

# Beam-based approach to verify nuclear hyperfine mixing effect

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Electromagnetic coupling between a nucleus and its bound electron(s) can induce mixing of nuclear levels—a phenomenon known as nuclear hyperfine mixing (NHM). This effect can greatly enhance light-nucleus interactions and is expected to play a key role in laser-based nuclear control. Although predicted in 1966, NHM has not yet been experimentally confirmed. Here we propose a beam-based scheme to test the NHM effect via collisions between hydrogenlike  $^{229}\text{Th}$  ions ( $^{229}\text{Th}^{89+}$ ) and a thin-layer target. This approach avoids the need for resonance tuning and enables detection of prompt ultraviolet photons as a clear experimental signature of NHM. To evaluate the excitation probability of the nuclear isomeric state, we develop, for the first time, nuclear excitation theories that incorporate NHM in both electron- and nucleus-induced processes.

**nuclear hyperfine mixing, thorium-229, electronic excitation, coulomb excitation**

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## 1 Introduction

In most cases, nuclear properties are influenced negligibly by bound electrons. However, notable exceptions have been predicted. For example, the radiative lifetime of the 8.27-eV isomeric state [1] in bare  $^{229}\text{Th}$  nuclei is on the order of 1000 s [2–4], but it shortens by five orders of magnitude to approximately 0.01 s in hydrogenlike ions ( $^{229}\text{Th}^{89+}$ ) [5–9]. For another example, the radiative lifetime of the 2.329-keV isomeric state of the  $^{205}\text{Pb}$  nucleus is approximately 900 s in bare nuclei ( $^{205}\text{Pb}^{82+}$ ), but it reduces by four orders of magnitude to 0.032 s in boronlike ions ( $^{205}\text{Pb}^{77+}$ ) [10, 11].

These dramatic shortenings in lifetime are attributed to nuclear hyperfine mixing (NHM), in which nuclear levels are mixed due to electromagnetic interaction with the

bound electron(s). This effect was first predicted in 1966 by Lyuboshitz et al. [12] and was also referred to as nuclear spin mixing in hyperfine fields [13]. The bound electron(s) greatly enhance the coupling strength between the nucleus and the vacuum light field, thereby accelerating radiative decay. Conversely, NHM can also enhance laser-induced nuclear excitation [14], potentially enabling access to highly nonlinear light-nucleus interactions when combined with intense light sources such as petawatt infrared lasers, synchrotrons, or x-ray free-electron lasers [15–17]. NHM is thus expected to play a key role in the emerging field of laser-nuclear physics.

Despite its importance, NHM has not been experimentally verified in the nearly 60 years since its theoretical proposal. One possible approach is laser-resonant excitation of  $^{229}\text{Th}^{89+}$  ions [18]. However, this method faces significant challenges due to the sizable uncertainties in the transition

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energies ( $\sim 10$  meV) [1, 9] and is applicable only to  $^{229}\text{Th}$  [19]. Note that the transition energy of  $^{229}\text{Th}^{89+}$  differs from the reported values measured for other  $^{229}\text{Th}$  ions doped in crystals [2-4, 20, 21]. It is therefore desirable to develop alternative approaches that do not rely on resonance conditions and are applicable to nuclei with higher transition energies.

## 2 Beam-based approach

In this work, we propose a beam-based method to verify the NHM effect without requiring resonance scanning. As illustrated in Figure 1(a), we consider a beam of  $^{229}\text{Th}^{89+}$  ions passing through a thin  $^{12}\text{C}$  target. During this interaction, the  $^{229}\text{Th}$  nucleus can be excited to its isomeric state via collisions with target electrons or nuclei. In the presence of NHM, the excited nuclei ( $^{229\text{m}}\text{Th}^{89+}$ ) decay rapidly—within  $\sim 0.01$  s—after the ions exit the target. In contrast, without NHM, the lifetime exceeds 1000 s. Thus, detecting the prompt decay photons within a short time window downstream provides a direct and unambiguous signature of the NHM effect.

Quantifying the nuclear excitation probability requires theoretical developments beyond existing theories for electronic [22-25] and Coulomb excitation [26, 27], which neglect NHM. In this work, we develop new theoretical frameworks for both excitation mechanisms that explicitly incorporate strong hyperfine mixing. This represents a key methodological advancement and enables quantitative predictions for the proposed experiment.

In addition to nuclear excitation, collisional ionization and recombination processes occur during the beam-target interaction, leading to changes in the charge state of the projectile ions. We use the ETACHA4 code [28] to model the evolution of the charge-state distribution of  $^{229}\text{Th}$  ions during the collision. Figure 1(b) shows the relative abundances of the dominant charge states for an incident  $^{229}\text{Th}^{89+}$  beam with an energy of 100 MeV/u. The  $^{12}\text{C}$  target has a density of  $2.26 \text{ g/cm}^3$ , so a thickness of  $5.65 \text{ mg/cm}^2$  corresponds to  $25 \text{ }\mu\text{m}$ . As the ions traverse the target, the fraction of

$^{229}\text{Th}^{89+}$  decreases, while the populations of lower charge states ( $^{229}\text{Th}^{88+,87+}$ ) increase. After passing through the target, approximately 61.3% of the ions remain in the  $^{229}\text{Th}^{89+}$  state.

The isomeric excitation probability is obtained by summing over all possible ionic charge states and integrating over the target depth:

$$P_{\text{exc}} = \sum_q \int f_q(x) [n_e \sigma_{e,q}(x) + n_C \sigma_{C,q}(x)] dx, \quad (1)$$

where  $f_q(x)$  is the fraction of the  $^{229}\text{Th}^{q+}$  ion at depth  $x$ ,  $n_e$  ( $= 6.80 \times 10^{23} \text{ cm}^{-3}$ ) and  $n_C$  ( $= 1.13 \times 10^{23} \text{ cm}^{-3}$ ) are the densities of electrons and  $^{12}\text{C}$  nuclei in the target, and  $\sigma_{e,q}$  and  $\sigma_{C,q}$  are the isomeric excitation cross sections due to target electrons and nuclei, respectively, for each charge state  $q$ . To account for energy loss in the target, we use the SRIM software [29] to calculate the stopping power. For an incident  $^{229}\text{Th}^{89+}$  beam with an energy of 100 MeV/u, the energy decreases slightly to 98.8 MeV/u after traversing the  $25\text{-}\mu\text{m}$ -thick  $^{12}\text{C}$  target. In the rest frame of the  $^{229}\text{Th}^{89+}$  ion, the energy of a target electron is given by  $\varepsilon = (m_e/m_{\text{ion}})\varepsilon_{\text{ion}}$ , where  $m_e$  and  $m_{\text{ion}}$  are the mass of the electron and the  $^{229}\text{Th}$  ion, respectively.

## 3 Electronic excitation with NHM

We turn to the rest frame of  $^{229}\text{Th}^{89+}$  and consider its excitation by an incoming electron beam. From Fermi's golden rule, the cross section is calculated as (atomic units are used,  $\hbar = m_e = |e| = 1$ ):

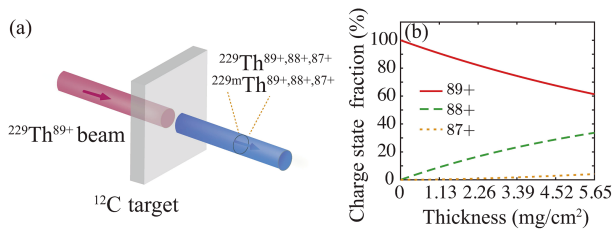
$$\frac{d\sigma_{e,89}}{d\Omega} = \frac{\varepsilon_f \varepsilon_i}{4\pi^2 c^4} \frac{p_f}{p_i} |\langle f | H_{\text{en}} + H_{\text{ee}} | i \rangle|^2, \quad (2)$$

where  $p_{i,f}$  is the initial or final momentum of the incident electron, and  $\varepsilon_{i,f} = \sqrt{p_{i,f}^2 + c^4}$  is the corresponding energy. The Hamiltonians  $H_{\text{en}}$  and  $H_{\text{ee}}$  describe the interactions between the incident electron and the nucleus, and between the incident electron and the bound electron, respectively. Note that existing theories of electronic excitation of the  $^{229}\text{Th}$  nucleus only include the  $H_{\text{en}}$  term [22-24], whereas in the current case the  $H_{\text{ee}}$  term must be included.

For low-energy nuclear transition,  $H_{\text{en}}$  can be written as [30]

$$H_{\text{en}} = \sum_{\tau L} \mathcal{M}^{(\tau L)} \cdot T^{(\tau L)}, \quad (3)$$

where  $\mathcal{M}^{(\tau L)}$  and  $T^{(\tau L)}$  are multipole operators of rank  $L$  for the nucleus and for the incident electron, respectively. The electric or magnetic multipoles are distinguished by  $\tau = E$  or



**Figure 1** (Color online) (a) Illustration of nuclear excitation during the collision between a  $^{229}\text{Th}^{89+}$  beam and a layer  $^{12}\text{C}$  target. (b) Evolution of charge state fraction  $f_q$  as a function of the  $^{12}\text{C}$  target thickness. The incoming beam energy is 100 MeV/u.

$M$  [11].  $H_{ee}$  can also be given as a multipole summation

$$H_{ee} = \sum_{Lm} \int \rho_1(\mathbf{r}_1) \rho_2(\mathbf{r}_2) C_m^{(L)}(\hat{\mathbf{r}}_1) C_m^{(L)*}(\hat{\mathbf{r}}_2) \frac{r_{>}^L}{r_{>}^{L+1}} d\mathbf{r}_1 d\mathbf{r}_2 \\ + \sum_{Lm} \frac{1}{c^2 L(L+1)} \int j_1(\mathbf{r}_1) \cdot \mathbf{L}_1 C_m^{(L)}(\hat{\mathbf{r}}_1) \\ \times j_2(\mathbf{r}_2) \cdot \mathbf{L}_2 C_m^{(L)*}(\hat{\mathbf{r}}_2) \frac{r_{>}^L}{r_{>}^{L+1}} d\mathbf{r}_1 d\mathbf{r}_2, \quad (4)$$

where  $\rho_1(\mathbf{r}_1)$  [ $\rho_2(\mathbf{r}_2)$ ],  $j_1(\mathbf{r}_1)$  [ $j_2(\mathbf{r}_2)$ ], and  $\mathbf{L}_1$  ( $\mathbf{L}_2$ ) are the charge density, current density, and angular momentum operators for the incident (bound) electron, respectively, and  $r_{>}$  ( $r_{<}$ ) denotes the greater (smaller) of  $r_1$  and  $r_2$ . Here,  $C_m^{(L)}(\hat{\mathbf{r}}) = \sqrt{4\pi/[L]} Y_{Lm}(\hat{\mathbf{r}})$  in which  $[L]$  is a shorthand notation for  $2L+1$ .

The initial and final states in Eq. (2) are

$$|i\rangle = \mathcal{A}\{|F_g M_g\rangle |\chi_i^+\rangle\}, \\ |f\rangle = \mathcal{A}\{|F_e M_e\rangle |\chi_f^-\rangle\}, \quad (5)$$

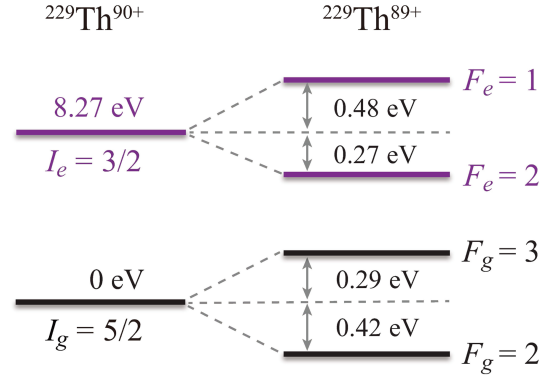
where  $\mathcal{A}$  is an antisymmetrization operator that exchanges the incident electron with the bound electron.  $|\chi_i^+\rangle$  and  $|\chi_f^-\rangle$  are the initial and final states of the free electron, respectively.  $|F_g M_g\rangle$  and  $|F_e M_e\rangle$  are eigenstates of the ion associated with the nuclear ground state and the isomeric state, respectively, as illustrated in Figure 2. Note that the eigenstates of the ion are dressed hyperfine states including the NHM effect [11, 31], and they are given by:

$$|F_g M_g\rangle \approx |I_g J_g; F_g M_g\rangle + b_g |I_e J_g; F_g M_g\rangle, \\ |F_e M_e\rangle \approx |I_e J_g; F_e M_e\rangle + b_e |I_g J_g; F_e M_e\rangle, \quad (6)$$

where  $b_\mu$  ( $\mu = g, e$ ) is the mixing coefficient [6, 8, 9, 15, 16],  $I_g = 5/2$  ( $I_e = 3/2$ ) denotes the spin of the nuclear ground (isomeric) state, and  $J_g = 1/2$  is the angular momentum of the bound  $1s$  electron. On the right hand side of the above equations,  $|I_\mu J_\mu; F_\mu M_\mu\rangle$  is the hyperfine-coupled basis with total angular momentum  $F_\mu$ , formed by coupling the nuclear spin  $I_\mu$  and angular momentum of the bound electron  $J_g$  [11]. Hyperfine mixing exists between the two levels with total angular momentum  $F_e = 2$  or  $F_g = 2$ , since  $b_\mu$  is nonzero only in these cases. Specifically, from Eq. (3) of ref. [9], we have  $b_g = 0.048$  for  $F_g = 2$  and  $b_e = -0.048$  for  $F_e = 2$ .

In the relativistic distorted-wave approximation, the state of the free electron with momentum  $\mathbf{p}$  and spin projection  $\eta$  can be expanded in the following form [22]

$$|\chi_{p\eta}^\pm\rangle = \frac{4\pi}{p} \sqrt{\frac{\varepsilon + c^2}{2\varepsilon}} \sum_{km} i^l e^{\pm i\delta_k} \langle lm - \eta, \frac{1}{2}\eta | jm \rangle \\ \times Y_{lm-\eta}^*(\hat{\mathbf{p}}) \frac{1}{r} \begin{pmatrix} P_\kappa(r) \Omega_{km}(\hat{\mathbf{r}}) \\ i Q_\kappa(r) \Omega_{-km}(\hat{\mathbf{r}}) \end{pmatrix}, \quad (7)$$



**Figure 2** (Color online) Energy levels (not drawn to scale) of the nuclear ground state and the isomeric state for the bare nucleus ( $^{229}\text{Th}^{90+}$ ) and for the hydrogenlike ion ( $^{229}\text{Th}^{89+}$ ). The hyperfine splittings for  $F_e = 1$  and  $F_g = 3$  levels are adopted from ref. [9], while those for  $F_e = 2$  and  $F_g = 2$  are calculated using Eq. (4) in ref. [9]. Note that the eigenstates of the ion are dressed hyperfine states with the NHM effect included, as given in Eq. (6).

where  $\delta_k$  is the partial-wave phase shift and  $\Omega_{km}$  is the spherical spinor with quantum number  $\kappa$  and magnetic number  $m$  [32]. The total and orbital angular momentum numbers  $j$  and  $l$  are related to  $\kappa$  by  $\kappa = \pm(j + 1/2)$  for  $l = j \pm 1/2$ .  $P_\kappa(r)$  and  $Q_\kappa(r)$  are the large and small components of the solutions to the radial Dirac equation in a potential  $V(r)$ . They satisfy the normalization condition  $P_\kappa(r) \rightarrow \sin(kr + \delta'_k)$  as  $kr \rightarrow \infty$ , where  $\delta'_k$  denotes a phase.

The total electronic excitation cross section is obtained by substituting Eqs. (5), (6), and (7) into the differential cross section in Eq. (2), integrating the resulting expression, and then averaging over the initial states and summing over all final states:

$$\sigma_{e,89} = \frac{2\pi}{c^4 [F_g]} \frac{p_f \varepsilon_f + c^2}{p_i} \frac{\varepsilon_i + c^2}{p_f^2} \frac{1}{p_i^2} \\ \times \sum_{\kappa_f, \kappa_i, N} [N] |t_{fi}^{\text{en},N} + t_{fi,\text{dir}}^{\text{ee},N} - t_{fi,\text{exc}}^{\text{ee},N}|^2, \quad (8)$$

where  $N$  is the total angular momentum coupled by  $F_e$  and  $J_f$ . The matrix elements  $t_{fi}^{\text{en},N}$ ,  $t_{fi,\text{dir}}^{\text{ee},N}$ , and  $t_{fi,\text{exc}}^{\text{ee},N}$  describe the transition between the nucleus and the incident electron, the direct part and the exchange part of the transition between the bound and the incident electrons, respectively. Note that  $t_{fi,\text{dir}}^{\text{ee},N}$  and  $t_{fi,\text{exc}}^{\text{ee},N}$  arise from the NHM effect, without which these two terms would vanish. The explicit expressions of these matrix elements are given in the Supplemental Material.

## 4 Coulomb excitation with NHM

The derivation of the Coulomb excitation cross section with the NHM effect closely follows that of electronic excita-

tion, except that the antisymmetrization need not be performed and the Schrödinger wave functions are used to describe the incident  $^{12}\text{C}$  nuclei (which have zero spin, and non-relativistic velocity):

$$|\chi_p^\pm\rangle = \frac{4\pi}{p} \sum_{lm} i^l e^{\pm i\Delta_l} Y_{lm}^*(\hat{p}) \frac{1}{r} S_l(r) Y_{lm}(\hat{r}). \quad (9)$$

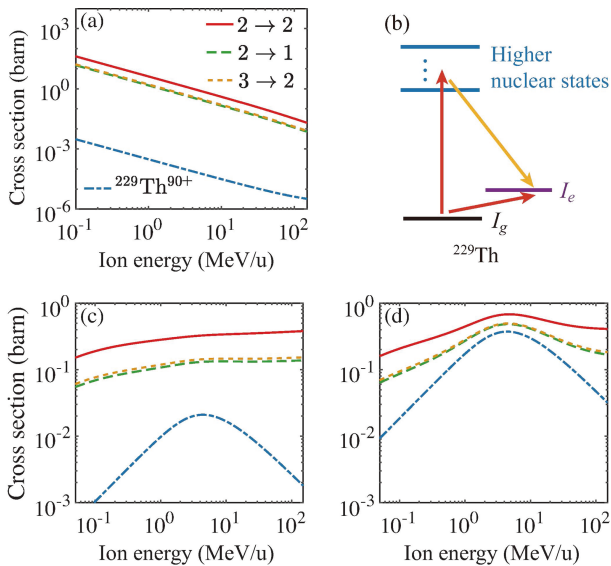
Thus the effective Coulomb excitation cross section is

$$\sigma_{C,89} = \frac{16\pi q_C^2 m_C^2}{p_{Ci}^2 p_{Cf}^2 [F_g]} \frac{p_{Cf}}{p_{Ci}} \sum_{l_f, l_i} |t_{fi}^{Cn} + t_{fi}^{Ce}|^2, \quad (10)$$

where  $q_C, m_C$  and  $p_C$  are the charge, reduced mass and momentum of the  $^{12}\text{C}$  nucleus, respectively. The matrix element  $t_{fi}^{Cn}$  corresponds to the transition induced by direct Coulomb excitation by the  $^{12}\text{C}$  nucleus, while  $t_{fi}^{Ce}$  corresponds to the contribution of the bound electron to the Coulomb excitation. The  $t_{fi}^{Ce}$  term originates from the NHM effect. The explicit expressions of  $t_{fi}^{Cn}$  and  $t_{fi}^{Ce}$  are given in the Supplemental Material.

## 5 Results and discussions

In the following calculations, we return to the laboratory frame and consider incident beam energies in the range of 0.05–150 MeV/u. Figure 3(a) shows the calculated cross sections for electronic excitation of  $^{229}\text{Th}^{89+}$ , compared with that



**Figure 3** (Color online) (a) Cross sections of electronic excitation of  $^{229}\text{Th}^{89+}$ , in comparison to that of  $^{229}\text{Th}^{90+}$ . For the former, three different  $F_g \rightarrow F_e$  channels are calculated, as labeled. (b) For Coulomb excitation by  $^{12}\text{C}$  nuclei, both direct ( $I_g \rightarrow I_e$ ) and indirect ( $I_g \rightarrow$  higher nuclear states  $\rightarrow I_e$ ) isomeric excitation channels are possible. (c) Cross sections of only the direct channel of Coulomb excitation. (d) Cross sections of all channels of Coulomb excitation.

for the bare nucleus  $^{229}\text{Th}^{90+}$ . For the hydrogenlike ion, three  $F_g \rightarrow F_e$  hyperfine transitions contribute:  $2 \rightarrow 2$ ,  $2 \rightarrow 1$ , and  $3 \rightarrow 2$  (Figure 2). The first two channels are relevant when the ion is initially in the  $F_g = 2$  ground state.

For  $^{229}\text{Th}^{89+}$ , the excitation cross sections lie in the range of 10<sup>2</sup> to 10<sup>-3</sup> barns in the considered energy range. In contrast, the cross section for the bare nucleus  $^{229}\text{Th}^{90+}$  is significantly smaller, ranging from 10<sup>-2</sup> to 10<sup>-6</sup> barns over the same energy interval. This demonstrates that the NHM effect enhances the isomeric excitation probability by three to four orders of magnitude. In both cases, the cross section decreases with increasing electron energy.

For Coulomb excitation by  $^{12}\text{C}$  nuclei in the considered energy range, in addition to direct excitation of the isomeric state, indirect excitation via higher-lying nuclear states is possible. These higher states can decay rapidly, usually within 100 ns, to the isomeric state (and also to the ground state, according to the corresponding branching ratios), as illustrated in Figure 3(b). The cross sections for direct isomeric excitation, calculated using Eq. (10), are shown in Figure 3(c). For the bare nucleus, the cross section lies between approximately 10<sup>-3</sup> and 10<sup>-2</sup> barns, peaking at around 5 MeV/u. The presence of the NHM effect in  $^{229}\text{Th}^{89+}$  enhances the cross section by one to two orders of magnitude.

Figure 3(d) shows the total Coulomb excitation cross sections, including both direct and indirect pathways. Since NHM has a negligible effect on the higher-lying nuclear states, the indirect excitation contribution in the  $^{229}\text{Th}^{89+}$  case is identical to that of the bare nucleus, as previously treated in ref. [27]. The considered higher nuclear states, together with the required nuclear parameters, are given in Table I of ref. [27]. When indirect channels are included, the relative enhancement due to NHM becomes less dramatic, but a moderate enhancement—typically by a factor of a few—remains observable.

With the electronic and Coulomb excitation cross sections determined, we return to the proposed scheme illustrated in Figure 1 and evaluate the isomeric excitation probability—that is, the probability of producing  $^{229\text{m}}\text{Th}^{89+}$  from each incident  $^{229}\text{Th}^{89+}$ . For an ion energy of 100 MeV/u, the relative electron energy is 54.5 keV, and the total electronic excitation cross section is 0.044 barn for  $^{229}\text{Th}^{89+}$  ( $2 \rightarrow 2$ : 0.032 barn,  $2 \rightarrow 1$ : 0.012 barn).  $^{229}\text{Th}^{87+}$  has a weaker NHM effect, and its yield is about an order of magnitude lower than that of  $^{229}\text{Th}^{89+}$  (Figure 1(b)). For  $^{229}\text{Th}^{88+}$ , which has no NHM, the cross sections are comparable to those of  $^{229}\text{Th}^{90+}$  and are roughly four orders of magnitude smaller than  $^{229}\text{Th}^{89+}$ . Although electronic excitation of higher nuclear states is possible in this case, the corresponding cross sections are three or more orders of magnitude smaller than the direct isomeric excitation cross section of  $^{229}\text{Th}^{89+}$  (see Supplemental Ma-

terial). Consequently, the contributions of  $^{229}\text{Th}^{88+,87+}$  and higher nuclear states to the isomeric excitation are negligible compared to those of  $^{229}\text{Th}^{89+}$ . Using Eq. (1), we compute the isomeric excitation probability from electronic excitation to be  $5.9 \times 10^{-5}$ .

For a relative ion energy of 100 MeV/u, the Coulomb excitation cross section is 0.047 barn for  $^{229}\text{Th}^{90+}$  and 0.60 barn for  $^{229}\text{Th}^{89+}$  ( $2 \rightarrow 2$ : 0.42 barn,  $2 \rightarrow 1$ : 0.18 barn). The excitation cross sections for the lower charge states are similar to the  $^{229}\text{Th}^{90+}$  case. Using Eq. (1), the excitation probability due to Coulomb interaction with  $^{12}\text{C}$  nuclei is  $1.4 \times 10^{-4}$ . Therefore, the total isomeric excitation probability is  $P_{\text{exc}} = 2.0 \times 10^{-4}$ , with contributions from electronic and nuclear Coulomb channels on the same order of magnitude. Considering that the final fraction of  $^{229}\text{Th}^{89+}$  is 61.3%, the effective excitation probability in the final  $^{229}\text{Th}^{89+}$  beam is about  $\tilde{P}_{\text{exc}} = 1.2 \times 10^{-4}$ . For comparison, we have also calculated the effective excitation probabilities at other beam energies and provided them in the Supplemental Material.

The anticipated experimental signature of the NHM effect is the spontaneous decay of the produced  $^{229\text{m}}\text{Th}^{89+}$  within 0.01 s following the beam-target interaction. Detection of emitted ultraviolet photons—centered around 8.27 eV, and more precisely at 7.71, 8.42, 8.46, and 9.17 eV as indicated by the level scheme in Figure 2—within this time window would serve as a direct confirmation of NHM. It should be noted that due to the relativistic Doppler effect, the photons with frequency  $\omega$  emitted into all solid angles by ions at 100 MeV/u are broadened in the laboratory frame to the range  $0.63\omega$ – $1.58\omega$  (see Supplemental Material). Therefore, in an actual measurement, the signal photons are distributed over an energy range of approximately 5–14 eV.

The expected ultraviolet photon signal can be estimated as  $N_{ph} = N_0 \tilde{P}_{\text{exc}} \eta$ , where  $N_0$  is the number of incident  $^{229}\text{Th}^{89+}$  ions,  $\tilde{P}_{\text{exc}}$  is the excitation probability given above, and  $\eta$  is the overall detection efficiency. For a storage ring with a circumference of approximately 100 m, about 1% of the emitted photons can be collected within a detection region of  $\sim 1$  m. After accounting for the Doppler broadening of the VUV photons, the effective quantum efficiency of a photomultiplier tube for the signal photons is about 8% [33] (see Supplemental Material). We therefore adopt an overall detection efficiency of  $\eta \approx 10^{-3}$ . It should be noted that the possible background photons in the proposed detection scheme mainly originate from two sources: radiation from the target region and radiation from the  $^{229}\text{Th}^{89+}$  ions. The former is emitted mainly near the target and can be suppressed by placing the detector away from this region. For the latter, the relevant VUV atomic transitions have lifetimes no longer than the nanosecond scale and therefore do not affect the detection of the signal from millisecond-scale nuclear decay.

The  $^{229}\text{Th}^{89+}$  beam can be produced via projectile-fragmentation reactions, for example by colliding a  $^{232}\text{Th}$  beam with a thick  $^9\text{Be}$  target. It has been estimated that approximately  $10^6$   $^{229}\text{Th}^{89+}$  ions/s can be generated from a  $^{232}\text{Th}$  beam of  $10^{10}$  particles/s [14]. Beam accumulation techniques may further increase the intensity by a factor of 20 [34], while deceleration from 400 to 100 MeV/u results in only a modest loss of beam intensity ( $\sim 30\%$ ) [35]. Assuming  $N_0 = 10^7$  particles/s, we estimate  $N_{ph} \approx 1 \text{ s}^{-1}$ , i.e., about one signal photon detected every one second of beam time. At next-generation facilities, such as the High-Intensity Heavy-Ion Accelerator Facility (HIAF) in China [36],  $N_{ph}$  is expected to increase by an additional order of magnitude. These estimations indicate that the proposed scheme is feasible with current or near-future experimental facilities.

In the absence of NHM, the isomeric excitation probability would be reduced by approximately one order of magnitude. This reduction arises because the electronic excitation would be suppressed by four orders of magnitude and thus become negligible, while the Coulomb excitation probability would decrease by one order of magnitude. Moreover, without NHM the radiative lifetime of  $^{229\text{m}}\text{Th}$  is on the order of 1000 s. Consequently, the detectable signal rate would be further reduced by five orders of magnitude. Under these conditions, virtually no decay photons would be expected within the same 0.01-s time window following the collision.

## 6 Conclusion

We have proposed a beam-based approach to verify the long-standing but unconfirmed NHM effect using  $^{229}\text{Th}^{89+}$  ions interacting with a  $^{12}\text{C}$  target. This approach avoids the need for resonance tuning and enables detection of prompt ultraviolet photons as a clear experimental signature. To support this scheme, we developed new theories for electronic and Coulomb excitation in the presence of NHM. Numerical results indicate that the excitation probability is within the detectable range using current or near-future heavy-ion facilities. Last but not least, this approach can be extended to nuclei with higher transition energies [10, 11, 19] that are inaccessible to current laser-resonant excitation techniques.

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**Conflict of interest** The authors declare that they have no conflict of interest.

**Data availability** The data that support the findings of this work are openly available ([https://figshare.com/articles/figure/Figures\\_for\\_the\\_article\\_Beam-Based\\_Approach\\_to\\_Verify\\_Nuclear\\_Hyperfine\\_Mixing\\_Effect\\_/2948-3342](https://figshare.com/articles/figure/Figures_for_the_article_Beam-Based_Approach_to_Verify_Nuclear_Hyperfine_Mixing_Effect_/2948-3342)).

### Supporting Information

The supporting information is available online at <http://phys.scichina.com> and <https://link.springer.com>. The supporting materials are published as submitted, without typesetting or editing. The responsibility for scientific accuracy and content remains entirely with the authors.

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