmetries. Then Dain goes on to show that for asymptotically flat spaces, for the case of approximate translational KID's $\mathcal{I}(N, N^i)$ can measure the *nonstationary energy* contained in the hypersurface $\Sigma$. To simplify his calculations Dain considered the time symmetric initial data ($K_{ij} = 0$) in three spatial dimensions. There are two crucial points to note about Dain's construction: first, one can show that for any asymptotically flat three manifold, the approximate KID equation has nontrivial solutions which are not KIDs; second, using integration by parts, one can convert the volume integral (6) to a surface integral. We shall discuss these in the next section, but let us first give another formulation of this invariant.

### A. Nonstationary energy via time-evolution equations

In Dain's construction, as is clear from the above summary, time evolution of the initial data has not played a role: in fact one only works with the constraints on the hypersurface. This fact somewhat obscures the interpretation of the proposed invariant as the nonstationary energy contained in the initial data. In what follows, we propose another formulation of this invariant with the help of the time evolution equations which makes the interpretation clearer. For this purpose let us consider the phase space variables to be the spatial metric $\gamma_{ij}$ and the canonical momenta $\pi^{ij}$; the latter can be found from the Einstein-Hilbert Lagrangian

$$\begin{aligned} \mathcal{L}_{\mathrm{EH}} &= \frac{1}{\kappa} \sqrt{-g}(R - 2\Lambda) \\ &= \frac{1}{\kappa} \sqrt{\gamma} N({}^{\Sigma}R + K_{ij}^2 - K^2 + \Lambda) \\ &\quad + \text{boundary terms} \end{aligned} \tag{10}$$

which are

$$\pi^{ij} := \frac{\delta \mathcal{L}_{\rm EH}}{\delta \dot{\gamma}_{ij}} = \frac{1}{\kappa}\sqrt{\gamma}(K^{ij} - \gamma^{ij}K). \tag{11}$$

Using the canonical momenta, it pays to recast the densitized versions of the constraints (4) for $T_{\mu\nu} = 0$ and setting $\kappa = 1$ as

$$\Phi_0(\gamma,\pi) := \sqrt{\gamma}(-{}^{\Sigma}R + 2\Lambda) + G_{ijkl}\pi^{ij}\pi^{kl} = 0,$$
$$\Phi_i(\gamma,\pi) := -2\gamma_{ik}D_j\pi^{kj} = 0, \tag{12}$$

where the *DeWitt metric* [11] $G_{ijkl}$ in $D$ dimensions reads

$$G_{ijkl} = \frac{1}{2\sqrt{\gamma}}\left(\gamma_{ik}\gamma_{jl} + \gamma_{il}\gamma_{jk} - \frac{2}{D-2}\gamma_{ij}\gamma_{kl}\right). \tag{13}$$

Ignoring the possible boundary terms, the ADM Hamiltonian density turns out to be a sum of the constraints as

$$\mathcal{H} = \int_{\Sigma} d^{D-1}x\langle \mathcal{N}, \Phi(\gamma,\pi)\rangle, \tag{14}$$

with $\mathcal{N}$ being the lapse-shift vector with components $(N, N^i)$ which play the role of the Lagrange multipliers; and the angle-brackets denote the usual contraction. Given an $\mathcal{N}$, the remaining evolution equations can be written in a compact form (the Fischer-Marsden form [12]) as

$$\frac{d}{dt}\begin{pmatrix}\gamma \\ \pi\end{pmatrix} = J\circ D\Phi^*(\gamma,\pi)(\mathcal{N}), \tag{15}$$

where the $J$ matrix reads

$$J = \begin{pmatrix} 0 & 1 \\ -1 & 0\end{pmatrix}. \tag{16}$$

The reason why the formal adjoint of the linearized constraint map $D\Phi^*(\gamma,\pi)$ appears in the time evolution can be seen as follows: the Hamiltonian form of the Einstein-Hilbert action

$$\mathcal{S}_{\rm EH}[\gamma,\pi] = \int dt \int d^{D-1}x(\pi^{ij}\dot{\gamma}_{ij} - \langle \mathcal{N}, \Phi(\gamma,\pi)\rangle), \tag{17}$$

when varied about a background $(\gamma,\pi)$ gives

$$D\mathcal{S}_{\rm EH}[\gamma,\pi] = \int dt \int d^{D-1}x(\delta\pi^{ij}\dot{\gamma}_{ij} + \pi^{ij}\delta\dot{\gamma}_{ij} - \langle \mathcal{N}, D\Phi(\gamma,\pi)\cdot(\delta\gamma,\delta\pi)\rangle). \tag{18}$$

Here the linearized form of the constraint map can be computed to be

$$D\Phi\begin{pmatrix}h_{ij} \\ p^{ij}\end{pmatrix} = \begin{pmatrix}\sqrt{\gamma}({}^{\Sigma}R^{ij}h_{ij} - D^iD^jh_{ij} + \triangle h) \\ -hG_{ijkl}\pi^{ij}\pi^{kl} + 2G_{ijkl}p^{ij}\pi^{kl} + 2G_{njkl}h_{im}\gamma^{mn}\pi^{ij}\pi^{kl} \\ -2\gamma_{ik}D_jp^{kj} - \pi^{jk}(2D_kh_{ij} - D_ih_{jk})\end{pmatrix}, \tag{19}$$

where $\delta\gamma_{ij} := h_{ij}$, $h := \gamma^{ij}h_{ij}$, $\delta\pi^{ij} := p^{ij}$ and $\triangle := D_kD^k$. We have used the vanishing of the constraints to simplify the expression. In (18) using integration by parts when necessary and dropping the boundary terms one arrives at the desired result

$$D\mathcal{S}_{\rm EH}[\gamma,\pi] = \int dt \int d^{D-1}x(\delta\pi^{ij}\dot{\gamma}_{ij} - \dot{\pi}^{ij}\delta\gamma_{ij} - \langle(\delta\gamma,\delta\pi), D\Phi^*(\gamma,\pi)\cdot\mathcal{N}\rangle), \tag{20}$$

where the adjoint constraint map appears in the process which reads

$$D\Phi^*\begin{pmatrix}N \\ N^i\end{pmatrix} = \begin{pmatrix}\sqrt{\gamma}({}^{\Sigma}R^{ij} - D^iD^j + \gamma^{ij}\triangle)N - N\gamma^{ij}G_{klmn}\pi^{kl}\pi^{mn} + 2NG_{klmn}\gamma^{ik}\pi^{jl}\pi^{mn} + 2\pi^{k(i}D_kN^{j)} - D_k(N^k\pi^{ij}) \\ 2NG_{ijkl}\pi^{kl} + 2D_{(i}N_{j)}\end{pmatrix}. \tag{21}$$

Setting the variation (20) to zero one obtains the evolution equations (15) or in more explicit form one has

$$\frac{d\gamma_{ij}}{dt} = 2NG_{ijkl}\pi^{kl} + 2D_{(i}N_{j)}, \tag{22}$$

and

$$\frac{d\pi^{ij}}{dt} = \sqrt{\gamma}(-{}^{\Sigma}R^{ij} + D^iD^j - \gamma^{ij}\triangle)N + N\gamma^{ij}G_{klmn}\pi^{kl}\pi^{mn} - 2NG_{klmn}\gamma^{ik}\pi^{jl}\pi^{mn} - 2\pi^{k(i}D_kN^{j)} + D_k(N^k\pi^{ij}). \tag{23}$$

Together with the constraints (12) these two tensor equations constitute a set of constrained dynamical system for a *given* lapse-shift vector an $(N, N^i)$. The constraints have a dual role: they determine the viable initial data and also generate time evolution of the initial data once the lapse-shift vector is chosen. As noted above, if $D\Phi^*(\gamma,\pi)(\mathcal{N}) = 0$, namely $\mathcal{N} = \xi$ is a Killing vector field then the time evolution is trivial. In particular this would be the case for a stationary Killing vector.

Consider now an $\mathcal{N}$ which is *not* a Killing vector, which means $D\Phi^*(\gamma,\pi)(\mathcal{N}) \neq 0$; and in particular directly from the evolution equations we can find how much

$D\Phi^*(\gamma,\pi)(\mathcal{N})$ differs from zero (or how much a given $\mathcal{N}$ fails to be a Killing vector) as

$$D\Phi^*(\gamma,\pi)(\mathcal{N}) = J^{-1}\circ\frac{d}{dt}\begin{pmatrix}\gamma\\ \pi\end{pmatrix}. \tag{24}$$

To get a number from this matrix, first one should note that the units of $\gamma$ and $\pi$ are different by a factor of $1/L$ and so a naive approach of taking the "square" of this matrix does not work. At this stage to remedy this, one needs the (adjoint) operator of Bartnik that we have introduced above: so one has

$$\begin{aligned}\mathcal{P}^*(\mathcal{N}) &:= \begin{pmatrix}1 & 0\\ 0 & D_m\end{pmatrix}\circ D\Phi^*(\gamma,\pi)(\mathcal{N})\\ &= \begin{pmatrix}1 & 0\\ 0 & D_m\end{pmatrix}\circ J^{-1}\circ\frac{d}{dt}\begin{pmatrix}\gamma\\ \pi\end{pmatrix},\end{aligned} \tag{25}$$

which yields $\mathcal{P}^*(\mathcal{N}) = (-\dot{\pi}, D_m\dot{\gamma})$. Since $\pi$ is a tensor density to get a number out of this vector, we further define

$$\tilde{\mathcal{P}}^*(\mathcal{N}) := \begin{pmatrix}\gamma^{-1/2} & 0\\ 0 & 1\end{pmatrix}\circ\mathcal{P}^*(\mathcal{N}). \tag{26}$$

Then the integral of $\tilde{\mathcal{P}}^*(\mathcal{N})\cdot\tilde{\mathcal{P}}^*(\mathcal{N})$ over the hypersurface yields

$$\begin{aligned}\mathcal{I}(\mathcal{N}) &= \int_\Sigma dV\tilde{\mathcal{P}}^*(\mathcal{N})\cdot\tilde{\mathcal{P}}^*(\mathcal{N})\\ &= \int_\Sigma dV\left(|D_m\dot{\gamma}_{ij}|^2 + \frac{1}{\gamma}|\dot{\pi}^{ij}|^2\right),\end{aligned} \tag{27}$$

where $|D_m\dot{\gamma}_{ij}|^2 := \gamma^{mn}\gamma^{ij}\gamma^{kl}D_m\dot{\gamma}_{ik}D_n\dot{\gamma}_{jl}$ and $|\dot{\pi}^{ij}|^2 := \gamma_{ij}\gamma_{kl}\dot{\pi}^{ik}\dot{\pi}^{jl}$. This is another representation of Dain's invariant which explicitly involves the time derivatives of the canonical fields. We have also not assumed that the cosmological constant vanishes, hence our result is valid for generic spacetimes. Note that this expression is valid for any $\mathcal{N}$ which is not necessarily an approximate KID, hence given a solution to the constraint equations and a choice of the lapse-shift vector, one can compute this integral. But the volume integral becomes a surface integral when $\mathcal{N}$ is an approximate KID which is the case considered by Dain. Observe that by construction, $\mathcal{I}(\mathcal{N})$ is a non-negative number. To get the explicit expression as a volume integral in terms of the canonical fields and not their time derivatives, one should plug the two evolution equations (22) and (23) to (27). The resulting expression is

$$\begin{aligned}\mathcal{I}(\mathcal{N}) = \int_\Sigma dV\{&|D_mV^{ij}|^2 + {}^\Sigma R_{ij}^2N^2 + (D_iD_jN)^2\\ &- 2{}^\Sigma R^{ij}ND_iD_jN + 2{}^\Sigma RN\triangle N + (D-3)\triangle N\triangle N\\ &+ 2Q\triangle N + Q_{ij}^2 + 2{}^\Sigma R_{ij}NQ^{ij} - 2Q^{ij}D_iD_jN\\ &+ 4D_mD_{(i}N_{j)}D^mD^{(i}N^{j)} + 4D_mD_iN_jD^mV^{ij}\},\end{aligned} \tag{28}$$

where

$$V^{ij} := \frac{2N}{\sqrt{\gamma}}\left(\pi^{ij} - \frac{1}{D-2}\pi\gamma^{ij}\right), \tag{29}$$

and

$$\begin{aligned}Q^{ij} :=& \frac{2N}{\gamma}\left(\pi^i_k\pi^{kj} - \frac{\pi\pi^{ij}}{D-2}\right) - \frac{N}{\gamma}\gamma^{ij}\left(\pi_{kl}^2 - \frac{\pi^2}{D-2}\right)\\ &- \frac{1}{\sqrt{\gamma}}D_k(N^k\pi^{ij}) + \frac{2}{\sqrt{\gamma}}\pi^{k(i}D_kN^{j)},\end{aligned} \tag{30}$$

and $Q := \gamma_{ij}Q^{ij}$. Equation (28) is our main result: given a solution, that is an initial data, one an compute this integral which measures the deviation from stationarity. We can also write (28) in terms of $\gamma_{ij}$ and the extrinsic curvature $K_{ij}$. For this purpose all one needs to do is to rewrite $V^{ij}$ and $Q^{ij}$ in terms of these variables. They are given as

$$V^{ij} = 2NK^{ij}, \tag{31}$$

and

$$\begin{aligned}Q^{ij} :=& 2N(K^i_kK^{kj} - KK^{ij}) - N\gamma^{ij}(K_{kl}^2 - K^2) - D_k(N^kK^{ij})\\ &+ \gamma^{ij}D_k(N^kK) + 2K^{k(i}D_kN^{j)} - 2KD^{(i}N^{j)}.\end{aligned} \tag{32}$$

Up to now we have not made a choice of gauge or coordinates. Let us now choose the Gaussian normal coordinates ($N=1$, $N^i=0$) on $\Sigma$ for which the integral reads

$$\begin{aligned}\mathcal{I}(\mathcal{N}) = \int_\Sigma dV\Bigg\{&\frac{4}{\gamma}\left(|D_m\pi^{ij}|^2 - \frac{D-3}{(D-2)^2}|D_m\mathrm{Tr}(\pi)|^2\right)\\ &+ {}^\Sigma R_{ij}^2 + \frac{4}{\gamma}{}^\Sigma R_{ij}\pi^{ik}\pi^j_k - \frac{4}{(D-2)\gamma}{}^\Sigma R_{ij}\pi^{ij}\mathrm{Tr}(\pi)\\ &- \frac{4}{\gamma}\Lambda\left(\mathrm{Tr}(\pi^2) - \frac{1}{(D-2)}(\mathrm{Tr}(\pi))^2\right)\\ &+ \frac{D-7}{\gamma^2}\left(\mathrm{Tr}(\pi^2) - \frac{1}{D-2}(\mathrm{Tr}(\pi))^2\right)^2\\ &+ \frac{4}{\gamma^2}\Bigg(\mathrm{Tr}(\pi^4) - \frac{2}{D-2}\mathrm{Tr}(\pi)\mathrm{Tr}(\pi^3)\\ &+ \frac{1}{(D-2)^2}(\mathrm{Tr}(\pi))^2\mathrm{Tr}(\pi^2)\Bigg)\Bigg\},\end{aligned} \tag{33}$$

where $\mathrm{Tr}(\pi) := \gamma_{ij}\pi^{ij}$ and $\mathrm{Tr}(\pi^2) := \pi^{ij}\pi_{ij}$ and so on. In terms of the extrinsic curvature, in the Gaussian normal coordinates, one has

$$\mathcal{I}(\mathcal{N}) = \int_\Sigma dV\{4|D_m K_{ij}|^2 + {}^\Sigma R_{ij}^2 + 4{}^\Sigma R_{ij}(K^{ik}K_k^j - KK^{ij}) + 4\Lambda(K^2 - K_{ij}^2) + 4K_{ij}K^{jl}K_{lm}K^{mi} - 8KK_{ij}K^{jl}K_l^i - 2(D-9)K^2K_{ij}^2 + (D-7)((K_{ij}^2)^2 + K^4)\}. \tag{34}$$

For a physically meaningful solution whose ADM mass and angular momenta are finite for the asymptotically flat case, or in the case of $\Lambda \neq 0$ whose Abbott-Deser [13] charges are finite, this quantity is expected to be finite and represents the nonstationary part of the total energy by construction. Observe that while the ADM momentum ($P_i = \oint_{\partial\Sigma} K_{ij} dS^j$) and angular momenta ($J^{jk} = \oint_{\partial\Sigma}(x^j K^{km} - x^k K^{jm})dS_m$) are linear in the extrinsic curvature given as integrals over the boundary, $\mathcal{I}(\mathcal{N})$ has quadratic, cubic and quartic terms in the extrinsic curvature in the bulk integral$\oint$.

Before we lay out the details of the above discussion, let us note that our final formula (28) can be reduced in various ways depending on the physical problem or the numerical integration scheme: for example, one can choose the maximal slicing gauge for which $\mathrm{Tr}(\pi) = K = 0$. If the problem permits time-symmetric initial data $\pi^{ij} = K^{ij} = 0$, then in this restricted case, $V^{ij} = Q^{ij} = 0$, and the integral (28) reduces to

$$\mathcal{I}(\mathcal{N}) = \int_\Sigma dV({}^\Sigma R_{ij}^2 N^2 + (D_i D_j N)^2 - 2{}^\Sigma R^{ij} N D_i D_j N + 2{}^\Sigma R N \triangle N + 4 D_m D_i N_j D^m D^{(i}N^{j)} + (D-3)\triangle N \triangle N).$$

Let us go back to (27) which was the defining relation of the invariant and try to write it as a boundary integral over the boundary of the hypersurface $\Sigma$. Then one has

$$\mathcal{I}(\mathcal{N}) = \int_\Sigma dV \tilde{\mathcal{P}}^*(\mathcal{N}) \cdot \tilde{\mathcal{P}}^*(\mathcal{N}) = \int_\Sigma dV \mathcal{N} \cdot \tilde{\mathcal{P}} \circ \tilde{\mathcal{P}}^*(\mathcal{N}) + \oint_{\partial\Sigma} dS n^k B_k, \tag{35}$$

which requires $\tilde{\mathcal{P}} \circ \tilde{\mathcal{P}}^*(\mathcal{N}) = 0$. This the approximate KID equation introduced by Dain [1] and $B_k$ is the boundary term to be found below. Note that our bulk integral (28) is more general and does not assume the existence of approximate symmetries.

## III. DETAILS OF THE CONSTRUCTION IN $D$ DIMENSIONS

### A. Boundary integral

The importance of the Einstein constraints (4) cannot be overstated: clearly the initial data is not arbitrary, one must solve these equations to feed the evolution equations; but, as importantly, the constraints also determine the evolution equations and they are related to the symmetries of the spacetime in a rather intricate way as we have seen above. One can consider the constraints (4) as the kernel of a map $\Phi$

$$\Phi\colon \mathcal{M}_2 \times \mathcal{S}_2^* \to \mathcal{C}^* \times \mathcal{X}^*, \tag{36}$$

where $\mathcal{M}_2$ denotes the space of the Riemannian metrics and $\mathcal{S}_2^*$ denotes the space of symmetric rank-2 tensor densities, $\mathcal{C}^*$ denotes the space of scalar function densities and $\mathcal{X}^*$ the space of vector field densities on the hypersurface $\Sigma$. We can express the constraint map explicitly as

$$\Phi\begin{pmatrix}\gamma_{ij}\\ \pi^{ij}\end{pmatrix} = \begin{pmatrix}\sqrt{\gamma}(2\Lambda - {}^\Sigma R) + \gamma^{-1/2}(\pi_{ij}^2 - \frac{\pi^2}{D-2})\\ -2\gamma_{ki}D_j\pi^{kj}\end{pmatrix}, \tag{37}$$

whose linearization can be found to be

$$D\Phi\begin{pmatrix}h_{ij}\\ p^{ij}\end{pmatrix} = \begin{pmatrix}\sqrt{\gamma}({}^\Sigma R^{ij} - D^iD^j + \gamma^{ij}\triangle)h_{ij}\frac{1}{\sqrt{\gamma}}\left(\gamma^{ij}\left(\frac{\pi^2}{D-2} - \pi_{ij}^2\right) + 2(\pi^{ik}\pi_k^j - \frac{\pi^{ij}\pi}{D-2})\right)h_{ij} + \frac{2}{\sqrt{\gamma}}(\pi_{ij} - \frac{\pi\gamma_{ij}}{D-2})p^{ij}\\ (\pi^{ij}D_k - 2\delta_k^{(i}\pi^{j)l}D_l)h_{ij} - 2\gamma_{k(i}D_{j)}p^{ij}\end{pmatrix}. \tag{38}$$

We can define a $2\times 2$ matrix as

$$D\Phi := \begin{pmatrix}\sqrt{\gamma}({}^\Sigma R^{ij} - D^iD^j + \gamma^{ij}\triangle) + \frac{1}{\sqrt{\gamma}}\left(\gamma^{ij}\left(\frac{\pi^2}{D-2} - \pi_{ij}^2\right) + 2(\pi^{ik}\pi_k^j - \frac{\pi^{ij}\pi}{D-2})\right) & \frac{2}{\sqrt{\gamma}}(\pi_{ij} - \frac{\pi\gamma_{ij}}{D-2})\\ \pi^{ij}D_k - 2\delta_k^{(i}\pi^{j)l}D_l & -2\gamma_{k(i}D_{j)}\end{pmatrix}, \tag{39}$$

such that

$$D\Phi\begin{pmatrix}h_{ij}\\ p^{ij}\end{pmatrix} = D\Phi\circ\begin{pmatrix}h_{ij}\\ p^{ij}\end{pmatrix}. \tag{40}$$

Defining [7]

$$\tilde{\mathcal{P}} \coloneqq D\Phi\circ\begin{pmatrix} \gamma^{-1/2} & 0 \\ 0 & -D^m \end{pmatrix}, \tag{41}$$

one finds

$$\tilde{\mathcal{P}} \coloneqq \begin{pmatrix} {}^{\Sigma}R^{ij} - D^iD^j + \gamma^{ij}\triangle + \frac{1}{\gamma}\left(\gamma^{ij}\left(\frac{\pi^2}{D-2} - \pi^2_{ij}\right) + 2\left(\pi^{ik}\pi^j_k - \frac{\pi^{ij}\pi}{D-2}\right)\right) & \frac{2}{\sqrt{\gamma}}\left(\frac{\pi\gamma_{ij}}{D-2} - \pi_{ij}\right)D^m \\ \frac{1}{\sqrt{\gamma}}(\pi^{ij}D_k - 2\delta^{(i}_k\pi^{j)l}D_l) & 2\gamma_{k(i}D_{j)}D^m \end{pmatrix}, \tag{42}$$

which is a map as

$$\tilde{\mathcal{P}}\colon\ \mathcal{S}_2\times\mathcal{S}_{1,2}\to\mathcal{C}\times\mathcal{X}, \tag{43}$$

where $\mathcal{S}_2$ denotes the space of covariant rank-2 tensors, $\mathcal{S}_{1,2}$ denotes the space of covariant rank-3 tensors which are symmetric in last two indices, $\mathcal{C}$ denotes the space of scalar function and $\mathcal{X}$ the space of vector fields on the hypersurface Σ.

The formal adjoint of $\tilde{\mathcal{P}}$-operator was defined in (26) via the (21) and it is a map of the form

$$\tilde{\mathcal{P}}^*\colon\ \mathcal{C}\times\mathcal{X}\to\mathcal{S}_2\times\mathcal{S}_{1,2}. \tag{44}$$

Working out the details, one arrives at

$$\tilde{\mathcal{P}}^*\begin{pmatrix} N \\ N^k \end{pmatrix} = \begin{pmatrix} N^{\Sigma}R^{ij} - D^iD^jN + \gamma^{ij}\triangle N + Q^{ij} \\ D_m(2D_{(i}N_{j)} + V_{ij}) \end{pmatrix}, \tag{45}$$

where $V^{ij}$ and $Q^{ij}$ were given (29), (30) respectively. We have used this expression in the previous section to find the bulk integral of the nonstationary energy. Now let us use this operator and its adjoint to find an expression on the boundary. For this purpose we need the following identity:

$$\int_\Sigma dV\begin{pmatrix} N \\ N^k \end{pmatrix}\cdot\tilde{\mathcal{P}}\begin{pmatrix} s_{ij} \\ s_{kij} \end{pmatrix} = \int_\Sigma dV\begin{pmatrix} s_{ij} \\ s_{kij} \end{pmatrix}\cdot\tilde{\mathcal{P}}^*\begin{pmatrix} N \\ N^k \end{pmatrix} + \oint_{\partial\Sigma} dSn^k\mathcal{B}_k, \tag{46}$$

with generic $s_{ij}\in\mathcal{S}_2$ and $s_{kij}\in\mathcal{S}_{1,2}$. After making use of (42) and (45), a slightly cumbersome computation yields the boundary term:

$$\begin{aligned}\mathcal{B}_k &= s_{kj}D^jN - ND^js_{kj} + ND_ks - sD_kN + 2N^iD^js_{jki} \\ &\quad - 2s_{kij}D^iN^j + \frac{2N}{\sqrt{\gamma}}\left(\frac{\pi}{D-2}s_{kj}{}^j - s_{kij}\pi^{ij}\right) \\ &\quad + \frac{1}{\sqrt{\gamma}}(\pi^{ij}s_{ij}N_k - 2s_{ij}N^i\pi^j_k),\end{aligned} \tag{47}$$

where $s = \gamma^{ij}s_{ij}$. Let us now assume a particular $s_{ij}$ and a particular $s_{kij}$ such that

$$\begin{pmatrix} s_{ij} \\ s_{kij} \end{pmatrix} \coloneqq \tilde{\mathcal{P}}^*\begin{pmatrix} N \\ N^k \end{pmatrix}, \tag{48}$$

which yields

$$\tilde{\mathcal{P}}\begin{pmatrix} s_{ij} \\ s_{kij} \end{pmatrix} = \tilde{\mathcal{P}}\circ\tilde{\mathcal{P}}^*\begin{pmatrix} N \\ N^k \end{pmatrix}. \tag{49}$$

Then (46) becomes

$$\int_\Sigma dV\begin{pmatrix} N \\ N^k \end{pmatrix}\cdot\tilde{\mathcal{P}}\circ\tilde{\mathcal{P}}^*\begin{pmatrix} N \\ N^k \end{pmatrix} = \mathcal{I}(\mathcal{N}) + \oint_{\partial\Sigma} dSn^k\mathcal{B}_k, \tag{50}$$

where $\mathcal{B}_k$ given in (47) must be evaluated with

$$s_{ij} = N^{\Sigma}R_{ij} - D_iD_jN + \gamma_{ij}\triangle N + Q_{ij} \tag{51}$$

and

$$s_{kij} = D_k(2D_{(i}N_{j)} + V_{ij}). \tag{52}$$

Equation (50) shows that generically $\mathcal{I}(\mathcal{N})$ cannot be written as an integral on the boundary of the hypersurface unless $\tilde{\mathcal{P}}\circ\tilde{\mathcal{P}}^*(\mathcal{N}) = 0$. In that case, the invariant reduces to

$$\mathcal{I}(\mathcal{N}) = -\oint_{\partial\Sigma} dSn^k\mathcal{B}_k. \tag{53}$$

Explicit computation shows that one has

$$\begin{aligned}\mathcal{B}_k &= \frac{N^2}{2}D_k{}^{\Sigma}R + N^{\Sigma}R_{kj}D^jN - D_kD_jND^jN \\ &\quad - (D-3)D_kN\triangle N + (D-2)ND_k\triangle N \\ &\quad + 4N^i\triangle D_{(k}N_{i)} - 4D_kD_{(i}N_{j)}D^{(i}N^{j)} + b_k,\end{aligned} \tag{54}$$

where

$$\begin{aligned} b_k &:= Q_{kj}D^jN - ND^jQ_{kj} + ND_kQ - QD_kN + 2N^i\Delta V_{ki} \\ &- 2D_kV_{ij}D^iN^j + \frac{1}{\sqrt{\gamma}}\frac{2N\pi}{D-2}(2D_kD_iN^i + D_kV) \\ &- \frac{2N\pi^{ij}}{\sqrt{\gamma}}(2D_kD_iN_j + D_kV_{ij}) \\ &+ \frac{1}{\sqrt{\gamma}}(\pi^{ij}N_k - 2N^i\pi^j_k) \\ &\times (N^\Sigma R_{ij} - D_iD_jN + \gamma_{ij}\Delta N + Q_{ij}). \end{aligned} \tag{55}$$

In the Gaussian normal coordinates the boundary integral reads

$$\begin{aligned} \mathcal{I}(\mathcal{N}) = \oint_{\partial\Sigma} dSn^k\bigg(\left(D-\frac{5}{2}\right)D_kK^2_{ij} + \left(\frac{7}{2}-D\right)D_kK^2 \\ + 2K^{lj}D_jK_{lk}\bigg). \end{aligned} \tag{56}$$

Another physically relevant case is the time symmetric asymptotically flat case for which the boundary integral reduces to

$$\begin{aligned} \mathcal{I}(\mathcal{N}) = \oint_{\partial\Sigma} dSn^k(D_kD_jND^jN + (D-3)D_kN\Delta N \\ - (D-2)ND_k\Delta N - 4N^i\Delta D_{(k}N_{i)} \\ + 4D_kD_{(i}N_{j)}D^{(i}N^{j)}). \end{aligned}$$

In the most general form $N$ and $N^i$ should satisfy the fourth order equations $\tilde{\mathcal{P}}\circ\tilde{\mathcal{P}}^*(\mathcal{N}) = 0$ which explicitly read

$$\tilde{\mathcal{P}}\circ\tilde{\mathcal{P}}^*\begin{pmatrix} N \\ N^i \end{pmatrix} = \begin{pmatrix} (D-2)\Delta\Delta N - {}^\Sigma R_{ij}D^iD^jN + N(\frac{1}{2}\Delta{}^\Sigma R + {}^\Sigma R^2_{ij}) + 2{}^\Sigma R\Delta N + \frac{3}{2}D_i{}^\Sigma RD^iN + Y \\ 4D^j\Delta D_{(k}N_{j)} + Y_k \end{pmatrix} = 0, \tag{57}$$

where

$$\begin{aligned} Y :=^\Sigma R^{ij}Q_{ij} - D^iD^jQ_{ij} + \Delta Q + \frac{2}{\sqrt{\gamma}}\left(\frac{\pi\gamma^{ij}}{D-2} - \pi^{ij}\right) \\ \times\Delta(2D_iN_j + V_{ij}) + \left(\frac{2}{\gamma}\left(\pi^{ik}\pi^j_k - \frac{\pi\pi^{ij}}{D-2}\right)\right. \\ \left. - \frac{\gamma^{ij}}{\gamma}\left(\pi^2_{kl} - \frac{\pi^2}{D-2}\right)\right)(N^\Sigma R_{ij} - D_iD_jN + \gamma_{ij}\Delta N + Q_{ij}) \end{aligned} \tag{58}$$

and

$$\begin{aligned} Y_k := \frac{1}{\sqrt{\gamma}}(\pi^{ij}D_k - 2\delta^i_k\pi^{jl}D_l) \\ \times (N^\Sigma R_{ij} - D_iD_jN + \gamma_{ij}\Delta N + Q_{ij}) + 2D^i\Delta V_{ik}. \end{aligned} \tag{59}$$

### B. The approximate KID equation in $D$ dimensions

Following the $D = 4$ discussion of Dain [1] let us now study the approximate KID equation (57) in $D$ dimensions. It is easy to see that it is a fourth order elliptic operator for $D > 2$. This follows by computing the leading symbol: for this purpose let us consider the higher order derivative terms and set $D_i = \zeta_i$ and $|\zeta|^2 = \zeta^i\zeta_i$. Using (57), the leading symbol of operator reads

$$\sigma[\tilde{P}\circ\tilde{\mathcal{P}}^*](\zeta)\begin{pmatrix} N \\ N_i \end{pmatrix} = \begin{pmatrix} (D-2)|\zeta|^4N \\ 4|\zeta|^2\zeta^j\zeta_{(k}N_{j)} \end{pmatrix}. \tag{60}$$

For a nonzero covector $\zeta$, if $\sigma$ is an isomorphism (here a vector bundle isomorphism), then the operator is elliptic. For the first component, this requires $D \neq 2$ and for the second component contraction with $\zeta^k$ yields

$$|\zeta|^4\zeta^kN_k = 0. \tag{61}$$

Assuming $D \neq 2$ one has $\zeta^kN_k = 0$. Inserting it back in the second component one obtains

$$|\zeta|^4N_k = 0, \tag{62}$$

so for $|\zeta|^2 \neq 0$, the leading symbol is injective and the operator $\tilde{P}\circ\tilde{\mathcal{P}}^*$ is elliptic for $D > 2$.

### C. Asymptotically flat spaces

Consider the initial data set $(\Sigma, \gamma_{ij}, \pi^{ij})$ for the vacuum Einstein field equations with $n > 1$ asymptotically Euclidean ends: this is to avoid bulk simplicity and allow black holes. There exists a compact set $\mathcal{B}$ such that $\Sigma\backslash\mathcal{B} = \sum_{k=1}^n \Sigma_{(k)}$, where $\Sigma_{(k)}$, $k = 1, \ldots, n$ are open sets diffeomorphic to the complement of a closed ball in $\mathbb{R}^{D-1}$. Each asymptotic end $\Sigma_{(k)}$ admits asymptotically Cartesian coordinates. We consider the following decay assumptions, for $D > 3$, which are consistent with finite ADM mass and momenta:

$$\gamma_{ij} = \delta_{ij} + o(|x|^{(3-D)/2}), \tag{63}$$

$$\pi^{ij} = o(|x|^{(1-D)/2}), \tag{64}$$

where $\delta_{ij} = (++\cdots+)$. Note that $\delta_{ij} = \mathcal{O}(1)$ and beware of the small $o$ and the big $\mathcal{O}$ notation. One can compute the following decay behavior for the Christoffel connection

$$^{\Sigma}\Gamma^k_{ij} = o(|x|^{(1-D)/2}), \tag{65}$$

and the curvatures

$$^{\Sigma}R^k{}_{lmn} = o(|x|^{-(1+D)/2}), \qquad ^{\Sigma}R_{ij} = o(|x|^{-(1+D)/2}), \qquad ^{\Sigma}R = o(|x|^{-(1+D)/2}). \tag{66}$$

### D. KIDs in $D$ dimensions

Let $(\Sigma, \gamma_{ij}, \pi^{ij})$ denote a smooth vacuum initial data set satisfying the decay assumptions (63), (64). Let $N, N^i$ be a smooth scalar field and a vector field on $\Sigma$ satisfying the KID equations. Then generalizing the $D = 4$ result of [9], the behavior of all the possible solutions were given in [10] which we quote here.

(1) There exits an antisymmetric tensor field $\omega_{\mu\nu}$, such that

$$N - \omega_{0i}x^i = o(|x|^{(5-D)/2}), \qquad N^i - \omega^i{}_j x^j = o(|x|^{(5-D)/2}). \tag{67}$$

(2) If $\omega_{\mu\nu} = 0$, then there exists a vector field $\mathcal{U}^\mu$, such that

$$N - \mathcal{U}^0 = o(|x|^{(3-D)/2}), \qquad N^i - \mathcal{U}^i = o(|x|^{(3-D)/2}). \tag{68}$$

(3) If $\omega_{\mu\nu} = 0 = \mathcal{U}^\mu$ then one has the trivial solution $N = 0 = N^i$. Both $\omega_{\mu\nu}$ and $\mathcal{U}^\mu$ are constants in the sense that they are $\mathcal{O}(1)$ whenever they do not vanish.

Case 1 above corresponds to the rotational Killing vectors while case 2 corresponds to the translational ones we shall employ the latter.

We explained in Sec. II that solutions of the $D\Phi^*(N, N^i) = 0$ yield spacetime Killing vectors. It is not difficult to see that the modified equation $\tilde{\mathcal{P}}^*(N, N^i) = 0$ yields only the Killing vectors for the case of translational KIDs (63), (64). Here is the proof: $\tilde{\mathcal{P}}^*(N, N^i) = 0$ implies

$$N^{\Sigma}R_{ij} - D_iD_jN + \gamma_{ij}\Delta N + Q_{ij} = 0, \tag{69}$$

$$D_m(2D_{(i}N_{j)} + V_{ij}) = 0. \tag{70}$$

If one assumes $(N, N^i)$ decay as in (68) we have $D_{(i}N_{j)} = o(|x|^{(1-D)/2})$; and $V_{ij} = o(|x|^{(1-D)/2})$, then

$$2D_{(i}N_{j)} + V_{ij} = o(|x|^{(1-D)/2}) \tag{71}$$

vanishes at infinity; and since it is covariantly constant, it must vanish identically

$$2D_{(i}N_{j)} + V_{ij} = 0. \tag{72}$$

Together with the first component of $\tilde{\mathcal{P}}^*(N, N^i) = 0$ we get the formal adjoint of the linearized constraint map, namely $D\Phi^*(N, N^i) = 0$. We can conclude that if $\tilde{\mathcal{P}}^*(N, N^i) = 0$ then $(N, N^i)$ solve the KID equations.

### E. Approximate KIDs in $D$ dimensions

Generalizing Dain's $D = 4$ result, let us search for translational solutions of the approximate Killing equation[2]

$$\tilde{\mathcal{P}}\circ\tilde{\mathcal{P}}^*\begin{pmatrix} N \\ N^i \end{pmatrix} = 0 \tag{73}$$

as a deformation of the KIDs $(X, N^i)$ in the following form:

$$N = \lambda\varphi + X, \qquad N^i = N^i, \tag{74}$$

where the function $\varphi$ is to be found, $\lambda$ is a constant. KIDs decay as

$$X - \mathcal{U}^0 = o(|x|^{(3-D)/2}), \tag{75}$$

$$N^i - \mathcal{U}^i = o(|x|^{(3-D)/2}). \tag{76}$$

Inserting the ansatz (74) into the approximate KID equation (73), one gets

$$\tilde{\mathcal{P}}\circ\tilde{\mathcal{P}}^*\begin{pmatrix} \varphi \\ 0 \end{pmatrix} = -\tilde{\mathcal{P}}\circ\tilde{\mathcal{P}}^*\begin{pmatrix} X \\ N^i \end{pmatrix} = 0, \tag{77}$$

or more explicitly

$$\tilde{\mathcal{P}}\circ\tilde{\mathcal{P}}^*\begin{pmatrix} \varphi \\ 0 \end{pmatrix} = \begin{pmatrix} (D-2)\Delta\Delta\varphi - {}^{\Sigma}R_{ij}D^iD^j\varphi + \varphi(\frac{1}{2}\Delta^{\Sigma}R + {}^{\Sigma}R_{ij}^2) + 2^{\Sigma}R\Delta\varphi + \frac{3}{2}D_i{}^{\Sigma}RD^i\varphi + Y \\ Y_k \end{pmatrix} = 0. \tag{78}$$

[2]We work in a given asymptotic end and not the clutter the notation we do not denote the corresponding index referring to the asymptotic end.

For such a $\varphi$, the bulk integral (28) becomes

$$\mathcal{I}(\mathcal{N}) = \lambda^2 \int_\Sigma dV\{|D_m V^{ij}|^2 + {}^\Sigma R_{ij}^2 \varphi^2 + (D_i D_j \varphi)^2 - 2{}^\Sigma R^{ij} N D_i D_j \varphi + 2{}^\Sigma R \varphi \Delta\varphi + (D-3)\Delta\varphi\Delta\varphi + 2Q\Delta\varphi + Q_{ij}^2 + 2{}^\Sigma R_{ij}\varphi Q^{ij} - 2Q^{ij} D_i D_j \varphi\}, \tag{79}$$

where

$$V^{ij} = 2\varphi K^{ij}, \tag{80}$$

and

$$Q^{ij} = 2\varphi(K^i_k K^{kj} - K K^{ij}) - \varphi\gamma^{ij}(K_{kl}^2 - K^2). \tag{81}$$

The boundary form for the asymptotically flat case follows similarly

$$\mathcal{I}(\mathcal{N}) = -\lambda^2 \oint_{\partial\Sigma} dS n^k\{-D_k D_j \varphi D^j \varphi - (D-3) D_k \varphi \Delta\varphi + (D-2)\varphi D_k \Delta\varphi + Q_{kj} D^j \varphi - \varphi D^j Q_{kj} - 2\varphi K^{ij} D_k V_{ij}\}, \tag{82}$$

where we used $K_{kl}^2 - K^2 = {}^\Sigma R = 0$ on the boundary.

## IV. CONCLUSIONS

Using the Hamiltonian form of the Einstein evolution equations as given by Fischer and Marsden [8], we constructed an integral that measures the nonstationary energy contained in a spacelike hypersurface in $D$ dimensional general relativity with or without a cosmological constant. This integral was previously studied by Dain [1] who used the Einstein constraints but not the evolution equations. The crucial observation is the following: the critical points of the first order Hamiltonian form of Einstein equations correspond to the initial data which possess Killing symmetries, a result first observed by Moncrief [5]. Hence, our vantage point is that the failure of an initial data to possess Killing symmetries is given by the evolution equations, namely nonvanishing of the time derivatives of the spatial metric and the canonical momenta. Then manipulating the evolution equations, one arrives at the integral (28). Once an initial data is given, one can compute this integral, which by construction, vanishes for stationary spacetimes.

## APPENDIX: ADM SPLIT OF EINSTEIN'S EQUATIONS IN $D$ DIMENSIONS

For the sake of completeness let us give here the ADM split of Einstein's equations and all the relevant tensors. Using the $(D-1)+1$ dimensional decomposition of the metric given as (1) we have:

$$g_{00} = -N^2 + N_i N^i, \qquad g_{0i} = N_i, \qquad g_{ij} = \gamma_{ij}, \tag{A1}$$

and

$$g^{00} = -\frac{1}{N^2}, \qquad g^{0i} = \frac{1}{N^2} N^i, \qquad g^{ij} = \gamma^{ij} - \frac{1}{N^2} N^i N^j. \tag{A2}$$

Let $\Gamma^\mu_{\nu\rho}$ denote the Christoffel symbol of the $D$ dimensional spacetime

$$\Gamma^\mu_{\nu\rho} = \frac{1}{2} g^{\mu\sigma}(\partial_\nu g_{\rho\sigma} + \partial_\rho g_{\nu\sigma} - \partial_\sigma g_{\nu\rho}) \tag{A3}$$

and let ${}^\Sigma\Gamma^k_{ij}$ denote the Christoffel symbol of the $D-1$ dimensional hypersurface, which is compatible with the spatial metric $\gamma_{ij}$ as

$${}^\Sigma\Gamma^k_{ij} = \frac{1}{2}\gamma^{kp}(\partial_i \gamma_{jp} + \partial_j \gamma_{ip} - \partial_p \gamma_{ij}). \tag{A4}$$

Then a simple computation shows that

$$\Gamma^0_{00} = \frac{1}{N}(\dot{N} + N^k(\partial_k N + N^i K_{ik})) \tag{A5}$$

and

$$\Gamma^0_{0i} = \frac{1}{N}(\partial_i N + N^k K_{ik}), \qquad \Gamma^0_{ij} = \frac{1}{N} K_{ij}, \qquad \Gamma^k_{ij} = {}^\Sigma\Gamma^k_{ij} - \frac{N^k}{N} K_{ij} \tag{A6}$$

and

$$\Gamma^i_{0j} = -\frac{1}{N} N^i(\partial_j N + K_{kj} N^k) + N K_j{}^i + D_j N^i \tag{A7}$$

and also

$$\Gamma^i_{00} = -\frac{N^i}{N}(\dot{N} + N^k(\partial_k N + N^l K_{kl})) + N(\partial^i N + 2N^k K_k{}^i) + \dot{N}^i + N^k D_k N^i. \tag{A8}$$

Starting with the definition of the $D$ dimensional Ricci tensor

$$R_{\rho\sigma} = \partial_\mu \Gamma^\mu_{\rho\sigma} - \partial_\rho \Gamma^\mu_{\mu\sigma} + \Gamma^\mu_{\mu\nu}\Gamma^\nu_{\rho\sigma} - \Gamma^\mu_{\sigma\nu}\Gamma^\nu_{\mu\rho} \tag{A9}$$

one arrives at

$$R_{ij} = {}^\Sigma R_{ij} + K K_{ij} - 2K_{ik} K^k_j + \frac{1}{N}(\dot{K}_{ij} - N^k D_k K_{ij} - D_i D_j N - K_{ki} D_j N^k - K_{kj} D_i N^k), \tag{A10}$$

where ${}^{\Sigma}R_{ij}$ denotes the Ricci tensor of the hypersurface. The remaining components can also be found to be

$$R_{00} = N^i N^j R_{ij} - N^2 K_{ij} K^{ij} + N(D_k D^k N - \dot{K} - N^k D_k K + 2N^k D_m K^m_k) \tag{A11}$$

and

$$R_{0i} = N^j R_{ij} + N(D_m K^m_i - D_i K). \tag{A12}$$

The scalar curvature can be found as

$$R = {}^{\Sigma}R + K^2 + K_{ij}K^{ij} + \frac{2}{N}(\dot{K} - D_k D^k N - N^k D_k K). \tag{A13}$$

Under the above splitting the cosmological Einstein equations

$$R_{\mu\nu} - \frac{1}{2} g_{\mu\nu} R + \Lambda g_{\mu\nu} = \kappa T_{\mu\nu} \tag{A14}$$

split in to constraints and evolution equations in local coordinates. The momentum constraints read

$$N(D_k K^k_i - D_i K) - \kappa(T_{0i} - N^j T_{ij}) = 0, \tag{A15}$$

via the Hamiltonian constraint becomes

$$N^2({}^{\Sigma}R + K^2 - K^2_{ij} - 2\Lambda) - 2\kappa(T_{00} + N^i N^j T_{ij} - 2N^i T_{0i}) = 0. \tag{A16}$$

On the other hand the evolution equations for the metric and the extrinsic curvature become

$$\frac{\partial}{\partial t}\gamma_{ij} = 2NK_{ij} + D_i N_j + D_j N_i, \tag{A17}$$

$$\frac{\partial}{\partial t}K_{ij} = N(R_{ij} - {}^{\Sigma}R_{ij} - KK_{ij} + 2K_{ik}K^k_j) + \mathcal{L}_{\vec{N}} K_{ij} + D_i D_j N, \tag{A18}$$

where $\mathcal{L}_{\vec{N}}$ is the Lie derivative along the shift vector.


[1] S. Dain, A New Geometric Invariant on Initial Data for the Einstein Equations, Phys. Rev. Lett. **93**, 231101 (2004).

[2] R. Arnowitt, S. Deser, and C. Misner, The dynamics of general relativity, Phys. Rev. **116**, 1322 (1959); Canonical variables for general relativity, Phys. Rev. **117**, 1595 (1960); in*Gravitation: An Introduction to Current Research*, edited by L. Witten (Wiley, New York, 1962).

[3] J. A. V. Kroon and J. L. Williams, Dain's invariant on non-time symmetric initial data sets, Classical Quantum Gravity **34**, 125013 (2017).

[4] B. P. Abbott *et al.* (LIGO Scientific and Virgo Collaborations), Observation of Gravitational Waves from a Binary Black Hole Merger, Phys. Rev. Lett. **116**, 061102 (2016).

[5] V. Moncrief, Spacetime symmetries and linearization stability of the Einstein equations. I, J. Math. Phys. (N.Y.) **16**, 493 (1975).

[6] R. Beig and P. T. Chruściel, Killing initial data, Classical Quantum Gravity **14**, A83 (1997).

[7] R. Bartnik, Phase space for the Einstein equations, Commun. Anal. Geom. **13**, 845 (2005).

[8] A. E. Fischer and J. E. Marsden, The Einstein evolution equations as a first-order quasi-linear symmetric hyperbolic system. I, Commun. Math. Phys. **28**, 1 (1972).

[9] R. Beig and P. T. Chruściel, Killing vectors in asymptotically flat spacetimes. I. Asymptotically translational Killing vectors and rigid positive energy theorem, J. Math. Phys. (N.Y.) **37**, 1939 (1996).

[10] P. T. Chrusciel and D. Maerten, Killing vectors in asymptotically flat space-times. II. Asymptotically translational Killing vectors and the rigid positive energy theorem in higher dimensions, J. Math. Phys. (N.Y.) **47**, 022502 (2006).

[11] B. S. DeWitt, Quantum theory of gravity. 1. The canonical theory, Phys. Rev. **160**, 1113 (1967).

[12] A. E. Fischer and J. E. Marsden, Linearization stability of the Einstein equations, Bull. Am. Math. Soc. **79**, 997 (1973).

[13] L. F. Abbott and S. Deser, Stability of gravity with a cosmological constant, Nucl. Phys. **B195**, 76 (1982).

T08

# Sub-GeV dark matter search at ILC beam dumps

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Light dark matter particles may be produced in electron and positron beam dumps of the International Linear Collider (ILC). We propose an experimental setup to search for such events, the Beam-Dump eXperiment at the ILC (ILC-BDX). The setup consists of a muon shield placed behind the beam dump, followed by a multi-layer tracker and an electromagnetic calorimeter. The calorimeter can detect electron recoils due to elastic scattering of dark matter particles produced in the dump, while the tracker is sensitive to decays of excited dark-sector states into the dark matter particle. We study the production, decay and scattering of sub-GeV dark matter particles in this setup in several models with a dark photon mediator. Taking into account beam-related backgrounds due to neutrinos produced in the beam dump as well as the cosmic-ray background, we evaluate the sensitivity reach of the ILC-BDX experiment. We find that the ILC-BDX will be able to probe interesting regions of the model parameter space and, in many cases, reach well below the relic target.
In this talk, we show the capability of the ILC-BDX, based on Ref. [1].

[1] K. Asai, S. Iwamoto, M. Perelstein, Y. Sakaki and D. Ueda, [arXiv:2301.03816 [hep-ph]].

T09

# The potential of MeerKAT for indirect dark matter searches: Abstract for Dark Side of the Universe 2023

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The nature of dark matter remains a major hiatus in the modern understanding of cosmology and astrophysics. In the pursuit of this mystery, the potential for radio-frequency indirect dark matter detection has long been known in the literature [1, 2, 3, 4, 5, 6]. However, the potential of MeerKAT, currently the world's best radio interferometer, has been only minimally discussed [7, 8, 9, 10, 11, 12]. In this work we will present the potential non-observation constraints on WIMPs and axion-like particles from the MeerKAT instrument in a variety of environments, from dwarf galaxies to galaxy clusters, identifying the properties of ideal search targets and discussing some prominent objects like Omega Centauri. These are performed via simulation of MeerKAT observations using the Stimela [13] software suite.

[1] S. Colafrancesco, S. Profumo, & P. Ullio, A&A, 455, 21 (2006).
[2] S. Colafrancesco, S. Profumo, & P. Ullio, Phys. Rev. D, 75, 023513 (2007).
[3] E. Storm, T. Jeltema, S. Profumo, & L. Rudnick, ApJ, 768, 106 (2013).
[4] G. Beck, S. Colafrancesco & P. Marchegiani, JCAP, **02**, 032C (2015).
[5] M. H. Chan, Phys. Rev. D, **96**, 043009 (2017).
[6] A. Caputo, M. Regis, M. Taoso, & S. J. Witte, JCAP, 2019, 027 (2019).
[7] G. Beck, Galaxies, 7(1), 16 (2019).
[8] E. Pinetti et al, JCAP, **07**, 044 (2020).
[9] A. Ayad & G. Beck, JCAP, **03**, 005 (2022).
[10] N. Lavis & G. Beck, arXiv: 2208.06813 (2022).
[11] Y. Zhou et al, Phys. Rev., D, **106**, 083006 (2022).
[12] G. Beck et al, Astropart. Phys., **148**, 102821 (2023).
[13] S. Makhathini, PhD thesis, Rhodes University, South Africa, available via http://hdl.handle.net/10962/57348 (2018).

T10

# WIMP and FIMP multi-component dark matter

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We consider the case of multi-component dark matter where one of the dark matter candidate is a WIMP, the other a FIMP. We discuss the interplay of the freeze-out and freeze-in production mechanism and the relative contribution to the final relic density within an extension of the standard model which can explain both neutrino masses and dark matter. This model can feature long-lived particles that can alter the successful BBN predictions for the abundance of light elements. We discuss possible probes of the model in ongoing direct, indirect and collider experiments, as well as in future detectors such as MATHUSLA [1].

[1] G. Bélanger, S. Choubey, R. Godbole, S. Khan, M. Mitra, A. Roy, JHEP. **11**, 133 (2022).

T11

# Gravitational wave probes of baryogenesis and dark matter scenarios

T12

# Effects of General Neutrino Interactions on Cosmic Neutrino Background Detection at PTOLEMY

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The standard big bang theory, along with cosmic microwave background (CMB), predicts the existence of a cosmic neutrino background (C$\nu$B). While the discovery and subsequent precision measurement of CMB help us understanding the Universe back till the time of recombination, the C$\nu$B remained undetected till date. The C$\nu$B, which consists of the relic neutrinos that decoupled at the time of neutrino decoupling when the Universe was about one second old, if detected can help us probe back even deeper than CMB. PTOLEMY is one of the proposed experiment that will try to detect C$\nu$B by capturing electron neutrino on a 100 g tritium target. In the standard model (SM) of particle physics neutrinos interact weakly via charged and neutral current interactions. In the low-energy limit these interactions be written as effective four-fermion operators. In the most general case the neutrinos can have exotic interactions which are known as general neutrino interactions (GNI). GNIs can be thought of as low-energy effective description of many well-motivated BSM scenarios. In this talk we will focus on the effect of GNIs in the detection prospects of PTOLEMY. We also show how the differential electron spectrum is sensitive to the finite experimental resolution, mass of the lightest neutrino eigenstate, the strength of these interactions and the ordering of neutrino mass.

[1] I. K. Banerjee, U. K. Dey, N. Nath, S. S. Shariff, arXiv:2304.02505.

T13

# Probing primordial non-gaussianity with the multi-tracer technique.

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In the near future, cosmic large-scale structure will be mapped with increasing detail by the next generation of observational facilities operating at various wavelengths (radio, optical, infra-red) and exploiting various techniques. Meanwhile, theory and simulations are increasing in sophistication in their ability to describe large-scale structure. These advances could potentially allow greater constraints into cosmological parameters and theories of galaxy evolution. Constraining $f_{\rm NL}$ values provides important information about the mechanisms that generated the primordial non-Gaussian fluctuations and the physics of the early universe. Observations of the CMB, large-scale structure, and galaxy clustering have been used to place limits on $f_{\rm NL}$. However, current constraints are still far from the target precision needed to discriminate between different models of the early universe. We review the current state of $f_{\rm NL}$ constraints and the ongoing efforts to improve their accuracy, challenges, limitations of different observational methods, and the potential for future experiments to significantly improve our understanding of $f_{\rm NL}$ and the physics of the early universe.

[1] S. Jolicoeur, R. Maartens, S. Dlamini, **Eur. Phys. J. C** , (2023) 83:320.

T14

# Synergies between Rubin Observatory and the Roman Space Telescope

T15

# A direct detection view of the NSI landscape

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The next generation of dark matter (DM) direct detection (DD) experiments are becoming sensitive to the scattering of solar neutrinos. It will be the first time coherent elastic neutrino-nucleus scattering (CEvNS) will be detected from incident solar neutrinos, as opposed to spallation sources. Simultaneously, it will provide a complementary measurement of solar neutrinos via elastic neutrino-electron scattering.
In this talk I will present the results of [1] highlighting the implications of these novel signals as a new means of testing neutrino non-standard neutrino interactions (NSI). For this purpose, I will introduce a convenient parameterization of the NSI parameter space that allows for an easy visualisation of the complementarity of different experimental sources. I will then contrast the sensitivities of DM DD experiments on NSIs with previous results from spallation sources and neutrino oscillation experiments, highlighting the importance of including DD experiments in future global fits of the NSI landscape.

[1] Dorian W. P. Amaral, David Cerdeno, Andrew Cheek, Patrick Foldenauer, submitted to JHEP, arXiv:2302.12846 [hep-ph] (2023).

T16

# Dark matter heating vs vortex creep heating in old neutron stars

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Old isolated neutron stars have been gathering attention as targets to probe Dark Matter (DM) through temperature observations [1]. DM will anomalously heat neutron stars through its gravitational capture and annihilation process, which predicts $T_s \simeq (1-3) \times 10^3$ K for $t \gtrsim 10^6$ years. If we find even colder neutron stars in the infrared telescope, such as the James Webb Space Telescope [2], we may constrain DM-nucleon scattering cross section [3].

This scenario, however, assumed that there is no relevant heating source for old neutron stars [4, 5]. To examine this assumption, we studied the creep motion of vortex lines in the neutron superfluid of the inner crust as the heating mechanism [6, 7]. This mechanism is inherent in the structure of neutron stars and is expected to be universal. Therefore, to constrain DM physics through temperature observations, this mechanism may cause serious background.

To address this contamination for dark matter heating, we characterized the quantitative impact of vortex creep heating on the predicted surface temperature. First, we examined the numerical evaluation of the creep rate, a radial velocity of the vortex line, based on the model proposed in Ref. [8]. We concluded that the steady creep motion is realized in the old isolated neutron stars whose temperatures are recently observed. Second, we specified the predicted temperature by inputting the vortex-nucleon interactions [9, 10, 11]. Finally, we compared this prediction with the observed surface temperature of old isolated neutron stars and revealed the current status of the vortex creep heating in light of observations. In this talk, we will conclude the significance of DM heating against vortex creep heating within the current uncertainties from both nuclear and astrophysical sides.

[1] C. Kouvaris, Phys. Rev. D **77** (2008), 023006.
[2] J. P. Gardner, J. C. Mather, M. Clampin, R. Doyon, M. A. Greenhouse, H. B. Hammel, J. B. Hutchings, P. Jakobsen, S. J. Lilly and K. S. Long, *et al.* Space Sci. Rev. **123** (2006), 485.
[3] M. Baryakhtar, J. Bramante, S. W. Li, T. Linden and N. Raj, Phys. Rev. Lett. **119** (2017) no.13, 131801.
[4] D. G. Yakovlev, K. P. Levenfish and Y. A. Shibanov, Phys. Usp. **42** (1999), 737-778.
[5] D. Page, J. M. Lattimer, M. Prakash and A. W. Steiner, Astrophys. J. Suppl. **155** (2004), 623-650.
[6] M. A. Alpar, D. Pines, P. W. Anderson, and J. Shaham, Astrophys. J. **276** (1984) 325-334.
[7] N. Shibazaki and F. K. Lamb, Astrophys. J. **346** (1989) 808-822.
[8] B. Link, R. I. Epstein, G. Baym, Astrophysical Journal, **403**, no. 1, 285-302.
[9] P. Donati and P. M. Pizzochero, Nucl. Phys. A **742** (2004), 363-379.
[10] P. Avogadro, F. Barranco, R. A. Broglia and E. Vigezzi, Nucl. Phys. A **811** (2008), 378-412.
[11] S. Seveso, P. M. Pizzochero, F. Grill, and B. Haskell Monthly Notices of the Royal Astronomical Society **455** (2015).

T17

# Dark Matter, Muon g-2 and lepton flavor violation in SUSY-GUT theories.

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Supersymmetry is still an attractive extension of the Standard Model, since it provides a natural framework for unification of the fundamental interactions, as well as suitable Dark Matter candidates. In addition, SUSY contributions can explain certain discrepancies between experimental values and SM predictions for observables such as the muon g-2, and predict new physics, including lepton flavor violation. In our work we expand the results of [1] (where we have shown how SUSY contributions can solve the muon g-2 problem while predicting DM candidates in agreement with cosmological observations) to also explain LFV processes like BR($\mu \to e\gamma$), on the range of the experimental searches[2]. We also identify the properties of the lightest SUSY particle in several DM scenarios and show the complementarity between LFV and searches for SUSY particles at the LHC.

[1] M. E. Gomez, Q. Shafi, A. Tiwari and C. S. Un,
"Muon $\mathbf{g-2}$, neutralino dark matter and stau NLSP,"
Eur. Phys. J. C **82** (2022) no.6, 561 [arXiv:2202.06419 [hep-ph]].

[2] M. E. Gomez, S. Lola, Q. Shafi and C. S. Un,
"Dark Matter, Muon g-2 and lepton flavor violation in SUSY-GUT theories."
Work in progress.

T18

# Cosmology with Fast Radio Bursts

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Fast radio bursts (FRBs) are very short and bright transients visible over extragalactic distances. Their origin is still a mystery, but since the radio pulse undergoes dispersion caused by free electrons along the line of sight, FRBs can be used to probe the distribution of baryons on cosmological scales. In this talk, I will give a brief overview of this rapidly evolving field. Recently, together with my collaborators, I used the small sample of currently known events to measure the Hubble constant, and constraints are expected to improve considerably in the near future as more and more data from the ongoing search programs becomes available. Future large FRB samples provide exciting opportunities not only to understand the physical mechanism, but also to use their dispersion to measure the cosmic electrons on largest scales and to look for new physics such as primordial non-Gaussianity or deviations from the equivalence principle. The unique perspective FRBs offer on the baryon distribution allows to study feedback processes on cosmological scales and promises great complementarity with upcoming large-scale structure surveys.

T19

# New constraints on the dark matter-neutrino and the dark matter-photon scattering cross-sections

T20

# Modelling the radio - γ-ray emission components of Jetted AGNs

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**Abstract**

In this work, the *Fermi*-LAT multiband selected sample of active galactic nuclei (AGN) was used to investigate the emission spectra of blazars and radio galaxies. We computed the broadband emission spectra from the low-energy components of radio to X-ray, radio to γ-ray and the high energy component of X-ray to γ-ray bands. Our findings from the distributions of the continuous spectra clearly indicate that radio galaxies form the tail of the distributions in the low energy components, overlapping in a well-determined range (up to 4 orders of magnitude) in the high energy spectrum. A two-sample Kolmogorov-Smirnov test (*K-S* test) of the computed spectra showed that radio galaxies differ from blazars in the low energy components while there is no clear difference between them in the high energy component, which implies that high energy emissions in radio galaxies and blazars may be as a result of the same emission mechanism. There is a regular sequence of distributions on the continuous spectra planes for radio galaxies and blazar subsamples. Simple linear regression analyses yield significant positive correlations ($r \geq 0.60$) within the low energy components. This upturns into anti-correlation ($r >$ - 0.60) in the high energy component. These results are not only consistent with unified scheme for blazars and radio galaxies but also show that the emission mechanisms of these sources are similar.

**Keywords:** active galactic nuclei – blazars: - radio galaxies: emission spectrum

T21

# Emulating large-scale structure formation at the field level with styled neural networks

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Surveying the large-scale structure of the universe will yield an enormous amount of high quality data for constraining cosmology and potentially detecting new physics. However, extracting the maximum amount of information from this dataset and using it to its full potential requires fast and accurate methods of simulating cosmic structure formation in the nonlinear regime. Normally this is achieved with computationally expensive N-body simulations, which are too slow to use directly for inference. In this talk I will present the results from a new field-level emulator for large-scale structure formation that is trained to map the linear perturbations of the early universe to the nonlinear outcomes of cosmological N-body evolution. The emulator is made of a convolutional neural network, augmented with style parameters that capture both cosmology and redshift dependence. The model is autodifferentiable by construction, and the redshift dependence allows for time derivatives to be computed during training, thus the model is training on the full phase-space distribution of the N-body particles. The cosmology dependence allows the model to act effectively as an ensemble of CNNs each trained on simulations with different cosmological backgrounds, making the model an emulator for structure formation that is autodifferentiable with respect to both initial conditions, cosmological parameters, and redshift. The emulator achieves percent-level accuracy down to nonlinear scales of $k \sim 1\ \mathrm{Mpc}^{-1}\ h$, and can be used for both fast, accurate generation of a large number of mock catalogs and for field-level inference.

T22

- Abstract -

# $30\circ$-ish Directional Modulation Anomaly in the Cosmic Microwave Background

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We discovered that the CMB spatial modulation anomaly is more severe on a $30\circ$-ish spherical cap than on a hemisphere. We have studied this fact by dividing the CMB sphere by caps with different sizes and measured standard deviation as well as two point correlation function of temperature fluctuations. Both statistical measures show that a spherical cap with a radius of about 30 degrees has more power than its supplementary. We also looked for finer structure by tiling the CMB sky by stripes which help us to look for higher multipoles by expanding the standard deviation map by the Legendre polynomials. The result shows, for our resolution, the few first higher multipoles are as significant as the dipole. To make this observation more clear we suggest to name it directional modulation anomaly instead of dipole modulation anomaly.

T23

# Projections for the impact of the Legacy Survey of Space and Time (LSST) telescope's detection of new Milky Way satellite dwarf galaxies on the indirect detection of dark matter


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The nature of dark matter remains one of the most intriguing unsolved problems in modern physics. We have developed a Monte Carlo code to predict the number of new dwarf galaxies that are likely to be discovered by the Legacy Survey of Space and Time (LSST) telescope. We incorporated the results from the Sloan Digital Sky Survey (SDSS), The Dark Energy Survey (DES), and the PanSTARRS telescopes to ensure that our code utilized a realistic sample of Milky Way-like halos. For certain sets of the input simulated halos, we reliably reproduce the numbers and distributions of the actual Milky way dwarfs for many of the quantities measured by the telescopes, within an acceptable level of uncertainty. We utilize an appropriate set of these Monte Carlo models to predict statistical distributions of the number of dwarfs detected, absolute magnitude, surface brightness, half-light radius, and J-factors for the indirect detection of dark matter for the new dwarfs that are likely to be discovered by LSST. The final step will be to translate these new dwarf galaxies into a range of improvements on the Fermi-LAT's ability to discover/constrain the properties of dark matter. This step utilizes the indirect detection of the dark matter particles through annihilation to standard model particles that eventually produce gamma rays that could be detectable at the Fermi-LAT.

T24

Short Abstract
by
Mponeng Kopana

HI intensity mapping (21cm maps) trace the large-scale structure in the Universe, and contain a wealth of information about the origins and evolution of that structure, which in turn provide measurements and tests of the standard cosmological model. In particular, 21cm maps can provide accurate measurements of redshift space distortions generated by galaxy peculiar velocities. They quantify the growth rate of structure, which is a powerful probe of the cosmological model and the nature of gravity itself. We employ and develop models of 21cm intensity maps and their power spectra, as expected from HIRAX and PUMA, in order to forecast the precision with which we can measure the growth rate. We will include nonlinear effects in redshift space, in addition to telescope beam and foreground effects. These are interferometer modes of 21cm intensity mapping will be investigated. Furthermore, we investigate the combination of 21cm intensity maps with contemporaneous galaxy surveys, such as those planned for the Euclid and DESI telescopes. Such synergies can suppress some of the systematics that affect each survey, as well as improving on the precision and accuracy of cosmological measurements made by the separate surveys.

T25

# Dipole Cosmology: The Copernican Paradigm Beyind FLRW

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Abstract submitted for the *DSU2023 (The Dark Side of the Universe) — (smr 3863)* to be held from *10th July 2023* to *14th July 2023*.

We discuss the dipole cosmological principle, the idea that the Universe is a maximally Copernican cosmology compatible with a cosmic flow. It serves as the most symmetric paradigm that generalizes the FLRW ansatz, in light of the increasingly numerous (but still tentative) hints that have emerged in the last two decades for a non-kinematic component in the CMB dipole. Einstein equations in our ”dipole cosmology” are still ordinary differential equations – but instead of the two Friedmann equations, now we have four. The two new functions can be viewed as an anisotropic scale factor that breaks the isotropy group from $SO(3)$ to $U(1)$, and a ”tilt” that captures the cosmic flow velocity. The result is an axially isotropic, tilted Bianchi V/VII$_h$ cosmology. We note that multiple fluid components with independent flows can be realized in this set up. This fact does not seem to have been appreciated in early works on “tilted” Bianchi models, but it makes dipole cosmology a viable setting for very conventional model building with fluid mixtures, as in FLRW. We present a dipole $\Lambda$CDM model which has radiation and matter with independent flows, with (or without) a positive cosmological constant. A remarkable feature of models containing radiation (including dipole $\Lambda$CDM) is that the $relative$ flow between radiation and matter can increase at late times, which can contribute to the CMB dipole. We emphasize the significance of this observation for late time cosmic tensions and late time flows.

T26

# The impact of dark matter self-scatterings on its relic abundance

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Elastic self-scatterings do not change the number of dark matter particles and as such have been neglected in the calculation of its relic abundance. In this work we highlight the scenarios where the presence of self-scatterings has a significant impact on the effectiveness of annihilation processes through the modification of dark matter momentum distribution. We study a few example freeze-out scenarios involving resonant and subthreshold annihilations, as well as a model with an additional source of dark matter particles from the decays of a heavier mediator state. Interestingly, when the calculation is performed at the level of dark matter momentum distribution function, we find that the injection of additional energetic dark matter particles onto the thermal population can lead to a *decrease* of its final relic abundance.

[1] A. Hryczuk, M. Laletin, Phys.Rev.D **106** (2022) 2, 2, 2204.07078

T27

# Similarity, trends and halo-to-halo scatter in a first 'observed' dark matter phase-space in nearby Milky Way like galaxies

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In 2013, the velocity distribution function (VDF) of Dark Matter (DM) in the Milky Way (MW) was directly estimated starting from the rotation curve (RC) for the first time [1] in a self-consistent manner. The obtained VDFs were in broad agreement with N-body simulation counterparts but inconsistent with the widely used Maxwell-Boltzmann form, known as the "Standard Halo Model" (SHM). In a following work, simulations showed that while the phase-space density of DM haloes has a uniform form, it exhibits halo-to-halo scatter in shape [2]. Deviations of the VDF from Maxwell-Botzmann have important implications for direct detection experiments [3, 4].

In this talk, we [5] shows results using high-quality rotation curves of approximately 100 spiral galaxies identified as MW-like galaxies. We use this sample to study the halo-to-halo scatter and similarity in the full 6D phase-space density of DM particles in real galaxies for the very first time. We give an updated form of the VDF and pseudo-phase-space density that best describe the isotropic and constant anisotropic VDFs obtained from these galaxies. The VDFs at the solar radius of the MW-like sample are used to study the associated scatter in the rate and exclusion limit calculations for direct detection experiments.

[1] P. Bhattacharjee, S. Chaudhury, S. Kundu, S. Majumdar, PRD, **8**, 87 (2013)
[2] Y. Mao, L. E. Strigari, R. H. Wechsler, H. Wu, O Hahn, ApJ, **35**, 764 (2013)
[3] S. Mandal, S. Majumdar, V. Rentala, R. Basu-Thakur,PRD, **100**,023002 (2019)
[4] A. Pandey, S. Majumdar, V. Rentala, JCAP (submitted), (2023)
[5] M. Manju, S. Majumdar (to be submitted) (2023)

T28

# The DarkSide-20k experiment: prospects and current status

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DarkSide-20k (DS-20k) by the Global Argon Dark Matter Collaboration will explore the WIMP hypothesis by looking for WIMP-nucleon elastic scattering with a dual-phase time projection chamber (TPC) detector filled with 50 tonnes of low-radioactivity liquid argon extracted from underground sources. The detector will allow the experiment to reach a sensitivity to cross sections equal to 6.3 x $10^{-48}$ $cm^2$ (90% C.L. exclusion limit for a WIMP mass of 1 $TeV/c^2$ for a 200 t yr exposure in the 20 t fiducial volume). DS-20k will be equipped with low-radioactivity SiPM-based photodetectors with high photon-detection efficiency and a neutron veto system based on Gd-loaded acrylic panels for neutron moderation and capture. These specifications were chosen to satisfy the ultra-low background goal which has been demonstrated already by the DarkSide Collaboration with its previous experiment, Darkside-50. DS-20k is being built in the Hall-C of the INFN Laboratori Nazionali del Gran Sasso, shielded by 3800 m (water-equivalent) of rocks of the Gran Sasso massif. In this talk, the current status of the experiment construction will be reported, with insights on the topics of photo-electronics, argon procurement and production of the TPC.

T29

# Is Cosmology Self-Regulating?

T30

# DES Year-3 Cosmological Constraints and Comparison with Planck, KiDS, and HSC: is there a $\sigma_8$ tension?

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and

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We review the cosmological results using the full 5000 deg$^2$ of imaging data from the Dark Energy Survey (DES), performing an analysis combining three two-point correlation functions: cosmic shear using 100 million source galaxies; galaxy clustering; and galaxy-galaxy lensing, the cross-correlation of source galaxy shear with lens galaxy positions. Their combination yields constraints in fair agreement with the prediction of the model favoured by the Planck 2018 CMB data, although with a somewhat lower value for the amplitude of matter fluctuations, $\sigma_8$. We then compare our results with those from the two other current large imaging surveys, KiDS and HSC. In a very recent paper, DES and KiDS have jointly reanalysed their cosmic shear data, finding good agreement when using a common framework, but some relevant differences when using their particular modelling choices, pointing to the need to further understand non-linear structure formation and some astrophysical systematics. As the analyses stand now, the combination of DES, KiDS, and HSC provides a stringent constraint on $\sigma_8$ in agreement with Planck at the two-standard-deviation level.

T31

# Rescuing leptogenesis parameter space of inverse seesaw

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The generic inverse seesaw (ISS) framework is unable to produce an adequate amount of lepton asymmetry, which consequently fails to generate the observed baryon asymmetry of the Universe (BAU). This happens due to mainly two reasons, (i) cancellation of the lepton asymmetries between/among the pseudo-Dirac pairs, and (ii) strong wash out caused by the inverse decays. In this work, we have looked at the possible ways to achieve successful leptogenesis in the ISS framework. We have emphasized on two pathways, by enhancing the lepton asymmetry and/or by reducing the washout, to accomplish so. We have considered a non-degenerate right handed neutrino mass spectrum, which results into a larger order of lepton asymmetry that is sufficient to account for the observed BAU. For the implementation of the second idea, we have assumed the presence of a non-standard cosmological era in the pre-BBN epoch, that results into a faster expansion of the Universe, thereby reducing the washout by several orders of magnitude. We analyse the above two cases separately and find them to work remarkably well in order to rescue the ISS parameter space of leptogenesis (baryogenesis). Most interestingly, we find that the leptogenesis viable parameter space of the generic ISS model yields the branching ratio for $\mu \to e\gamma$ LFV decay, which is very much close to the present and future sensitivity as set by the MEG and MEG-II expectations.

[1] A. Mukherjee, N. Narendra, arXiv: 2204.08820.

T32

# Directional detection of dark matter boosted by cosmic-rays from direction of the Galactic center

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Dark matters with MeV- or keV-scale mass are difficult to detect with standard direct search detectors. However, they can be searched for by considering the up-scattering of kinetic energies by cosmic-rays [1]. Since dark matter density is higher in the central region of the Galaxy, the up-scattered dark matter will arrive at Earth from the direction of the Galactic center. Once the dark matter is detected, we can expect to recognize this feature by directional direct detection experiments. In this study [2], we simulate the nuclear recoils of the up-scattered dark matter and quantitatively reveal that a large amount of this type of dark matter is arriving from the direction of the Galactic center.

[1] T. Bringmann and M. Pospelov, Phys. Rev. Lett. **122**, no.17, 171801 (2019) doi:10.1103/PhysRevLett.122.171801 [arXiv:1810.10543 [hep-ph]].
[2] K. I. Nagao, S. Higashino, T. Naka and K. Miuchi, [arXiv:2211.13399 [astro-ph.CO]].

T33

# Electroweak Loop Contributions to the Direct Detection of Wino Dark Matter

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[d]*Tsung-Dao Lee Institute & School of Physics and Astronomy, Shanghai Jiao Tong University, Shanghai 200240, China*

Electroweak loop corrections to the matrix elements for the spin-independent scattering of cold dark matter particles on nuclei are generally small, typically below the uncertainty in the local density of cold dark matter. However, as shown in this paper [1], there are instances in which the electroweak loop corrections are relatively large, and change significantly the spin-independent dark matter scattering rate. An important example occurs when the dark matter particle is a wino, e.g., in anomaly-mediated supersymmetry breaking (AMSB) and pure gravity mediation (PGM) models. We find that the one-loop electroweak corrections to the spin-independent wino LSP scattering cross section generally interfere constructively with the tree-level contribution for AMSB models with negative Higgsino mixing, $\mu < 0$, and in PGM-like models for both signs of $\mu$, lifting the cross section out of the neutrino fog and into a range that is potentially detectable in the next generation of direct searches for cold dark matter scattering.

[1] J. Ellis, N. Nagata, K. A. Olive and J. Zheng, [arXiv:2305.13837 [hep-ph]].

# T34

Current observations of the universe suggest that the cosmic energy budget is dominated by dark energy (DE) which accounts for approximately 70% of the energy budget and a further 25% of this energy distribution is from the dark matter (DM) distribution. Cosmic microwave background (CMB) radiation experiments have already provided a wealth of information constraining the current cosmological model. Galaxy surveys have also provided valuable information on the evolution of the universe at later times. Next-generation experiments will probe these tracers to much higher precision providing unprecedented constraints on the various cosmological parameters of the ΛCDM model and testing various exotic models.
We investigate the ability of future large-scale structure surveys including HI intensity mapping surveys to constrain cosmological model parameters, dark energy equation of state parameters, and modified gravity models. In particular, we study how the HI surveys can be combined with other large-scale structure probes such as CMB lensing and galaxy surveys to place even tighter constraints on these models.

T35

# On the Inflationary Production of Light Dark Photon Dark Matter

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[3] *Kavli Institute for the Physics and Mathematics of the Universe (Kavli IPMU, WPI),UTIAS, The University of Tokyo, 5-1-5 Kashiwanoha, Kashiwa, Chiba, 277-8583, Japan*

The properties of dark matter are largely unknown under the current observational results, and light dark matter is a unique parameter space that is yet to be explored and may alleviate controversial small-scale issues of dark matter. We study the possibility of producing a light dark photon dark matter during the inflationary phase. The dark photon with a large wavelength can be efficiently produced owing to an effective violation of its conformal invariance and becomes non-relativistic before the time of matter-radiation equality. To meet the constraint from the isocurvature perturbation, the production is localized in time and consequently exhibits a localized spectrum. We further discuss the detectability of the remnant gravitational-wave signature from the dark-photon production by upcoming observational missions.

T36

# Thermodynamics of a rotating and non-linear magnetic-charged black hole in the quintessence field

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Here, we purpose a way to analyse the thermodynamics of a rotating and non-linear magnetic-charged black hole with quintessence. Especially, we compute various thermodynamic quantities of the black hole, such as mass, temperature, potential provided from the magnetic charge, and the heat capacity. This leads us to a better analysis of how quintessence modifies the behaviour of the black hole. Moreover, we study second-order phase transitions of this black hole, analysing the plot of its heat capacity. Then, we show that the black hole mass would have a phase of decrease. From the behaviour of the heat capacity, we point out that the black hole undergoes a second-order phase transition, which is shifted towards higher values of entropy as we increase the rotating parameter $a$ or the magnetic parameter $Q$. However, we have found that when we increase the quintessence parameter $c$, the second-order phase transition is instead shifted to lower entropy values. This presentation is based on [1].

[1] R. Ndongmo, S. Mahamat T. Bouetou and T. Kofane, , Physica Scripta, **96**(9), 095001 (2021).

T37

# Matter power spectrum in Scalar-tensor theory of gravity

T38

# Initial Performance of AMoRE-II Muon Detector

**Pendo B. Nyanda**[1] , **Yeongduk Kim**[1,2] , **Jaison Lee**[2] , **Gowoon Kim**[2]

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AMoRE-II is an experiment to search for neutrinoless double beta (0νββ) decay of $^{100}$Mo in molybdate crystals ($CaMoO_4$ and $Li_2MoO_4$) with a cryogenic system. It will be conducted in the 1,000m deep underground laboratory, Yemilab in Jeongseon, South Korea. To achieve the desired sensitivity of AMoRE-II, the muon induced background in AmoRE-II should be below $10^{-5}$ counts/keV/kg/year (ckky). Two types of muon veto detectors, plastic scintillators and water Cherenkov detectors have been designed, constructed and installed at Yemilab. The muon detectors's initial performance test results and the first measurement of the muon rate will be presented.

T39

# Higgs Spectrum Is Non-thermal after Inflation: Primordial Condensate vs Stochastic Fluctuation

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In this talk, I will present our recent work [1]: *"Non-thermal Higgs Spectrum in Reheating Epoch: Primordial Condensate vs. Stochastic Fluctuation."* The abstract is the following:

Since electroweak symmetry is generally broken during inflation, the Standard Model Higgs field can become supermassive even after the end of inflation. In this paper, we study the non-thermal phase space distribution of the Higgs field during reheating, focusing in particular on two different contributions: primordial condensate and stochastic fluctuations. We obtain their analytic formulae, which agree with the previous numerical result [2]. As a possible consequence of the non-thermal Higgs spectrum, we discuss perturbative Higgs decay during reheating for the case it is kinematically allowed. We find that the soft-relativistic and hard spectra are dominant in the decay rate of the stochastic fluctuation and that the primordial condensate and stochastic fluctuations decay almost at the same time.

## References

[1] K. Kaneta and K.-y. Oda, *Non-thermal Higgs Spectrum in Reheating Epoch: Primordial Condensate vs. Stochastic Fluctuation*, 2304.12578.

[2] K. Kaneta, S.M. Lee and K.-y. Oda, *Boltzmann or Bogoliubov? Approaches compared in gravitational particle production*, *JCAP* **09** (2022) 018 [2206.10929].

T40

# Increased sensitivity of strong gravitational lensing to populations of dark matter subhaloes

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In the search for the nature of dark matter, strong gravitational lenses are becoming increasingly important. With a technique called gravitational imaging, the presence (or absence) of otherwise invisible dark matter haloes can be uncovered. Haloes in the vicinity of the lens galaxy are called subhaloes. The cold dark matter (CDM) model predicts a large population of low mass subhaloes, increasing in number inversely to mass. In warm dark matter (WDM) models, the formation of subhaloes below a certain mass, $M_{\rm hm}$, is suppressed relative to CDM. This mass is inversely proportional to the dark matter particle mass. By measuring the properties of these low mass subhaloes in strong lens galaxies, gravitational imaging can constrain the underlying dark matter model.

We previously introduced a machine learning method for estimating the sensitivity of strong lens images to dark matter subhaloes and applied this to simulated data for the Euclid survey [1]. We found that subhaloes with mass larger than $M = 10^{8.8\pm0.2} M_\odot$ could be detected at $3\sigma$ in Euclid data. However, to probe WDM models that have not already been ruled out, subhalo detections are necessary with $M < 10^8 M_\odot$.

In this work, we show that detections in this previously forbidden mass range may be possible, when the effect of large populations of subhaloes are accounted for. We use a simulated dataset of $100$ Einstein rings observed at HST resolution and S/N typical of previous strong lens observation campaigns. We compute the subhalo sensitivity for this dataset using our previous method, which considers each detection an isolated event [1]. We find that individual $5\sigma$ detections are possible for subhaloes larger than $M = 10^{8.3\pm0.1} M_\odot$. Assuming a cold dark matter (CDM) scenario, we find that strong lenses of this quality should produce subhalo detections $4.5\%$ of the time, similar to the frequency of detections in real HST observations [2].

We then attempt to detect subhaloes in realisations of this dataset in images where populations of subhaloes are present. By comparing the frequency of detections made on data with populations, versus that predicted by the sensitivity to individual detections, we can quantify the increased sensitivity due to population effects. We find that, for CDM, detections are made $6.5\%$ of the time, much more frequently than when the subhaloes are isolated objects. This indicates that substructures smaller than the individual detection limit become detectable when a large population of them are present.

We also perform this comparison in WDM scenarios, parametrised by $M_{\rm hm}$. As $M_{\rm hm}$ increases, a greater number of low mass subhaloes are removed from the population. We find that WDM models with fewer subhaloes below $M = 10^8 M_\odot$ lack the detection excess found in CDM, and the detection frequency for images with populations matches the prediction for individual detections. Using these data to infer the dark matter model from the number of detections, we show that a biased result is obtained for CDM, where a model with $50\%$ more substructure is preferred over the truth. In the WDM case, the input model can be correctly recovered.

This shows that strong lensing observations are more sensitive to dark matter substructure than previously thought. If the effect of the population can be accounted for in modelling methods, constraints on dark matter models from strong lensing can become more competitive.

[1] O'Riordan C. M., Despali, G., Vegetti, S., Moliné, Á., Lovell, M., MNRAS **521**, 2342 (2023).
[2] Vegetti, S., Koopmans, L. V. E., Auger, M. W., Treu T., Bolton, A., S. MNRAS **442**, 2017 (2014).

T41

# Cosmology using 1-loop EFT power spectrum for HI intensity mapping

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We explore the information from HI power spectrum in redshift space using 1-loop EFT power spectrum model and HIRAX survey. This work is an extension of our recent paper [1] on Cosmological constraints from the power spectrum and bispectrum of 21cm intensity maps. We marginalise over bias and nuisance parameters. We notice from the results that CMB Planck 2018 constraints can be improved if we add non-linearity (EFT).

[1] D. Karagiannis, R. Maartens, L. Randrianjanahary, JCAP. **11**, 003 (2022).

T42

# Asymmetric Mediator in Scotogenic Model

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**, Yasutaka Takanishi(b,c)§and Masato Yamanaka(b)¶**

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## Abstract

The scotogenic model is the Standard Model (SM) with Z2 symmetry and the addition of Z2 odd right-handed Majorana neutrinos and SU(2)L doublet scalar fields. We have extended the original scotogenic model by an additional Z2 odd singletscalar field that plays a role in dark matter. In our model, the asymmetries of the lepton and Z2 odd doublet scalar are simultaneously produced through CP-violating right-handed neutrino decays. While the former is converted into baryon asymmetry through the sphaleron process, the latter is relaid to the DM density through the decay of SU(2)L doublet scalar that is named “asymmetric mediator”. In this way, we provide an extended scotogenic model that predicts the energy densities of baryon and dark matter being in the same order of magnitude, and also explains the lowenergy neutrino masses and mixing angles.Asymmetric mediator in scotogenic model. It is based on ref.[1].

[1] Phys.Lett.B 836 (2023) 137627 .

T43

# Muon g – 2 anomaly in left-right model with inverse seesaw or Modified Hybrid Inflation, Reheating and Stabilization of the Electroweak Vacuum

T44

# The type-II seesaw mechanism in alternative gauged $U(1)_X$ and dark sector

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We propose an extra $U(1)_X$ model with an alternative charge assignment for right-handed neutrinos. The type-II seesaw mechanism by a triplet Higgs field is promising for neutrino mass generation because of the alternative charge assignment. With the minimal Higgs field for $U(1)_X$ with the charge 1, right-handed neutrinos are candidates for Dirac dark matter (DM) and dark radiation (DR). We have derived and imposed the LHC bound, the DR constraint, and the bound from DM direct searches in the wide range of parameter space.

[1] N. Okada and O. Seto, Phys. Rev. D **105**, 123512 (2022).

T45

# Pseudo-Scalar Fields in Early Universe

T46

# An alternative method for measuring accretion rate in Galaxies

T47

# Gravitational Positivity for Dark Gauge Bosons